

DOE/ER/40150--81

Proceedings

DE89 014423

CEBAF/SURA

1986 Summer Workshop

Continuous Electron Beam Accelerator Facility

Newport News, Virginia

June 23-27, 1986

Edited by: Franz Gross and Ralph Minehart

Sponsored by: CEBAF/SURA

and

The Department of Energy

Endorsed by: The National Science Foundation

Continuous Electron Beam Accelerator Facility

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November 1986

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Preface

The 3rd CEBAF Summer Workshop was held at a time of growing enthusiasm for the CEBAF project. Everyone was expecting that Congress would approve initial construction funds (\$17M was approved in October); the first four commercially fabricated superconducting cavities had been tested and all exceeded specifications (the number is now six); J. Dirk Walecka had accepted the position of Scientific Director and his arrival at CEBAF was eagerly anticipated (he sent the letter of greeting to the Workshop participants printed on page ix); the Research Division had published its first major report on planning for the scientific program (the RPAC report distributed to all the users); and J. Mougey and V. Burkert had recently joined the Research Division staff. In this atmosphere of anticipated funding and growth, a record number of 131 physicists attended the Workshop. Also attending the Workshop were eleven graduate students who made up the first class of the program of Hampton University Graduate Studies (HUGS) at CEBAF.

It was only a few days before the Workshop was to begin that it became clear that the attendance would be so large. Since CEBAF has no seminar room which can hold more than 80-100 people, it was necessary at the last minute to shift the location of the plenary sessions to neighboring Christopher Newport College, which graciously loaned us the use of their beautiful new auditorium. We thank CNC for allowing us to use their facilities; their help and generosity is typical of the support given to CEBAF by the people and communities in the Hampton Roads area. (Both the cities of Hampton and Newport News supported the HUGS at CEBAF program with generous financial contributions which made it possible to get the program underway.) The last minute changes put an unusually heavy burden on our Support Staff; Ramona Rickman, Administrative Assistant for the Research Division, and Chris Bach and Maria Thornton. They undertook this burden so skillfully that most Workshop participants were completely unaware of the near crises. We are all grateful to them for maintaining a calm, user friendly environment.

The organization of this volume parallels the program of the 1986 Workshop; it consists of summary talks on the status of CEBAF, invited talks on theoretical and experimental subjects related to the CEBAF program, and working group meetings. The summary talks published here have been revised to include developments subsequent to the Workshop, and are current as of the end of

October. There were five working groups, reflecting the organization of the RPAC report, chaired by invited participants:

- A Single Nucleon Emission - A. Gerard (Saclay)
- B Two Nucleon Emission - J. Lightbody (NBS)
- C Baryon Spectroscopy - N. Mukhopadhyay (RPI)
- D Studies of Strange Nuclear Systems - E. Hungerford (Houston)
- E Studies in the x-Scaling Region - K. Van Bibber (LLNL)

These working group meetings are very valuable; they have become the heart of the CEBAF Summer Workshop series, and will again be a major part of the 1987 Summer Workshop (to be held June 22-26). Reports of the working groups written by their chairman are included in the third section of these proceedings, together with Ingo Sick's Workshop Summary. Finally, we were pleased with the variety of contributions received this year, and they are included in the last section.

The Users Group held its third annual meeting on June 25. Outgoing Chairman Ralph Minehart presided, and the prize for the user who made the most valuable contribution to CEBAF programs during the previous year went to Warren Buck (Hampton University) for his organization of the HUGS program. The new chairman, Claude Williamson, reminded everyone of the important role the Users must play in the definition of the CEBAF experimental facilities which must be made during the coming year.

The discussions at the Workshop were a valuable contribution to planning for the CEBAF scientific program, and we thank all the participants for coming and sharing their ideas with us.

Franz Gross
CEBAF and College of
William and Mary

Ralph Minehart
University of Virginia

Workshop Editors

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Letter to the Workshop Participants from the Scientific Director

Professor Franz Gross
Acting Associate Director for Research
CEBAF
12070 Jefferson Avenue
Newport News, Virginia 23606

June 16, 1986

Dear Franz,

I would like to ask you, on my behalf, to welcome the participants in the CEBAF 1986 Summer Workshop. Because of commitments made some time ago, I will not be able to attend the workshop this summer. The earliest I can get to CEBAF is July 1, 1986. I plan to spend the month of July in Newport News, honor the commitments I have made for August, and then assume the position as Scientific Director on a full-time basis starting September 1, 1986.

The program you have arranged for the workshop looks interesting and wide-ranging. The topics covered are at the forefront of the field, and the material will be of immediate use in planning the future program at CEBAF. It is my intention to use the material produced at this workshop as a base for the development of the scientific program at CEBAF. In this regard, I see the continued recruiting of an outstanding scientific staff as the first priority. The next priority is the development of the basic experimental equipment for the laboratory.

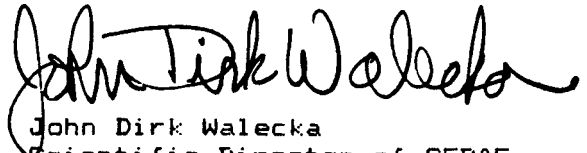
The true base for the success of this project is outstanding people with strong drive and motivation to understand the properties and behavior of the nucleus. This base of support must come from the permanent scientific staff, from nuclear physicists in the SURA institutions, and ultimately, from the national nuclear physics community whose members will be the users of this national facility. My biggest regret in missing the workshop is that I will not be able to meet the people there, to renew old acquaintances and establish new ones, and to get the participant's views and advice on how to make this project a special success. Rest assured that I will make every effort to solicit this advice after I arrive. Perhaps in the meantime, Franz, you will gather some of these ideas at the workshop and transmit them to me.

Herman Grunder and his staff will deliver a magnificent instrument. In my opinion, the scientific goal of this facility is to study the structure and behavior of the nuclear many-body system, its quark substructure, and the strong and electroweak interactions governing the properties of this fundamental form of matter. It is an exciting time for nuclear physics. The past few years have seen new insights into nuclear structure that raise as many questions as they answer. In addition, our whole picture of the nucleus has been revolutionized by the recent advances in high-energy physics. There is an entirely new underlying set of degrees of freedom in this fundamental condensed matter system. CEBAF will be a national users facility at the forefront of the field. It will provide an unmatched tool for the next generation of nuclear physicists. It is my goal

that CEBAF should be not only a superb experimental facility, but also one of the world's centers for nuclear physics research.

Again, Franz, please give both my welcome to the workshop participants, and my regrets at being unable to attend this one. I hope the workshop is a great success. I eagerly look forward to studying the proceedings. And finally, my heartfelt thanks to you, to Ramona, and to all the staff for the hard work that went into organizing this important endeavor.

Sincerely yours,

A handwritten signature in black ink, reading "John Dirk Walecka". The signature is fluid and cursive, with the first name "John" and last name "Walecka" being more prominent than the middle name "Dirk".

John Dirk Walecka
Scientific Director of CEBAF
As of September 1, 1986

Overview and Status of the CEBAF Project*

Summary Description

The Continuous Electron Beam Accelerator Facility will be a superconducting recirculating cw electron accelerator operating with 100% duty factor at 200 μA and at energies of 0.5 GeV to 4 GeV.

Acceleration will take place in two antiparallel linac segments, with recirculation beamlines for four orbits. Each linac segment will contain 200 CEBAF-Cornell superconducting accelerating cavities with an energy gain per segment variable up to at least 0.5 GeV, or 1 GeV per orbit. A central helium liquefier will supply liquid helium at 2.0 K to insulated cryomodules containing eight cavities each. An injector, an instrumentation and control system, and an rf power system will also be provided. The extraction system will direct simultaneous beams at different (but correlated) energies to each of three end stations.

Beam Parameters

Key performance objectives are for the machine to provide a high-energy, high-duty-factor, high-intensity electron beam of sufficient beam quality to allow high-resolution experiments. Table 1 displays baseline beam parameters for meeting these objectives.

Table 1
Baseline Beam Parameters

Energy	$0.5 \leq E_0 \leq 4.0 \text{ GeV}$
Duty factor	100%
Total average beam current	$1 \mu\text{A} \leq I_{\text{ave}} \leq 200 \mu\text{A}$
Number of simultaneous beamlets in accelerating structure	≤ 4
Number of simultaneous extracted beams	≤ 3
Beam emittance ($\pi\epsilon$)	$\leq 10 \text{ nm } \pi \text{ rad (0.5--4 GeV)}$
Beam energy spread (4σ)	$\Delta E/E \leq 4 \times 10^{-4}$

The design preserves the potential for significant upgrades in energy beyond and outside the scope of the present project.

* Information from *CEBAF Design Report* and *CEBAF Design Handbook*.

Detailed Description

The following discussion makes general reference to Figure 1, a schematic view of the accelerator, and to Figure 2, a plan view of the CEBAF site as it will appear in 1992.

The major components of the linac, listed by accelerator function, are:

- The injector.
- Linac segments I and II (the north and south segments, respectively) where acceleration occurs, with voltage gain in each segment variable up to at least 0.5 GeV.
- The beam transport elements making up the spreaders, recombiners, matching sections, and recirculation arcs.
- Finally, connected to the accelerator by the extraction elements, the beamlines of the beam switchyard (BSY) leading to the experimental stations.

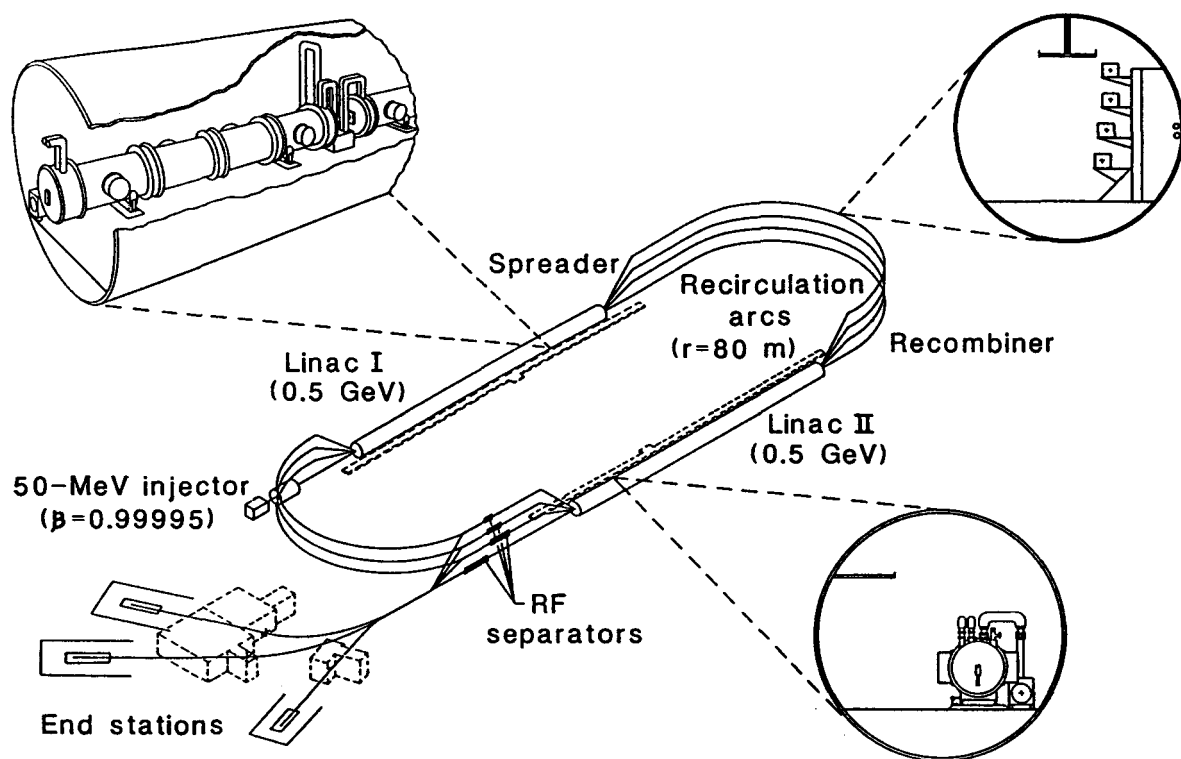


Figure 1 Schematic overview of the CEBAF accelerator.

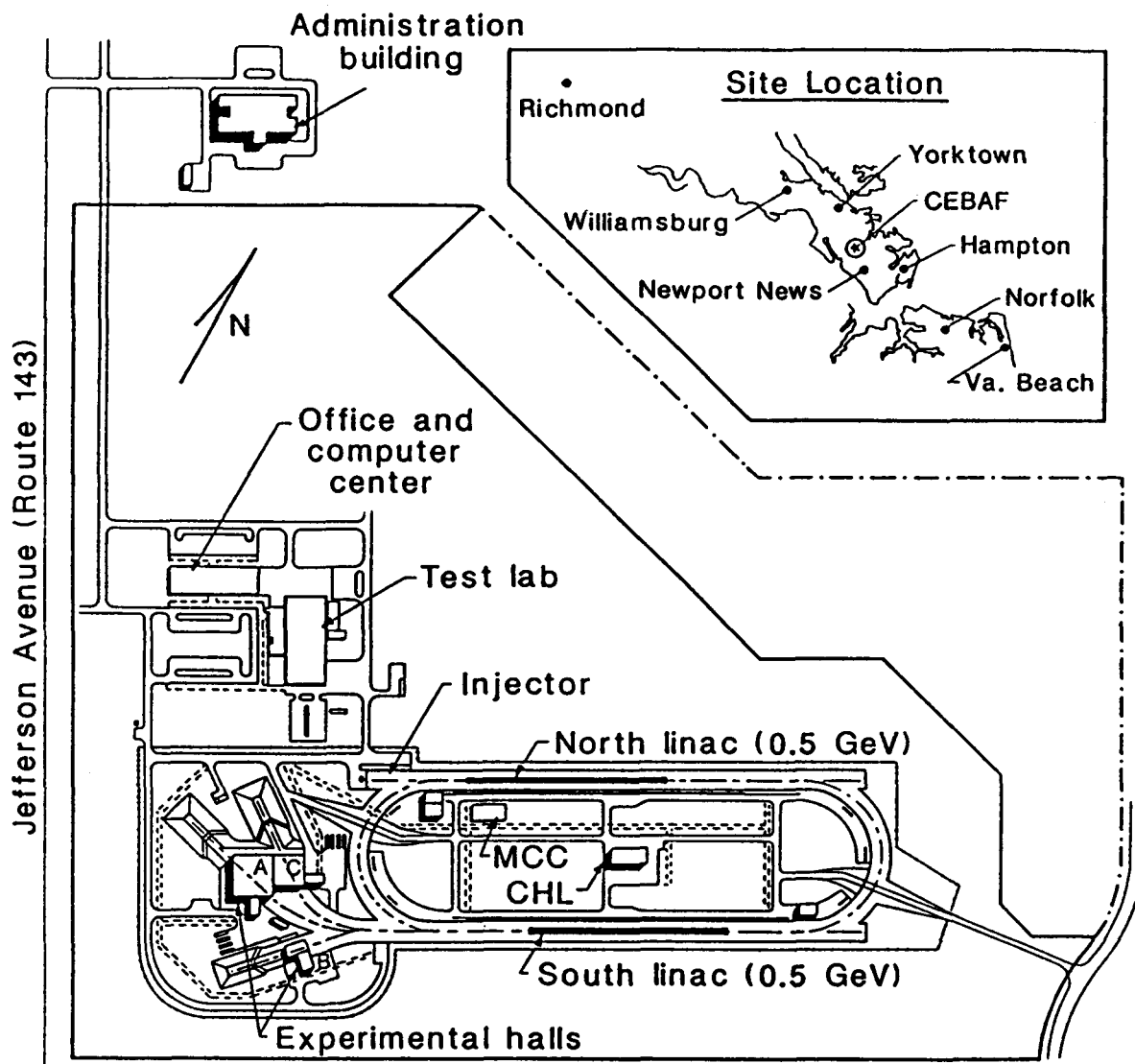


Figure 2 Plan view of CEBAF.

This site plan shows CEBAF as it is planned to look in 1992. CEBAF's present work takes place in the administration building, with plans for expansion soon into the rf superconductivity test laboratory, the other pre-existing structure on the site. The injector, the accelerator, and the beamlines to the end stations will be enclosed in tunnels covered by earthen berms. MCC is the machine control center; CHL is the central helium liquefier.

Beam is injected at 50 MeV (nominal), $\beta = 0.99995$, sufficiently relativistic to slip in phase only $\simeq 1.9^\circ$ through the first linac, with almost no phase slip after that—an essential ingredient for the recirculation scheme. Up to four passes through the accelerator are allowed by the recirculation arcs that transport the beam from one linac segment to the next. Key design considerations for the recirculation arcs are high-quality optics (isochronicity, achromaticity, matching in all phase planes) and the avoidance of, or minimization of, quantum excitation effects. Their average radius has been chosen such that with a new magnet lattice, but within the same enclosure, beams of up to 16 GeV could be transported with high beam quality maintained.

Matching the three experimental areas, the injector creates three distinct bunch trains (bunches $1 + N \cdot 3$, $2 + N \cdot 3$, $3 + N \cdot 3$) that can be of individually adjusted current. The extraction system, based on rf separators, i.e. cavities in a deflecting mode, operating at 2.5 GHz, has the capability of dispatching beams of up to three different (although correlated) energies to the end stations.

In this four-pass system, four electron beams, each at a different energy, are simultaneously present on the same trajectory in the linac segments. Since each beam is at a different energy, each will require a separate recirculation path tuned to accommodate its particular energy. Vertically separated recirculation beamlines comprise these paths—four of them in the case of arc I (the east arc), and three in arc II (the west arc). Upon exiting a linac segment, the beams at different energies are spread apart and transported to the appropriate recirculation beamlines through the use of a *spreader*. At the end of the recirculation arc, the beams are transported from each recirculation beamline, recombined, and injected into the next linac through the use of a *recombiner*.

The path of a single electron bunch from injector to experimental area therefore comprises up to four acceleration cycles. After injection and before extraction, each of these passes is as follows:

- acceleration in linac segment I,
- spreading to the proper beamline in recirculation arc I,
- recirculation,
- recombining for injection into linac II,
- acceleration in linac II,

- spreading to the proper beamline in arc II,
- recirculation, and
- recombining for reinjection into linac I.

The sequence can be interrupted on any cycle for beam extraction following spreading into arc II. Extraction elements in each of the arc II beamlines allow extraction and transport via the beam switchyard to the experimental area. With this extraction system the accelerator can deliver continuous beams to each of the three end stations simultaneously, and can allow the experimental group in any end station some latitude in choosing its beam's energy. That is, if the maximum energy being delivered to any of the end stations is $E_{\max}$, one beam will use $E_{\max}$, another beam can use $E_{\max}$, $3/4 E_{\max}$, $1/2 E_{\max}$, or $1/4 E_{\max}$, and the third beam can use either $E_{\max}$ or whichever of the three fractions of $E_{\max}$ the second beam is *not* using. Any of the three selected energies can be delivered to any of the three end stations.

The extraction elements, time-dependent deflecting devices called rf separators, use the two systems shown in Figure 3. For beam extraction at cycles 1, 2, and 3, they allow one out of every three passing bunches to be steered into the extracting beamline rather than into the recirculation beamline. A fourth rf separator, capable of generating up to three beams, is used in the extraction line after the fourth cycle when either two or three end stations require $E_{\max}$.

Energies between 500 MeV and 4 GeV can be achieved by lowering the energy gains per linac segment and extracting bunches of the desired energy after the applicable number of cycles. For example, a beam of 1.6 GeV can be obtained through use of an energy gain of 0.4 GeV per segment, with extraction after the second pass—that is, after two orbits through the pair of linac segments.

Cavity Prototyping Program

CEBAF's five-cell, 0.5-meter-long niobium cavity (Figure 4) was developed at Cornell University's Newman Laboratory of Nuclear Studies. Working with the Cornell laboratory, CEBAF is engaged in an industrial prototyping program preparatory to larger-scale production of these cavities and their cryostats. The goals are to qualify industry to produce these key components, and also to optimize cavity performance. So far eleven cavities have been ordered from five vendors, and four cryo-units from three vendors. Six cavities from four vendors have been fabricated and tested. During 1987, industrially fabricated cavities will be assembled in pairs

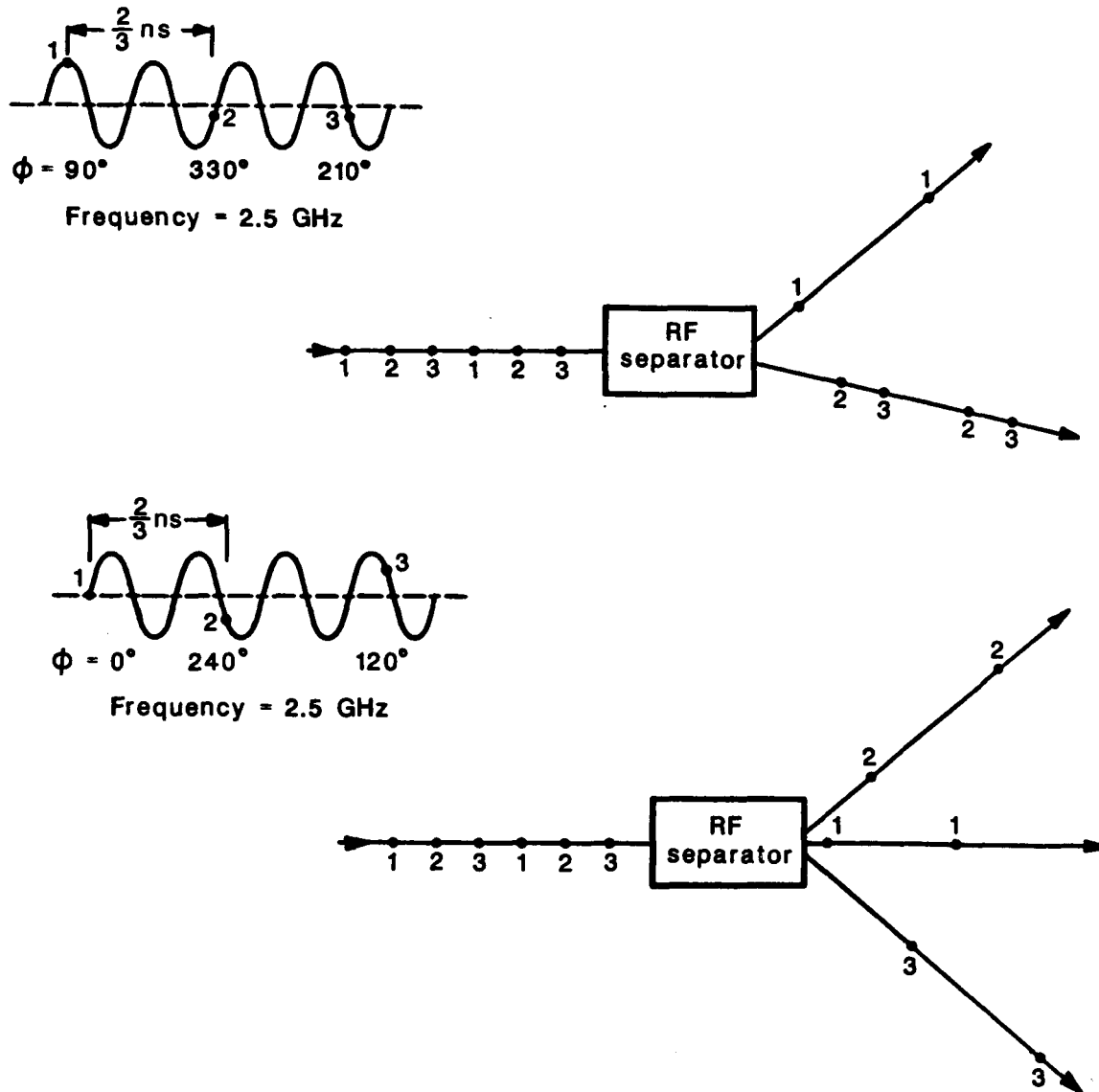


Figure 3 Deflection of microbunches by rf separators.

The system depicted at the top of the figure is phased so that one of the microbunches receives a deflection opposite to the other two; this system can therefore split the available current into two separate beams, one of which may be extracted while the other is recirculated for further acceleration. The system at the bottom involves the same principle, but the phase of the microbunches relative to the rf has been adjusted to provide for the generation of three distinct beams. The CEBAF extraction scheme is based upon these two systems.

in cryo-units, and then in a cryomodule of four cryo-units—the actual operational configuration—for further testing. It is expected that much of this work will be done at the superconducting rf test laboratory being prepared at the CEBAF site (Figure 2). The cryomodule is shown in perspective and in cross section in Figure 1; a cavity pair is shown in Figure 5. The operating frequency of the cavity is 1497 MHz.

The six cavities produced so far by industry have exceeded design specifications substantially (Figure 6). To meet specifications, the CEBAF accelerator must achieve a gradient of 5 MV per meter of active accelerating structure, with a Q_0 of 2.4×10^9 at 2 K. In gradient, the six cavities have achieved from 6.1 to 10.5 MV/m; since the cavities will be individually controlled, high gradients like these can be exploited for high-energy operation. Q_0 measurements have ranged from 5 to 7.5×10^9 .

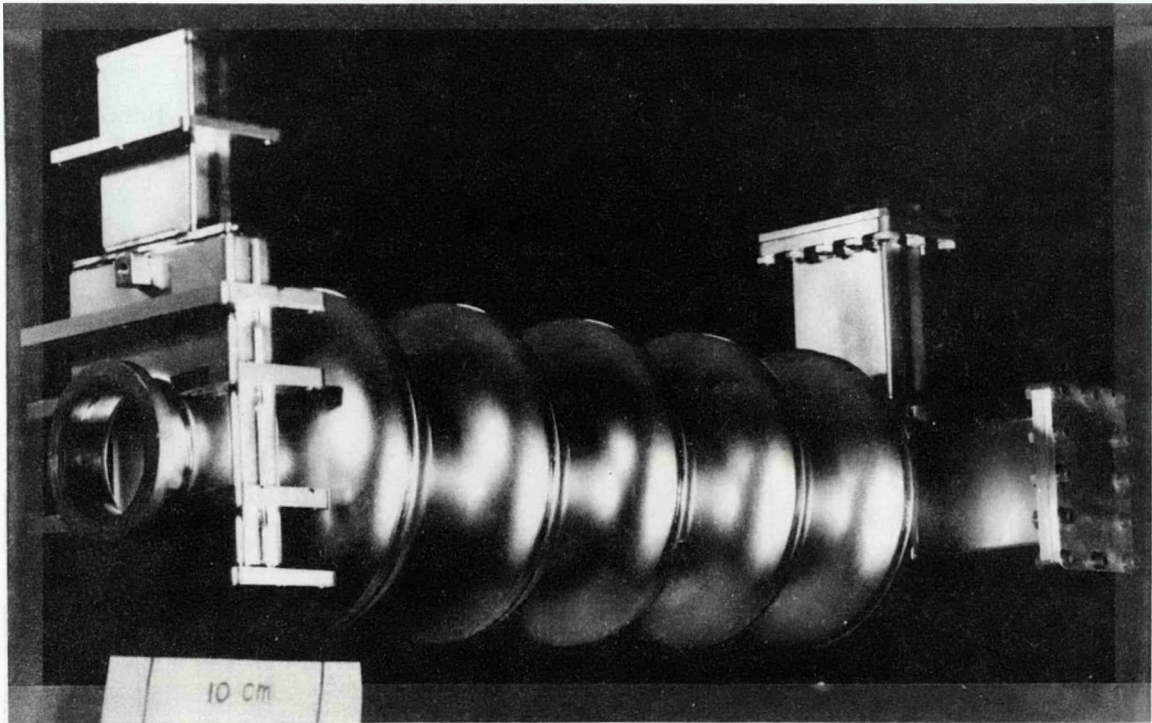


Figure 4 The CEBAF-Cornell niobium superconducting accelerating cavity.

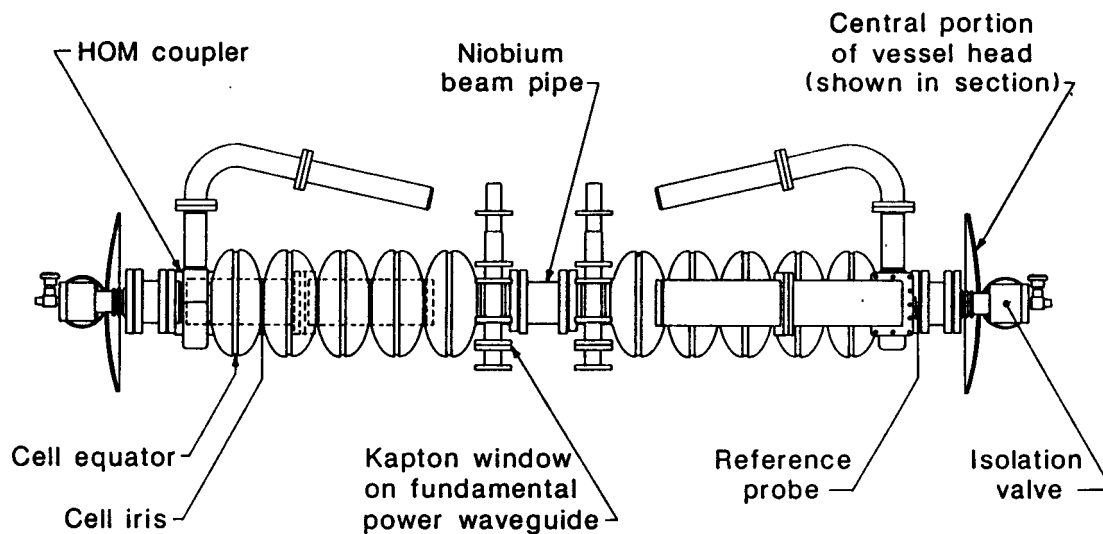


Figure 5 Diagram of a cavity pair.

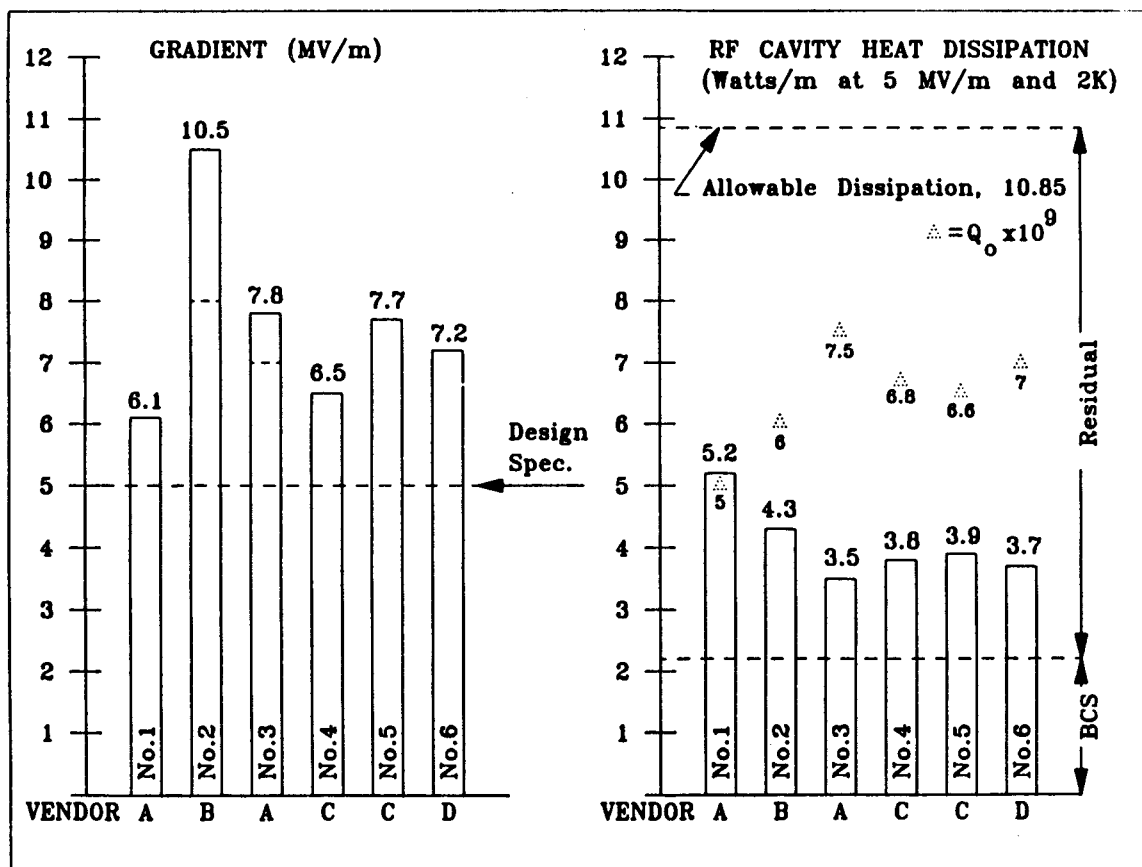


Figure 6 Performance of prototype cavities.

Project Status

With a detailed design in hand, and with initial construction money appropriated by Congress for fiscal year 1987, CEBAF is ready for groundbreaking. The design and specifically its key technology, rf superconductivity, have been thoroughly reviewed and approved by the Department of Energy and by several scientific panels as well. During the summer of 1986 a comprehensive *Design Handbook* was prepared; this two-volume manual will serve as an ongoing reference during, and after, construction. The federal budget contains a \$17 million appropriation for commencement of construction, which will thus begin early in 1987.

Total estimated cost for the project is \$236 million; Figure 7 breaks this down into its four major components. *Conventional construction* refers to the tunnels, buildings, and other structures and facilities that house and support the accelerator and the research equipment. *Contingency* refers to funds held in reserve.

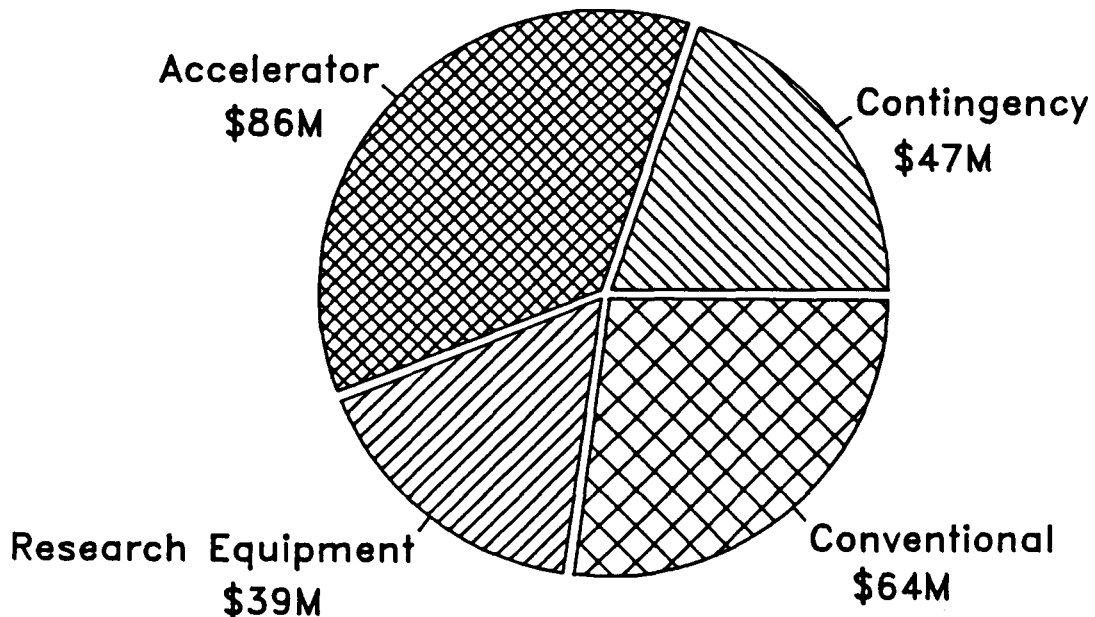


Figure 7 Components of CEBAF's total estimated cost.

Meanwhile the CEBAF staff continues to grow, from 66 in December of 1985 to nearly 100 projected by the end of 1986. It is expected that the staff will number around 290 when the laboratory is ready for physics research in 1992.

CEBAF Scientific Program

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1. Introduction and Popular Overview

The exciting challenge facing nuclear physicists in the years ahead is to understand precisely how nuclei and nuclear matter are built up out of quarks and gluons, nucleons and mesons, and to understand how the effective forces in nuclei arise from the fundamental strong forces describable by Quantum Chromodynamics (QCD) and the fundamental electroweak forces describable by the standard model. The nuclear physicist today finds himself in a situation similar to that of the condensed matter physicist; he knows the fundamental force laws and the basic constituents and seeks to explain the structure and behavior of macroscopic systems constructed from these elementary constituents. In both fields, the path to be followed to achieve such an understanding is far from clear. Knowing the fundamental force law is a significant advantage, but collective phenomena not immediately suggested by the basic force laws can be expected (such as superconductivity in condensed matter physics, and one pion exchange forces in nuclear physics) and a major experimental and theoretical program is needed to discover and understand these phenomena.

The principal scientific mission of CEBAF is to carry out these studies on cold (or normal) nuclear matter. The electron, with its penetrating point-like interaction is an ideal probe of cold nuclear matter; by independently varying the energy and three-momentum transferred by the virtual photon, it is possible to either leave

the cold nucleus undisturbed by the probe, or to highly excite it and study individual constituents which emerge from the interaction. The principle scientific mission of the relativistic heavy ion collider (RHIC), planned to be constructed at Brookhaven National Laboratory, is to carry out these studies on hot nuclear matter. By smashing two high energy heavy ions together, nuclear matter at high temperature and energy density is obtained, and it should be possible to create a plasma of quarks and gluons. With both facilities it will be possible to understand the behavior of nuclear matter under a wide variety of conditions. The need for both of these new facilities has been recognized by the Nuclear Science Advisory Committee (NSAC)¹.

Specific experimental programs currently planned for CEBAF are addressed in detail in Part 2. In this section, issues which the CEBAF experimental and theoretical program is designed to address are discussed in broad popular terms. Among the questions which CEBAF will address are:

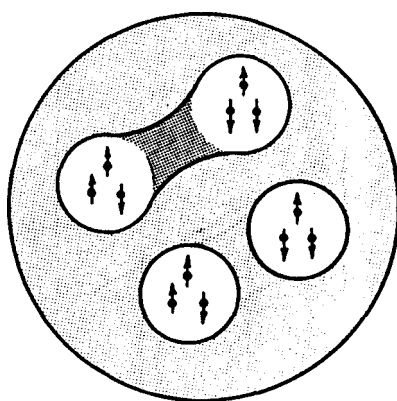
- How is the structure of the nucleon modified in the nuclear medium? Does the nucleon swell?
- What is the nature of two-nucleon subsystems in the nuclear medium? Are nucleon-nucleon correlations large? Are there important contributions from independent 6-quark channels which do not overlap the NN subspace (dibaryons)?
- What is the precise role of excited baryons in the nuclear medium? Specifically, what is the free baryon spectrum, and how important are these baryons to the structure of cold nuclei? How is the structure of N^* 's and Δ 's modified by the nuclear medium? What are the N^*-N and $\Delta-N$ forces?
- What is the role of strange and other heavy quarks in the nuclear medium?
- How are elementary process, such as e-p scattering modified by the nuclear medium? How do quarks hadronize in the nuclear medium?
- How can the dynamics of quark confinement be best described and modeled (solitons, bags, fluxtubes, potentials)? How is this confinement modified by the presence of other nucleons, such as occurs in the nuclear medium?

- What degrees of freedom are important to the description of the nuclear wave function at short distances (high momentum transfer), and how large are its high momentum components?
- What are the "smoking gun" signatures (if any) of new phenomena related to quark degrees of freedom in nuclei?

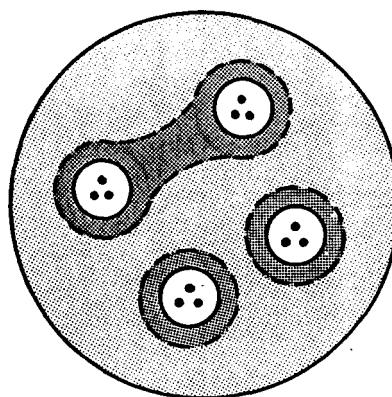
A broad program of both theoretical and experimental work will be necessary before many of these questions can be answered.

If there are new phenomena related to the underlying quark degrees of freedom, their appearances and importance can be expected to depend strongly on the effective quark confinement radius. When nucleons are far apart, their constituent quarks remain with the parent nucleon. In this case the nuclear force is transmitted by the pion, a collective mechanism not currently understood in the language of QCD, and new phenomena related to the mingling of quarks should be absent. As nucleons are brought closer together, the confining forces may weaken, allowing the quarks to move from one nucleon to another, and new phenomena related to this mingling may occur. Clearly, the critical distance at which this happens will be related to the effective confinement radius of quarks inside of a nucleon.

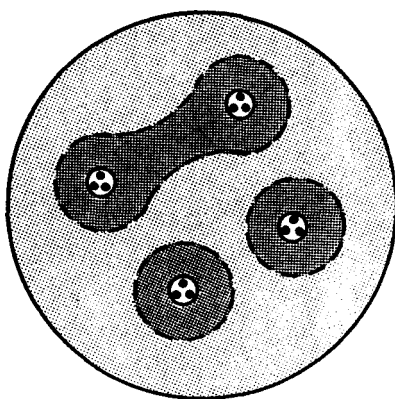
Figure 1 shows three different views of the ${}^4\text{He}$ nucleus, which are appropriate for different values of the confinement radius. In the first case, the confinement radius is large -- comparable to the radius of the nucleon itself. The meson cloud around the nucleon is absent, or very insignificant. Here the long range pion force still would play a role, but shorter range meson exchange forces would be replaced by quark interchange mechanisms. Signatures of explicit quark degrees of freedom should be visible even at comparatively low excitation energies and low momentum transfers, and the success of conventional nuclear models based on hadronic degrees of freedom with meson exchange would be an accident, or a consequence of duality (see below). This picture is suggested by the original MIT bag model².



(a)



(b)



(c)

Figure 1 Cartoons of the ${}^4\text{He}$ nucleus for different values of the effective confinement radius: (a) large (quarks occupy the entire volume of the nucleon), (b) intermediate (quark core with a thick meson "skin"), (c) small (nucleons are mostly mesons with a tiny quark core). Dark shading - pion (and other meson) clouds. Light shading - normal vacuum. White - QCD vacuum.

The other extreme possibility is that the confinement radius is very small, and that the nucleon is largely a meson cloud with the valence quarks confined inside of a tiny volume of nearly zero radius, as illustrated in Figure 1(c). This picture is suggested by the $N_c \rightarrow \infty$ limit of QCD (N_c is the number of colors), in which, to lowest order in N_c^{-1} , baryons are built from topological soliton solutions of effective meson field equations³, such as the Skyrmin model⁴. It is also suggested by relativistic meson theory, in which the structure of the nucleon would arise primarily from its meson cloud. In this case, it is unlikely that effects from the quark core will be visible, and nuclear structure is likely to be fully understood in terms of meson interactions, with the possibility that non-linear meson interactions, such as those treated by Skyrmin models, will play a role.

Perhaps the most likely possibility is that the confinement radius will be intermediate in size, as illustrated in Figure 1(b). This picture is suggested by the cloudy bag model⁵, which has been quite successful phenomenologically. In this case, explicit effects from the quark core may appear only at the high momentum transfers which can be studied at CEBAF.

The previous discussion assumes that there are new phenomena related to quark degrees of freedom. It is possible that there are no smoking gun signatures, i.e. new phenomena which cannot be explained by meson theory. This would be the case if an exact duality between quark-gluon descriptions and meson-nucleon descriptions were to hold. That this might be the case is suggested by the observation⁶ that any color singlet 6-quark state can be expressed in terms of a sum of states consisting of two color singlet 3-quark states. The three quark states in the sum may be highly deformed, corresponding to highly excited baryons, so that even in this case it could be more efficient to use a 6-quark basis consisting of a few states rather than a basis consisting of a large number of excited 3-quark baryons. But this observation raises the possibility that, if the number of meson and baryon states included

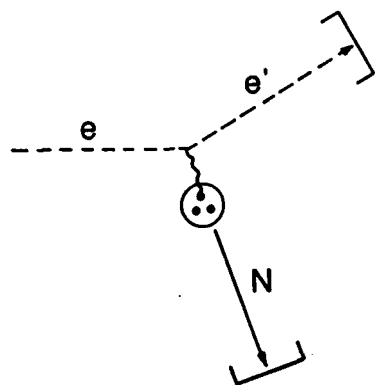
in a effective meson theory is enlarged sufficiently, it may be possible to describe all short range phenomena in terms of meson theory. Theoretical models which allow for this duality are needed.

While most of the questions raised in the beginning of this section will require considerably more thought before they can be formulated in sufficiently precise terms to be tested, the first three questions can be captured in a cartoon which will serve to give a popular focus to the CEBAF program. The cartoon, Figure 2, shows a ${}^4\text{He}$ nucleus being probed by a high energy electron. At the moment it is being probed, it is in a state where one of its nucleons is swollen, one is in an excited Λ or N^* state, and two have overlapped to form a 6-quark state. In Figure 2(b), the electron strikes one of the quarks in the swollen nucleon, knocking the nucleon from the nucleus. Both the electron and the nucleon are detected in an $(e, e'N)$ experiment, and the cross section, when compared to that from a free nucleon (illustrated in Figure 2(a)) will reflect the altered charge distribution of the swollen nucleon. Alternatively, if the electron strikes the 6-quark state (Figure 2(c)), it could lead to an enhanced $2N$ emission, which, if the two nucleons are highly correlated, could be detected in two spectrometers, one of which has a large acceptance. Finally, striking the Λ or N^* (Figure 2(d)), could lead to a final $2\pi N$ state, and all of these fragments can be detected in a Large Acceptance Spectrometer (LAS), with a nearly 4π acceptance as shown. These three experiments can be thought of as typical of the programs to be carried out in the three experimental halls described in Section 3; each requires its own unique complement of experimental equipment.

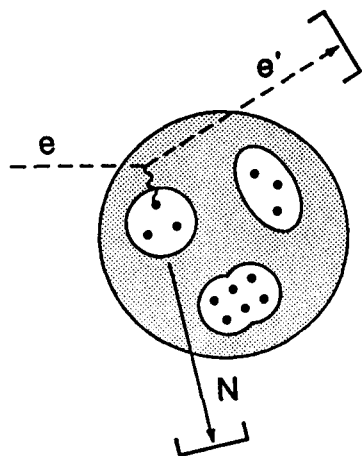
2. Highlights of Experimental Programs Planned for CEBAF

Highlights of seven experimental programs which have served to define the CEBAF experimental equipment will be discussed in this section. These are:

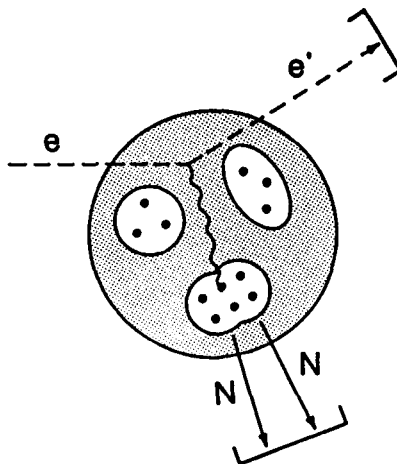
- Single nucleon emission - $(e, e'N)$
- Two-nucleon emission - $(e, e'2N)$



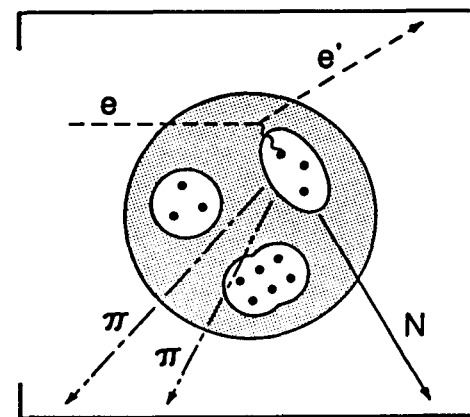
(a)



(b)



(c)



(d)

Figure 2 Cartoons representing electron scattering. In each case the electron emits a virtual photon which is absorbed by a quark. (a) Scattering from a free nucleon followed by detection of the outgoing electron and nucleon in coincidence. The other three figures all show scattering from ^4He , with the electron striking different constituents. (b) The electron strikes a bound nucleon, swollen by the interactions with the nuclear medium, and the outgoing nucleon and electron are detected in coincidence. (c) The electron strikes a six-quark state, which emerges as two nucleons detected in coincidence with the electron. (d) The electron strikes an excited nucleon, which shatters into two pions and a nucleon, all of which are detected in a LAS.

- Study of excited baryons
- Hypernuclei - (e,e'K)
- Charge distribution of the neutron and deuteron
- Studies in the deep inelastic region
- Studies of the electroweak interaction

The first six of these programs were reviewed briefly in the RPAC report⁷; the latter has been discussed in the original SURF proposals⁸, and in the 1981 report of the Bates Study Group⁹. CEBAF contributors to RPAC include V. Burkert, F. Gross, B. Mecking, J. Mougey, and R.R. Whitney.

2.1 Single Nucleon Emission

The original motivation for measuring single nucleon emission in coincidence with electron scattering is that it could give a direct measurement of the square of the momentum space nuclear wave function, called the single particle density. As an introduction to coincidence measurements, and to see how this comes about, assume that meson exchange effects and final state interactions are small (this is often not the case -- see the discussion below). In this case, the electromagnetic interaction from the simplest nucleus, the deuteron, can be calculated from the two impulse approximation diagrams shown in Figure 3. If magnetic scattering is neglected (for simplicity) the differential cross section obtained from these diagrams is

$$\frac{d^5\sigma}{d\Omega dE' d\Omega_{p_1}} \simeq \sigma_M |F_p(Q^2)\psi(p_2^2) + F_n(Q^2)\psi(p_1^2)|^2 \quad (1)$$

where the scale of the overall interaction is set by the Mott cross section,

$$\sigma_M = \left(\frac{2\alpha E' \cos\theta/2}{Q^2} \right)^2, \quad (2)$$

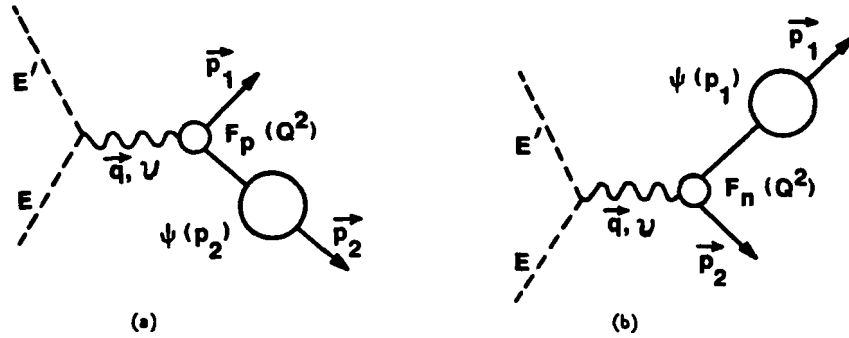


Figure 3 Impulse approximation diagrams for the $d(e, e'p)$ experiment.

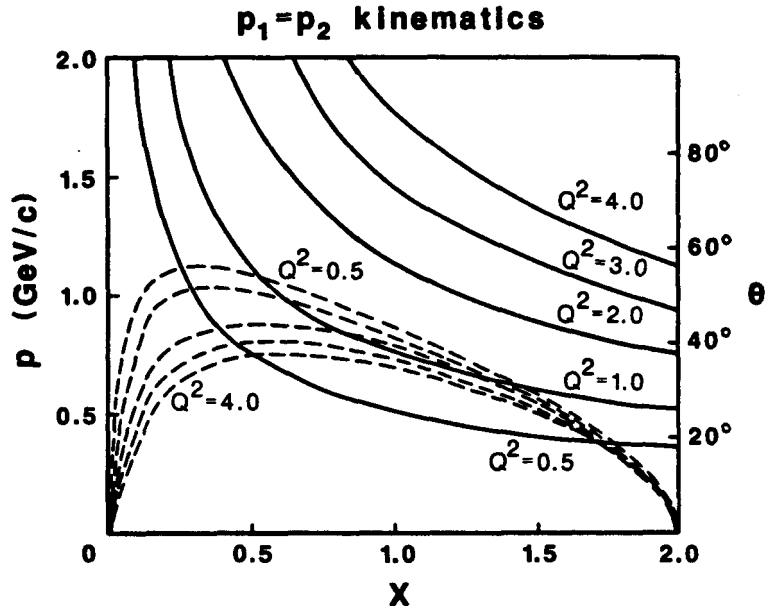


Figure 4 The dependence of nucleon momentum, p , and angle, θ (equal to the angle between p_1 and q -- also one-half the opening angle between the two nucleons), for the symmetric kinematics described in the text. The solid lines show the variation of p for constant Q^2 , to be read using the scale on the left, and the dashed lines are θ , to be read on the right.

with Q^2 the 4-momentum transferred by the electron

$$Q^2 = q^2 - \nu^2 = 4EE' \sin^2 \frac{1}{2}\theta \quad , \quad (3)$$

$F_p(Q^2)$ and $F_n(Q^2)$ are the form factors of the bound proton and neutron, respectively, and ψ is the wave function of the deuteron (which can be defined relativistically¹⁶ by relating it to the d-n-p vertex function in a covariant way). To extract the square of the wave function from (1), it is helpful to choose symmetric kinematics so that $|\vec{p}_1| = |\vec{p}_2| = p$. Then

$$\frac{d^5\sigma}{d\Omega \, dE' \, d\Omega_p} \simeq \sigma_M |F_p(Q^2) + F_n(Q^2)|^2 |\psi(p^2)|^2 \quad . \quad (4)$$

By varying Q^2 and the Bjorken scaling variable

$$x = \frac{Q^2}{2M\nu} \quad (5)$$

independently for this symmetric kinematics, it is possible to measure the p^2 variation of this expression for a fixed Q^2 , or the Q^2 variation for a fixed p^2 , as illustrated in Figure 4, and hence to measure the bound form factor dependence and wave function dependence independently.

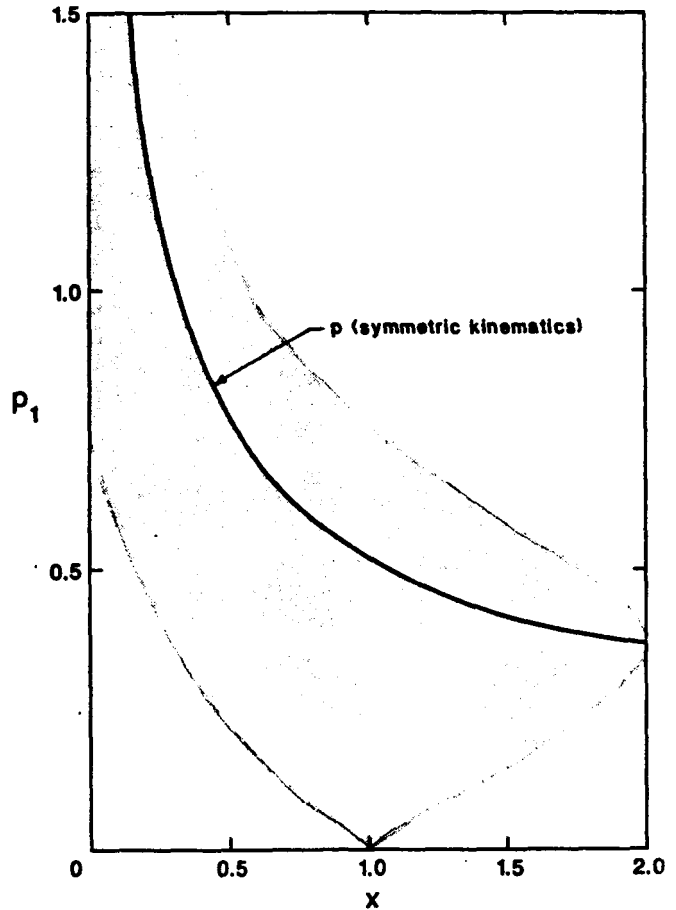
For comparison, the same simple theory gives the following expression for the single arm inclusive (e, e') cross section:

$$\frac{d^3\sigma}{d\Omega \, dE'} \simeq \sigma_M \int d\Omega_{p_1} |F_p(Q^2)\psi(p_2^2) + F_n(Q^2)\psi(p_1^2)|^2 \quad . \quad (6)$$

Since the outgoing nucleon is not detected, the cross section is dominated by that part of the region of integration where the wave function is the largest, which occurs for S states when the relative momentum is small. Figure 5 shows, for a fixed $Q^2 = 0.5 \text{ (GeV/c)}^2$, how the magnitude of $\vec{p}$ varies as the integral over its direction

in Eq. 6 is carried out. Note that for all $x \neq 2$, values of $|\vec{p}_1|$ occur which are considerably smaller than p given by symmetric kinematics. In particular, at the quasielastic peak, defined by the condition $x = 1$, the integral is dominated by contributions from regions where $\vec{p}_1^2$ or $\vec{p}_2^2 \approx 0$, regardless of the value of Q^2 . Coincidence experiments allow the experimenter to fix the momentum of the wave function at a particular value, and study the wave function in regions where it is small. In this way coincidence measurements give important new information not obtainable from single arm measurements.

Figure 5 Range of variation of $|\vec{p}_1|$ (or $|\vec{p}_2|$) under the integral in Eq. (6) for $Q^2 = 0.5 \text{ (GeV/c)}^2$ and different values of x . For comparison, the value of p fixed by symmetric kinematics is also shown.



Recent high resolution measurements¹¹ of $(e, e'p)$ from Pb at NIKHEF, give beautiful confirmation of the shell structure of nuclei (see Figure 6). A resolution $\Delta E \approx 100 \text{ keV}$ was obtained in the experiments, permitting individual nuclear levels to be resolved. Other measurements from Saclay on the deuteron¹² and ^3He (ref. 13) are shown in Figures 7, 8 and 9. The deuteron measurements are very

sensitive to the D state, and the ^3He measurements show the potential this program has for giving detailed information about three body wave functions. Even studies of the few body system require high resolution; separation of the two- and three-body final states requires the resolution of $\Delta E \approx 1$ MeV, obtained in this experiment and shown in Figure 8.

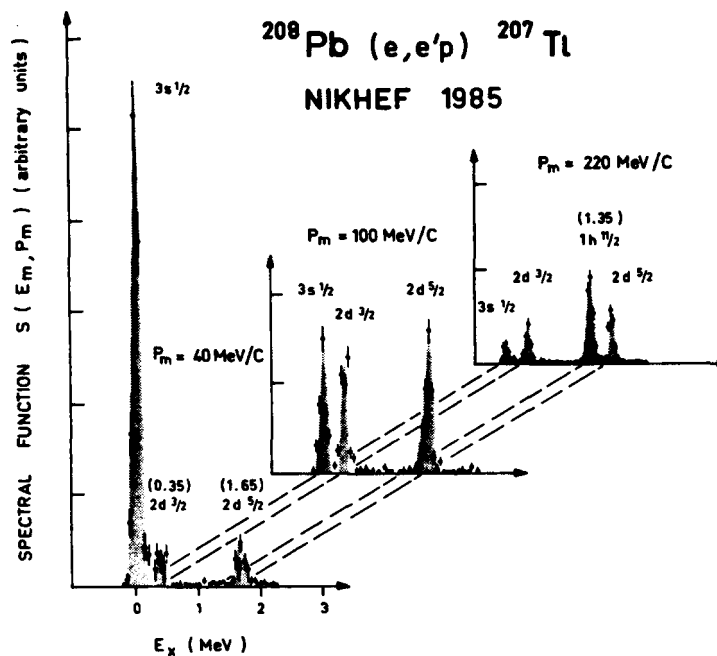
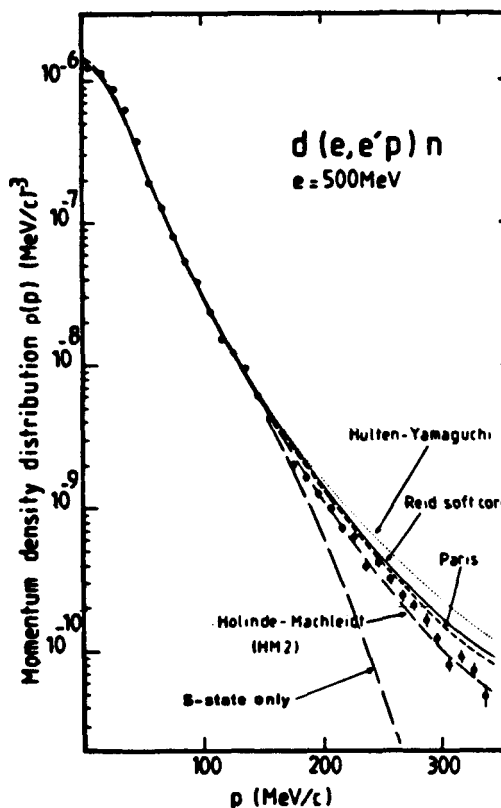


Figure 6 Recent NIKHEF data as plotted in Reference 11. High resolution permits different states of the final ^{207}Tl nuclear system, corresponding to knockout of protons from different shells in the lead nucleus, to be distinguished.

In practice, final state interactions and meson exchange currents will make important contributions to the cross section, altering significantly the ease with which quantitative conclusions about the nuclear wave function can be extracted from these experiments. For example, Arenhovel¹⁴ has shown that such effects vary from a few percent at $p^2 \approx 0$ to 100% at $p^2 \approx 300$ MeV for the deuteron measurements shown in Figure 7, and for other cases these effects can be considerably larger. This complicates the theoretical interpretations of these experiments and requires that

meson exchange effects, final state interactions, and contributions of the initial nuclear wave functions all be treated consistently.

Figure 7 Data from the Saclay measurement described in Ref. 12 compared with four theoretical models for the deuteron momentum density distribution. A dashed line showing the S-state contribution only, has been added to the figure.



At CEBAF, the $(e, e'N)$ program will be extended to higher nucleon momenta p and higher momentum transfer Q^2 , and the goal will be to study high momentum components of nuclear (especially few body) wave functions, study the three nucleon system fully and use such data to learn about 3 body forces, study the form factor of bound nucleons at high Q^2 (to see if nucleons are swollen as discussed in Section 1), and to study neutron distributions through the $(e, e'n)$ process. The latter cross sections are expected to be 10 times smaller than $(e, e'p)$ cross sections, and are much harder to measure because of the inefficiency of neutron detectors, but are expected to be every bit as rewarding as the $(e, e'p)$ program. They will require the full capability of CEBAF with its high current and high duty factor.

The general form of the $(e, e'N)$ cross section has been given by Donnelly¹⁵. If the recoil polarization of the nucleon is

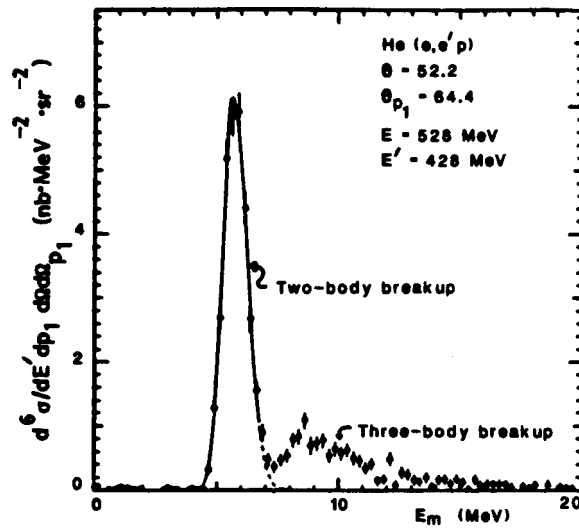


Figure 8 Distribution of missing energies, showing how the two-body $p + d$ final state can be separated from the three-body $p + p + n$ final state.

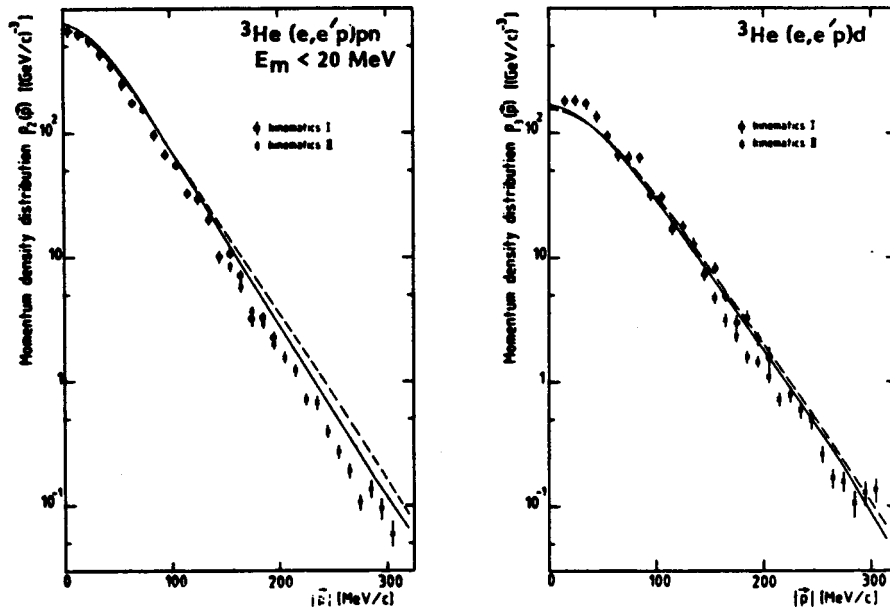


Figure 9 Momentum density distribution for the two-body and the three-body sectors in inelastic electron scattering from ${}^3\text{He}$.

undetected, and the target is unpolarized, the cross section can be written

$$\begin{aligned} \frac{d^6\sigma}{d\Omega dE' d^3p_1} = \frac{\sigma_M}{M_T} \left\{ v_L W_L^{(1)} + v_T W_T^{(1)} + v_{TT} W_{TT}^{(1)} \cos 2\phi_1 \right. \\ \left. + v_{TL} W_{TL}^{(1)} \cos \phi_1 + h v_{TL'} W_{TL'}^{(1)} \sin \phi_1 \right\} \end{aligned} \quad (7)$$

where ϕ_1 is the out-of-plane scattering angle of the nucleon (as shown in Figure 10), h is the helicity of the incident electron, and the kinematic factors are

$$\begin{aligned} v_L &= \rho^2 \quad 0 \leq \rho = Q^2/q^2 \leq 1 \\ v_T &= \frac{1}{2} \rho + \tan^2 \theta/2 \equiv \frac{1}{2} \rho/\epsilon \\ v_{TT} &= -\frac{1}{2} \rho \\ v_{TL} &= -\frac{1}{\sqrt{2}} \rho \sqrt{\rho + \tan^2 \theta/2} = -\frac{1}{2} \rho \sqrt{(1 + \epsilon) \rho/\epsilon} \\ v_{TL'} &= -\frac{1}{\sqrt{2}} \rho \tan \theta/2 = -\frac{1}{2} \rho \sqrt{(1 - \epsilon) \rho/\epsilon} \end{aligned} \quad (8)$$

where ϵ is the longitudinal polarization of the virtual photon, and the five W_i are structure functions that depend on four variables, which can be taken to be Q^2 , ν , p_1 , and θ_1 , as shown in Figure 10. An alternative form for Eq. (7) is:

$$\begin{aligned} \frac{d^6\sigma}{d\Omega dE' d^3p_1} = \Gamma \left\{ \sigma_T + \epsilon \sigma_L + \epsilon \sigma_{TT} \cos 2\phi_1 \right. \\ \left. + \sqrt{\epsilon(\epsilon + 1)} \sigma_{TL} \cos \phi_1 + h \sqrt{\epsilon(1 - \epsilon)} \sigma_{TL'} \sin \phi_1 \right\} \end{aligned} \quad (9)$$

where Γ is the virtual photon flux produced by the scattered electron

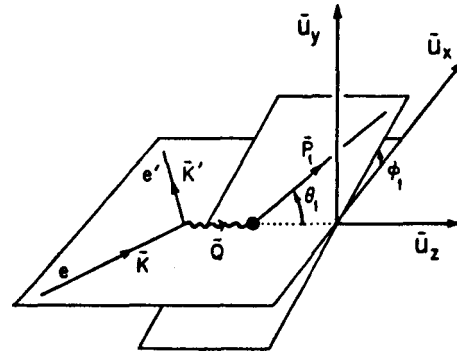
$$\Gamma = \frac{a}{2\pi^2} \left(\frac{E'}{E} \right) \frac{|\vec{q}|}{q^2 (1 - \epsilon)} = \frac{\sigma_M \rho |\vec{q}|}{4\pi^2 a \epsilon} \quad (10)$$

and

$$\begin{aligned} \sigma_T &= \frac{2\pi^2 a}{M_T |\vec{q}|} W_T^{(1)} \equiv C W_T^{(1)} \\ \sigma_L &= 2\rho C W_L^{(1)} \\ \sigma_{TT} &= - C W_{LL}^{(1)} \\ \sigma_{TL} &= - \sqrt{\rho} C W_{TL}^{(1)} \\ \sigma_{TL'} &= - \sqrt{\rho} C W_{TL'}^{(1)} \end{aligned} \quad (11)$$

Each of these five structure functions or partial cross sections are sensitive to different physical processes. The longitudinal structure function W_L is sensitive to charge distributions and nuclear wave functions, W_T is sensitive to magnetic moments, meson exchange currents, and Δ processes, and $W_{TL'}$ is zero without final state interactions. An experimental goal for CEBAF is to separately measure all five of these functions, and to thereby fully determine this important process. To measure W_{TT} and $W_{TL'}$ requires measurements out of the (e, e') reaction plane, but the others can be separated by making measurements at different electron scattering angles (different values of ϵ) and by observing the scattered proton on both the "left" (where $\cos \phi_1 = 1$) and "right" (where $\cos \phi_1 = -1$) sides of the three-momentum transfer $\vec{q}$.

Figure 10 Exclusive-1 electron scattering. Here the $\bar{z}$ -axis is chosen to be along $\bar{q}$ and the electron scattering occurs in the xz -plane. The particle detected in coincidence with the electron has momentum $\bar{p}_1$, whose direction is given by angles (θ_1, ϕ_1) in this coordinate system.



2.2 Two-Nucleon Emission

Two-nucleon knockout, $(e, e' 2N)$, can be expected to be sensitive to the pair distribution of nucleons in a nucleus, and to reflect the presence of localized six-quark states as discussed qualitatively in Section 1. Deuteron knockout, a special case, should reflect the existence of pre-existing deuterons or ^3S correlations. However, experience with pion absorption processes where two-nucleon knockout has been extensively studied¹⁶ suggests that final state interactions are very important. One mechanism which seems to occur is Λ excitation, followed by de-excitation through the $\Lambda N \rightarrow 2N$ process leading to two-nucleon emission. It is also possible that such processes can generate two Λ 's leading to three- or four-nucleon emission¹⁷. Very little is known about the electromagnetic two-nucleon knockout process; but studies with electrons, where it is possible to independently vary the energy momentum transfer, and polarization, and to separate a variety of structure functions sensitive to different physical mechanisms (as discussed in Section 2.1) should give much insight into this important process. CEBAF will be an ideal facility at which to carry out such studies.

The $2p$ knockout reaction is thought to be a promising way to study two-nucleon correlations, since exchange current contributions are strongly suppressed and the $2p$ system has no dipole moment¹⁸.

The transverse cross section is therefore expected to be quite small, and the longitudinal cross section should be sensitive to the correlations. However, calculations by Laget¹⁸ suggest that this process is very sensitive to final state interactions. Figure 11 shows the longitudinal (σ_L) and transverse (σ_T) cross sections as a

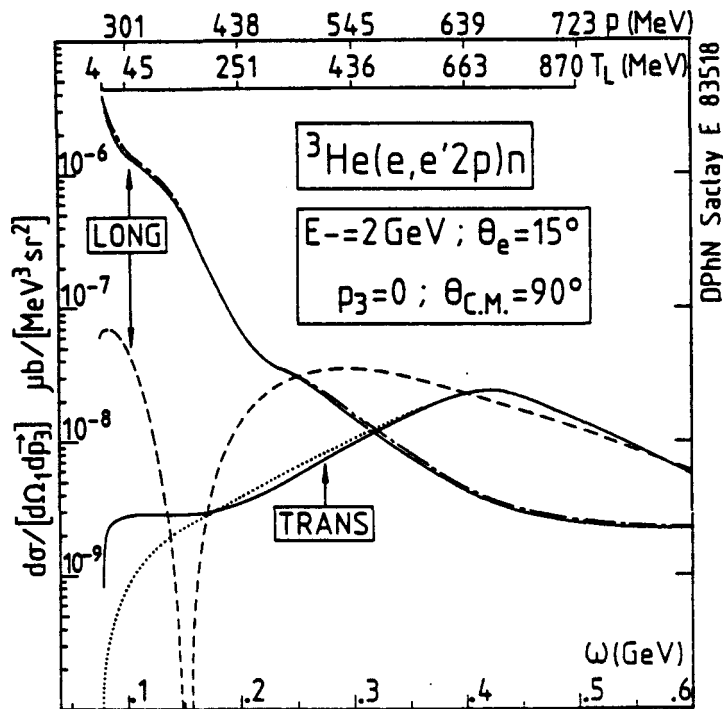


Figure 11 The excitation functions for the transverse and longitudinal cross sections as defined in Eq. (9). Symmetric kinematics with the proton pair initially at rest are assumed. The magnitude of each proton momentum p , and the relative kinetic energy T_L are shown along the top axis. Solid curves: full calculations including all final state interactions. Dashed curves: impulse approximation (plane waves in final state) only. Dotted curves: impulse approximation with meson exchange currents.

function of the incoming virtual photon energy ω for the case of symmetric kinematics described in Section 2.1 and Figure 4. In the

impulse approximation, σ_T is zero and σ_L shows the oscillatory structure of the initial state correlations clearly. When final state interactions are included, however, the clear connection with initial state correlations is washed out, σ_L can change its value by more than an order of magnitude, and σ_T is no longer zero. More calculations are needed, but it is clear that initial state correlations are only a part of what is happening at lower energies. Things may be simpler at the higher energies accessible to CEBAF where final state interactions are expected to be smaller.

The $(e, e'np)$ cross section is expected to be larger than the $(e, e'2p)$ cross sections by at least an order of magnitude, providing a strong incentive for the development of a neutron detection capability to go with this program.

As in the discussion presented in Section 2.1, it may be useful to vary the Q^2 dependence and the p^2 dependence of the cross sections independently. This can be done by using the symmetric kinematics described there, and varying both Q^2 and x while holding p constant, or p and x while hold Q^2 constant.

2.3 Study of Excited Baryons

Measurements of the spectrum of free baryons, together with their electromagnetic and hadronic decay modes, give information essential for constructing and testing realistic quark models of hadrons. Such models are important to nuclear physics for two reasons. First, they provide a bridge between QCD and nuclear structure by providing information and insight needed for the development of models of nuclei based on their underlying quark structure. Secondly, even if the construction of such models is trivial because the effective confinement radius is small and meson theories work for nuclear physics, it is necessary to fully understand the properties of the excited baryons which must necessarily play a role in such a meson theory.

It is also of equally great interest to study the propagation and decay of excited baryons in a nuclear medium, in order to determine N^*N interaction strengths (where N^* will be used to denote all excited baryons, including the Δ and Δ^* 's) and the effect of the nuclear medium on the structure of the N^* themselves.

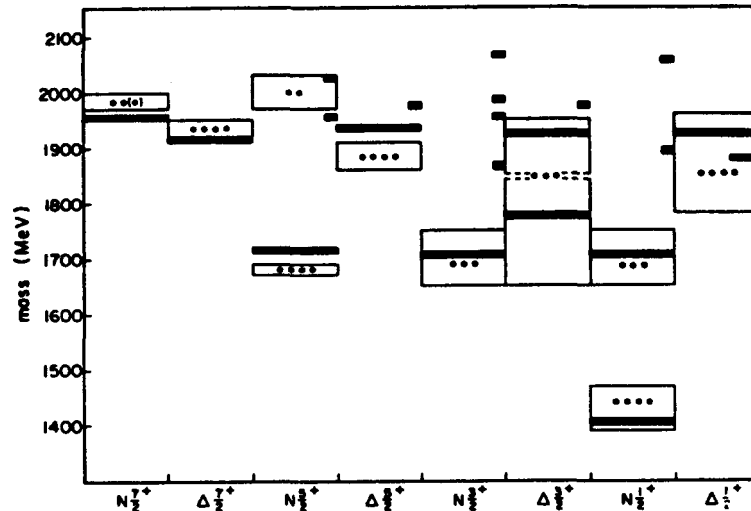
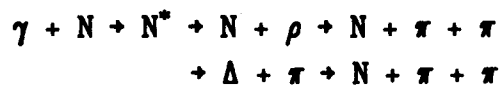


Figure 12 A chart of excited baryons (taken from Ref. 19). The open boxes with asterisks inside correspond to resonances which have been seen in πN scattering (the greater the number of asterisks the more certain are the observations). The solid bars are predictions. The short solid solid bars label states which are not expected to couple strongly to the πN channel, but could be seen in photoproduction experiments.

Figure 12 shows a table of the predicted and observed lowest lying baryon resonances¹⁹, all of which can be studied at CEBAF. Note that many have not been seen. This may be due to the fact that they do not appear to couple strongly to the πN channel²⁰, which has been the major experimental means of studying baryon resonances. Such resonances can be studied at CEBAF through the reactions



These studies will require a Large Acceptance Spectrometer (LAS) capable of detecting and identifying π 's, K 's, and N 's over a large

solid angle (80% of 4π). Initial studies will probably be limited to photoproduction measurements, where the goal will be to find the missing states, and to determine their mass and decay widths into γN , πN , and $2\pi N$. Such studies will be complicated by the fact that the widths of the resonance are expected to be as broad (at least 50 MeV) as the spacing between them, so that they strongly overlap (as shown in Figure 13 (and Figure 19 in Section 2.6 below). They will have to be separated by a partial wave analysis using a number of the observables which can be measured, including those which depend on polarizations.

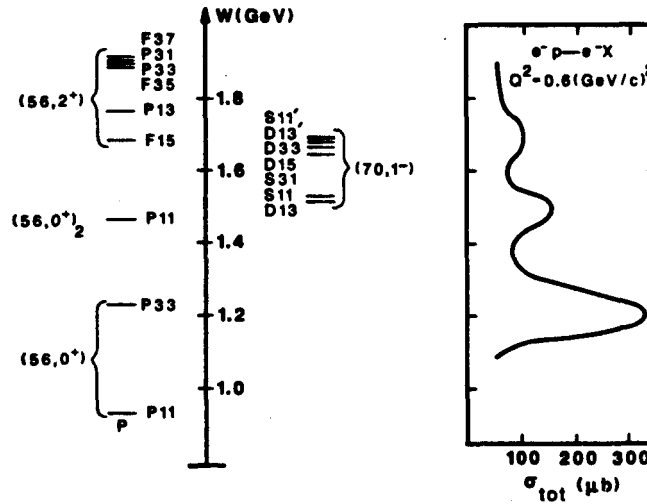


Figure 13 The best known nucleon resonances with masses below 2 GeV (left) and the inclusive electron scattering cross section at $Q^2 = 0.6 \text{ GeV}^2$ (right). The numbers in brackets indicate the assignments to multiples within the group $SU(6)_w$.

CEBAF can also study well established resonances in greater detail than has been possible up until now. The Roper $P_{11}(1440)$ is an interesting case. While its coupling to real photons is reasonably well known, it has hardly been seen in electron scattering. What data exists suggest that this resonance has a strong coupling to longitudinal photons²¹, a result not obtainable from quark models.

New measurements of the effect of the Roper on electro-production of pions are needed to clarify these issues. It turns out²² that the production of π^0 's is particularly sensitive to the Roper if the proton target is polarized in a direction transverse to the virtual photon in the electron scattering plane. The expected asymmetry for the process $\vec{p}(e, e'p)\pi^0$ is shown in Figure 14. Note that $\phi_1 = \pi/2$; the proton must be detected out-of-plane. The VAS spectrometers were designed by V. Burkert for measurements such as these²³, and with them this program could be readily carried out at CEBAF with a state-of-the-art polarized target.

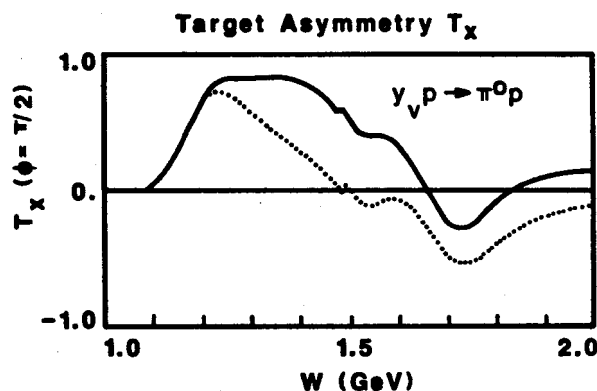


Figure 14 The target asymmetry $T_x(\phi = \pi/2)$ for single π^0 production from protons at $Q^2 = 1 \text{ GeV}^2$, and its sensitivity to the $P_{11}(1440)$ excitation strength. The prediction is based on a recent analysis of DESY and NINA data²¹. Solid line: with P_{11} ; dotted line: without P_{11} .

2.4 Hypernuclei - $(e, e'K)$

One of the novel programs proposed for CEBAF is the study of the excitation and decay of hypernuclei by the $(e, e'K)$ reaction. Hypernuclei are interesting because they contain one strange valence quark, which can be thought of as an impurity in the nuclear system. Study of hypernuclei should give insight into both hyperon-nucleon interactions and nuclear structure.

One system which illustrates how the study of hypernuclei can shed light on fundamental questions is ${}^4\text{He}$ and ${}^5\text{He}$ (ref. 24). The ${}^4\text{He}$ nucleus has 2 neutrons and 2 protons; or 6 valence up and 6 valence down quarks, which completely fill the lowest shell in either the language of nucleons or quarks. Adding a fifth nucleon (neutron) forces all three quarks to occupy the next highest shell, while adding a Λ would require that only the two non-strange quarks occupy the next highest shell, but would allow the strange quark into the lowest shell. In contrast, if the Λ were truly elementary, it could "fully" occupy the lowest shell. This suggests that the behavior of ${}^5_\Lambda\text{He}$ may be sensitive to the underlying quark structure of the Λ , which in this system seems to satisfy a partial exclusion principle. Studies of the level structure and properties of Λ in ${}^5\text{He}$ may give insight into the underlying quark dynamics.

The $(e, e'K)$ reaction has several advantages over the more traditional (K^-, π^-) reaction as a means of producing hypernuclei²⁵. The electromagnetic process excites both natural and unnatural parity states. Furthermore, since K^+ (or K^0) and γ both interact weakly with the nuclear system, these particles can penetrate (or escape from) deep inside of the nuclear medium. This means that the simple impulse mechanism, illustrated in Figure 15, is more likely to describe the reaction than in the (K^-, π^-) case; distorted wave calculations give corrections of a few 10's of a percent for the $(e, e'K)$ process as opposed to factors of 10 for (K^-, π^-) . It also means that it is possible with $(e, e'K)$ to convert a proton into Λ deep inside a large nucleus (such as Pb) opening up an entirely new area of study. A disadvantage of the $(e, e'K)$ process is the energy-momentum mis-match, which means that the momentum transferred to the nucleus cannot be made zero, as in the (K^-, π^-) reaction. However, at high energy the momentum transferred in the $(e, e'K)$ process can be reduced to a few hundred MeV, favoring transitions in which the angular momentum of the particle is changed by a few units, so that the electromagnetic process naturally favors conversion of a proton in an outer shell into a Λ in an S state.

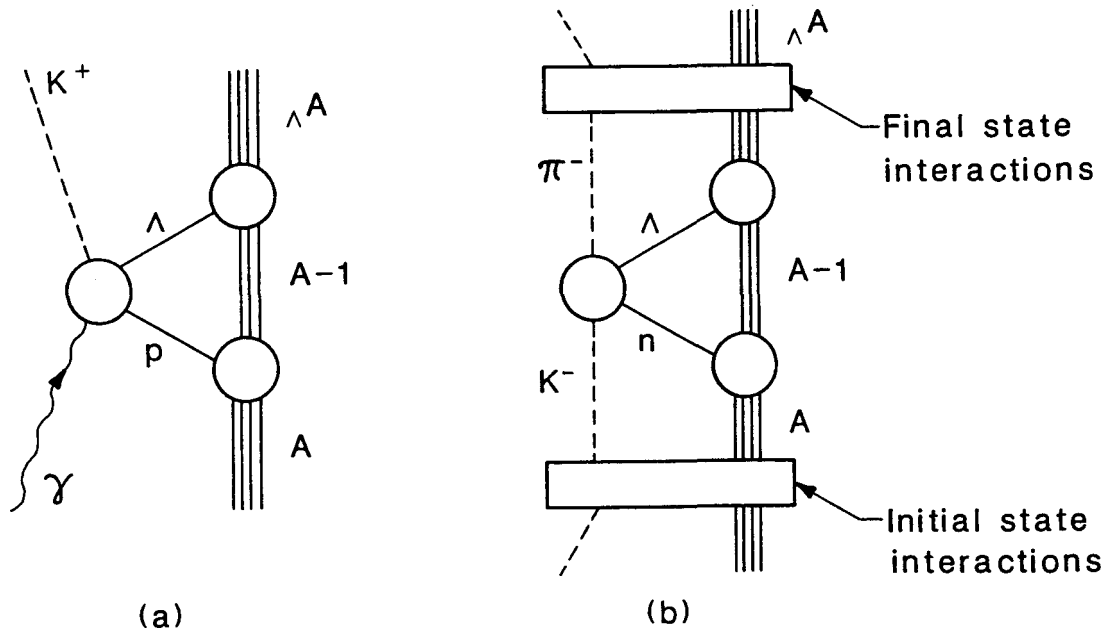


Figure 15 Diagrams showing single-step formation of hypernuclei through $(e, e'K)$ and (K^-, π^-) reactions. In the (K^-, π^-) case, initial and final state interactions are very strong and cannot be ignored.

Mecking²⁶ has found that the best way to achieve high resolution and reasonable counting rates in $(e, e'K)$ is to the work in the forward direction, where the electron scattering angle is less than 1° and the kaon angle is $\leq 5^\circ$. A possible experimental setup is shown in Figure 16. High resolution is achieve by using an incident electron energy of about 2 GeV, so the outgoing electron has low momentum and can be measured to high precision in a special split-pole spectrometer. The forward kaon can be measured in a high resolution spectrometer, similar (or identical perhaps) to those under consideration for Hall A. If high beam currents and a thin target are used, resolutions of the order of a few 100 keV are possible, and can be improved as the incident beam resolution is reduced below the design value of 10^{-4} . This method is superior to using tagged photons. The resolution expected is 10 times better than that obtainable with K^- beams.

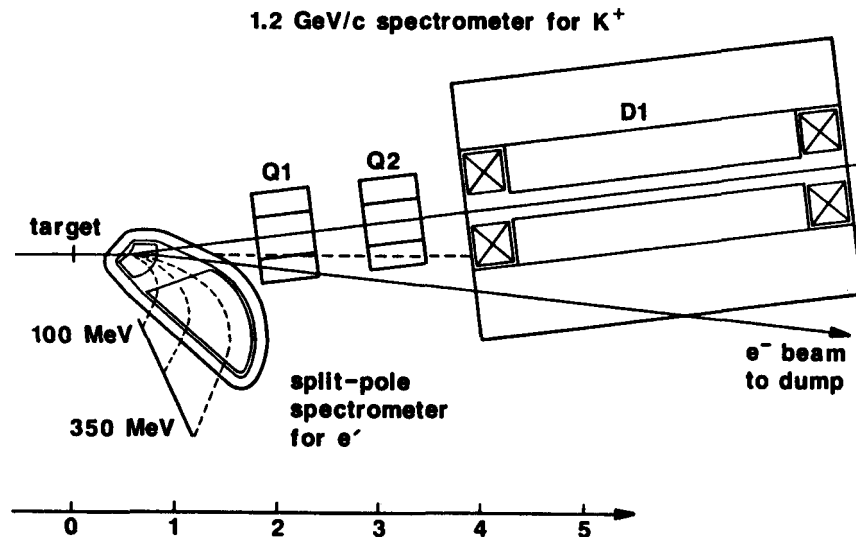


Figure 16 Possible experimental layout for the $(e, e' K^+)$ program discussed in the text.

2.5 Charge Distribution of the Neutron and Deuteron

The charge form factor of the neutron, G_{En} , is of fundamental importance both because it is sensitive to the distribution of the quarks in its interior, and because it must be taken into account when extracting nuclear structure information from all electron scattering data. Until recently, it has been believed that G_{En} is small, safely ignorable in all experimental analyses, but a recent attempt by Gari to fit all existing nucleon form factor data²⁷ has lead to speculation that it might be large at $Q^2 \gtrsim 1 \text{ (GeV/c)}^2$. If it is large, it may be possible to measure it by separating longitudinal and transverse response functions from $d(e, e'n)p$ scattering. In any case, use can be made of the CEBAF polarized electron beam to determine G_{En} in one of two ways:

- the transfer of the electron polarization to the recoiling neutron can be measured in a $d(\vec{e}, e' \vec{n})p$ experiment²⁸.
- the target asymmetry for a target polarized transverse to $\vec{q}$ in the electron scattering plane can be measured in a $\vec{d}(\vec{e}, e'n)p$ experiment²⁹. (Alternatively, G_{En} may also be measured using polarized ^3He targets³⁰.)

In either case, the experimentally measured quantity is proportional to

$$T = \frac{-2 \sqrt{\tau(1+\tau)} G_M G_E \tan \theta/2}{G_E^2 + G_M^2 + 2(1+\tau)\tau G_M^2 \tan^2 \theta/2} \quad (12)$$

where $\tau = Q^2/4M^2$. Hence both the magnitude and sign of G_E can be obtained directly.

Currently, G_{En} is completely unknown for $Q^2 \gtrsim 1 \text{ (GeV/c)}^2$, and the measurements at low Q^2 have large systematic and statistical errors, and are sensitive to the model used for the deuteron wave function, as shown in Figure 17. Precise measurements of G_{En} at both low and high Q^2 are very much needed, and can be done at CEBAF.

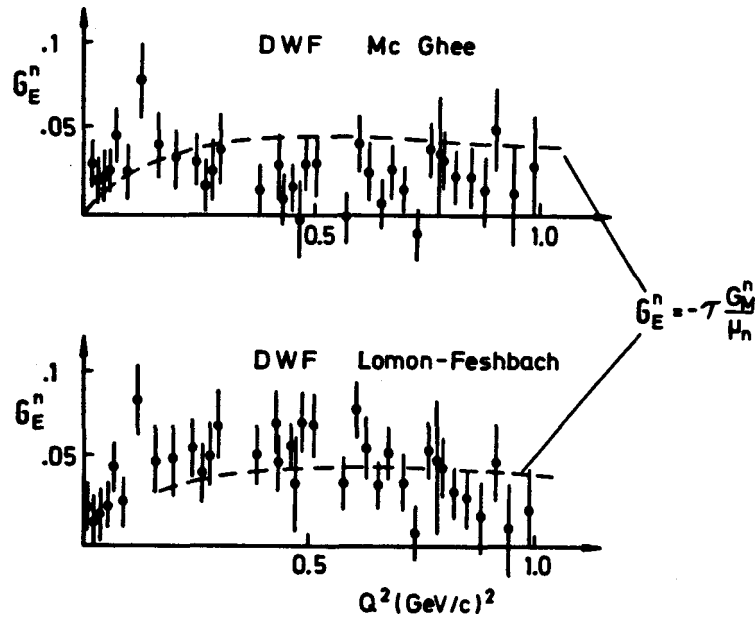


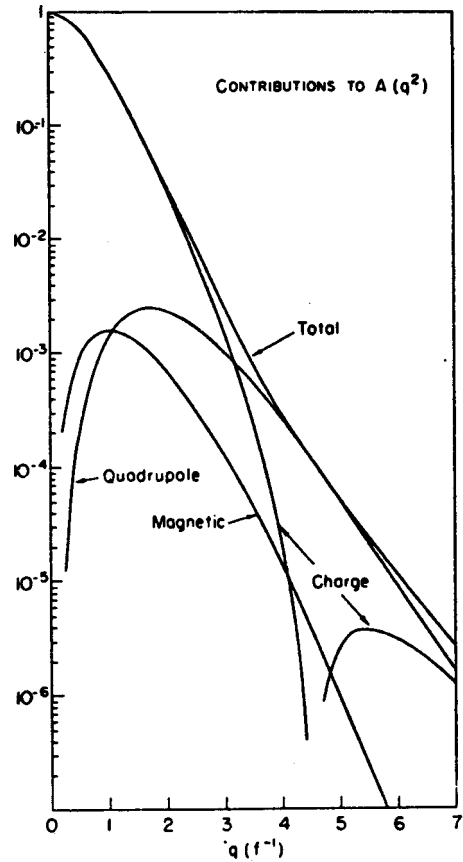
Figure 17 The electric form factor of the neutron derived from elastic electron deuteron scattering for two models of the deuteron wave functions. The results indicate a strong dependence on wave function.

The deuteron charge form factor, G_C , is also currently unknown. Beautiful high Q^2 measurements of the $A(Q^2)$ and $B(Q^2)$ structure functions exist, but from these it is only possible to measure

$$G_C^2 + \frac{8}{9} \eta^2 G_Q^2 = A(Q^2) - \frac{B(Q^2)}{2(1 + \eta)} \quad (13)$$

where G_Q is the quadrupole form factor and $\eta = Q^2/4M_d^2$. Unfortunately, much of the nuclear structure is washed out by this combination of form factors, as shown in Figure 18. A separate determination of G_C , particularly near its predicted zero, would be a sensitive indicator of short range structure, relativistic and meson exchange effects, and quark degrees of freedom.

Figure 18 Contributions to the A structure function of the deuteron from the charge, G_C^2 , quadrupole, $8/9\eta^2 G_Q^2$, and magnetic terms for a typical model. Note the zero in G_C^2 near $Q^2 \approx 0.85 \text{ (GeV/c)}^2$ is completely obscured by quadrupole terms.



As for the neutron, a vector polarization transfer can be measured with a polarized electron beam, or, alternatively,

asymmetries due to tensor polarization can be measured with unpolarized electron beams²⁸. From these asymmetries G_C and G_Q can be separated. Measurements of T_{20} have already been carried out at Bates and Bonn, but at momentum transfers well below the region where G_C^2 has its minimum³¹. A new experiment is planned for Bates, which should probe T_{20} near $Q^2 = 1 \text{ (GeV/c)}^2$, a very interesting region. At CEBAF, better measurements at lower and higher momentum transfers will be possible.

2.6 Studies in the Deep Inelastic Region

Inclusive measurements of electron and muon scattering, (e, e') and (μ, μ') , have provided some of the best information about quarks. When analyzed in conjunctions with neutrino and antineutrino charge changing reactions (ν, μ^-) and $(\bar{\nu}, \mu^+)$, they give an explanation of the data in terms of lepton scattering from individual quarks, yielding an almost model independent determination of the number of valence quarks, the distribution functions for valence and ocean quarks, the percentage of momentum carried by gluons, and the rms charge carried by the quarks³². And the EMC effect, which has generated such interest, gives information about quark distributions in nuclei, regardless of which explanation eventually prevails.

In view of the success of the quark description of inclusive scattering in the deep inelastic region, it is particularly interesting to see whether or not some of this success can be extended to the description of the behavior of exclusive channels such as $(e, e'\pi)$, $(e, e'2\pi)$, $(e, e'K)$ and other processes. How is the total cross section built up out of these individual channels, and how does this composition change as one moves deeper into the scaling region? Specific questions that could be addressed in an initial program of CEBAF experiments are the study of the A dependence of single hadron production, measurement of the ratio of positive to negative hadron production which may be sensitive to differences in densities of ocean $u\bar{u}$, $d\bar{d}$, and $s\bar{s}$ quarks, and studies of the Q^2 , x , high $p_\perp$ and ϕ dependencies of single hadron

production. A list of specific experiments of interest was developed at this Workshop by the working group on this topic and is summarized in Karl Van Bibber's report⁸⁸.

CEBAF's 4 GeV design energy was chosen to permit a complete study of the proton (and hence nucleus) inelastic structure functions from the region where resonances and other effects of collective quark structure are visible, up to the onset of deep inelastic scattering, where the structure functions become a function only of the Bjorken x-scaling variable, Eq. (5), and inclusive scattering is described by incoherent scattering from point-like quarks. Figure 19 shows the kinematic region which can be explored by a 4 GeV and a 6 GeV accelerator. The 4 GeV accelerator can cover the transition region, as was intended, but cannot carry on a very extensive program in the deep inelastic region. Extending the energy to 6 GeV, which appears to be a likely possibility if the superconducting cavities continue to exceed specifications (as described in the overview report with this proceedings), will considerably expand the possible program of deep inelastic studies, and further increases in energy will help still more.

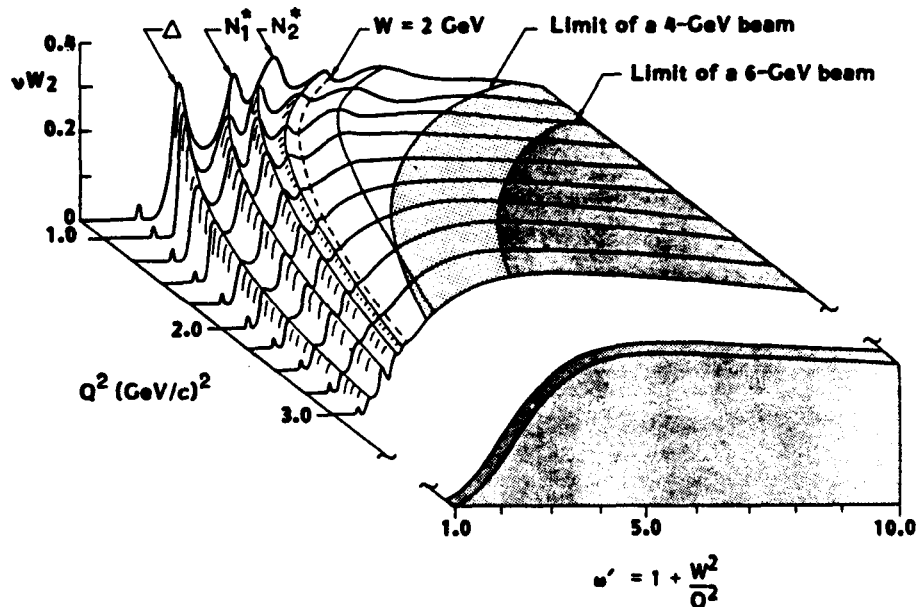


Figure 19 The proton inelastic structure function νW_2 shown as a function of Q^2 and the scaling variable w' . The regions which can be studied by a 4-GeV accelerator and a 6-GeV accelerator lie to the left of the limits shown.

2.7 Studies of the Electroweak Interaction

The weak neutral current can be studied by measuring the parity violating asymmetries

$$A = \frac{\sigma_+ - \sigma_-}{\sigma_+ + \sigma_-} \quad (14)$$

for the scattering of electrons with positive and negative helicity. This asymmetry measures the interference between the electromagnetic interaction and the parity violating parts of the weak neutral interaction, which arises from the exchange of a Z_0 boson. Assuming the electromagnetic interaction is known, precision measurements of A can be converted into precision measurements of the coupling constants of the Z boson to the constituent quarks in various nuclear targets. By selecting the quantum numbers of the initial and final hadronic states, and choosing different electron scattering angles, different combinations of these coupling constants can be isolated and measured. While the standard model gives all of these coupling constants in terms of only one parameter, $\sin^2\theta_W$, which has been measured to about 5%, it is important to test this model to the highest possible precision. Tests of semi-leptonic processes inside the nuclear medium, where the strong interactions are strong and confining, are particularly valuable.

Elastic scattering from the proton and Δ^+ production both have large cross sections, and have been proposed as a test of the electroweak theory which could be carried out at CEBAF³⁴. The general interaction of the Z_0 with u and d quarks can be expressed in terms of four constants, α , β , γ and δ . Only γ will be measured in the Bates asymmetry experiment on carbon, and the SLAC experiment was dominated by α ³⁵. Elastic measurements from the proton will give information about all 4 constants, but are more sensitive to γ and β , depending on the scattering angle. The reaction $e p \rightarrow e' \Delta^+$ depends on α and β only, and this measurement would therefore help to separate α , β and γ .

CEBAF's 100% duty factor will be helpful to this program by permitting instantaneous rates to be lowered to a level where individual pulses can be counted³⁴. The asymmetries to be measured are of the order of 10^{-5} , and hence systematic errors must be lowered to the 10^{-7} level to permit the 1% measurements required for precision tests. If the systematic errors can be reduced sufficiently, and special targets capable of taking high beam currents are available, it is estimated³⁴ that data for the needed 1% statistical accuracy can be accumulated in 100's of hours.

3. Experimental Requirements for the Planned Program

CEBAF will be equipped with three experimental halls, designated A, B, and C. While the mission of each of these halls is being currently worked out, and will probably undergo further evolutionary changes in the months ahead, the major function of each hall has already emerged, at least in broad outline, and will be summarized briefly. For a detailed discussion of each of the Halls, with a description of the equipment planned or under consideration, see the following talk by J. Mougey.

- Hall A will be devoted to studies where high momentum and angular resolution are needed. This will include single nucleon knockout $A(e, e'N)A-1$, and studies of hypernuclei $A(e, e'K)_{\Lambda}A$, where a resolution of ~ 200 keV is needed to separate closely spaced nucleon levels. Studies of the $(e, e'N)$ reaction from few body systems, which require resolutions of about 1 MeV, will also have to be done in this hall. It is desirable that the spectrometers in Hall A be able to move over the widest possible angles, and be capable of making out-of-plane measurements so that the individual structure functions can be fully separated. It should also be possible to use polarized targets, and detect neutrons in Hall A.
- Hall C will be devoted to the study of correlated, few particle final states in cases where it is desirable to trade

high resolution for large angular and momentum acceptance. It seems likely that resolutions of a few to 10 MeV are the best that will be achievable in this hall, but angular acceptances of up to 50 msr and momentum bites of a factor of 2-3 will be gained as a consequence, at least for the proposed VAS hadronic spectrometer. This hall is the likely home of $(e, e'2N)$ measurements, and will also be ideal for the study of baryon resonances in cases where specific signatures are sought after, as in the case of the Roper described in Section 2.3. As in Hall A, motion of spectrometers through the widest possible angular range, out-of-plane capability, ability to detect neutrons, and ability to use polarized targets are all needed. In addition, this may be the hall where a program of triple coincidence measurements is most likely to begin, and measurements of G_{En} and G_C/G_q carried out.

- Hall B will be devoted to the study of uncorrelated, many particle final states, where large nearly 4π acceptance is essential and 1% resolution is acceptable, and to tagged photon studies where large acceptance is needed to compensate for a low photon flux. It also will be an ideal place to initiate new programs (such as two-nucleon knockout) where a broad exploration of the phase space is needed before focusing in detail on specific processes or choices of kinematics. It will probably be the major home for studies of excited baryons and exclusive processes in the deep inelastic, x-scaling region. To achieve the needed 4π acceptance, a large acceptance spectrometer (LAS) is being developed. Finally, to allow for the possibility that parity studies may be carried out in Hall B, it will be designed, together with Halls A and C, to handle the full beam intensity.

To design and build equipment to carry out this exciting program, to keep within the budget allocated to experimental equipment, and to have the equipment ready by 1992 will be a major challenge. CEBAF looks toward its user Community to help and to contribute to this exciting endeavor.

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Overview of the Experimental Equipment for CEBAF

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1. Introduction

In nuclear physics, electromagnetic probes are known to provide precise and quantitative information, owing to the well-known and relatively weak electromagnetic interaction. Thus, the physics program at CEBAF will rely heavily upon the quality of a set of high-precision instruments, mainly magnetic spectrometers and detectors.

The planned experiments, implying mostly coincidence measurements, call for sets of two or three large angular and momentum acceptance spectrometers around the same pivot, or nearly 4π detectors with magnetic analysis. Other specific requirements come from the use of long liquid or gas targets, polarized beam and/or targets, neutral particle detection and measurements involving multiparticle detection in non-coplanar kinematics.

The CEBAF project include three experimental halls, designated A, B, and C, being fed simultaneously by continuous beams of different but correlated energies and independently controlled intensities. Each hall will be designed to match specific physics requirements. Hall A, designed for high resolution experiments (~ 1 to 0.1 MeV) will house two, possibly three, high resolution (10^{-4} or better) spectrometers, with reasonably large acceptancies. A nearly 4π magnetic spectrometer is planned in Hall B for photonuclear and low luminosity electronuclear

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reactions. Specific set-ups, like the one for parity violation experiments, can also be installed in Hall B. Hall C will be devoted to double and triple arm experiments with moderate (~ 10 MeV) resolution.

Preliminary designs for the spectrometers have been performed during the 1985 Summer Study Group Meeting and are described in the report on Research Program at CEBAF (RPAC)^[1] and in the CEBAF Design Report^[2]. Since then, improvements on the optics designs, studies of alternative solutions as well as definition of basic technological options have led to the present designs, still being improved. Modifications made during the 1986 Summer Study Meeting are included in the versions which will be briefly described hereafter. A summary of the main parameters for the CEBAF spectrometers are given in Table 1.

Contributors to the designs include P. Brindza, V. Burkert, J. LeRose, B. Mecking, J. Mougey, R. Whitney (CEBAF), I. Blomqvist (Lund), R. Neuhausen (Mainz), H. deVries (NIKHEF, Amsterdam), G. Fife, M. Finn, C. Perdrisat and V. Punjabi (William & Mary), and A. Saha (UVa). The designs of the magnet elements have been worked out during the April 1986 Magnet Workshop^[3].

2. End Station A Spectrometers

An important part of the CEBAF physics program deals with completely exclusive experiments in which the nuclear final state has to be fully specified. Examples of such experiments are:

- $(e, e'p)$ reactions. Single nucleon densities. Elementary ep interaction in nuclei.
- $(e, e'd)$, $(e, e'^3\text{He})$, ... to specific final states.
- $(e, e'\pi^{+/-})$ coherent production in nuclei.

Table 1
Preliminary Designs Parameters for CEBAF Spectrometers

	Type	Maximum Momentum (GeV/c)	Momentum Acceptance (%)	Momentum Resolution	Solid Angle (msr)	Angular Range
Hall A Electron	QQDDQQ horizontal	4. (6)	10	$< 10^{-4}$	11.	$10^\circ + 130^\circ$
Hall A Hadron Solution I	QQDD vertical	3.	15	$< 10^{-4}$	11.	$10^\circ + 160^\circ$
Hall A Hadron Solution II	QQDMD vertical	1.2	15	$\sim 3 \cdot 10^{-5}$	35.	$20^\circ + 150^\circ$
Hall C electron	QQDQ	5.5	30	$\sim 3 \cdot 10^{-4}$	6.6	$10^\circ + 160^\circ$
Hall C Hadron I	Dipole (VAS I)	3.5	> 50	$\sim 3 \cdot 10^{-3}$	50.	$16^\circ + 150^\circ$
Hall C Hadron II	Dipole (VAS II)	1.7	> 50	$\sim 3 \cdot 10^{-3}$	70.	$25^\circ + 150^\circ$
Hall B	Large Acceptance Spectrometer		Toroidal field, max. strength 8 superconducting coils			1T
			Solid Angle			80% of 4π
			Momentum Resolution			$\sim 10^{-2}$

- $(e, e' K^+)$ hypernuclear physics.

Typical values for the required missing mass resolution range from ~ 1 MeV in light systems ($d \leftrightarrow np$ separation) to ~ 100 keV in heavy nuclei, or for hypernuclear spectroscopy. Even if it is not required for isolating nuclear levels, high resolution may be very useful in exclusive experiments to improve signal/noise (true/accidental) ratio, like when measuring very low cross sections, or for separation of small (interference) structure functions. This could apply to experiments like $d(e, e' p)n$, ${}^4\text{He}(e, e' p){}^3\text{H}$, $p(e, e' p)\pi^0, \dots$

From kinematics and cross section evaluations of some typical experiments, mainly $(e, e' p)$, the following major requirements for end station A spectrometers have been obtained.

Electron spectrometer: Should be designed for a maximum momentum of 4 GeV/c, with possibility of being upgraded up to 6 GeV/c. Should have large momentum ($\Delta p/p \sim 10\%$) and angular ($\Delta\Omega \sim 10\text{msr}$) acceptancies. Momentum resolution below 4 GeV/c should be better than 10^{-4} , with possibility to reach a few 10^{-5} within reduced acceptance. Should be able to reach small ($\lesssim 10^\circ$) forward angles.

Hadron spectrometer: Maximum momentum of 3 GeV/c. Momentum acceptance larger than 10%, solid angle larger than 10 msr. Momentum resolution NOV 11, 1986 $\delta p/p \sim 10^{-4}$ or better. Small ($\lesssim 10^\circ$) forward angle capabilities. Short optical path length ($\lesssim 15$ m) to allow K or π detection with acceptable efficiency.

The two spectrometers should allow to operate at high luminosity values ($\sim 10^{38} \text{cm}^{-2} \text{sec}^{-1}$). Use of extended targets ($\gtrsim 5$ cm along beam direction) or polarized targets should be made possible. Out-of-plane measurement capabilities are also needed. Provision should be made for future implement of a third spectrometer, with $P_{\text{max}} \lesssim 1$ GeV/c, larger momentum and angular acceptancies, and

shorter optical length. Such spectrometer would have to be optimized for K detection, but could also be used for backward electron detection, and in triple coincidence measurements. Neutron detection capabilities should also be implemented.

As discussed in the RPAC report, focusing, essentially "hardware-type" spectrometers have been chosen as the only way to achieve the resolution goals. However, as the complete kinematics have to be determined, with angular resolutions of the order of 1msr to achieve the desired missing mass resolution, software corrections are possible, and are envisaged for obtaining the ultimate resolution. The designs rely on combinations of a few modular elements which should simplify magnet studies, give more flexibility, and optimize costs. Iron-dominated dipoles with rectangular pole faces, uniform fields ($B_{max} \lesssim 1.7T$) and superconducting coils have been chosen as dispersive elements. The focusing ones will be $\cos 2\theta$ -type superconducting quadrupoles. As they allow larger field gradients than iron-dominated ones, the large acceptancies can be achieved without increasing the overall length. These choices have been discussed in detail during the CEBAF Spectrometer Magnet Workshop, April 7-11, 1986^[3]. The electron spectrometer with particle trajectories in the horizontal plane, as i) a larger acceptance in the vertical plane is preferable due to the rapid variation of the cross section with scattering angle and ii) a larger acceptance in the transverse plane fits better with an iron-dominated dipole gap through a quadrupole doublet (Karl Brown's 3:1 rule). The hadron spectrometer has to be vertical in order to determine the target reaction point when operating with extended targets.

The spectrometer arrangement in End Station A is shown in Figure 1 and a layout of the electron spectrometer is given in Figure 2a. It is a QQDDQQ essentially symmetric design. The main characteristics are listed in Table 2. Point-to-point imaging in both planes, to second order, is realized by a sextupole field component built

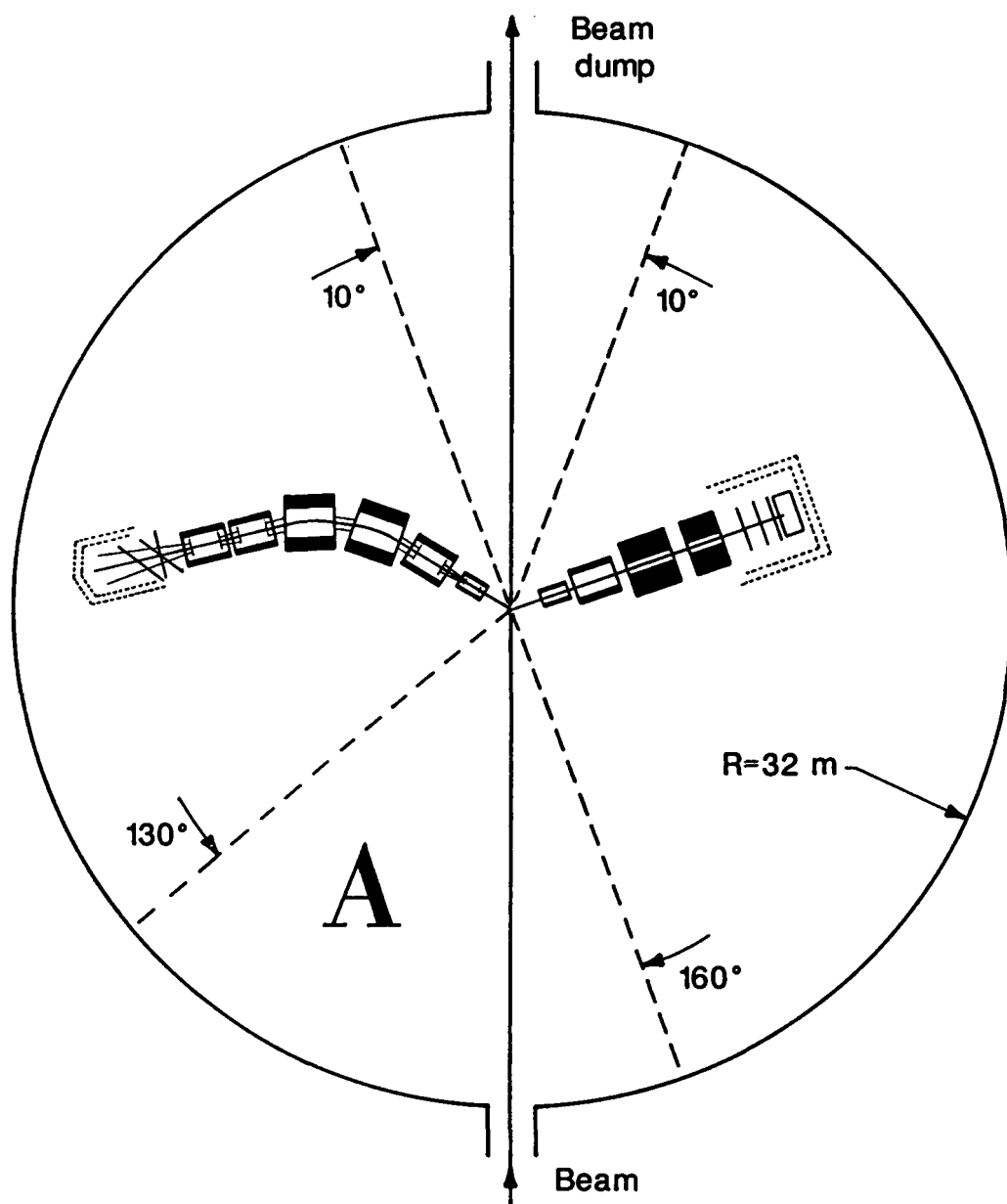


Figure 1 Spectrometer arrangement in End Station A.

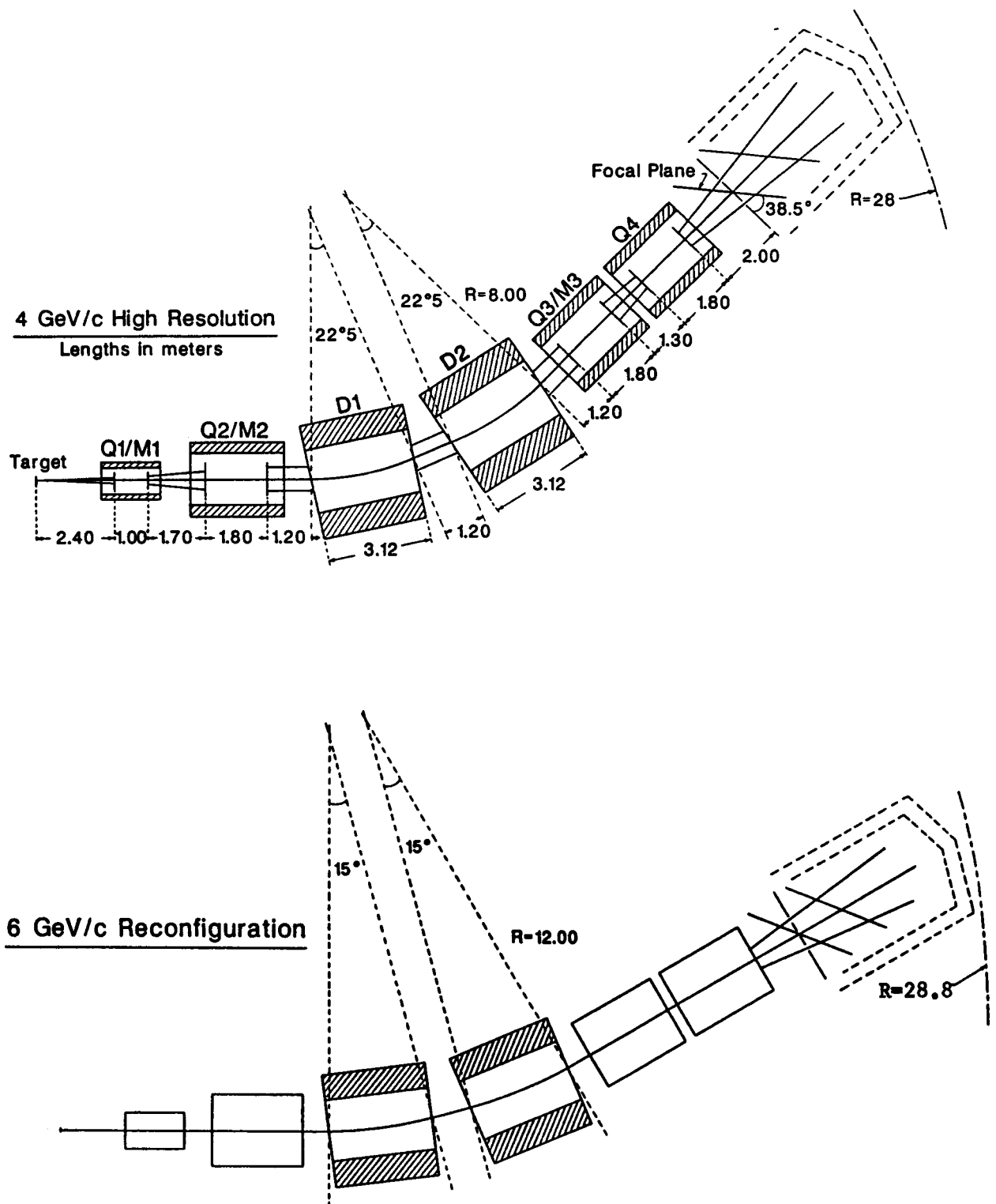


Figure 2 Sketch of the High Resolution Electron Spectrometer for End Station A. a) 4 GeV/c configuration b) 6 GeV/c configuration.

Table 2
Main Parameters of End Station A Spectrometer Designs

		Electron Spectrometer	Hadron Spectrometer (Solution I)	Hadron Spectrometer (Solution II)
Type		QQDDQQ	QQDD	QQDMD
Bending		horizontal, 45° (30°)*	vertical, 60°	vertical, 120°
Maximum momentum (GeV/c)		4, upgradable to 6	3	1.2
Momentum acceptance (%)		±5	±7.5	±10
Momentum resolution		~ 5•10 ⁻⁵	~ 10 ⁻⁴	~ 5•10 ⁻⁵
Maximum solid angle (msr)		10.8	10.8	20 to 35**
Angular acceptance (mr)	horiz.	±30	±90	±95
	vert.	±90	±30	±65
Momentum dispersion (cm/%)		9.9 (6.6)*		7.0
Linear magnification	horiz.	-1.06	-6.40	
	vert.	0.99	-1.15	-0.34
Angular magnification	horiz.	-0.94	-0.16	
	vert.	1.01	-0.87	-2.97
Focal plane angle	horiz.	51.5°		45°
	vert.	51.5°		50°
Maximum field in dipoles (T)		1.7 (1.7)*	1.7	1.8
Maximum gradient in quads (T/m)		8.9 (11.9)*	6.2	2.7
Maximum field in quads (T)		2.4 (3.5)*	1.6	0.8
Maximum sextupole strength (T/m ²)		0.12 (0.18)*		0.5

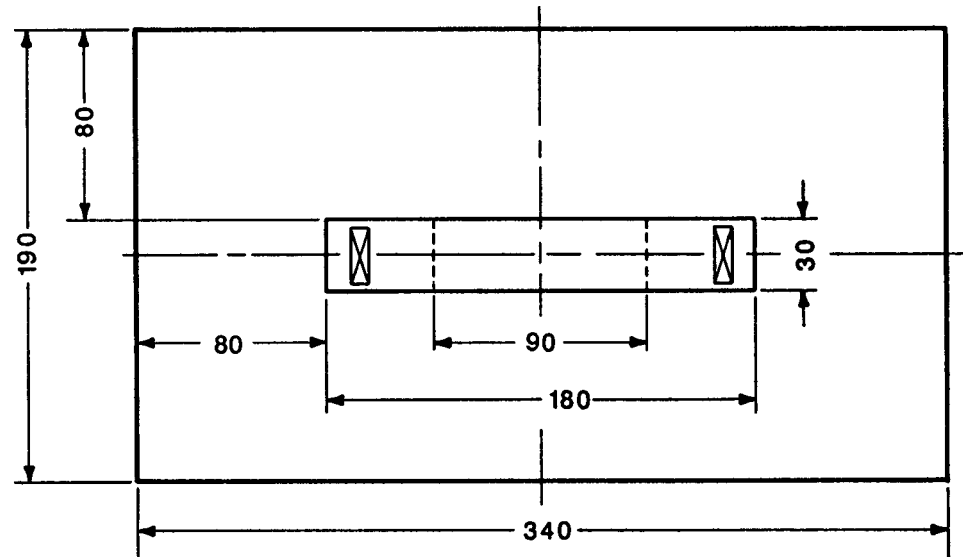
* Numbers in parenthesis are for 6 GeV/c maximum momentum.

** 20 msr with cos2θ quadrupoles, 35 msr with iron-dominated ones.

inside the element Q3/M3. In the middle of the drift space between the two 3.12m dipoles, the first order conditions $\langle \theta/\theta \rangle = \langle y/\phi \rangle = 0$ are realized. The cross-over in the transverse plane ensures relatively small dipole gaps (30 cm) notwithstanding the large acceptancies, while the other condition makes optimum use of the pole face areas to build up resolving power. With a total bending angle of 45° , the first order resolving power $R = |\langle x|\delta \rangle / \langle x|x \rangle| / \bar{x}_0$ is 47000 for a beam spot size of $\bar{x}_0 = 0.2$ mm. Without changing distances between the various elements, the spectrometer could be reconfigured (Fig. 2b) to a 30° bending angle device, allowing to increase the momentum up to 6 GeV/c. With the same total optical length, the overall length from the pivot is changed by only 0.8 m. In such configuration, the resolving power decreases to $R = 31000$.

The dipole element is sketched in Figure 3. The use of superconducting coils has been shown to be costs effective^[3]. Although it leads to slightly larger dimensions and weights, the racetrack coil H-type solution may be preferable to the saddle coil window-frame one, when superconducting coils are considered. As said before, the quadrupole elements (Figure 4) are current-dominated, $\cos 2\theta$ -type ones, using superconducting coils and warm iron yokes. Appropriate coil geometry (3 sectors or 5 layers) will allow to eliminate the dominant higher order multipoles, then to make the most effective use of the quadrupole aperture. A field gradient homogeneity of $\sim 10^{-3}$ is required within $\sim 2/3$ of the inner coil diameter. To keep the cylindrical geometry and ensure a minimum forward angle of 10° for the spectrometer, the front quadrupole Q_1 has a smaller size ($\phi = 0.5$ m inner coil diameter instead of $\phi = 1$ m) and a cold bore at liquid nitrogen temperature. In addition, a notch is cut within the iron yoke and the front part of the cryostat to leave room for the beam pipe. The approximate total weight of Q_1 is 7 tons, and 27 tons for the other quadrupoles. Additional layers of conductors powered independently can be used to produce sextupole or higher order

D1-D2,
Window-frame
solution



D1-D2,
H-type
solution

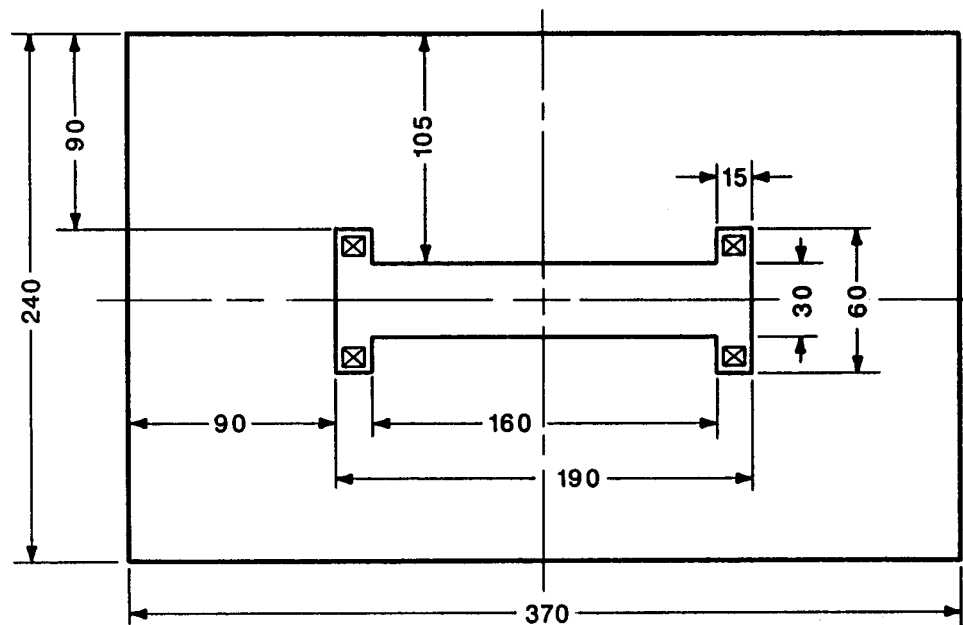


Figure 3 Sketch of the dipole element module. The length of the magnet is approximately 3 m.

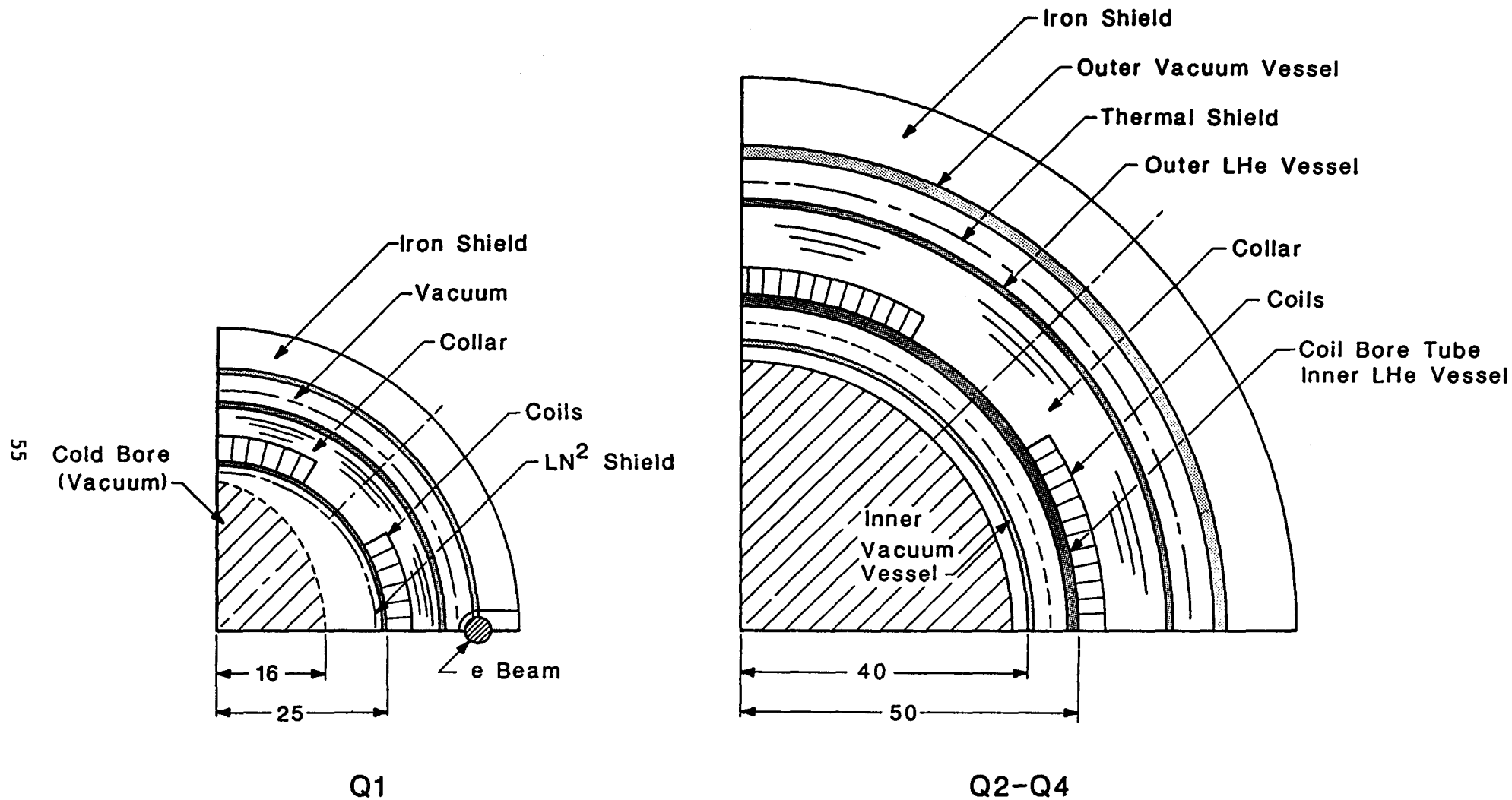


Figure 4 Schematic sections through the $\cos 2\theta$ superconducting quadrupole elements.

multipole fields which could be tuned either to correct second or higher order aberrations in the spectrometer, or to compensate for unwanted high order multipole contributions, coming in particular from fringe fields which have important effects in such open geometries. Although the $\cos 2\theta$ quadrupole configuration leads to higher needs of ampere-turns, higher value of stored energy, and more space requirements along the axis, it has been preferred to the Panofsky-type geometry sketched in the RPAC report (see the report of the quadrupole working group in Ref. 3). Two- and three-dimensional field configuration calculations are underway to determine optimal coil and iron geometry, end effects and force distribution on iron and conductors.

A layout of the 3 GeV/c hadron spectrometer is shown in Figure 5a. It consists in a QQDD vertical bending device, 18 m long, making use of the same basic modules as the electron spectrometer. The first order optics (point-to-point focusing in both planes) is sketched in Figure 5b, showing the beam envelope for a ± 2 cm target extension in the transverse direction. Although momentum acceptance is larger ($\Delta p/p = \pm 7.5\%$), the solid angle is the same as for the electron spectrometer to ensure a minimal forward angle of 10° . The possibility to achieve larger solid angles by sliding the entrance doublet towards the target is being studied.

An alternate solution for the hadron spectrometer, combining high resolution and large acceptancies at reduced maximum momentum capability of 1.2 GeV/c has been worked out by R. Neuhausen (Ref. 1, p. 11.31). An updated version of the QQDD design from Ref. 1 is shown in Figure 6. The entrance doublet can either make use of similar $\cos 2\theta$ quadrupoles as for the other spectrometers, or of iron-dominated quadrupoles, as the maximum field at the outer radius is 0.8T. The second solution leads to a larger solid angle (see Table 2) but an open geometry is required for the front quadrupole to allow comparable forward angles of $\sim 20^\circ$. To reach smaller angles (14°), the spectrometer may be pulled back (0.6 m) and used with a reduced

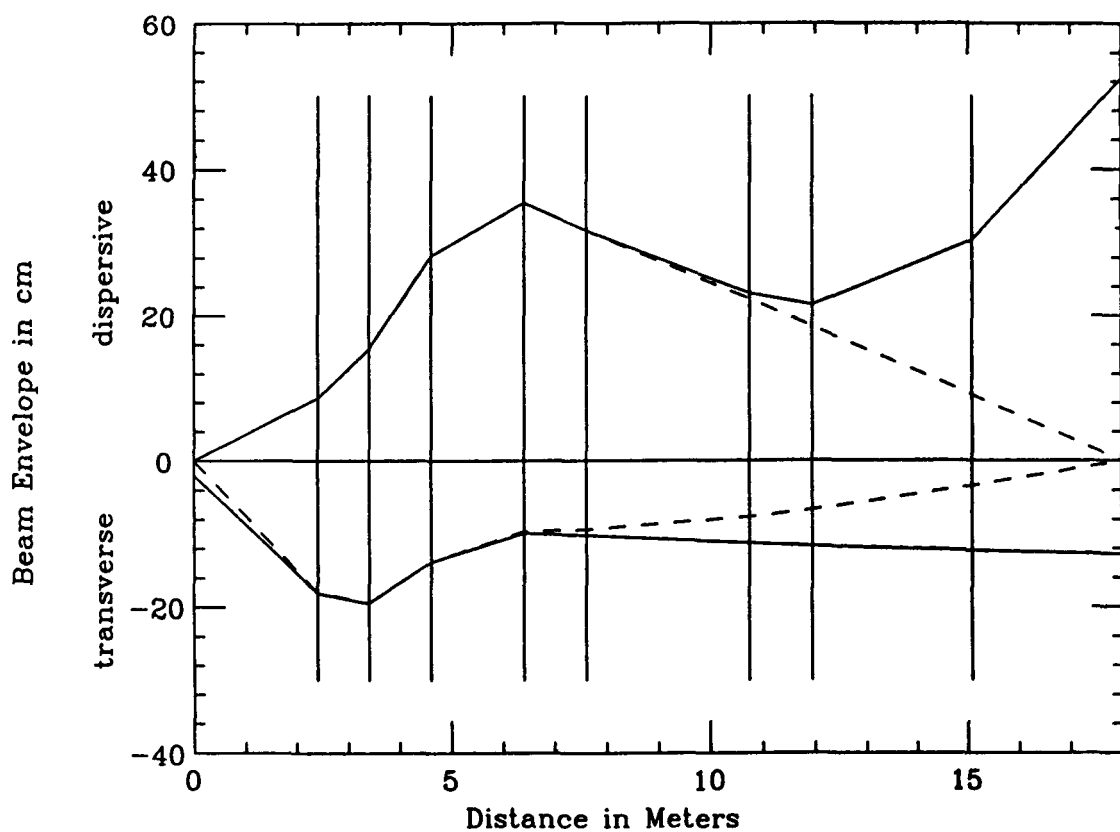
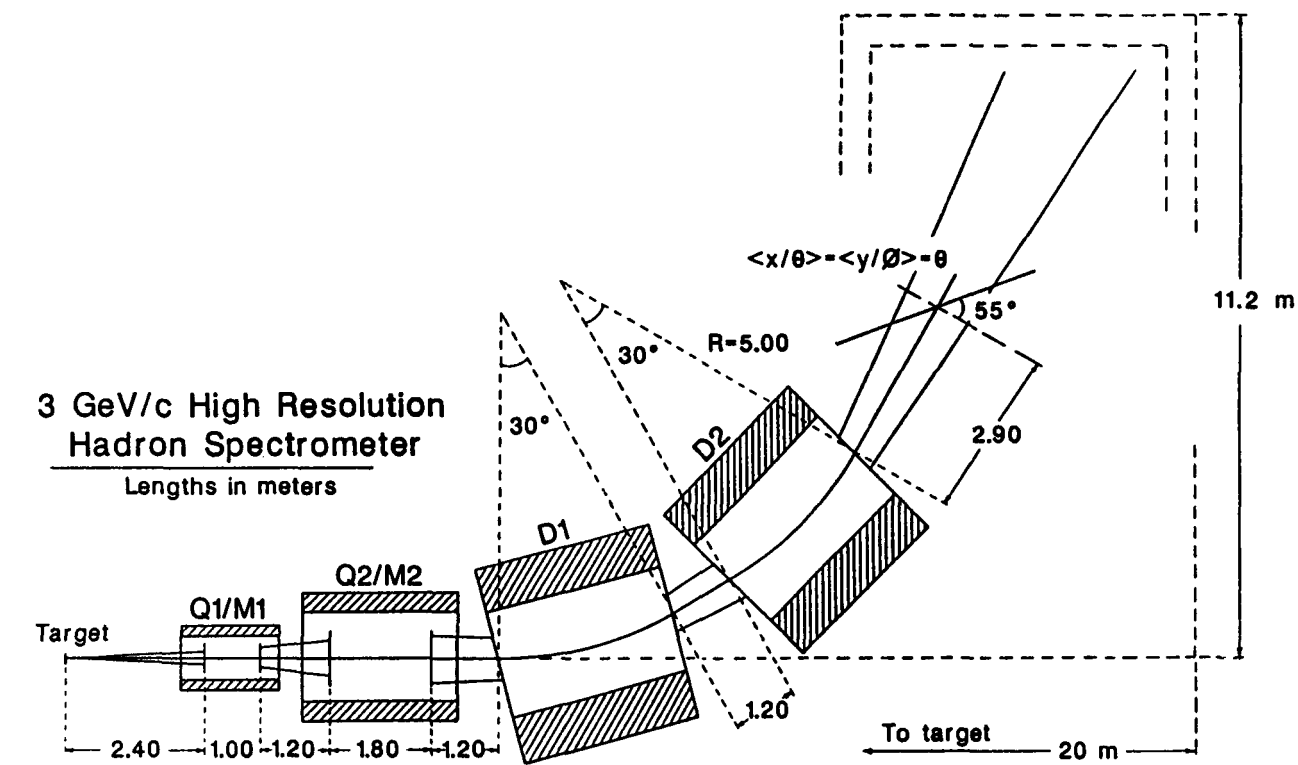


Figure 5 a) Layout of the spectrometer b) First order beam envelope for a point (dashed-line) and extended (± 2 cm, solid line) target.

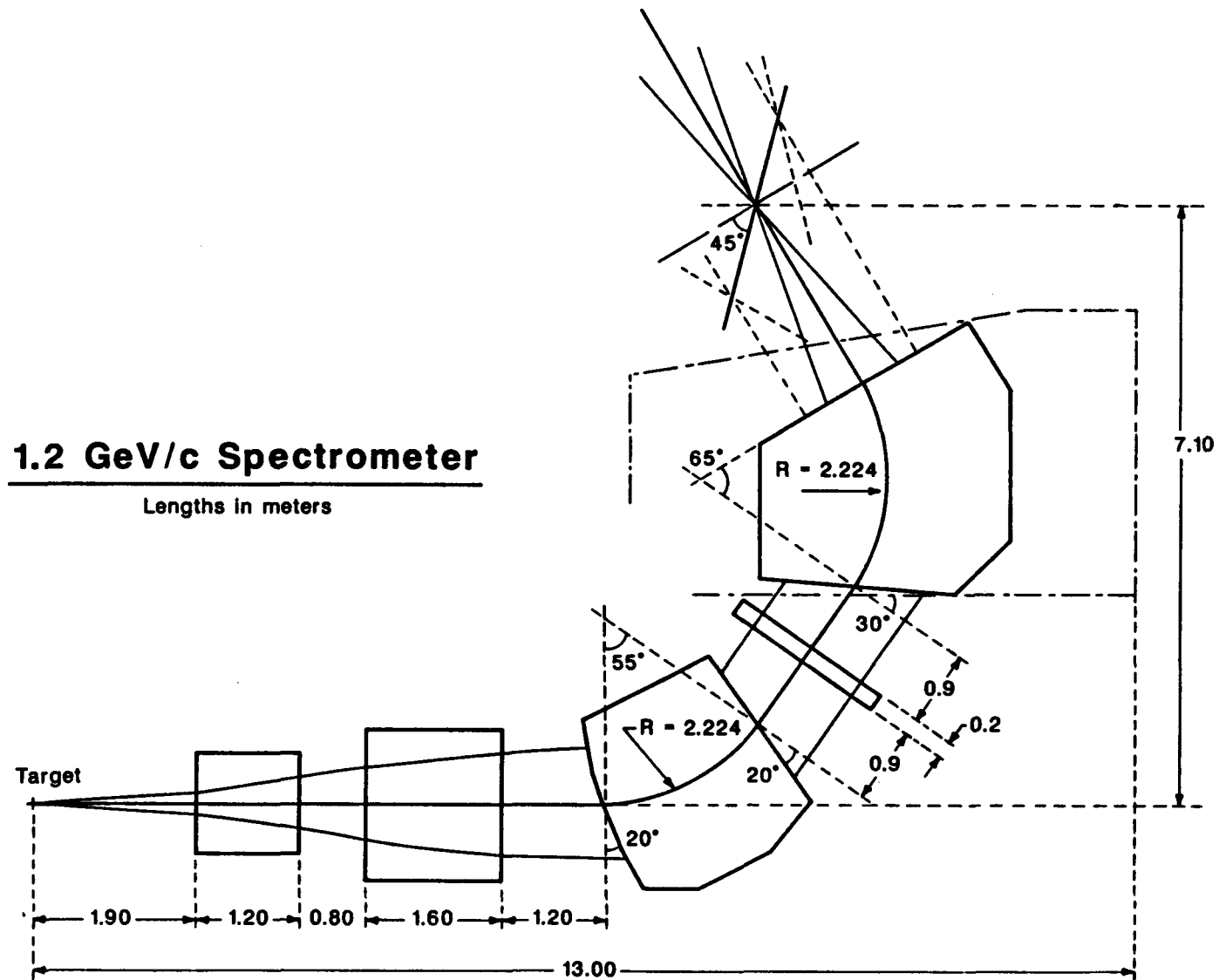


Figure 6 Sketch of the 1.2 GeV/c High Resolution Spectrometer.

solid angle (25msr). A multipole located between the two dipoles is used to restore the first order optics and focal plane angle of 45° .

Solutions making use of high field ($\sim 3\text{T}$), $\cos\theta$ combined function dipoles have also been investigated. Although they lead to interesting short designs with large acceptancies, their feasibility is more problematic, considering the combination of large bore dimensions and high field homogeneity requirements.

As mentioned earlier, the spectrometer arrangement should enable to perform measurements in which the hadron is detected out of the scattering plane. In that case, the relevant angles (see Figure 7) are the angle α between the scattering plane and the photonuclear reaction plane, for which a few discrete values, preferably including 90° , are needed, and the angle ϕ between the virtual photon and the emitted hadron, which should be selected independently at any value. The solution of moving one spectrometer out of the plane is clearly unpractical (weight ~ 400 tons, length ~ 20 m). A very wide vertical aperture for one spectrometer leads to complicated optics, and extreme difficulties for realizing the front quadrupole. The only practical solution appears to be moving the primary beam in the vertical plane (angle ϕ) direction. Assuming $\alpha = 90^\circ$ the relation between y and the kinematical angles θ and ϕ is

$$\tan\phi = \frac{e \sin \phi (\tan^2\phi + \sin^2\theta)^{1/2}}{q \sin \theta}$$

Provided the measurements are done at forward scattering angles ($\theta \lesssim 30^\circ$), values of ϕ below 20° seem sufficient to span a substantial range in ϕ in current (e,e'p)-like kinematics. Note however that ϕ has to vary continuously. The method (see Fig. 8) implies to move up and down two bending magnets. Disadvantages are dispersion produced on target, interference problems with spectrometers at forward angles, and the need for sufficiently long magnets to generate the beam

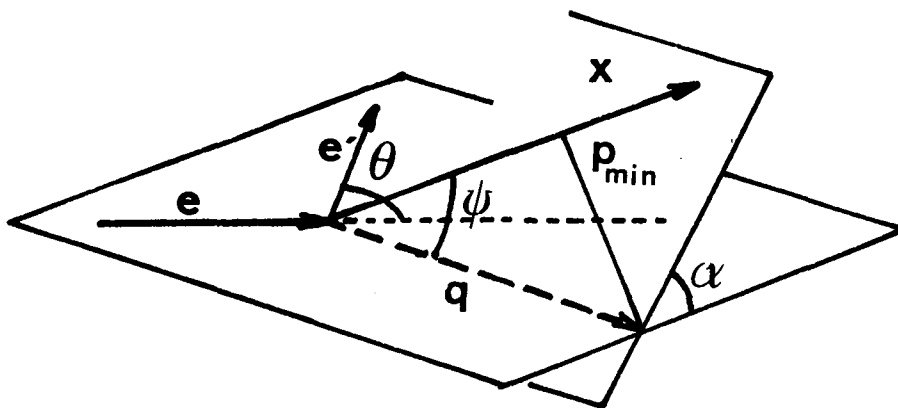


Figure 7 Kinematics for $(e, e'X)$ reactions

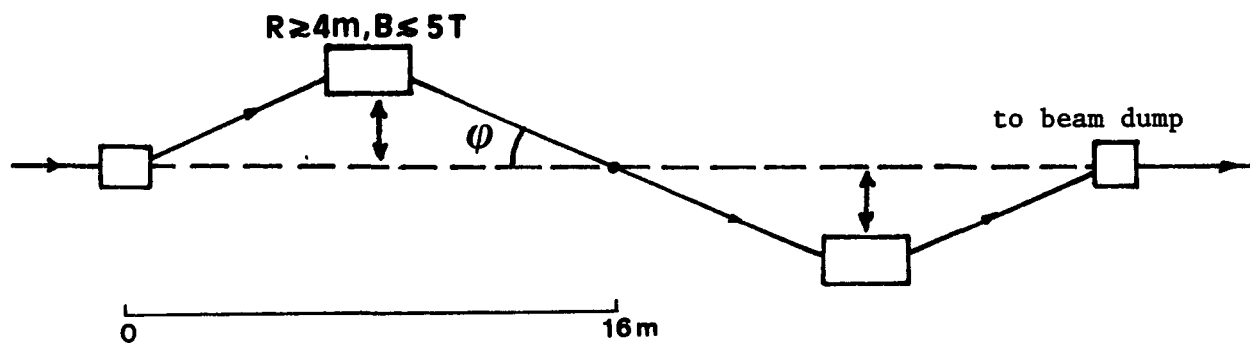


Figure 8 Steering the beam for out-of-plane measurements

deflection. Indeed, bending radii cannot be too small ($\gtrsim 2$ m) due to fluctuations in synchrotron radiation losses which otherwise would affect the energy resolution at high energy.

3. End Station C Spectrometers

The CEBAF physics program includes a large set of experiments for which high resolution is not needed, but high luminosities are required, which leads to a set of moderate resolution ($\sim 10^{-3}$) spectrometers. Examples of such experiments are:

(e,e'N) reactions in the continuum

(e,e'2N) reactions

- resonance transitions from free and bound nucleons through the reactions (e,e'N), (e,e' π), (e,e'N π)...
- measurements of nucleon (neutron) and deuteron form factors.

Some of those programs imply polarized targets, then very low luminosities, so that very large acceptancies are crucial. This is also true for triple coincidence measurements. Moreover, out-of-plane measurements are also needed to separate the various contributions to the coincidence cross sections.

The spectrometer arrangement currently planned for end station C is shown in Figure 9. The electron spectrometer (Figure 10) is a QQDQ design using basically the same elements as the two high resolution spectrometers. The maximum momentum of 4 GeV/c at 1.7T in the dipole can be raised to 5.5 by increasing the field to 2.1T. With a large gap of 0.45 m, and a 6.4 msr solid angle, this spectrometer can have 30% momentum acceptance and handle long (± 5 cm) targets. Its momentum resolution is $\sim 3 \cdot 10^{-4}$, and the parallel-to-point focusing condition in the transverse (horizontal) plane allow a very accurate determination of the scattering angle. Having similar entrance optics as the high resolution electron spectrometer, it can also reach small ($\sim 10^\circ$) forward angles.

Spectrometer Arrangement

Endstation C

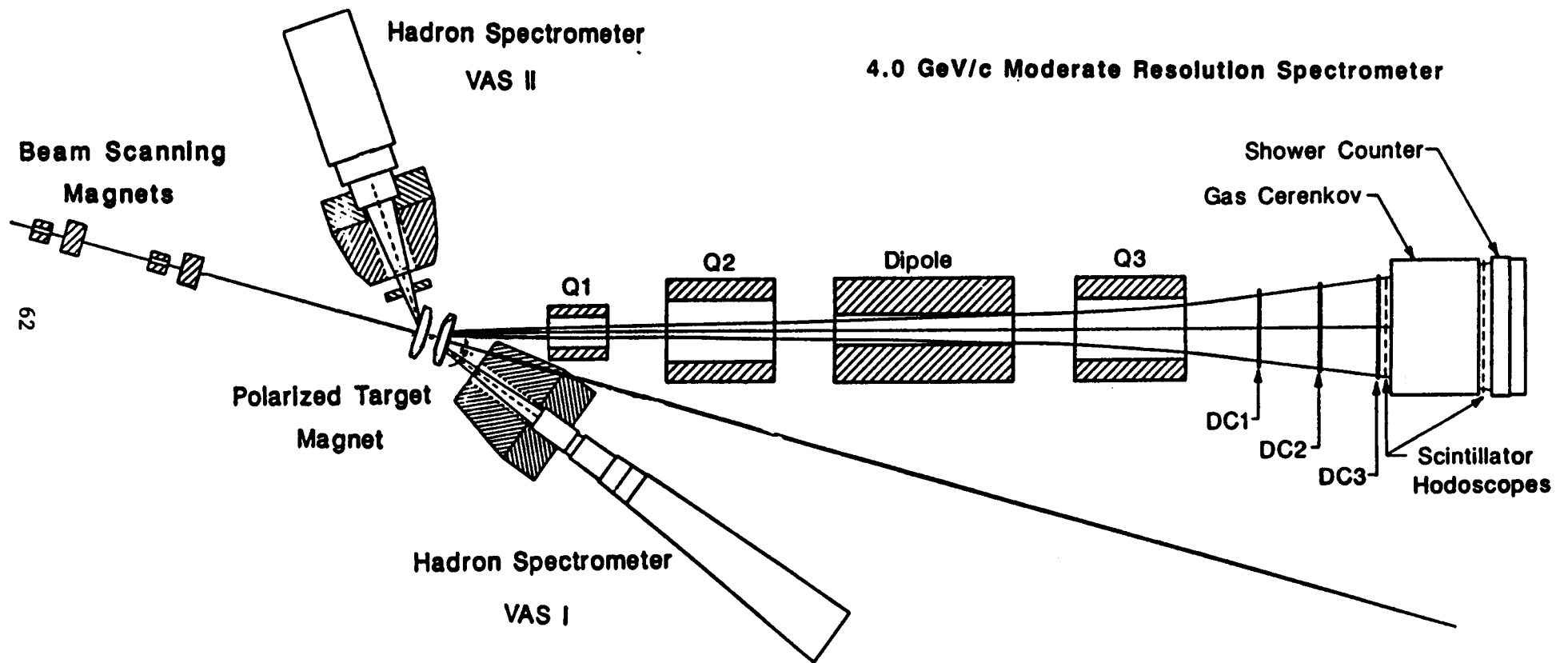


Figure 9 Spectrometer arrangement in End Station C.

4.0 GeV/c Moderate Resolution Spectrometer

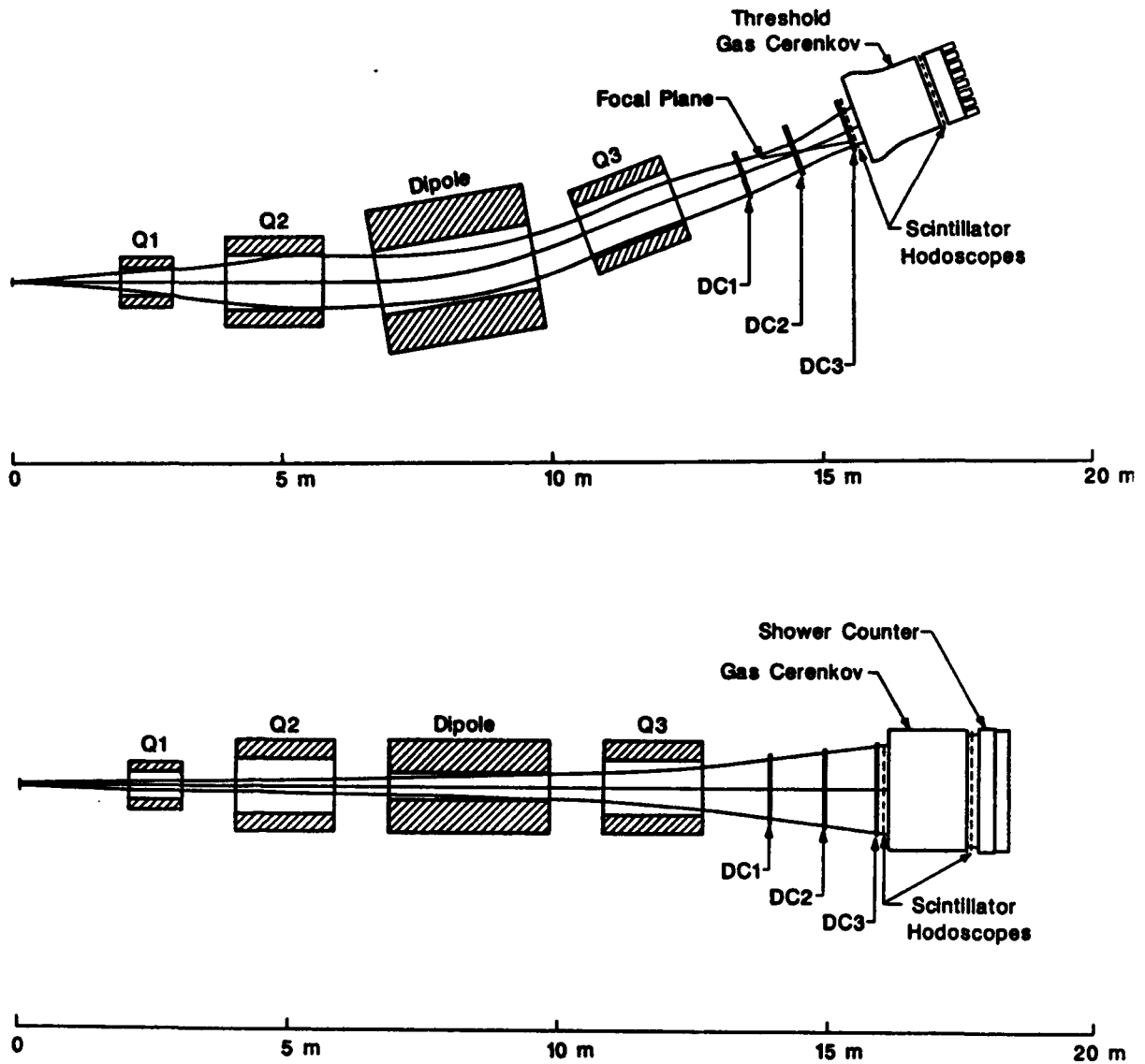


Figure 10 Moderate Resolution Electron Spectrometer for End Station C. The maximum momentum capability can be increased to 5.5 GeV/c by increasing the field in the dipole to 2.1T.

For hadron detection, the proposed solution has been derived by V. Burkert (Ref. 1, p. 11.54) from a similar one used at the Bonn 2.5 GeV electron accelerator in electroproduction experiments^[4]. Called Variable Acceptance Spectrometer (VAS), it consists in a single, uniform field, dipole with 35 cm gap, followed by a set of xy drift chambers and scintillation hodoscopes for particle identification and trajectory reconstruction. Figures 11 and 12 show two versions of it called VAS I and VAS II respectively.

VAS I has a maximum momentum capability of ~ 4 GeV/c. Using collimators to avoid the scintillation detectors to be in direct view of the target, luminosities larger than 10^{36} should be achievable, with adequate shielding around the detector system. Figure 13 shows the angular and momentum acceptancies at the maximum field value of 1.8T. Particle identification is obtained through time-of-flight measurement using microbunch timing, with the electron arm used to define which microbunch is involved. It has been shown that π/K separation up to ~ 1.5 GeV/c and $(\pi/K)/p$ separation up to ~ 2.5 GeV/c can be achieved. Momentum resolution is about $3 \cdot 10^{-8}$, and the maximum angular range is 20° to 160° .

As such a device is substantially lighter than the multi-element spectrometers, (total weight ~ 60 tons) one can envisage to move it for out-of-plane measurements. The first idea (Ref. 1, p. 11.54) was to raise it vertically, reverse the field and reconfigure the detector system accordingly. A preferred method, allowing continuous variation of the out-of-plane angle, would be to rotate the dipole around a horizontal axis going through the target and perpendicular to the dipole symmetry plane. The detector system, and possibly the shielding, would have to follow.

For coincidence experiments involving three charged particle, a third spectrometer would be necessary. As a second hadron spectrometer, VAS II can analyze particle momenta up to ~ 1.7 GeV/c with a maximum field of 1T, and have similar solid angle and

VAS II

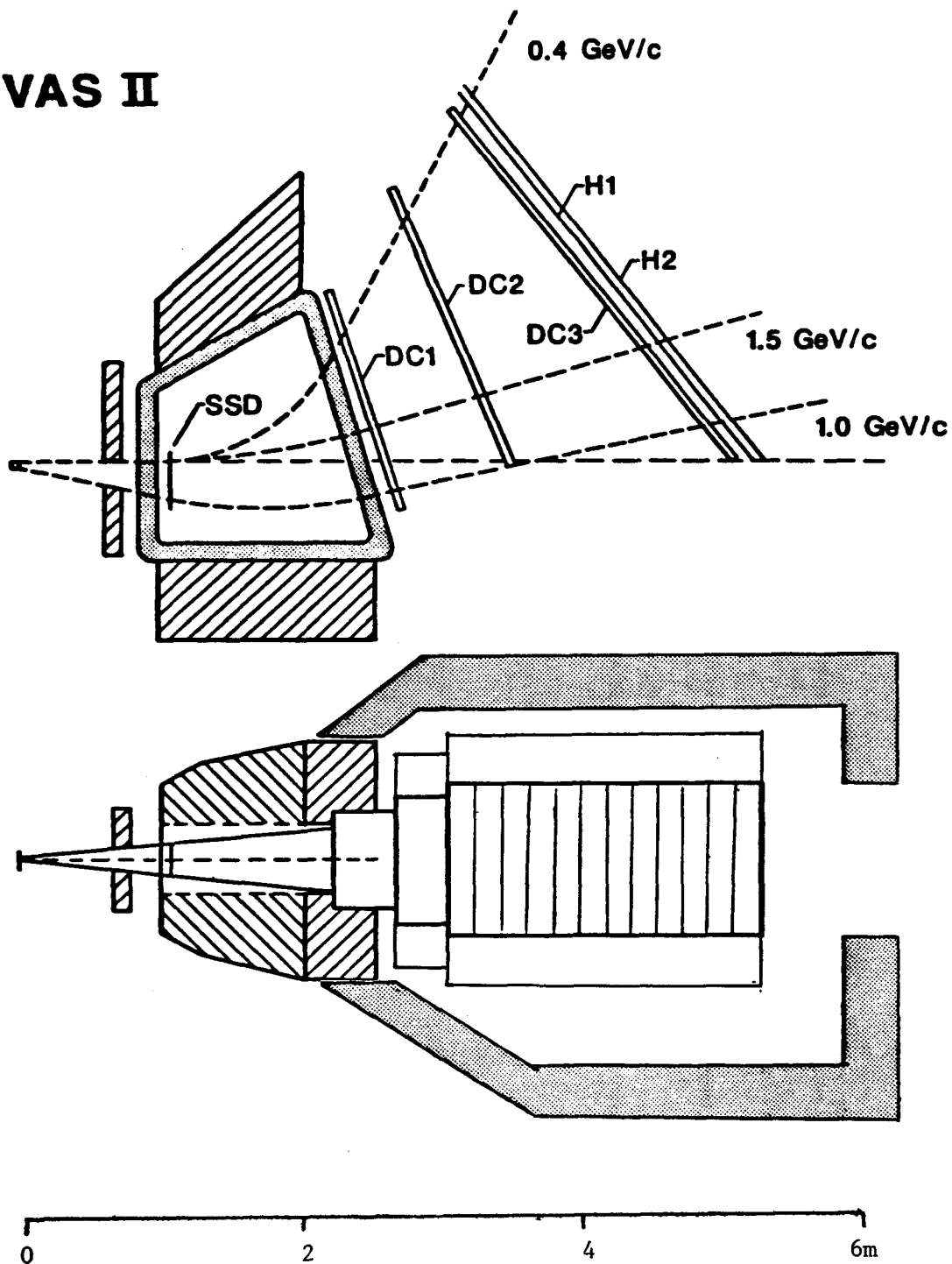


Figure 12 Variable Acceptance Spectrometer VAS II.

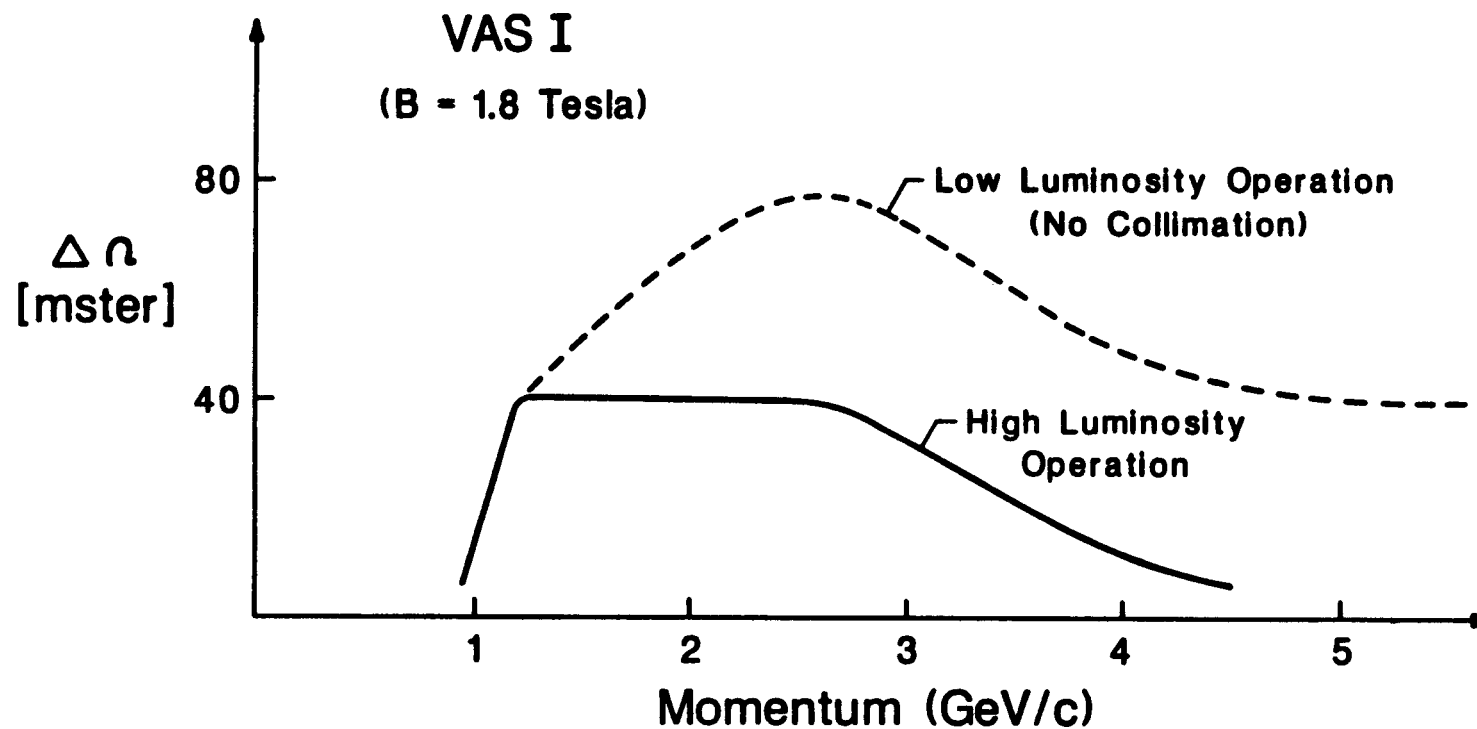


Figure 13 Solid angle as a function of the particle momentum in VAS I operated at 1.8T, with (full line) and without (dashed line) collimator to prevent direct detector exposure.

momentum resolution. Its total weight is about 40 tons. Neutron detection in End Station C is also under study.

4. End Station B Experimental Equipment

In addition to experiments that require high resolution magnetic spectrometers (with relatively small solid angles) there are two important classes of experiments at CEBAF that can only be performed with (or will at least benefit from) a large acceptance spectrometer (LAS):

1. The detection of multi-hadron final states, like in the photo- (or electro-) excitation of the nucleon resonances in both elementary and nuclear production ($\gamma N \rightarrow N^* \rightarrow \pi\pi N$, $\gamma D \rightarrow \Delta\Delta \rightarrow NN\pi\pi$, $\gamma A \rightarrow \pi NX$, $\gamma A \rightarrow NNX$).
2. Measurements at limited luminosity. The limitation can be due to the target (e.g. polarized target with low radiation resistance, polarized gas targets, active targets for the detection of low range recoils) or due to the beam (e.g. all experiments using a tagged bremsstrahlung or other secondary beams independent of the number of particles in the final state).

The physics program results in the following requirements for a general purpose large acceptance spectrometer:

1. Homogeneous coverage of a large angular and momentum range for charged particles (magnetic analysis), photons (total absorption counters) and neutrons.
2. Good energy and angular resolution (for all particles).
3. Good particle identification properties.
4. No transverse magnetic field at the beam axis (to avoid sweeping (e^+, e^-) pairs into the detector).
5. Large field free space around the target to allow for the installation of complicated targets (cryogenic, polarized, track

sensitive, etc.) or auxilliary equipment (nucleon polarimeter, γ -converter).

6. High count rate capability. The operation in a tagged photon beam ($N_\gamma = 10^7/\text{sec}$) should be easy; it will be more challenging to operate the detector in the high background environment caused by a reasonable intensity electron beam (e.g. $1 \mu\text{A}$ on a 1 mg/cm^2 target + luminosity $\approx 10^{38}/\text{cm}^2 \cdot \text{sec}$).

The requirements are best satisfied by a toroidal magnet (see Figure 14 for a perspective view) consisting of 8 coils arranged around the beam line to produce essentially a magnetic field in ϕ -direction. Charged particles are tracked by drift chambers; scintillation counters are used for the trigger and for time-of-flight; photons are detected by shower counters. A description of the main features of the detector is given below.

1. Magnetic field. The field is generated by a toroidal magnet (4 m diameter, 4 m length) consisting of 8 super-conducting coils. Each coil carries a current of 500 kA. The coil shape is given in Figure 15. The inner part of the coil is circular to avoid transverse defocusing effects. In the forward direction the coil comes close to the beam to generate a high $\int B_\phi dl$ for the high momentum forward going particles. The radial dependence of B_ϕ is shown in Figure 16 for 3 different values of z . The forces that pull the coils towards the axis (about 50 tons/coil) are handled by using a warm support structure around the outside of the coils. About 20% of the y -range is obstructed by the vacuum chamber.
2. Particle tracking. In each of the eight segments (see Figure 17) charged particles are tracked by planar wire chambers. The segments are instrumented independently to achieve high count rate capability. The estimated momentum resolution is $\approx 1\%$ in the momentum range of interest; the initial ϕ and θ of the track can be determined to $\leq 5 \text{ mrad}$.

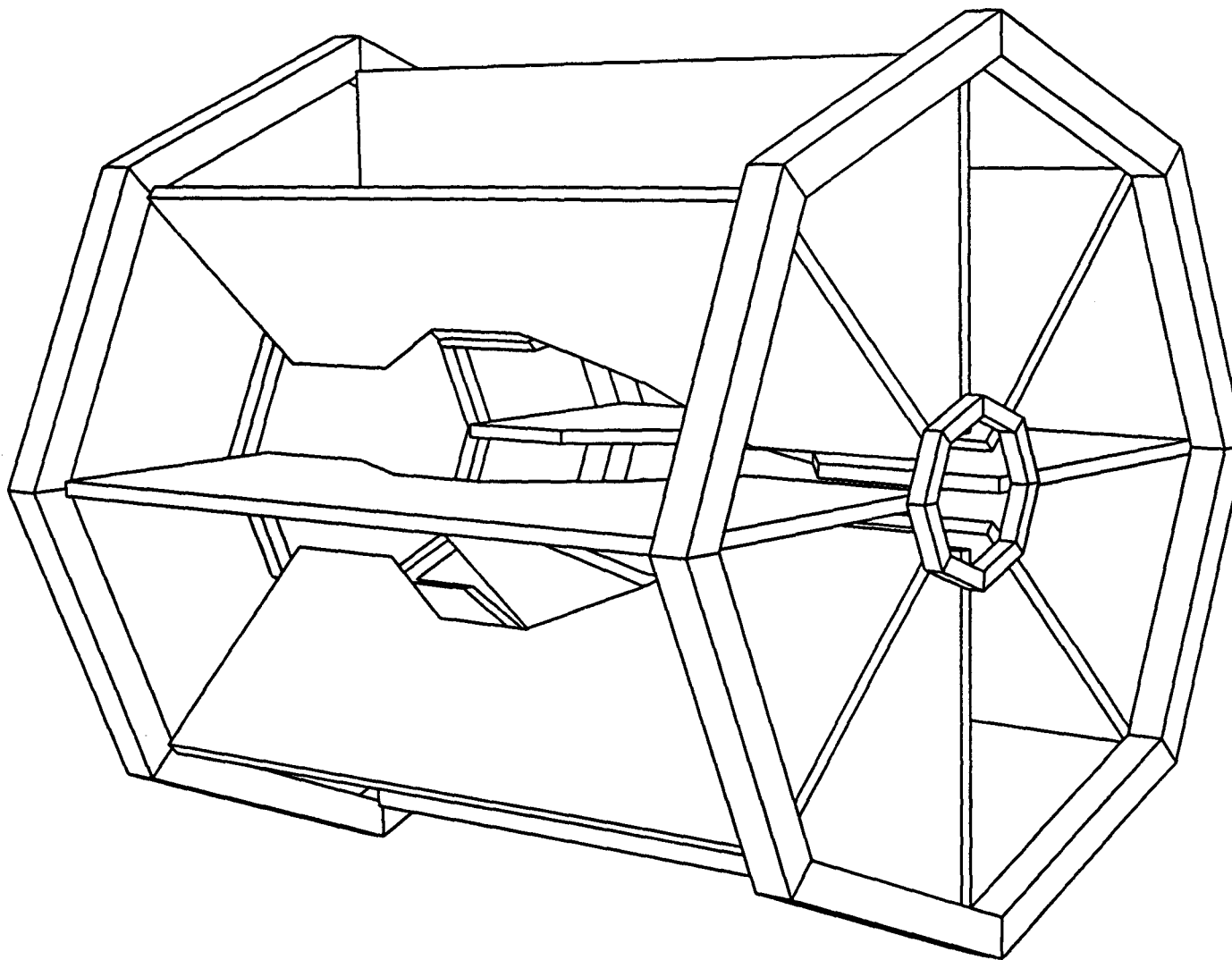


Figure 14 Perspective view of the Large Acceptance Spectrometer.

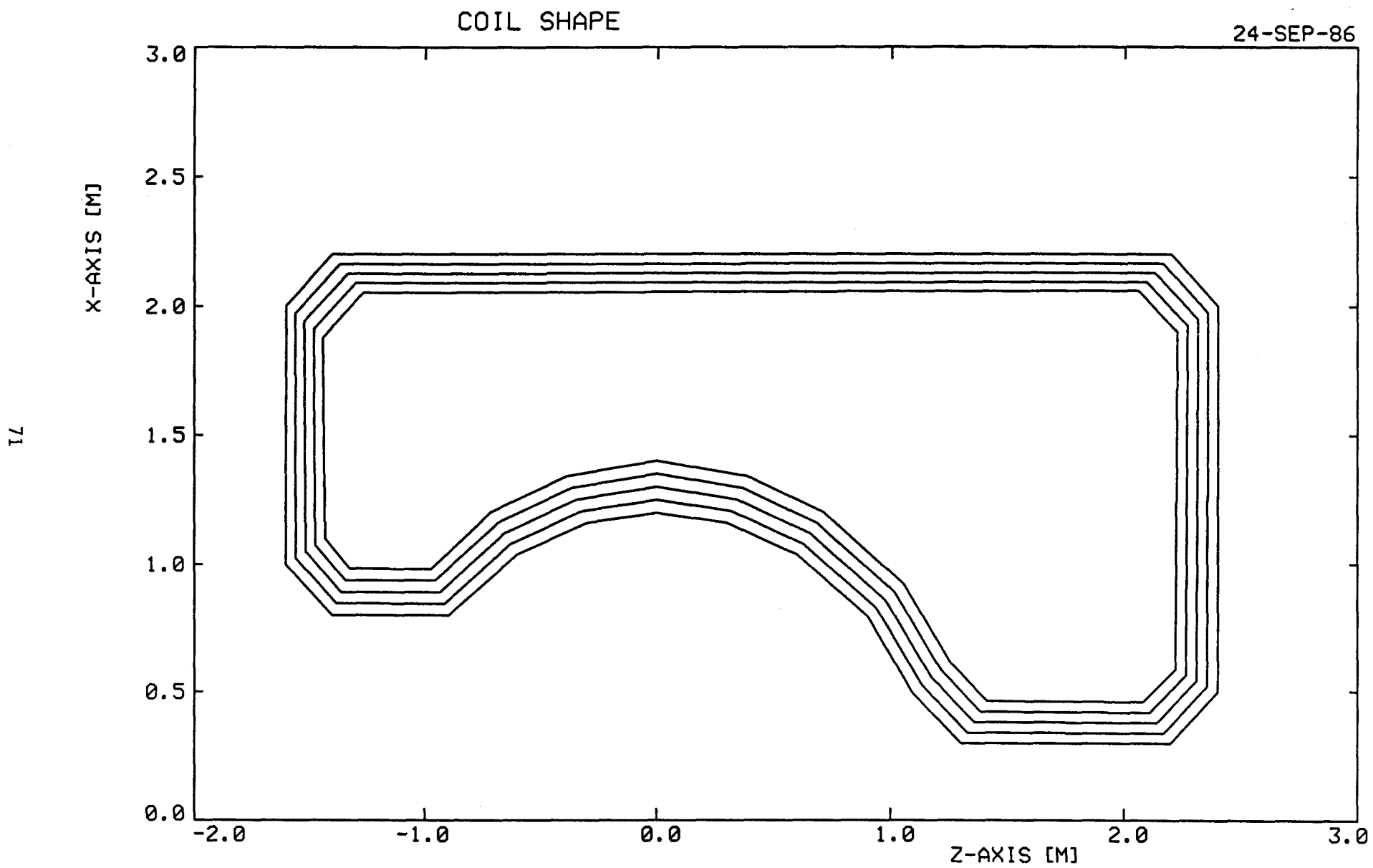


Figure 15 LAS coil shape.

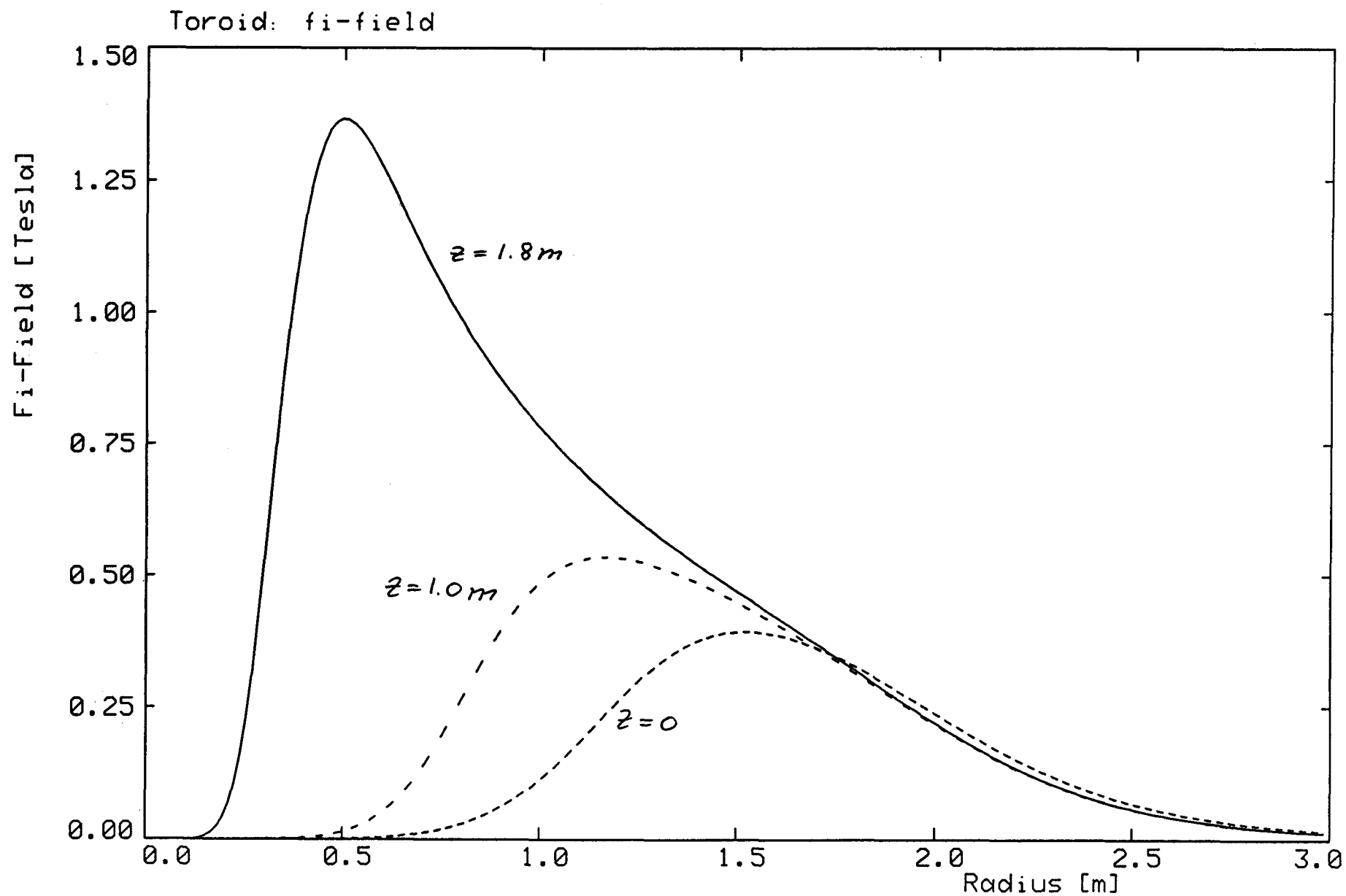


Figure 16 Radial dependence of the ϕ -component of the magnetic field in the symmetry plane of one octant, at different distances along the beam from the target location.

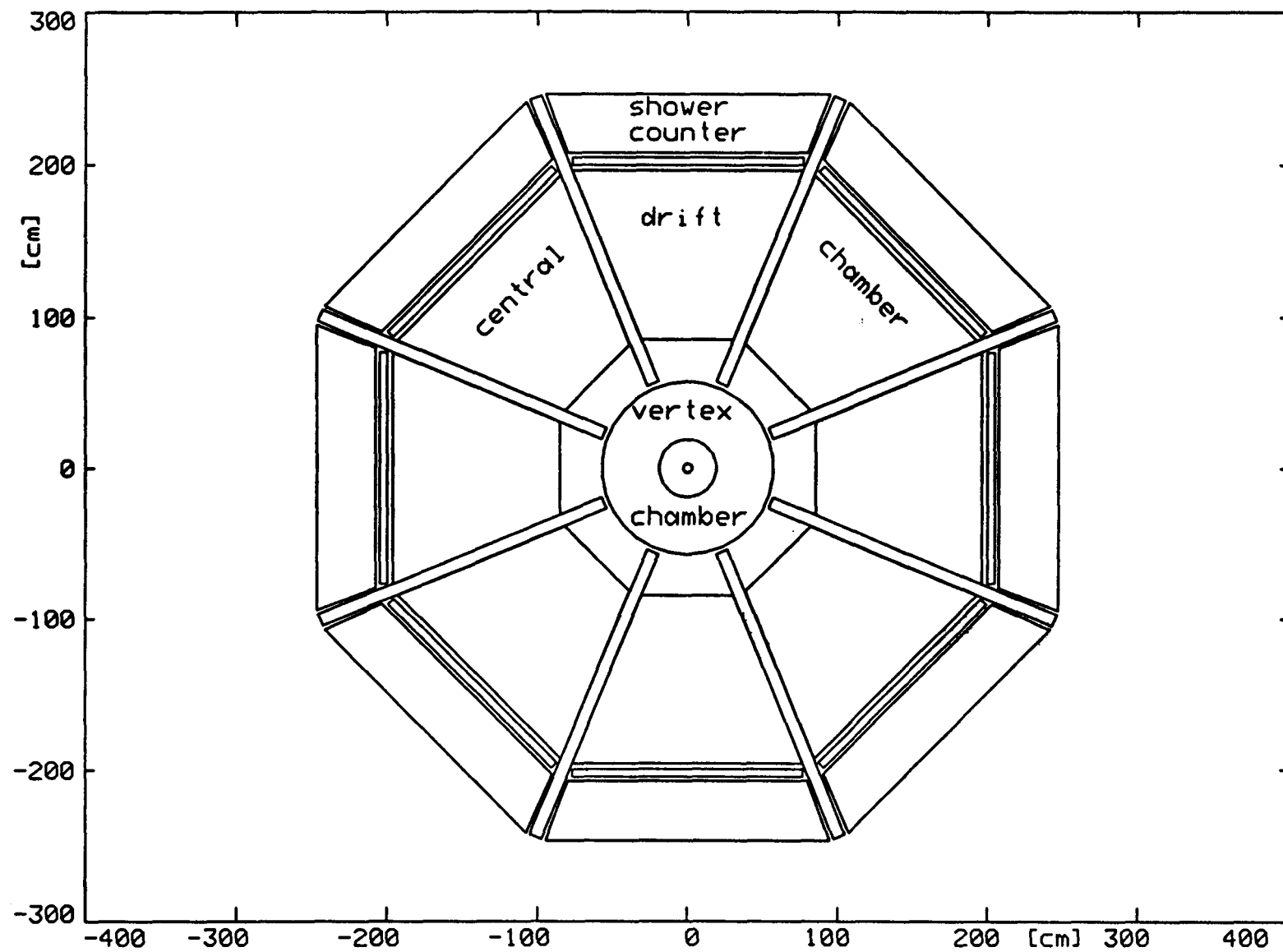


Figure 17 Detector arrangement in the LAS.

3. Scintillation counters. The outer planar drift chambers are covered by scintillation counters. In each segment there are 8 barrel counters, each about 400 cm long, 20 cm wide, and 5 cm thick. The counters are viewed by phototubes at both ends for improved timing and position resolution. Each endcap segment is covered by 4 counters viewed by a single photomultiplier. The scintillation counters serve the double purpose of providing the trigger and the time-of-flight information. Also, a fraction of the high energy neutrons ($\approx 5\%$) will interact in the outer scintillation counters and will thus be detected.
4. Shower counter. The detector can be surrounded by shower counters for the detection of showering particles, especially for high energy photons from Compton scattering, π^0 and η decays, etc.
5. Particle identification. The combination of momentum and time-of-flight is used for the identification of charged particles. Pions and kaons can be separated up to 1.5 GeV/c, kaons and protons up to 2.5 GeV/c. π/μ and μ/e separation can be achieved by using the pulse height in the shower counter in addition.

The performance of the LAS can be predicted using Monte-Carlo methods to generate multiple particle final states and to test particle tracking and reconstruction. As an example, the probability to reconstruct the complete $\pi^+\pi^-p$ final state from the $F35(1975) \rightarrow \pi^+\pi^-p$ decay has been found to be 30% (most of the events are lost because one of the 3 particles hits a coil).

In addition to low intensity electron and photon beams, the LAS is used for experiments with secondary kaon, pion and muon beams. It is complementary to the high resolution magnetic spectrometers: by using the standard spectrometers for high luminosity and the LAS for low luminosity experiments a wide range physics phenomena can be covered at CEBAF.

5. Conclusion

The spectrometers and detectors which have been described, although still in a preliminary stage of designs, will represent the major pieces of experimental equipment in the three CEBAF end stations. Choices between alternate solutions which have been presented here will have to be done soon as the end station sizes and configurations are planned to be frozen by April 1987. Studies are underway towards better understanding and optimization of field configurations using 2D and 3D codes, more detailed optics studies using ray-tracing codes, more detailed definition of the magnet technology, and more precise cost and schedule evaluations. To realize a better adaptation of the spectrometer arrangements to a broad physics program, typical multiparticle experiments have to be designed with some detail. Specific problems like out-of-plane measurements, triple arm experiments and use of (essentially) non-magnetic devices clearly need more work.

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FROM THE STANDARD MODEL TO SUPERSTRINGS

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INTRODUCTION

In the last fifteen or twenty years there has been incredible progress in our understanding of the elementary particles and their interactions. We now have a theory - the standard model of the strong, weak and electromagnetic interactions - which is almost certainly a correct description of Nature to some degree of approximation. With the possible exception of the anomalous positron events from Darmstadt, the standard model either predicts or is compatible with all known facts. It describes the physics that has been observed under normal terrestrial and astrophysical conditions, and at existing accelerators up to energies of about 100 GeV. Despite the great success of the standard model, however, it is certainly not the ultimate theory of Nature. This is because it is too complicated: it has many free parameters, and there are many aspects of Nature which are input and not explained in any fundamental way. Therefore, there is almost certainly something more fundamental. In this talk I will describe one particularly promising line of development of ideas for physics beyond the standard model. This line incorporates grand unified theories (GUTs), in which the strong, weak, and electromagnetic interactions are all components of a single underlying unified interaction. A further development involves supersymmetry and supergravity. Supersymmetry is a particular type of symmetry in which fermions and bosons are related to each other. Supersymmetry was introduced mainly to help explain why the weak interaction scale can be so much smaller than the Planck scale (the scale at which quantum gravity becomes important). Most supersymmetric models do not explain the origin of this scale, but at least they stabilize it with respect to higher order corrections. In addition to this technical but very important problem, supersymmetry brings gravity into the theory in a non-trivial way. Unfortunately, it does not explain how to generate a sensible quantum theory of gravity. Finally, in the last couple of years there has been a great deal of excitement about superstring theories. These very speculative theories suggest that physics beyond the Planck scale may be very different than what we see at low energies. In particular, the basic entities are extended one-dimensional objects called strings rather than point-like particles. Superstring theories apparently lead to finite quantum theories of gravity, and, moreover, they give the hope of being the ultimate theory of everything. However, we are very far from knowing whether Nature indeed chooses these very elegant theories.

THE STANDARD MODEL

The standard electroweak theory is a gauge field theory. Gauge

field theories are mathematically consistent, renormalizable field theories based on an underlying symmetry principle called gauge invariance. Gauge invariance predicts the existence of spin-1 gauge bosons which mediate the underlying interactions. The gauge symmetry not only predicts the existence of these gauge bosons, but it prescribes the form of their interactions with each other and with the fermions and other particles of the theory, up to a constant known as the gauge coupling constant. In the standard model the basic symmetry is described by the $SU_3 \times SU_2 \times U_1$ group. This is the group of transformations of the fundamental fields of the theory under which the equations of motion are left invariant. A gauge invariance means that the symmetry transformation can be performed independently at different space-time points. The strong interactions are associated with the SU_3 group (with gauge coupling constant g_s) of transformations of the different colors of quarks (and antiquarks) into each other; red and green quarks can be interchanged, for example, without changing the form of the equations of motion. The electroweak interactions are based on the $SU_2 \times U_1$ group, with two coupling constants g and g' . The charged current weak interactions are associated with transformations of up quarks into down quarks, and neutrinos into electrons. The neutral current and electromagnetic interactions are associated with phase transformations on the fields of all of the quarks, anti-quarks, leptons and anti-leptons.

THE STRONG INTERACTIONS (QCD)¹

It is now generally believed that the strong interactions are based on a gauge theory known as quantum chromodynamics (QCD). This is a gauge theory based on the group $SU(3)$ which describes transformations of the three colors (red, green, and blue) of quarks into one another. The $SU(3)$ group has 8 generators; therefore, there are 8 massless gauge bosons or gluons associated with it. The underlying force involves the emission of a gluon by a quark. The quark then changes into a quark of a different color but the same flavor and electric charge; the gluon can then be absorbed by another quark. The amplitude for emission or absorption is proportional to the strong coupling constant g_s . In addition, there are self-interactions amongst the gluons because, unlike the analog of electrodynamics, the gluons themselves carry color charge. There are three and four point vertices, the form of which are also prescribed by the gauge symmetry and have strengths related to the strong coupling constant. The

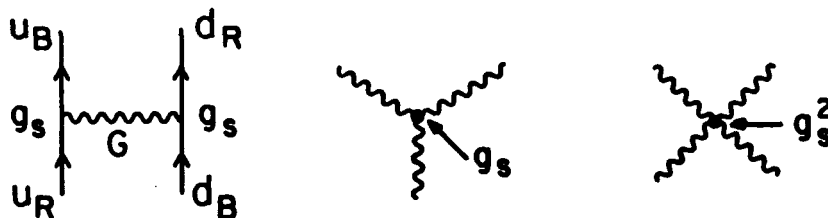
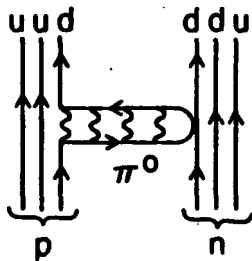


Fig. 1. Fundamental QCD interactions.

strong coupling constant is large (i.e. of order one) and for that reason the strong interactions are strong. It also means that the force is non-perturbative. It is very hard to do calculations, and therefore QCD has not been subjected to the same degree of rigorous quantitative testing as the other interactions. It is thought, though not rigorously proved, that the large value of the coupling constant and the existence of the gluon self-interactions lead to quark confinement (i.e. the non-observation of isolated quarks). It is also believed that all colored objects are confined and unable to propagate freely. This is the reason that the strong force is short range even though the gluon is massless.

QCD has been extremely successful qualitatively. It has all the right general features, correct limits, and so forth. For example, a) QCD incorporates the one-boson exchange potential model of the strong force in the appropriate limit, with the exchange of a pion or other meson interpreted as a higher order effect.



b) Current algebra and $SU_3 \times SU_3$ chiral symmetry are naturally incorporated into the theory. The success of the soft pion theorems is associated with the smallness of the up and down quark masses.

c) QCD naturally explains parity conservation to order α .

d) The approximate scaling observed in deep inelastic electron-nucleon scattering is incorporated in the theory. No other realistic theory has this property.

Fig. 2. The one pion exchange diagram.

Because of the strength of the QCD coupling constant it is very hard to do the necessary calculations to subject the theory to rigorous, quantitative tests; however, to the extent that such tests have been carried out, it has been generally successful. In particular, the logarithmic scaling violation observed in deep inelastic scattering of electrons, muons and neutrinos from nucleons is described in a satisfactory way.

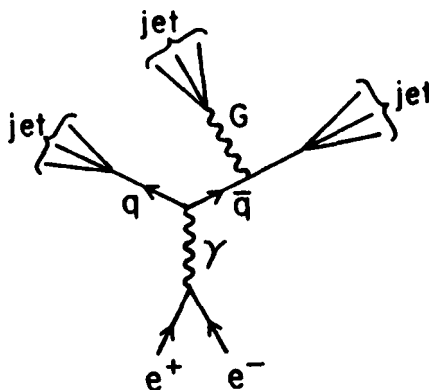


Fig. 3. Jet production in e^+e^- annihilation.

Secondly, hard scattering and jet production observed in e^+e^- , $p\bar{p}$, and pp scattering is consistent with the QCD predictions, although again it is hard to do the calculations in detail.

There is a great deal of current activity, both theoretical and experimental. Much of the ultimate motivation is to provide the final definite proof that QCD is the correct theory. For example, there is much theoretical effort trying to develop non-perturbative calculational techniques; e.g. lattice calculations of the hadronic spectrum, phase transitions, attempts

to prove quark confinement, etc. Another area of interest both theoretically and experimentally is the search for glueballs, which are predicted bound states of gluons only. There is much effort on the possibility of new forms of matter, such as quark matter or quark-gluon plasma, which could occur in conditions of sufficient density and energy. Finally, there is a lot of continuing activity in jet physics, deep inelastic scattering, heavy quark systems and so forth.

QCD is probably correct. It has all the right limits, and those tests that we can make are successful. However, the most compelling argument is that there just doesn't seem to be any competing candidate within the general framework of renormalizable field theories.

ELECTROWEAK INTERACTIONS²

The standard model incorporates the electroweak theory of Glashow, Weinberg and Salam. This is based on the gauge group $SU_2 \times U_1$,

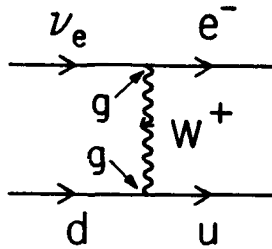


Fig. 4. The charged current interaction.

which has gauge couplings g and g' for the two subgroups. The SU_2 transformations relate up and down quarks, neutrinos and electrons, and so forth, whereas the neutral generator of SU_2 and the U_1 generator involve phase transformations on all of the quarks and leptons. The charged current weak interactions are associated with the exchange of a charged boson called the W . The Fermi constant is related by

$$G_F/\sqrt{2} \approx g^2/8 M_W^2.$$

In addition, there is a neutral current interaction associated with the exchange of a massive electrically neutral boson, called the Z . The couplings of the Z involve g and also $\sin^2 \theta_W = g'^2/(g^2 + g'^2) = 1 - M_W^2/M_Z^2$. Finally, the electromagnetic interactions (quantum electrodynamics) is incorporated in the theory as well. The amplitude

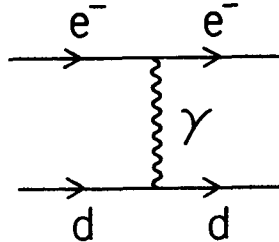
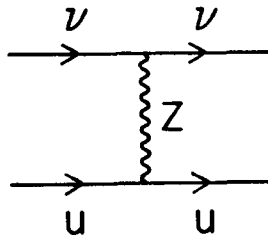


Fig. 5. The neutral current and electromagnetic interactions.

for photon emission or absorption is proportional to the electric charge of the positron, given by $e = g \sin \theta_W$.

The $SU_2 \times U_1$ model is a renormalizable field theory. It incorporates the Fermi theory of the charged current weak inter-

actions (as modified later by others to include parity violation and strangeness changing effects). In addition, it predicted the existence of the neutral current interaction. In the 1970's the neutral current interaction was observed and in the last decade it has been intensively studied in many types of reactions and over a very wide range of momenta (and therefore distances). For example, the

parameter $\sin^2 \theta_W$ has been measured in such diverse reactions as parity violation in atoms, neutrino-electron elastic scattering, polarization asymmetries in electron and muon scattering from hadrons, deep

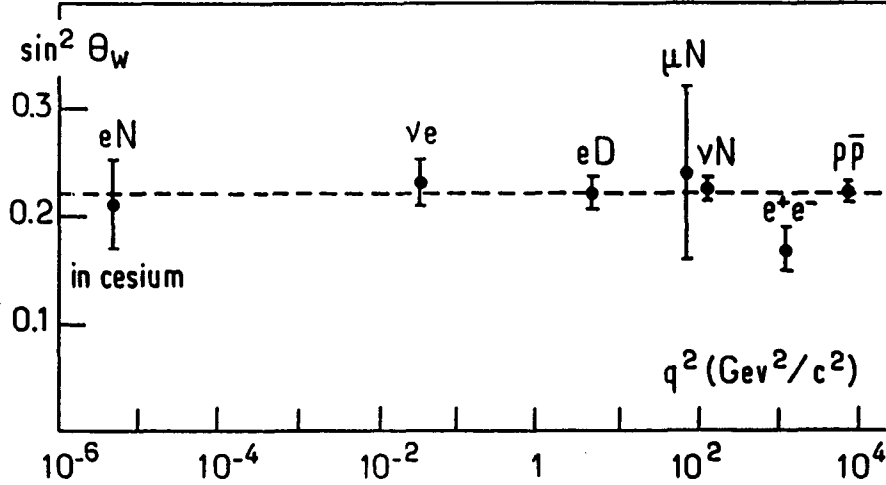


Fig. 6. $\sin^2 \theta_W$ determinations from a variety of processes. From M. A. Bouchiat et al.³

inelastic neutrino scattering, forward-backward asymmetries in electron-positron annihilation, and indirectly in the measured masses of the W and Z bosons. The consistency of these results indicates that the $SU_2 \times U_1$ model is approximately correct down to a distance scale of 10^{-16} cm, which corresponds to about 100 GeV. The value of the weak angle determined from these measurements is $\sin^2 \theta_W = 0.223 \pm .004$ (the quoted error is experimental only; there is an additional theoretical error of approximately 0.005. This average represents all of the available data as of the summer of 1985). In addition to the prediction of the neutral current interaction, the $SU_2 \times U_1$ model predicted the existence of the W and Z bosons and of their masses. The W and Z were observed by the UA1 and UA2 groups at the CERN $\bar{p}p$ collider a few years ago. By now their masses and other properties have been measured in detail. The measured masses are compared with the theoretical prediction in Table 1. The agreement is remarkable.

	UA1	UA2	Theory
M_W (GeV)	$83.1^{+1.3}_{-0.8} \pm 3$	$81.2 \pm 1.1 \pm 1.3$	$A/\sin \theta_W = 82.0 \pm 1.7$ (77.2 without rad. corr.)
M_Z (GeV)	$93.0 \pm 1.6 \pm 3$	$92.5 \pm 1.3 \pm 1.5$	$2A/\sin 2\theta_W = 93.0 \pm 1.4$ (88.2 without rad. corr.)

Table 1. Measured and predicted values for M_W and M_Z . $A \equiv (\pi\alpha/\sqrt{2}G_F)^{1/2} (1-\Delta r)^{-1/2}$, where $\Delta r = 0.0696 \pm 0.002$ is due to radiative corrections (other corrections affect the value of $\sin^2 \theta_W$ extracted from neutral current experiments).

CURRENT WORK

There is still a great deal of effort on the standard electro-weak model. Some of the major activities are in the following directions:

- a) A great deal of experimental work is going on to continue studying the neutral and charged current interactions to a high degree of precision, and to measure the masses and properties of the W and Z to higher accuracy. One of the motivations for this work is that one would like to test the theory at the level of the radiative corrections. At present the measured value of the W mass is about two standard deviations away from the prediction that one would have without including the corrections and it is in perfect agreement with the prediction that includes the radiation corrections. One is therefore on the verge of being able to test these effects, and with the advent of LEP and SLC it should be possible to make these tests precise and conclusive; unfortunately, it will probably never be possible to test the radiative corrections to the same degree of precision that has been done with quantum electrodynamics. Another motivation for these precision tests is to search for and put limits on possible new physics beyond the standard model, such as on the possible existence of additional heavy W and Z gauge bosons. At present the limits on additional neutral bosons are rather weak. For bosons with the couplings predicted by some of the more popular extensions of the standard model, one could tolerate masses as low as 100-200 GeV. At present the most stringent limits come from the indirect neutral current constraints, but as soon as the Fermilab Tevatron starts running the direct production limits will quickly become more important. The most popular theory predicting additional charged W's are the $SU_2 \times SU_2 \times U_1$ models. In such models the additional charged bosons couple to right-handed rather than to left-handed currents. Because of this new type of coupling it turns out that the extra W's can give very important contributions to the K_L - K_S mass difference. It is therefore likely that any such additional W's should be more massive than 2-3 TeV. However a great deal of theory goes into this limit and it should therefore be regarded as suggestive, but not rigorous.
- b) Another area of activity is the study of flavor mixing between quarks. This refers to the determination of the mixing angles which govern the rate of weak decays of heavy quarks into lighter quarks. These mixings are free parameters in the standard model; however, we would like to know them in detail because they are useful for computing new weak processes and also because they may give us some insight into the origins of the fermion mass.
- c) CP violation. We are still not certain whether CP violation is due to higher order weak interactions or whether it requires the introduction of new physics or new interactions.
- d) Rare Decays, such as $K_L \rightarrow \mu e$. The search for rare decays, such as those which violate quark or leptonic flavor, is very important. The predicted rates for such decays in the standard model (via higher order effects) are very low or zero. However, most extensions of the standard model predict one or another type of rare decay at some level. Therefore it is important to continue such searches as far as

is feasible as a probe of new physics.

e) The Top Quark. The observed decays of the bottom quark indirectly indicate that it must have a partner, the top quark. So far, the t has not been discovered. It must be more massive than about 23 GeV because of direct searches at PETRA. Theoretical arguments strongly suggest that the t should be less than about 200 GeV.

f) The Higgs Scalar. The masses of the W and Z bosons require that the $SU_2 \times U_1$ gauge symmetry be spontaneously broken down to the U_1 of electromagnetism. The simplest mechanism for accomplishing this breaking is the Higgs mechanism. One introduces a scalar field called the Higgs field. If the lowest energy state (the vacuum) has a non-zero value for this Higgs field, then the symmetry will be spontaneously broken. As the W and Z propagate, they constantly interact with this field and acquire effective masses. The quantum of this field is known as the Higgs scalar. The Higgs has not been observed - it is the last important ingredient of the $SU_2 \times U_1$ which needs to be verified. Unfortunately, the theory does not predict the Higgs mass. Plausible theoretical arguments suggest that it should be somewhere between 10 GeV and 1 TeV, but these are not rigorous. Unfortunately, the Higgs scalar is relatively weakly coupled, it is hard to produce, and it does not have a very dramatic signature. As we will see, the Higgs sector is not only the least understood, but is also the most unattractive feature of the standard model. It is possible that the Higgs mechanism is right, or it is possible that there is some other mechanism for spontaneous symmetry breaking. Investigations in this direction are one of the major motivations for the desire to build larger accelerators.

g) Neutrino Mass. One of the most interesting areas of investigation is the search for a possible non-zero mass for the neutrino. The neutrino is predicted to be massless in the minimal standard model, but almost all extensions predict a non-zero mass at some value. For some years the ITEP group in Moscow has reported evidence from tritium β decay for a non-zero mass for the electron neutrino somewhere in the 17 to 40 eV range. This result has not yet been confirmed by other experiments and in fact it is on the verge of being contradicted by the Zurich experiment. If true, however, the ITEP result would be extremely important, not only for particle physics but for cosmology as well. Relic neutrinos left over from the early Universe are expected to still fill space ($n_\nu \sim 50/\text{cm}^3$ for each neutrino type) and if the neutrino has even a tiny mass, it could dominate the energy density of the Universe and perhaps provide an explanation for the dark matter. There are various current neutrino mass experiments. In addition to Zurich, there are experiments at Los Alamos, Tokyo, Livermore and elsewhere which should confirm or refute the ITEP result in the near future. Results are eagerly awaited.

Another interesting development concerns the possibility of much smaller neutrino masses. Recently Mikheyev and Smirnov⁴ have developed an old idea of Wolfenstein that may well explain the missing solar neutrinos. The idea is that electron neutrinos produced at the core of the sun interact with electrons on their way out. For appropriate sets of parameters, an adiabatic conversion of the electron neutrinos into muon or tau neutrinos is possible, explaining the results of the

Davis experiment. This is a very elegant and plausible sounding solution to the solar neutrino problem. If true, it would suggest neutrino masses of order 10^{-2} eV or so for the ν_μ or ν_τ and even smaller for the ν_e . Unfortunately, it would be very difficult to verify neutrino masses of this range in the laboratory. Therefore further solar neutrino experiments, including gallium and heavy water experiments, are urgently needed.

THEORETICAL PROBLEMS

The standard $SU_3 \times SU_2 \times U_1$ model, combined with classical general relativity, is almost certainly correct at some level. However it contains far too much arbitrariness to be the final story of Nature. One way of seeing this is that the minimal version with massless neutrinos has 21 free parameters, and nobody believes that the ultimate theory of Nature could have so much arbitrariness. I have classified the problems or unexplained features of the standard model under five headings.

a) The Gauge Problem

The standard model gauge group is a complicated direct product of three factors with three independent coupling constants. Charge quantization, which refers to the fact that the magnitudes of the proton and of the electron electric charges are the same, is not explained and must be put into the standard model by hand.

b) Fermion Problem

The standard model involves a very complicated reducible representation for the fermions. Ordinary matter can be constructed out of the fermions of the first family (ν_e , e^- , u , d). However, in the laboratory we observe at least two additional fermion families - (ν_μ , μ^- , c , s), (ν_τ , τ^- , t , b) - which appear to be identical with the first except that the quarks and leptons are heavier. We have no fundamental understanding of why there should be these heavy families or how many there are. In addition, the standard model does not explain or predict the pattern of fermion masses, which are observed to vary over five orders of magnitude, and it does not explain the various mixings.

c) Higgs/Hierarchy Problem

The next problem is a fine-tuning problem. The spontaneous breakdown of the $SU_2 \times U_1$ symmetry in the standard model is accomplished by the introduction of a Higgs field. Consistency requires that the mass of the Higgs be not too much different from the weak interaction scale; that is, it should be equal to the W mass within one or two orders of magnitude. The problem is that there are higher order corrections involving loop effects (or even tree-level corrections in grand unified theories) which change (renormalize) the value of the square of the Higgs mass by $\delta m_H^2/M_W^2 \gtrsim 10^{24}-10^{34}$. Therefore the bare value of m_H^2 in the original Lagrangian must be adjusted or fine-tuned to 24 to 34 decimal places to achieve a satisfactory theory. Such a fine-tuning

is possible, but extremely unattractive.

d) Strong CP Problem

Another fine-tuning problem is the strong CP problem. It is possible to add an additional term $(\theta/32\pi^2)g_s^2 \bar{\psi}\psi$ to the QCD Lagrangian. This term breaks P, T, and CP invariance. Limits on the electric dipole moment of the neutron require θ to be less than 10^{-9} . However, weak interaction corrections change or renormalize the lowest order value of θ by about 10^{-3} - that is, 10^6 times more than the total value is allowed to be. Again, one must fine-tune the bare value against the correction to a high degree of precision.

e) Graviton Problem

Finally, I come to the graviton problem. This really has several aspects. First of all, gravity is not unified with the other interactions in the standard model in a fundamental way. Secondly, even though general relativity can be incorporated into the model by hand, we have no idea how to achieve a mathematically consistent theory of quantum gravity: attempts to quantize gravity within the standard model framework lead to horrible divergences and a non-renormalizable theory.

Finally there is yet another fine-tuning problem associated with the cosmological constant. The vacuum energy density associated with the spontaneous symmetry breaking of the $SU_2 \times U_1$ model generates an effective renormalization of the cosmological constant $\delta\Lambda = 8\pi G_N \langle V \rangle$, which is about 50 orders of magnitude larger than the observed value. One must fine-tune the bare cosmological constant against the correction to this incredible degree of precision.

EXTENSIONS OF THE STANDARD MODEL

It is therefore almost certainly true that there must be new physics beyond the standard model. Some of the possible types of new physics that have been discussed extensively in recent years are shown in Table 2. I do not have time to discuss all of these, so I will concentrate on one line of development which I find particularly promising, namely grand unification, supersymmetry/supergravity, and superstrings.

GRAND UNIFIED THEORIES (GUTs)⁵

In the standard model the apparent symmetry of Nature at energies less than 100 GeV is $SU_3 \times U_1^{\text{elm}}$ (as well as classical gravity). At the electroweak unification scale (≈ 100 GeV) one sees a larger symmetry group, namely $SU_3 \times SU_2 \times U_1$ combined with classical gravity. In the standard model this symmetry continues up to the Planck scale (10^{19} GeV), at which point we know that the theory breaks down because quantum gravitational effects become important. In grand unified theories one embeds the $SU_3 \times SU_2 \times U_1$ standard model into a larger group, which is not a direct product of factors, such as SU_5 , SO_{10} or E_6 . There is therefore a single underlying gauge coupling constant and a single underlying interaction. In a very real sense all of the

Extension	Typical Scale (GeV)	Motivation	Comments
extra W's, Z's extra Fermions extra Higgs	10^2 - 10^{19}	-	Generally remnants of something else
Family Symmetry Composite Fermions	10^2 - 10^{19}	Fermion problem	No compelling models
Composite Higgs Composite W, Z	10^2 - 10^3	Higgs problem	No compelling models
Grand Unification	10^{14} - 10^{19}	Gauge problem	strong $\leftrightarrow$ electro- weak
Supersymmetry/ Supergravity	10^2 - 10^4	Higgs, graviton problems	Fermion $\leftrightarrow$ boson
Kaluza-Klein	10^{19}	Graviton	(spin-0) $\leftrightarrow$ (spin-1) $\leftrightarrow$ $\leftrightarrow$ (spin-2) extra dimensions
Superstrings	10^{19}	all?	Gauge + Gravity Fermion $\leftrightarrow$ boson extra dimensions

Table 2. Extensions of the standard model.

interactions are unified: they are merely different aspects of this larger underlying symmetry group. The larger group has additional generators and corresponding gauge bosons, which acquire a large mass of order of the grand unification scale. If the theory is observed or probed at an energy that is much higher than that scale, one sees the full symmetry; all interactions look alike and there is only one observed gauge coupling. However, at energies much lower than the unification scale the strong and electroweak interactions look very different because of symmetry breaking effects, and the effective (momentum-dependent) coupling constants split apart. Therefore we observe completely different features and strengths for the strong and electroweak interactions.

In the simplest grand unified theories one can predict the unification scale in terms of the observed coupling constants at low energy, using the theoretical expression for the momentum dependence of the effective couplings. From the ratio of electromagnetic to strong interaction couplings one predicts the unification scale $M_X \sim 1.3 \times 10^{14 \pm .5}$ GeV, where the uncertainty is mainly due to the fact that it is hard to determine the strong coupling constant precisely. One can also use the fact that there are three couplings which come together at M_X to predict the weak angle, which is defined as a ratio of coupling constants. In the simplest models one predicts $\sin^2 \theta_W = .218 \pm .007$, which is remarkable agreement with the experimental number $.223 \pm .004 \pm .005$. This agreement strongly suggests that grand unified theories may indeed have something to do with the real world. However, it is possible that the agreement is fortuitous.

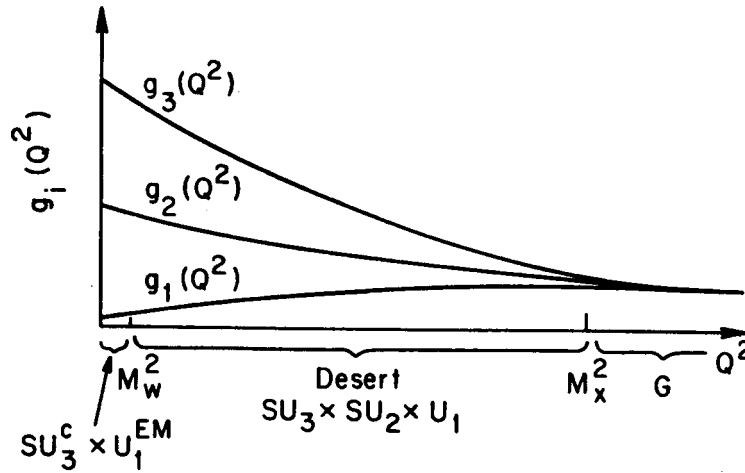


Fig. 7. The momentum dependent coupling constants $g_3 \equiv g_s$, $g_2 \equiv g$, and $g_1 \equiv \sqrt{5/3} g'$.

In addition to unifying the interactions, the additional symmetries within a grand unified theory relate quarks, antiquarks, leptons and antileptons. For example, the simplest grand unified theory - the SU_5 model of Georgi and Glashow - assigns the fermions of the first family to a reducible $5^* + 10$ dimensional representation:

$$\begin{array}{c}
 5^* \qquad \qquad \qquad 10 \\
 \begin{array}{c}
 W^\pm \updownarrow \left(\begin{array}{c} \nu_e \\ e^- \\ d^- \end{array} \right)_L \\
 \leftarrow X, Y \rightarrow
 \end{array}
 \quad
 \begin{array}{c}
 \left(\begin{array}{c} e^+ \\ u \\ d \end{array} \right)_L \\
 \leftrightarrow \quad \leftrightarrow \\
 X, Y
 \end{array}
 \end{array}$$

The neutrino and electron can be rotated into the anti-down quark by the new symmetry generators, as can the positron, the up and down and the anti-up quark. Because of this relation between the different types of fermions, the grand unified theories naturally explain charge quantization; that is, quark and lepton electric charges are related, and the fact that the electron and the proton have the same magnitude of electric charge is predicted by the theory. In addition, there are new gauge bosons associated with these extra symmetries. The SU_5 model contains new bosons which carry both color and electric charge, known as the X and Y bosons, with masses $M_X \sim M_Y \sim 1.3 \times 10^{14}$ GeV. These can mediate proton decay. For example, a diagram for $p \rightarrow e^+ \pi^0$ is shown in Figure 8. The proton lifetime can be estimated in the simplest SU_5 model: one has for the partial lifetime into $e^+ \pi^0$ the expression $\tau_{p \rightarrow e^+ \pi^0} \approx 6.6 \times 10^{28 \pm 0.7} (M_X / 1.3 \times 10^{14} \text{ GeV})^4$ yr, where the

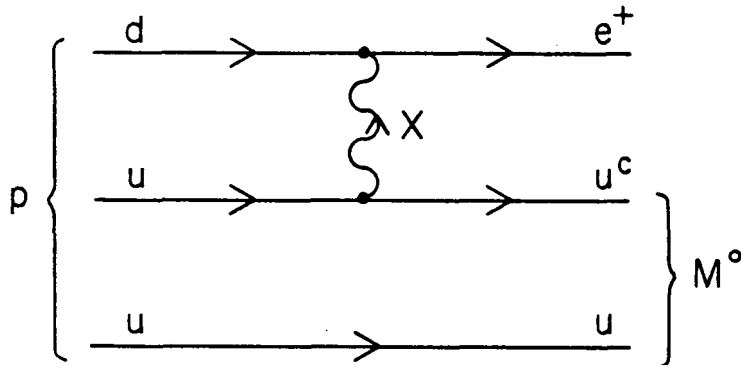


Fig. 8. A typical proton decay diagram.

uncertainty in the coefficient is due to the difficulty of computing the hadronic matrix elements of the quark operators derived from the diagram. For the expected range of unification scale, this implies $\tau_{p \rightarrow e^+ \pi^0} < 8.4 \times 10^{30} \text{ yr}$. However, the uncertainties in the calculation both of the coefficient and unification scale are large and difficult to quantify.

It is therefore useful to take a more conservative limit $\tau_{p \rightarrow e^+ \pi^0} < 1.4 \times 10^{32} \text{ yr}$ for the partial lifetime in SU(5) and similar models.

A number of experiments have been searching for proton decay for some time. There is no definitive evidence for proton decay and at present the lower limit is $\tau_{p \rightarrow e^+ \pi^0} > 3.3 \times 10^{32} \text{ yr}$. Therefore SU(5) and similar unified theories are excluded. However, grand unification itself is alive and well. Many variations and extensions of the simplest theories give longer lifetimes. For example, some models invoke extra gauge boson, fermion, or scalar thresholds; the new particles can lead to a larger M_X and a longer lifetime. Supersymmetric grand unified theories, for example, automatically lead to a larger unification scale. Some models invoke large mixing effects, which can suppress the decay amplitude, and some theories predict the existence of higher dimensional effective operators which have the ultimate effect of raising the unification scale. It should be cautioned, however, that many of the models which give a longer proton lifetime also lose much of their predictive powers for the weak angle.

To summarize the status of grand unified theories:

- a) Grand unified theories solve the gauge problem. The strong, weak and electromagnetic interactions are unified. Charge quantization is explained naturally and $\sin^2 \theta_W$ is correctly predicted, at least in the simplest theories.
- b) Another attractive feature of grand unified theories is that they can naturally explain the cosmological baryon asymmetry. It is an observational fact that our part of the Universe consists of matter and not antimatter: there is approximately one baryon for every 10^{10} microwave photons, but essentially no anti-baryons. In grand unified theories, this asymmetry does not need to be invoked as an initial condition on the Big Bang; rather, it can be generated dynamically in the first instant (the first 10^{-35} sec. or so) after the Big Bang, by baryon number violating interactions related to those which lead to proton decay.
- c) On the other hand, grand unified theories give no help with the fermion, strong CP, Higgs (the Higgs problem is even aggravated), or

graviton problems.

d) The non-observation of proton decay rules out the simplest grand unified theories, but many more complicated schemes are still possible.

e) There are other implications of grand unified theories as well. For example, all grand unified theories predict the existence of superheavy magnetic monopole states with masses $M_M \sim M_X/\alpha_G \sim 10^{16} \text{ GeV}$. These may have been produced during phase transitions during the early Universe and could still be left over as relics today. However, we have no firm prediction as to what the flux of these monopoles should be. Early estimates suggested that far too many of these monopoles should be left over and that they should contribute much more to the energy density of the present Universe than is allowed by observation. This monopole problem motivated the development of the inflationary cosmology. If the Universe did undergo a period of inflation subsequent to the grand unified period, there might be very low numbers of relic monopoles. Anything in between is also possible.

f) Another implication of the simplest grand unified theories is that they often predict the b quark mass in terms of the observed τ^- mass. However, the models which correctly predict m_b almost always have trouble with similar predictions for the d and s masses. However, small perturbations on such theories could solve the latter problem without significantly affecting the m_b prediction.

g) Finally, some grand unified theories (and even some analogous theories which are not truly grand unified) allow the oscillation of neutrons into anti-neutrons. The interactions which mediate such neutron to anti-neutron oscillations are generally driven by extra Higgs bosons in the theory.

In some ways then, the simplest grand unified theories are probably too naive. They have some successes, but there are many problems that they too leave unanswered. It is likely that some aspects of grand unification are correct, but they are almost certainly not the final story.

SUPERSYMMETRY (SUSY)⁶

Supersymmetry, which can be implemented either with or without grand unification, is motivated by the Higgs problem and by the graviton problem. It involves a new type of symmetry in which fermions are related to bosons. In order to solve the Higgs problem, this new symmetry must be broken at a scale not too much larger than the weak scale - that is between 10^2 and 10^4 GeV.

The primary motivation for supersymmetry is the Higgs problem. To generate the weak interaction scale by the Higgs mechanism, one needs to introduce a scalar field, the quantum of which (the Higgs boson) has a mass not too much larger than the weak scale. However, loop diagrams associated with Higgs fields, gauge fields, or fermions yield quadratically divergent renormalizations of the Higgs mass. One finds $\delta m_H^2 = 0 (\lambda, g^2, h^2) \Lambda^2$, where the cutoff Λ is interpreted as the Planck scale $G_N^{-1/2} \approx 10^{19}$ GeV or as the grand unification scale $M_X \approx 10^{14}$ GeV.

In the simplest type of theory with exact supersymmetry, fermions and bosons occur in degenerate pairs. For example, the spin-1/2 quark

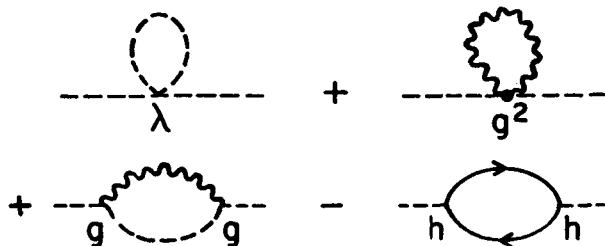


Fig. 9. Diagrams contributing to the renormalization of the Higgs mass.

there is a cancellation between the boson loops and the fermion loops, which occur with opposite sign, and the renormalization of the Higgs mass becomes zero.

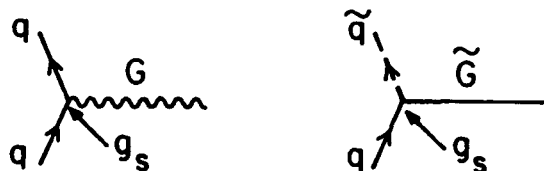


Fig. 10. Ordinary and new vertices.

finite though nonzero, yielding $\delta m_H^2 = 0(\lambda, g^2, h^2) |m_B^2 - m_F^2|$. Therefore, as long as the splitting between bosons and fermions is less than about 10 TeV, the Higgs problem is at least partially solved. There will be no unacceptable large renormalization of the Higgs mass and no large fine tunings will be necessary. I should emphasize that the origin of the weak scale and why it is so much smaller than the Planck and GUT scales is not explained in the simplest supersymmetric models. It is merely stabilized with respect to higher order corrections.

(q) has a partner, a spin-0 scalar-quark ($\tilde{q}$), and the spin-1 gluon (G) has as a partner a new spin-1/2 particle called the gluino ($\tilde{G}$). In addition, there are relations between the couplings of the particles and their new superpartners: the qqG vertex with strength g_s therefore implies the existence of a $qq\tilde{G}$ vertex with the same strength. If the supersymmetry is exact, then

The real world is not exactly supersymmetric. We do not have bosons and fermions in degenerate pairs. However, if there is an approximate supersymmetry which is broken by mass terms (and/or by cubic scalar interactions) then the renormalization effects remain

IMPLICATIONS OF SUPERSYMMETRY

The major prediction of supersymmetric theories is the existence of a large number of new particles. For various reasons, no known fermion-boson pairs can be identified as superpartners of each other. Therefore, it is necessary to assume a doubling of the spectrum; for every known fermion one must introduce a new boson and vice versa, as shown in Table 3. The new particles, denoted with a tilde, are generally given funny names. For example, associated with the W and Z there are new fermionic partners known as the wino and zino, respectively. Similarly, the photon and gluon have partners known as the photino and the gluino. The quarks and leptons have new spin-0 partners, known as scalar-quarks (or squarks) and scalar-leptons (or sleptons), and the Higgs (there must be at least two Higgs doublets in

$J = 1$	$J = 1/2$	$J = 0$
$W^\pm$	$\tilde{w}^\pm$ (wino)	
Z	$\tilde{z}$ (zino)	
γ	$\tilde{\gamma}$ (photino)	
G (gluon)	$\tilde{G}$ (gluino)	
	q_L, q_R (quarks)	$\tilde{q}_L, \tilde{q}_R$ (squarks)
	ℓ_L, ℓ_R (leptons)	$\tilde{\ell}_L, \tilde{\ell}_R$ (sleptons)
	$\tilde{H}_1, \tilde{H}_2$ (Higgsinos)	H_1, H_2 (Higgs)

Table 3. The low energy spectrum in the minimal supersymmetric model.

a supersymmetric model) have fermionic partners known as Higgsinos. The masses of these particles are very model dependent. All that we really know is that the masses should be larger than existing limits and less than 1-10 TeV. The photino, the zino, and the neutral Higgsinos are fermions which interact very weakly - very much like neutrinos. The signature of supersymmetric theories is therefore unobserved momentum. This is associated with the escape from the detector of the lightest supersymmetric particle (e.g. the photino), which is typically produced in the decays of heavier superpartners such as the gluino or scalar quarks. The existing limits from the nonobservation of such missing momentum events at the CERN collider are $m_{\tilde{q}}, m_{\tilde{G}} \geq 40$ GeV. The limits on the charged scalar leptons, the winos, and the charged Higgsinos are around 20 GeV because they have not been observed at the e^+e^- colliders PEP or PETRA. The limits on the photino, zino, and neutral Higgsinos are very weak; they could easily be as light as a GeV or so. In addition to the limits from direct searches at accelerators, there are a number of indirect limits on the new particles due to the fact that higher order diagrams involving these particles could lead to CP violation, to large parity violating effects, or to rare decays, such as $\mu \rightarrow e\gamma$.

SUSY - GUTs

One can also consider supersymmetric grand unified theories. In such theories one has supersymmetric partners not only for the ordinary light particles, but also for the superheavy gauge and Higgs bosons introduced in GUTs. An important feature of supersymmetric GUTs are that all of the new light superparticles, that is the scalar quarks, scalar leptons, Higgsinos, gauginos, etc., modify the theoretical predictions for the momentum dependence of the effective coupling constants and therefore for $\sin^2\theta_W$ and M_X . The $\sin^2\theta_W$ is about .02 larger than in standard GUTs, but this is still probably within an experimentally acceptable range. The unification scale is about 30 times larger than in the standard grand unified theory. The proton

	Ordinary Gut	SUSY-GUT	Experiment
$M_X(\text{GeV})$	1.3×10^{14}	5×10^{15}	-
$\sin^2 \theta_W$	~ 0.22	~ 0.24	$0.223 \pm 0.004 \pm 0.005$
$\tau_{p \rightarrow e^+ \pi^0}(\text{yr})$	$6.6 \times 10^{28 \pm 2}$	10^{35}	$\geq 3 \times 10^{32}$

Table 4. Predictions for M_X , $\sin^2 \theta_W$, and $\tau_{p \rightarrow e^+ \pi^0}$ (via X, Y exchange) in ordinary and supersymmetric GUTs.

lifetime due to ordinary heavy gauge boson exchange is therefore increased by a factor of approximately 10^6 to an observable rate of about 10^{35} yr.

However, there are new proton decay mechanisms (so-called dimension-5 operators) which can lead to much more rapid proton decays. For example, a typical diagram involves the exchange of the fermionic partner of a superheavy color Higgs scalar. This diagram leads to a proton lifetime that is proportional to the square of the unification scale and not to the fourth power as in traditional GUTs. There are many suppression factors, but nevertheless in many supersymmetric GUTs the proton lifetime is predicted to be in the range 10^{26} to 10^{31} yrs. The preferred decay mode in most such models is $p \rightarrow \bar{\nu} K$, rather

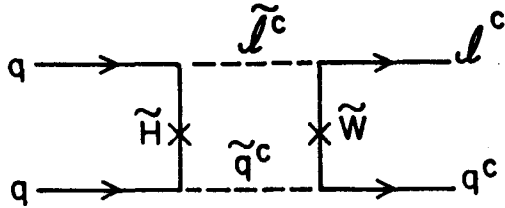


Fig. 11. A diagram leading to a dimension-5 operator.

than into $e^+ \pi^0$. Observational limits exclude proton lifetimes into this mode in this range, so many supersymmetric GUTs are in trouble. However, there are many unknown masses and parameters in the prediction, so one should keep pushing the proton decay experiments into these and other modes as far as possible.

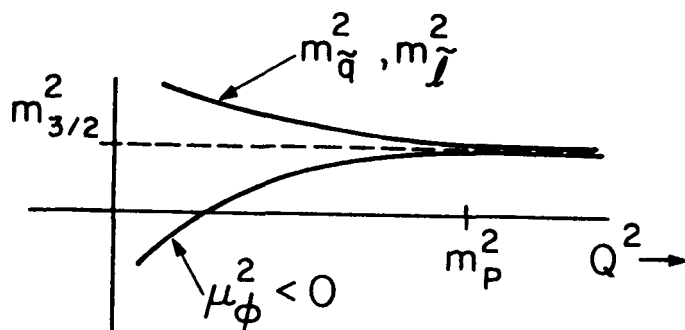
SUPERGRAVITY

So far I have been describing global supersymmetry in which the symmetry transformations between fermions and bosons are the same at all points in space and time. In complete analogy to gauge theories, in which symmetry transformations are allowed to take place independently at different space-time points, one can also consider local (gauge) supersymmetry in which the fermion-boson transformations can be performed independently at different space-time points. Such theories, known as supergravity theories, necessarily require the introduction of a spin-3/2 particle, called the gravitino, which is analogous to the spin-1 boson that is needed for a local gauge symmetry. In addition, the gravitino has a spin-2 partner, the graviton, which generates gravity. Supergravity therefore incorporates gravity in a

non-trivial way. However, it does not solve the problem of quantum gravity; supergravity theory is still non-renormalizable. The spontaneous breaking of supergravity generates a mass for the gravitino in much the same way that gauge bosons acquire a mass when an ordinary gauge symmetry is spontaneously broken.

SUSY BREAKING

We have seen that supersymmetry must be broken, with the typical breaking scale somewhere in the range from 10^2 to 10^4 GeV. The simplest possibility would be to simply introduce some symmetry breaking mechanism in this mass range. However, for a variety of reasons all such schemes fail. The so-called hidden sector models lead to an acceptable theory, but they are rather complicated. In these theories one introduces a set of particles which are almost totally decoupled from the ordinary particles and their superpartners. Supersymmetry breaking presumably occurs in this hidden sector and is transmitted only indirectly via weak couplings to the ordinary particles and their superpartners. This weak coupling may be via gravity or it may be via very high order gauge interactions. If the scale of symmetry breaking is F in the hidden sector, then the result is that the gravitino and all scalars in the ordinary sector acquire a mass of order $m_{3/2} \sim F^2/m_P \ll F$. Since the Higgs scalar has a positive mass² of this value, $SU_2 \times U_1$ is not spontaneously broken at tree level. However, large Yukawa couplings associated with the top quark mass may lead to $SU_2 \times U_1$ breaking via radiative corrections in such



models. The upshot is that the electroweak and effective supersymmetry breaking scales are naturally of the same order of magnitude, namely the gravitino mass. Since this magnitude must be around 10^2 to 10^4 GeV, one has that the hidden sector SUSY breaking₁₁ scale is of order 10^{11} GeV.

Fig. 12. Running scalar masses, leading to radiative $SU_2 \times U_1$ breaking in hidden sector models.

SUSY SUMMARY

- a) Supersymmetry is a very elegant idea.
- b) Supersymmetry helps to solve the Higgs Problem. It stabilizes the electroweak scale against unacceptably large radiative corrections as long as the effective breaking scale is less than about 10 TeV.
- c) Supersymmetry does not help to solve the fermion problem.
- d) Gravity is incorporated into supergravity theories, but quantum

gravity is still nonrenormalizable.

- e) There is not a shred of experimental evidence for supersymmetry. However, since the breaking scale cannot be arbitrarily large (and still solve the Higgs problem) it should be possible to produce and detect supersymmetric particles at new high energy supercolliders.
- f) Phenomenologically viable supersymmetric models are extremely complicated.

SUPERSTRINGS⁷

Superstrings are a very exciting development which may yield finite theories of all interactions with no free parameters. Superstring theories introduce new structure above the Planck scale. The basic idea is that one changes the fundamental entities in a superstring theory. Instead of working with (zero-dimensional) point particles as the basic quantities, one considers one-dimensional objects known as strings, which may be open or closed. The quantized vibrations of these strings lead to an infinite set of states, the spectrum of which is controlled by the string tension.

Strings were introduced in the early 1970's as a model of the strong interactions and of the hadronic states. However, the strong interaction strings never quite worked out, and since then they have been replaced by QCD as the most likely theory of the strong interactions. More recently, strings have re-emerged with the new interpretation that they are a theory of everything, not just the strong interactions. The string tension is now given by the square of the Planck scale $T \sim m_p^2$. When one probes or observes a string at energies much less

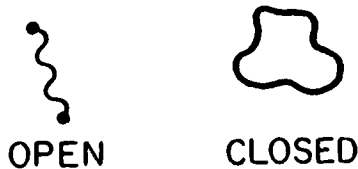


Fig. 13. Open and closed strings.

than the Planck scale, one sees only the "massless" modes and these represent ordinary particles. The physical size of the string is given by the inverse of the Planck scale $\approx 10^{-33}\text{cm}$, and at larger scales a string looks like a point particle.

There are various types of string theories. The most promising is the heterotic string, which is a theory with closed strings only in which the gauge symmetry charge densities are distributed along the length of the string. The basic interaction in the closed string theory involves the joining of two strings to form one or the breaking of one closed string into two. This should be compared to the three point vertex in ordinary field theory. The various string configurations correspond to ordinary matter particles such as fermions or their scalar particles, to gauge bosons and their partners, and to the graviton and the gravitino. String theories therefore incorporate both gauge interactions and gravity - they are unified. The enormously exciting development is that string theories apparently lead to completely finite theories of quantum gravity and of the other interactions; not just renormalizable but completely finite. No known ordinary field theory leads to a mathematically consistent, renormalizable theory of quantum gravity, let alone to a finite theory.

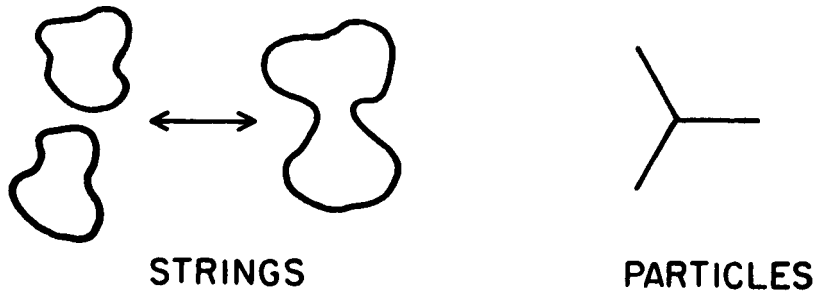


Fig. 14. Interactions in string and field theories.

An interpretation is that the string represents an infinite number of physical states and in field theory language one has an infinite set of diagrams and exchanges involving this hierarchy of states. The couplings and masses of these states are constrained in such a way that the divergences cancel, leading to a completely finite theory.

There are two important constraints on the string theory. The first is that the string theories are only consistent in ten dimensions of space and time, that is nine space and one time dimension. In different dimensionalities, one has unacceptable negative norm (ghost) states. The fact that we apparently live in four and not ten dimensions can be accounted for by the assumption that six of the dimensions compactify into a small manifold. That is, six of the dimensions curl up to a small closed manifold of radius $r_{\text{comp}} = 1/m_{\text{comp}} \approx 10^{-32 \pm 1} \text{ cm}$. This idea of compactification was borrowed from Kaluza-Klein theories, which are an alternative class of models that also gave some promise of unifying gravity (but which have run into serious problems). Another constraint is that string theories tend to be plagued with anomalies, which are pathologies or inconsistencies which can occur in a theory, leading to breaking of gauge and other symmetries that are needed for the consistency of the theory. In four-dimensional field theories, anomalies are generated by certain triangle diagrams, and one generally insists that the fermion content of a theory be such that the anomalies cancel. In a ten-dimensional string theory there are hexagon-type diagrams which lead to unacceptable anomalies. However, the anomalies cancel if the underlying gauge group is either $G = E_8 \times E_8$ or SO_{32} . The SO_{32} theories are not physically realistic, so the underlying gauge symmetry of a superstring theory is constrained to be the group $E_8 \times E_8$.

The scenario that one expects is the following: At high energies (large compared to the string tension or Planck scale) the theory of Nature would be an $E_8 \times E_8$ ten-dimensional string theory. However, six of the dimensions compactify into a small manifold at the compactification scale $m_{\text{comp}} \sim 10^{18 \pm 1} \text{ GeV}$. At energies much lower than m_{comp} one sees an effective four-dimensional particle theory. It is possible and expected that some of the gauge flux of the $E_8 \times E_8$ theory becomes trapped in the manifold so that one sees an effective gauge symmetry smaller than the original $E_8 \times E_8$ group. In order to have a chiral low energy supersymmetry below the compactification

scale, it is necessary that the extra six dimensions compactify into what is known as a Calabi-Yau manifold. In such theories the low energy gauge group is either $E_6 \times E_8$ or a subgroup of $E_6 \times E_8$. The E_8 factor corresponds to a hidden sector, while the E_6 factor (or the subgroup which survives the compactification) is associated with the ordinary particles. Because the E_6 is broken to a subgroup such as $SU_3 \times SU_2 \times U_1 \times G'$, the compactification scale plays the role of the GUT unification scale; compactification is a new method of spontaneous symmetry breaking that does not rely on an ordinary Higgs mechanism. Consistency of the entire approach requires $M_X \sim m_{\text{comp}} \sim 10^{18 \pm 1} \text{ GeV}$.

One of the most exciting features of the superstring is that essentially all low energy physics, that is all physics below the compactification scale, is in principle determined by the topology of the Calabi-Yau manifold. The nature of the low energy gauge group, the number of fermion families, and even the Yukawa couplings (and therefore the fermion masses and mixings) are in principle not free parameters but are determined by the topology of the compactification manifold. Unfortunately, there are many (10^4) possible Calabi-Yau manifolds. At present we are not capable of calculating which manifold is chosen or even why it is a Calabi-Yau manifold, why six dimensions are compactified, or what the compactification scale is. Therefore although everything is determined in principle, we are far from able to do the calculations.

CONSEQUENCES OF STRINGS

a) Proton decay is probably unobservable in superstring theories. This is because the basic unification scale is typically of order 10^{18} GeV so that gauge boson mediated decays are unobservably slow. The necessary Yukawa interactions can in principle be calculated in terms of topological considerations, but we are unable to do that at present. Models typically have proton decays via Higgs that are either much too fast, in which case those models are excluded experimentally, or much too slow to observe. However, it is always possible that in some models proton decay may occur at an observable rate. One should keep looking.

b) Baryon Asymmetry. It is simply too early to make meaningful statements about such cosmological implications as the baryon asymmetry.

c) Most Calabi-Yau manifolds imply the existence of extra Z bosons (i.e., the symmetry group below the GUT scale generally contains an extra factor G' . This is typically U_1 or $U_1 \times U_1$, although there are other possibilities. These may be very light; e.g., in the range from 10^2 to 10^3 GeV . Present experimental limits are very weak, typically allowing extra Z 's as light as 100 to 200 GeV. One should certainly pursue the search for such bosons vigorously.

d) Matter fields. Most manifolds imply the existence of either three or four families of fermions and their scalar partners. (The simplest manifolds imply far too many families and are therefore excluded.) The fermions and their partners occur in 27-dimensional representations of the E_6 group. Therefore, it is predicted that for

each family there should be the ordinary fermions $\begin{pmatrix} u \\ d \end{pmatrix}_L$ $\bar{u}_L$ $\bar{d}_L$ $\begin{pmatrix} \nu_e \\ e \end{pmatrix}_L$ e_L^+

(and their CPT conjugates) as well as a class of exotic new fermions

$D_L, \bar{D}_L$ $\begin{pmatrix} E^+ \\ E^0 \end{pmatrix}_L$ $\begin{pmatrix} \bar{E}^0 \\ E^- \end{pmatrix}_L$ N_{1L}, N_{2L} . These exotic new fermions and their

scalar partners may be in the hundred GeV range, or possibly much more massive.

e) $\sin^2 \theta_W$. The predictions for the weak angle in superstring theories are similar to those in other supersymmetric GUTs, except that one must take the exotic particles into consideration. The correct value $\sin^2 \theta_W$ is therefore a constraint on model building and on assumptions as to whether there are intermediate thresholds in the desert, at which the extra Z's or exotic fermions could possibly acquire masses.

f) Neutrino mass. The neutrino mass turns out to be too large in most superstring theories. It is typically comparable to the charged lepton mass. There are possible ways of evading this problem, but nothing particularly elegant or attractive has emerged so far.

g) SUSY breaking. The origin of supersymmetry breaking in the superstring model is somewhat obscure. It is thought that a gluino condensate in the hidden E_8 sector may lead to the original supersymmetry breaking. However, the details of how this is transmitted to the observed sector are somewhat obscure and it seems that the superstring models do not directly imitate the superstring breaking schemes that have been previously worked out for non-string theories.

STRING SUMMARY

a) Superstrings are very elegant.

b) They yield a completely finite theory of gravity and of all other interactions with no free parameters; therefore, the entire structure of physics below the compactification scale is, in principle, determined.

c) However, the details of the low energy physics depend on the topology of the manifold. We are not, at present, able to calculate what the manifold is, the number of compactified dimensions, the compactification scale, etc. Therefore we are unable at present to predict the structure of the low energy world.

d) If the superstring idea is correct, it is likely that superstring models will incorporate some, but by no means all, aspects of the standard supergravity-GUT picture.

SUMMARY

a) The standard model of the strong, weak, and electromagnetic interactions is almost certainly approximately correct up to around 100 GeV.

b) However, the standard model with classical gravity involves at least 21 free parameters and is plagued by many unexplained features,

such as the gauge, fermion, Higgs/hierarchy, strong CP, and graviton problems. Therefore new underlying physics is necessary.

C) Grand Unified Theories solve some, but not all of the problems of the standard model. They elegantly explain the gauge problem and charge quantization. However, proton decay limits exclude the simplest GUTs. Some aspects of grand unification are probably true, but it is unlikely that a simple grand unified theory is the final story either.

d) Supersymmetry and supergravity models stabilize the weak scale and therefore solve much of the Higgs problem, but they do not in general explain it in the first place. Supergravity theories incorporate gravity, but they are still nonrenormalizable. Supersymmetric models do not help with the fermion problem.

e) Superstring models are conceivably the ultimate theory of everything. They yield finite theories of quantum gravity, and in principle they lead to a completely determined theory, with no arbitrariness or free parameters. However, we are very far from being able to do the calculations needed to test superstring theories or to predict the structure of the low energy world. If superstrings are correct, they probably incorporate some of the aspects of the supergravity-GUT picture, but probably not all. I don't want to give the impression that the superstring theories are obviously correct. It may well be that we are on the wrong track. They are very exciting and promising, but whether Nature chooses to use these elegant theories remains to be seen.

f) There are many other possible extensions of the standard model, such as Kaluza-Klein theories or models in which fermions, Higgs bosons, or gauge bosons are composites of still more elementary particles. It is not clear what direction physics beyond the standard model will ultimately take. In the end it will be an experimental question. In the future we expect to probe physics beyond the standard model with a new generation of accelerators and with new proton decay experiments, neutrino mass experiments, monopole searches, etc. The future promises to be very exciting indeed.

ACKNOWLEDGMENT

It is a pleasure to thank Franz Gross and the organizers of this meeting for their invitation and travel support.

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CURRENT TECHNOLOGY OF PARTICLE PHYSICS DETECTORS

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CEBAF 1986 Summer Workshop
Newport News, VA

June 23, 1986

This work performed under the auspices of the
U.S. Department of Energy

CURRENT TECHNOLOGY OF PARTICLE PHYSICS DETECTORS*

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I. New Facilities for New Physics

Detection of particles is the experimental tool for both nuclear and elementary particle physics. In particle physics, more so than in nuclear, the evolution of the field has been closely coupled to the development of improved techniques for detecting and measuring an ever-widening spectrum of particle properties. The continuing quest for higher collision energies, and correspondingly rarer and more complex phenomena, has led to an exponential growth in the capabilities of accelerators and detectors, and a field in which much of the research is concentrated in a relatively small number of big detector systems at large accelerator facilities. At present these facilities are epitomized by a number of colliding beam accelerators which are at various stages from planning to actual operation. These facilities and some of their characteristic parameters are listed in Table I. Not all of HEP is done at colliders, but these facilities are currently most influential in driving the technologies of particle detectors. Note that one of the facilities listed, RHIC, is a nuclear physics initiative, albeit with detector requirements similar in many ways to its HEP counterparts. At each of these facilities the experimental program involves several hundred physicists working with a relatively small number of large detector systems. These large detectors, ranging in cost from \$20 M to \$100 M, are actually integrated systems of many different detector technologies and components.

This research supported in part by the U.S. Department of Energy under Contract DE-AC02-76CH00016.

Table I.
High Energy Colliding Beam Facilities

<u>Accelerator</u>	SPS CERN	Tevatron Fermilab	TRISTAN KEK	SLC SLAC	LEP CERN	HERA DESY	RHIC BNL	SSC
<u>First Operation</u>	1981	1986	1986	1987	1989	1990	1992*	1994*
<u>Beams</u>	$\bar{p}p$	$\bar{p}p$	e^+e^-	e^+e^-	e^+e^-	ep	Au+Au	pp
<u>No. Interaction Regions</u>	2	2 large 2 small	4	1	4	3	6 (4 initially)	4
<u>Maximum beam Energy (GeV)</u>	400	1000	30	50	60	28/820 (e/p)	20,000 (100/nucleon)	20,000
<u>Luminosity ($\text{cm}^{-2} \text{sec}^{-1}$)</u>	1×10^{30}	1×10^{30}	2×10^{31}	1×10^{30}	1.6×10^{31}	1.5×10^{31}	1×10^{31} (protons) 4×10^{26} (Au)	1×10^{33}
<u>Event Rate (sec^{-1})</u>	5×10^4	5×10^4	20	1	20	10^3	10^3 (Au)	10^8

*Proposed dates; facilities not yet funded.

The operational meaning of the term "particle detection" has changed considerably over the past two decades as the available collision energy, and hence the resolving power of experiments has gone up several orders of magnitude. Not long ago the "elementary" particles of interest were the entries in an ever-longer list of leptons and hadrons. These were produced in few-body final states (i.e., less than 10) in which precise momentum measurements and some selection among pions, kaons and protons were the basic requirements for sorting out resonances, often in exclusive channels. Experiments with bubble chambers or straight-forward magnetic spectrometers were typical.

The discovery of charm, the understanding of quark confinement in the QCD framework and the development of what is now called the standard model of electro-weak interactions has brought great changes in the character of particle detection. The characteristic lifetimes of charmed mesons and baryons, with mean path lengths of a few hundred micro-meters, has spurred the development of tracking chambers with position resolution in the 5-50 μm range. The importance of jets has led to highly sophisticated systems of well-segmented calorimeters. The special role of leptons, the need to isolate single high energy electrons, muons and even neutrinos in an environment of high track density has had a significant influence on the development of track chambers (pattern recognition with many, closely spaced tracks; the need to cover very large areas), calorimetry (precise missing-energy measurements in "hermetic" detectors), and particle identification (development of transition radiation detectors and ring-imaging Cerenkov detectors, for example).

In the quest for a plasma of quarks and gluons with high energy collisions of heavy nuclei, the need to contend with extremely high track densities is not confined to small solid angles and rare events, but is the normal state of affairs in an environment where typical events may have many thousands of final-state particles. In this case it may be argued that with so many particles it needn't be essential to measure each one individually, but rather to measure

local fluxes of energy, as with calorimeters, treating the whole as a classical wave of hadrons. Still, it remains necessary to resolve certain special particles, such as leptons.

In the search for phenomena beyond the standard model, e.g., for Higgs particles and manifestations of supersymmetry, the detector systems must be capable of reconstructing the 4-vectors of quark and gluon jets and combining them to examine the effective mass of multi-jet systems. This is certainly a giant step beyond the traditional spectroscopy in which we search for resonance bumps in few-particle mass spectra.

Event rates play an important role in the technology of detectors. While the e^+e^- machines in Table I have relatively low interaction rates (a few per second, typically) the proposed SSC, with luminosity $10^{33} \text{ sec}^{-1} \text{ cm}^{-2}$, will have an event rate approaching 10^8 sec^{-1} .

The coupling between high energy and high luminosity is of course no accident. The purpose of high energy supercolliders is to extend upward the mass scale of particle states which mediate the observable phenomena. The higher the mass the smaller the cross section and hence the need for ever-higher interaction rates to accumulate statistically accurate samples. High rates and complex events are facts of life in future generations of high energy collider experiments.

II. Position Sensitive Detectors

The precise measurement of charged particle trajectories is a familiar requirement for momentum measurement in magnetic spectrometers. As noted above, today's detector systems often make special demands on tracking detectors, calling for extremes of multitrack resolution and/or rate capability, and in the special cases of high resolution "vertex detectors" we must seek the fundamental limits of position accuracy. These limits are set by thermal noise and the scattering inherent in the mechanism of signal deposition ^{1,2,3}.

The performance capabilities of several different techniques are summarized in Table II.

For many years multi-wire proportional chambers and drift chambers have been the work horses for both high energy and nuclear physics. Measurements of sub-millimeter accuracy can be made with very thin wire chambers covering large areas (many square meters), or volumes (cubic meters). An important feature of these detectors is their very high rate capability, which can exceed 10^6 particle hits per wire per second⁴.

Drift chamber principles have been applied to many ingenious detector geometries. The most important intrinsic limitations to the accuracy of these detectors are

- i) The relatively sparse distribution of ionization charge along the track ($\approx 300 \mu\text{m}$ between clusters of ionization electrons for a minimum ionizing track in atmospheric pressure gas), and
- ii) diffusion of the drifting electrons. The effect of diffusion in various gases is displayed in Fig. 1.

Typically, these effects impose limits on r.m.s. spatial resolution at the 150-200 μm level in large detectors. One of the most important avenues of detector development in recent years has been to extend the drift chamber concept to produce an electronic detector capable of giving track "images" with the visual quality typically associated with bubble and streamer chambers. The important difference between such truly visual detectors and the classical planar wire chamber is that in a bubble chamber, for example, each bubble gives a track point whose position is known in all three space dimensions, and there are many bubbles along each track, whereas in an array of wire chambers a measurement at each wire plane gives only the coordinate in the plane perpendicular to the wires, and the position along the wire is not known. In fact, methods have now been

Table II.

Position Sensitive Charged Particle Detectors

	Position Accuracy (μm)	Resolution Elements per cm^2	Maximum Instantaneous Rate per Element
Semiconductor			
Strip detector	5	200	10^6
CCD	10	10^5	10^3
Scintillating fiber	≥ 100	≤ 10	10^7
Drift chamber			
Typical	200	0.1	10^6
Limit	≤ 50	1	10^7

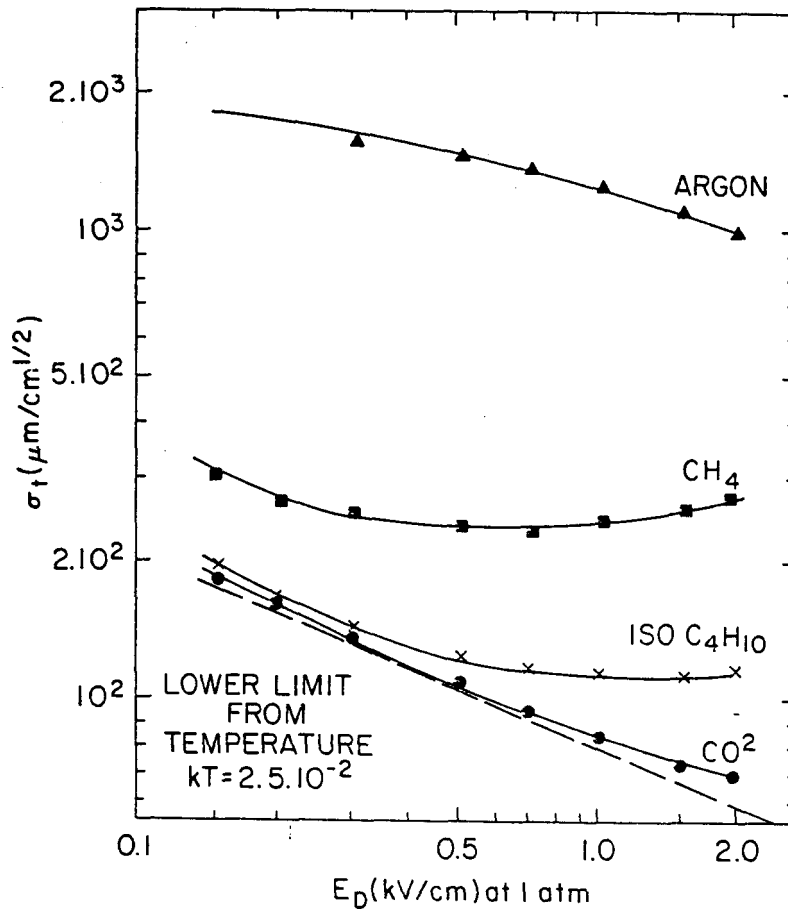


Fig. 1 Characteristic diffusion lengths, as function of electric field strength, for electrons in various gases at STP.

developed for localizing the position of the avalanche along the wire. One way of doing this employs resistive anode wires and read-out electronics to measure the pulse height at each end of every mode wire. The ratio of pulse heights at the two ends of a wire determines the position along the wire, with typical r.m.s. accuracy of about 1 cm. This technique is referred to as "charge division"⁵. Another approach is to measure the induced charge on segmented cathodes ("Pads") which are segmented along the wire direction. The centroid of the measured charge distribution on these pads locates the position along the anode wires and with a sufficiently large number of readouts this can be done with submillimeter accuracy⁶. Figure 2 illustrates the technique. Figure 3 shows an example of a large-volume "imaging" drift chamber⁷ and an event which it recorded. With such a device a volume of $\sim 10 \text{ m}^3$ is made fully sensitive with an effective cell size for track resolution of a few mm^3 and with a sufficiently small amount of material that tracks emerging from the detector have their trajectories virtually undisturbed.

Another direction being vigorously pursued in the development of improved tracking chambers involves efforts to significantly enhance the resolving power, both the position resolution and the ability to resolve pairs of nearby tracks. For wire chambers, specially designed electrode structures coupled with advances in electronics as well as advances in the knowledge of the physics of charge transport and collection in gases have led to a number of devices that approach the limits of resolution listed in Table II. An example of such techniques⁸ is shown in Fig. 4. Here the anode wire is shielded in such a way as to collect only the charge from a limited segment of the particle track, reducing the time-smearing effects of the geometric spread in arrival time as ionization charge drifts toward the wire. Also, the gas mixture used has CO_2 as its primary component: a gas for which electron diffusion is minimized, as seen in Fig. 1. Figure 5 illustrates a practical realization of such an electrode structure⁹, using wire electrodes at various potentials to achieve the channelling effect. This chamber has achieved r.m.s.

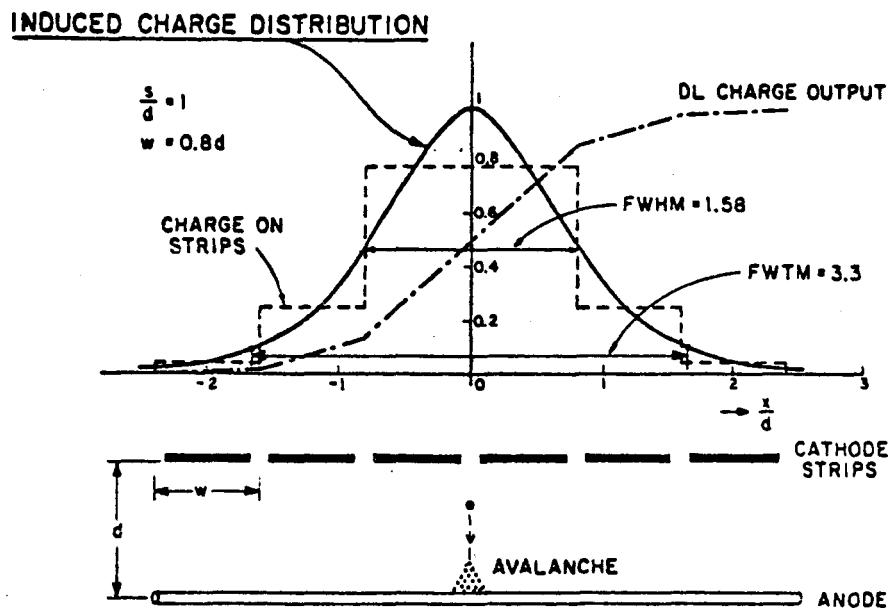


Fig. 2 Centroid finding by induced charge on a segmented cathode (Ref. 6).

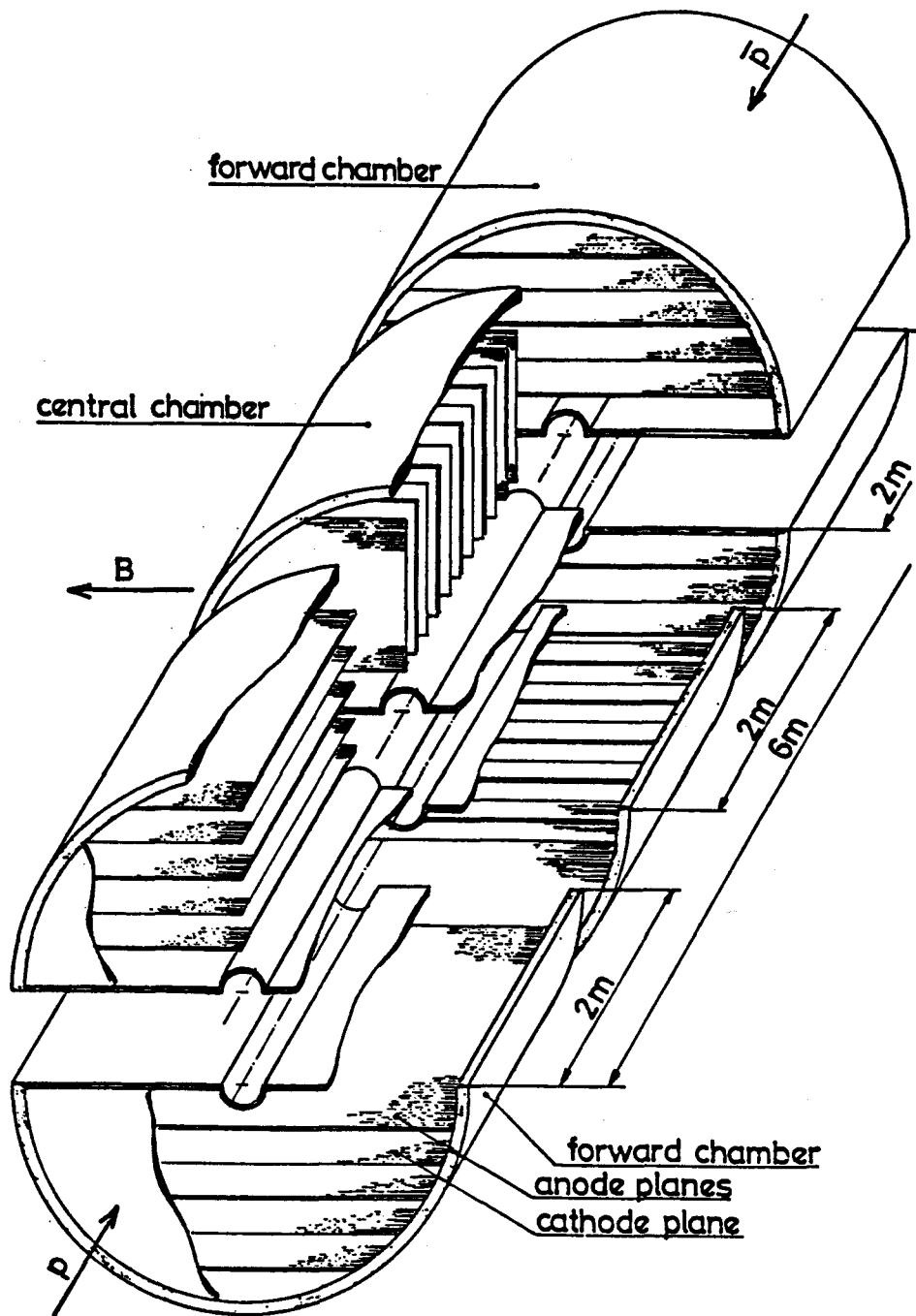


Fig. 3a The central track detector of the UA-1 dipole spectrometer (Ref. 7).

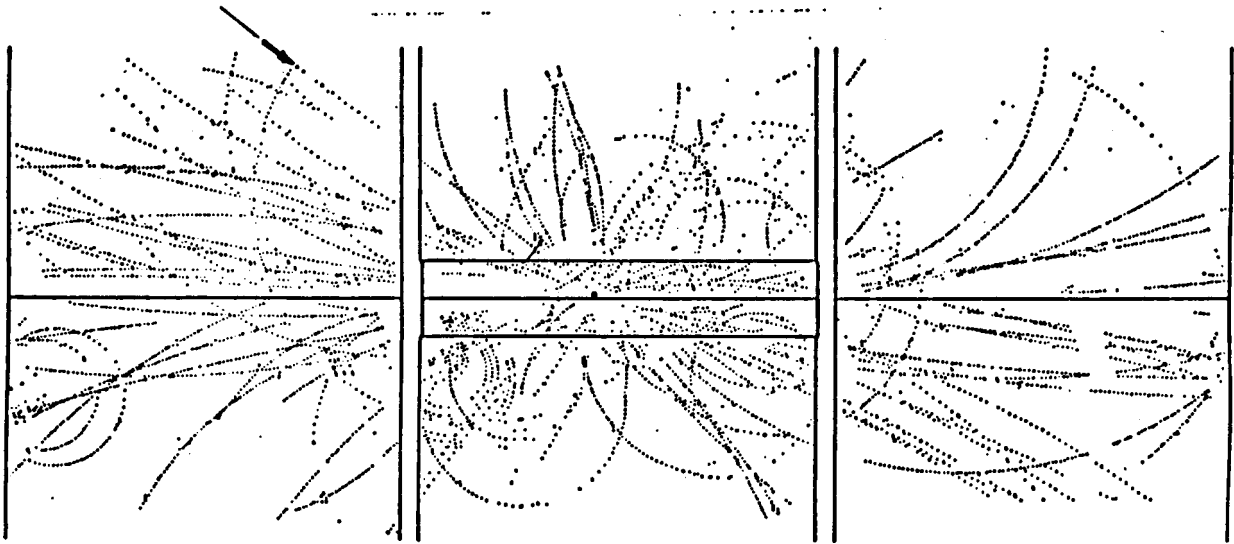


Fig. 3b A high energy $\bar{p}p$ event recorded in the UA-1 track detector.

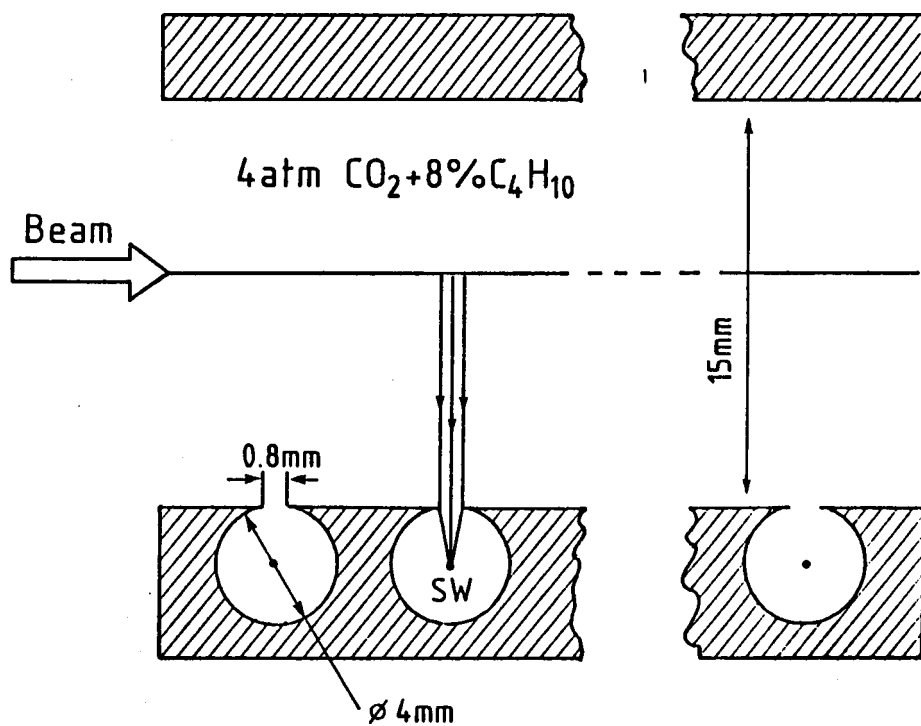
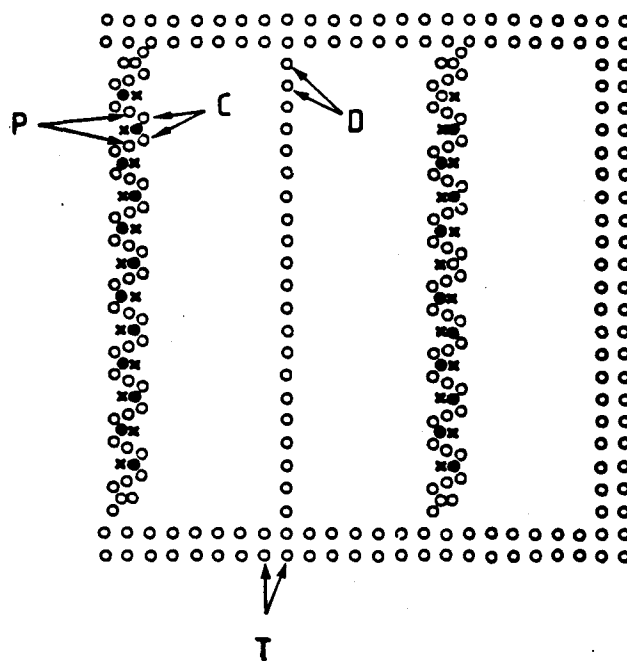


Fig. 4 Conceptual design for a "controlled geometry" anode structure (Ref. 8).



- × Guard wire (G)
- Anode wire
- Channelling wire (C)
- Potential wire (P)
- Electric field
- Termination wire (T)
- Drift field wire (D)

0 1 2 cm

Fig. 5 A practical realization of a controlled geometry anode structure (Ref. 9).

spatial resolution of $\leq 70 \mu\text{m}$, and double-track resolution $\sim 600 \mu\text{m}$. Some other ideas for achieving this kind of resolution in wire chambers can be found in Ref. 2.

Whereas wire chambers are pushed toward fundamental limits to achieve such resolution, a new class of detector has recently received broad attention for which the inherent spatial resolution is $\lesssim 10 \mu\text{m}$. The development of very high resolution position-sensitive semiconductor detectors^{10,11} has followed from years of experience in low energy nuclear physics experiments with unsegmented semiconductor detectors, plus the realization of large-scale integrated circuit techniques for putting many detector elements on a single silicon wafer.

The characteristic property of a single-crystal silicon semiconductor is the small separation between the conduction and valence band of electron energies. The mean energy deposited per electron-hole pair is 3.6 eV (cf. 26 eV per ion pair in argon gas). Figure 6 illustrates how a signal is developed as a charged particle traverses a silicon semiconductor detector element. For a relativistic (minimum-ionizing) particle the rate of energy loss is $\approx 3 \text{ MeV/cm}$, corresponding to 80 electron-hole pairs per micron of track length in silicon. This charge drifts in the applied field E , inducing a signal charge on the collection electrodes; there is no avalanche gain. Electrons and holes drift with nearly the same velocity; the hole mobility is the smaller and gives typical drift velocity of 1 cm/sec. Thus, as shown in Fig. 6, collection times can be $\lesssim 100 \text{ nsec}$ for detector thicknesses $\lesssim 1 \text{ millimeter}$.

The collection field E cannot be established with simple ohmic contacts between the silicon crystal and electrodes, as the finite resistivity of the semiconductor would result in a current flow ("leakage current") whose fluctuations would overwhelm the signal. Hence a diode structure is used. An n-type region with an excess of

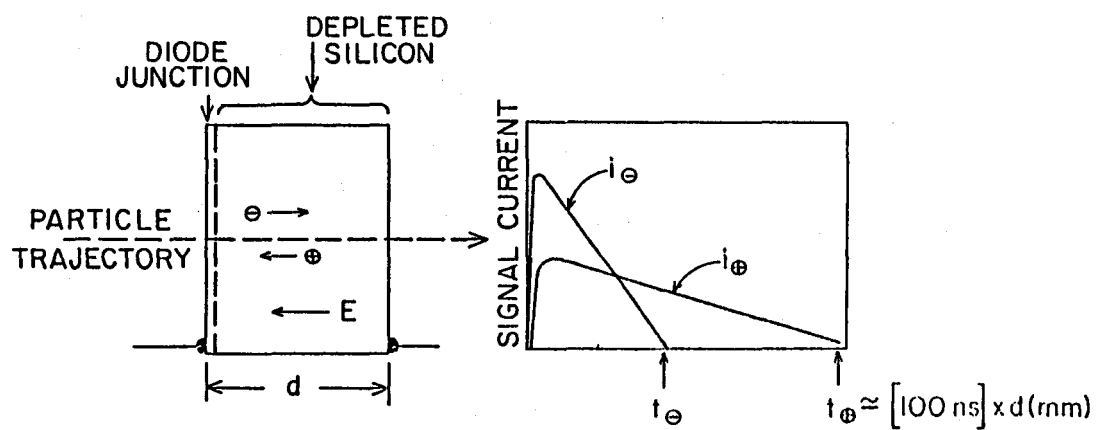


Fig. 6 Signal development in a semi-conductor diode detector.

free electrons is joined to a p-type region which has more holes. Near the junction the diffusion of charge establishes a region which is depleted of free charge carriers, blocking the current. The thickness of this depletion layer (X_D) is determined by the bias voltage (V) and the resistivity (ρ) of the semiconductor material. For silicon,

$$X_D \text{ (mm)} \approx 0.5\sqrt{\rho V}$$

with ρ in $K\Omega\text{-cm}$ and V in kV . The entire sensitive volume of a detector must be depleted. To achieve depletion layers of a few hundred microns thickness requires extremely pure silicon: typically a resistivity $\gtrsim 10 K\Omega\text{-cm}$ is required, corresponding to impurity levels of $\lesssim 10^{13}$ atoms per cm^3 (about 1 part in 10^{10} !).

The construction of a silicon microstrip detector is illustrated in Fig. 7a. Its application and readout geometry is much like that of a tiny wire chamber: a silicon wafer is subdivided into many strip electrodes, each of which acts as a diode detector element as described above. Space resolution figures $\sim 10\mu\text{m}$ are achieved with such detectors and, because of the fine granularity, high track densities can be tolerated. A serious limitation of these devices is that, in the simplest case, every strip must have its own chain of readout electronics. In some cases it is possible to interpolate among groups of strips, nevertheless it is generally true that a few square centimeters of detector may require square meters of fan-outs and amplifier cards for the readout. Several methods are now being explored to overcome this problem (see Radeka¹⁰). One approach is to integrate much of the front-end amplification and readout logic onto the detector chip; another is to develop charge coupled devices (CCD), such as are now common in video applications, as charged particle detectors. This approach is illustrated schematically in Fig. 7b. Figure 7c shows the principle of a "semiconductor drift chamber."

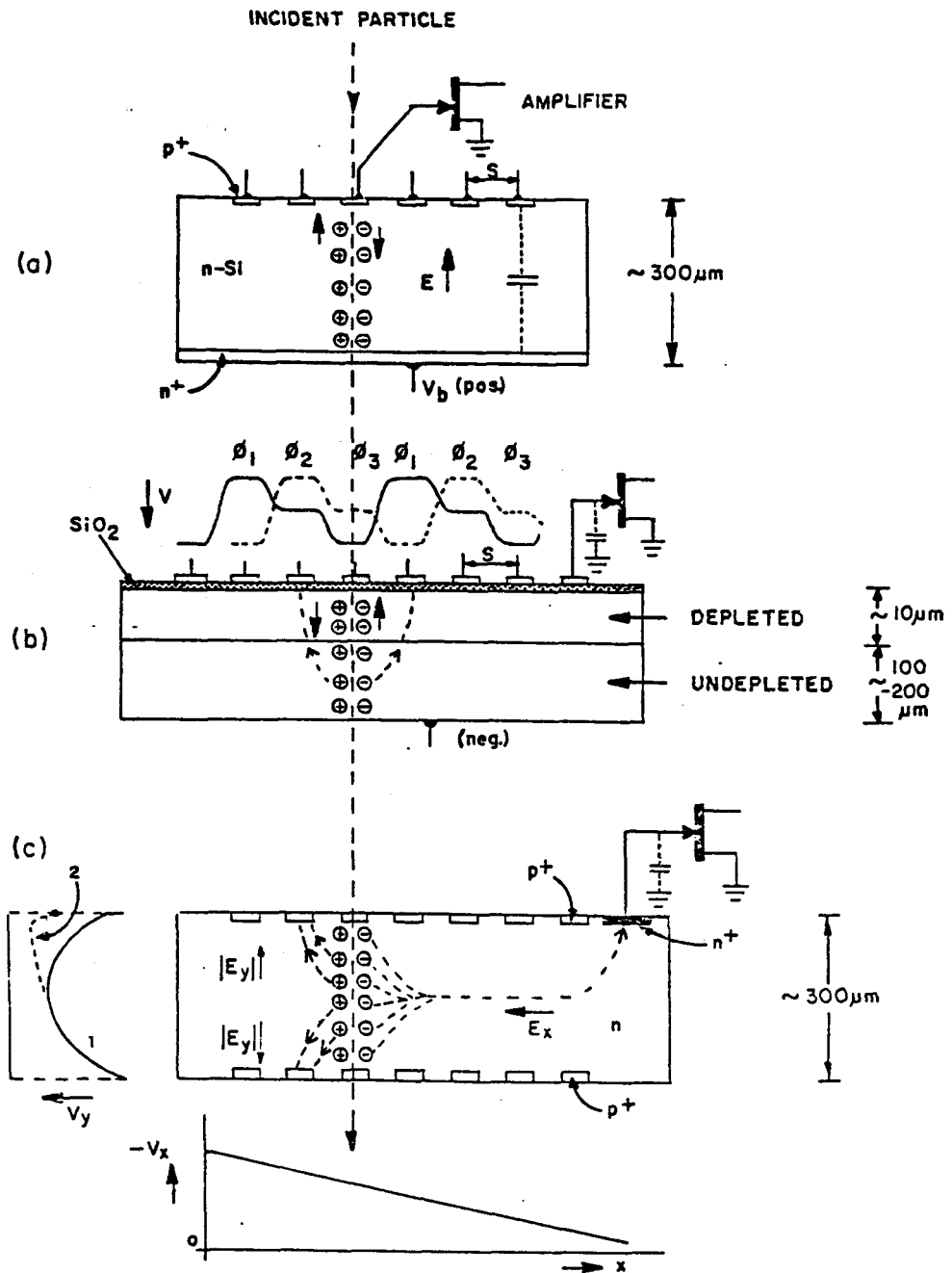


Fig. 7 Position-sensitive semi-conductor detectors¹⁰.
 (a) Microstrip detector with charge collection and measurement on individual strip electrodes.
 (b) Charge transfer device (CCD).
 (c) The concept of a semi-conductor drift chamber, in which a linearly varying potential in the x direction causes electrons to drift to the midplane and then to the positively biased n⁺ electrode.

Not surprisingly, the more one pushes toward very high resolution in position-sensitive detectors, the more difficult it becomes to cover large areas or volumes with such a detector. High resolution drift chambers in general require very fine, closely spaced wire electrodes with very stringent mechanical tolerances, and are difficult to construct if the wire lengths, for example, exceed a few tens of centimeters. Silicon detectors are limited by wafer technology to linear dimensions of only a few cm. The high density of readout makes the cost per unit area very high for such detectors.

III. Particle Identification

Traditionally, particle identification has referred to a determination of the particle mass and has applied to discrimination among pions, kaons, and protons. While this remains important, the identification of other types of particles, e.g., leptons, photons, and heavy flavor particles with characteristic lifetimes in the range 10^{-10} to 10^{-13} sec, and the identification of particles through the invariant mass of their decay products has become increasingly important. In the new physics of supercolliders the measurement of jet-jet invariant masses is expected to be a means of isolating Higgs bosons and other very massive particle states¹².

Among the more classical means of particle identification, those which involve the identification of photons are currently receiving a great deal of attention for high energy physics detectors. Highly segmented, high resolution detectors of single photons in the uv and x-ray range, capable of covering large areas, are being developed for ring-imaging Cerenkov detectors and the measurement of transition radiation.

The ring-imaging Cerenkov technique has been a long time in development, and is just now being realized in practical detectors

covering large solid angles. The concept is illustrated in Fig. 8. Here the cone of Cerenkov light emitted from a charged track is focussed into a ring by a spherical mirror. The diameter of the ring image is related to the Cerenkov angle, and hence to the velocity. The center of the ring gives the track position. The essential development to make such a detector feasible is a large area, highly segmented photon detector with sufficient sensitivity to produce many photons in a ring even for rather short lengths of radiator. Special uv sensitive wire chambers have been successfully developed for this purpose, an example of which¹³ is seen in Fig. 8. These detectors are made possible by the discovery of organic compounds with very low photoionization potential which can be added to the normal drift chamber gas to produce a high quantum efficiency for uv photons, in the range $E < 10$ eV, which are transmitted through calcium fluoride or fused quartz windows.

Recent developments in transition radiation detectors have improved both the sensitivity and the multitrack capability of these devices. These detectors are primarily used to discriminate electrons from hadrons. In the phenomenon of transition radiation, a highly relativistic particle moving in the discontinuous dielectric medium emits x-ray photons whose intensity is proportional to the Lorentz factor γ . A measurement of this radiation offers a means of particle identification in the ultra-relativistic regime ($\beta\gamma > 500$) where methods sensitive to the particle velocity are not useful. Many detailed studies have been carried out to optimize the radiator for such measurements. One such optimization¹⁴ resulted in a radiator consisting of ~1000 layers of lithium foil, each 50 μm thick, totaling about 2% of a radiation length. The radiator does not always consist of foils: volumes of closely-packed bits of carbon fiber have been successfully used as a radiator¹⁵, and means of creating low-density "plastic foams" have been considered. Figure 9 illustrates a configuration of radiator plus photon detector currently being used¹⁶. Here the wire chamber and its high speed readout electronics are designed to resolve individual clusters of

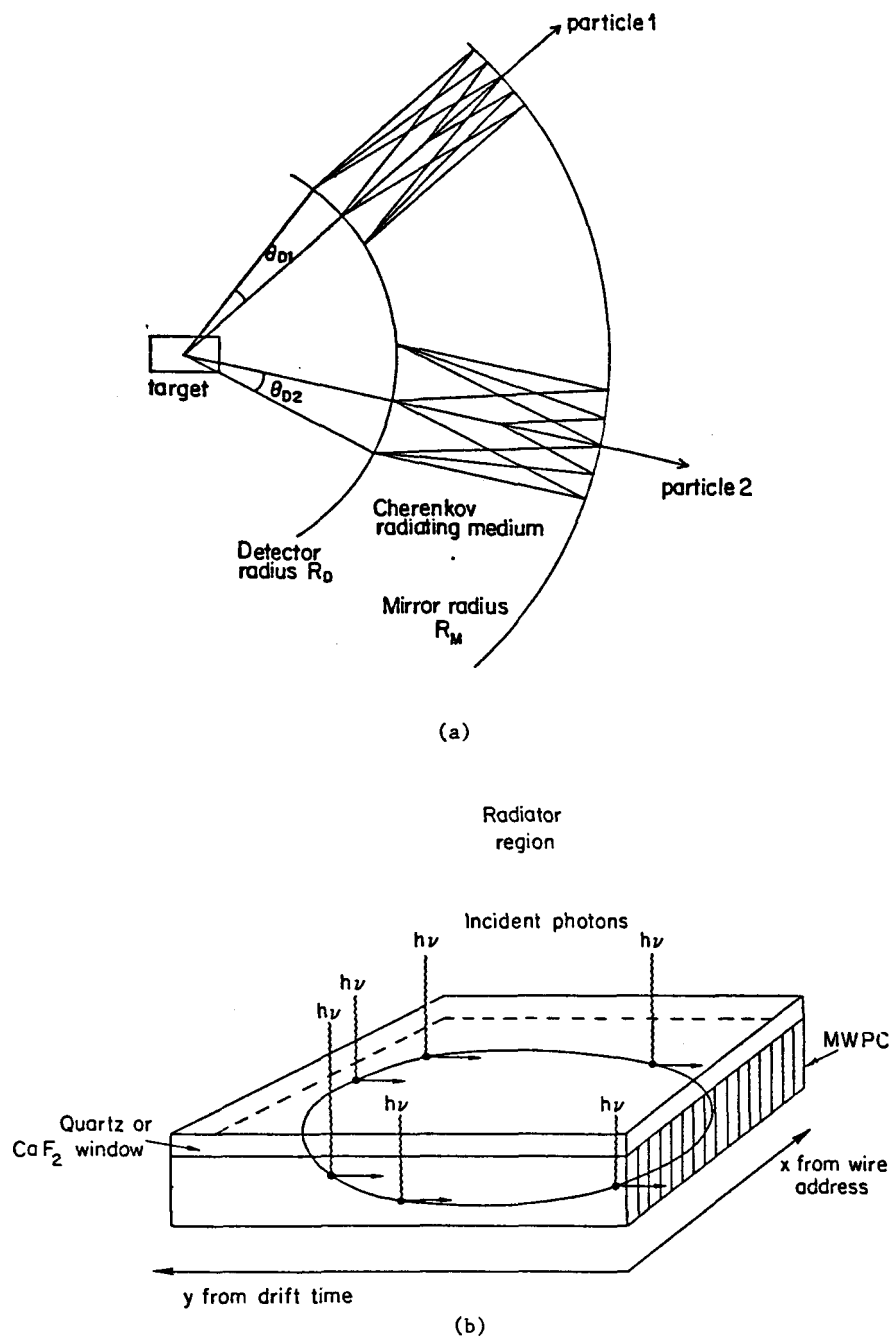


Fig. 8 (a) Diagram of the ring-imaging Cerenkov geometry. The outer spherical surface is a mirror and the inner surface a photon detector.

(b) A photon detector using drift chamber techniques. The photons are converted into photo-electrons by photo-ionization of gas molecules in the drift volume. (Ref. 13)

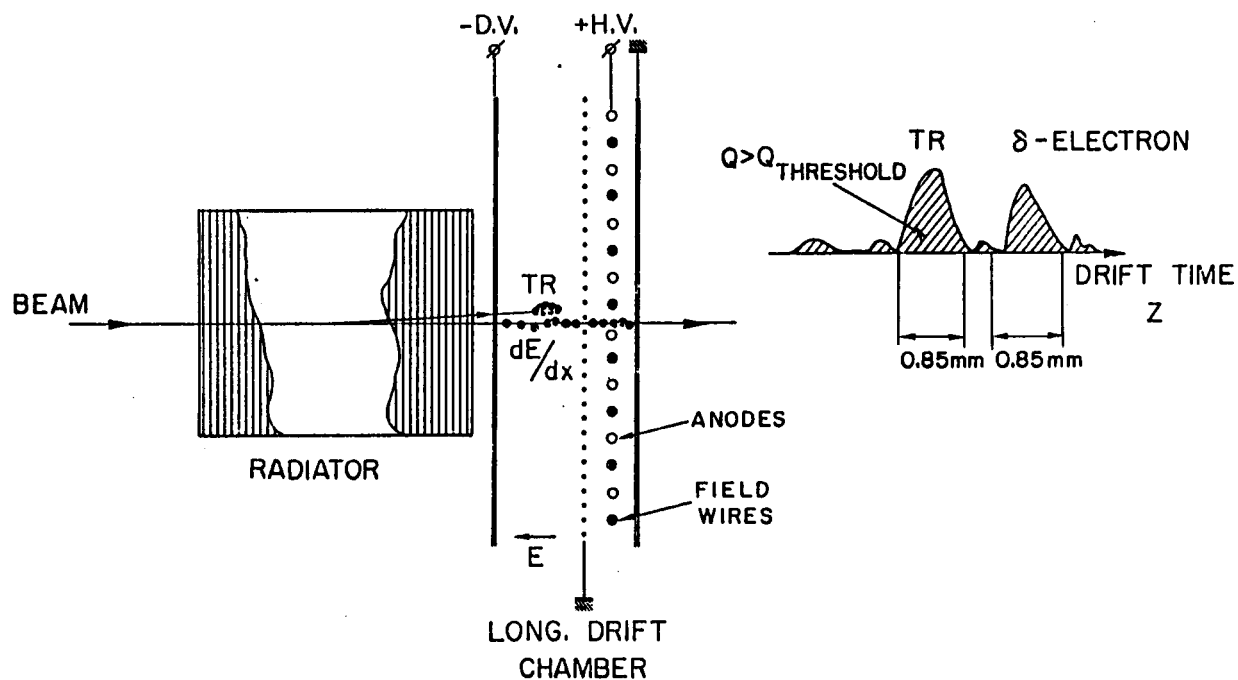


Fig. 9 Diagram showing transition radiator and drift chamber photon detector, the latter configured to detect individual ionization clusters. Also shown is the signal current waveform at the anode wire as a function of the drift time¹⁶.

ionization electrons in the wire chamber. In this way the dE/dx signals from charged tracks (small clusters) can be distinguished from the signals due to absorption of x-ray quanta (large clusters). The sensitivity to the TR signal is enhanced and the techniques discussed above for high resolution and multitrack capability in wire chambers can be brought to bear.

IV. Calorimetry

Calorimeters measure the energy and position of particles through total absorption¹⁷. Among the reasons why this technique is attractive for high energy physics are the following:

- i) It is the only way to measure neutral particles among the secondaries produced in high energy collisions.
- ii) The relative accuracy of the total energy loss measurement (a statistical process) improves with increasing E as $E^{-1/2}$. Therefore, with the increasing energy of accelerators, absorptive spectroscopy is becoming frequently more accurate than measurements with magnetic spectrometers -- especially if large solid-angle coverage is required.
- iii) The signal from these detectors is developed on a time scale of 10^{-8} to 10^{-7} sec and is therefore useful for very fast trigger decisions.

It is useful to group calorimeters according to the type of incident particle to be detected. Electrons and photons interact electromagnetically with the absorbing material through processes such as bremsstrahlung, pair production, and Compton scattering. The magnitudes of these cross sections set the scale for the size and depth of "electromagnetic calorimeters", the characteristic length being the radiation length, $X_0 \approx 180 A/Z^2$ (g cm^{-2}). Hadrons,

on the other hand, interact mainly through the strong force and hence the nuclear interaction length is the appropriate scale.

In recent years much has been learned about the properties of electromagnetic and hadronic showers, and there has been rapid development of instrumentation to realize the intrinsic resolution of calorimeters¹⁷.

In Table III we list the major contributions to the final resolution of practical detectors. For electromagnetic showers, the intrinsic fluctuations set a limit so small as to be seldom approached in practice. The most precise electromagnetic measurements are achieved with homogeneous shower counters, in which the entire volume responds with a measurable signal to the passage of charged particles. The familiar sodium iodide crystal is one such detector with excellent performance, in which resolution values close to the shower fluctuation limit have been obtained for energies ≈ 1 GeV. At higher energies, however, the measurement of electromagnetic showers in this and other detectors is dominated by various instrumental limitations.

Table IV summarizes the performance characteristics of a number of homogeneous shower detectors¹⁸. I have singled out 3 for particular scrutiny, two of which - Barium Fluoride and BGO - are of high current interest for high energy physics. The energy resolution obtained with BGO, and its time response, are very similar to sodium iodide. Its mechanical properties are much better than NaI, and it is nearly twice as dense. If we look at the costs, which I have written in below the table, we see that BGO costs ten times as much as NaI by volume. However, for a well-designed detector system, in which the segmentation of the e.m. shower detector is scaled to the shower size, the essential consideration is cost per cubic radiation length. In these terms BGO is actually cheaper than NaI at current prices.

Barium fluoride crystal has two scintillation lines, one of which has a characteristic decay time less than one nanosecond. Thus, unlike NaI or BGO, this detector could operate at the rates

Table III.

Limitations to energy resolution

Contributing mechanisms (add in quadrature)	Electromagnetic showers	Hadronic showers
Intrinsic shower fluctuations	Track length fluctuations: $\sigma/E \approx 0.005/\sqrt{E}$ [GeV]	Fluctuations in the mechanism of energy loss: $\sigma/E \approx 0.50/\sqrt{E}$ [GeV]; With compensation for nuclear effects: $\sigma/E \approx 0.20/\sqrt{E}$ [GeV]
Sampling fluctuations ^a	$\sigma/E \approx 0.04\sqrt{\Delta E/E}$	$\sigma/E \approx 0.09\sqrt{\Delta E/E}$
Instrumental effects	Noise and pedestal width: $\sigma/E \sim 1/E$ - determine minimum detectable signal; - limit low-energy performance. Calibration errors; non-uniformities: $\sigma/E \sim \text{constant}$ - limit high-energy performance.	
Incomplete containment of shower (energy leakage)	$\sigma/E \sim \log E$ For leakage fraction $\gtrsim$ few %: non-linear response and non-Gaussian "tail".	

^a ΔE = energy loss by a single charged particle in one sampling layer, measured in MeV; E = total energy, measured in GeV.

Table IV.

Properties and Performances of Homogeneous e.m. Shower Detectors

Detector type	<u>Nal(Tl)</u>	CsI(Tl)	<u>BaF₂</u>	<u>Bi₄Ge₃O₁₂</u>	Scintillating glass	Lead glass 55% PbO + 45% SiO ₂	Tl(HCO ₃)-liquid 'Helicon'	Liquid argon
Radiation length (cm)	<u>2.59</u>	1.86	<u>2.1</u>	<u>1.12</u>	~ 4	2.36	~ 1.9	14
Density (g/cm ³)	<u>3.7</u>	4.51	<u>4.9</u>	<u>7.13</u>	~ 3.5	4.08	~ 4.3	1.4
Detection mechanism	Scintillation	Scintillation	Scintillation (20% around 210 nm, 80% around 310 nm)	Scintillation	Scintillation	Cherenkov light	Cherenkov light	Ionization charge
Energy resolution (E in GeV)	$\sim 0.015 E^{-1/2} < 1$ $< 0.015 E^{-1/4} > 1$	Comparable to Nal(Tl)	Comparable to Nal(Tl)	Comparable to Nal(Tl)	$\sim 0.002 E^{-1/2}$	$\sim 0.04 E^{-1/2}$	Comparable to lead glass	$\geq 0.02 E^{-1/2}$
Principal limitation to $\sigma(E)$	Shower fluctuations optically non-uniform	Similar to Nal(Tl)	Light collection non-uniformities	Similar to Nal(Tl)	Photon statistics	Photon statistics	Photon statistics	Effect of shower fluctuation on electron collection
Signal ^{a)} (photo-el/GeV)	$\sim 10^3$	$\sim 5 \times 10^4$	$\sim 10^4$	$\sim 10^4$	Few $\times 10^3$	10^3	$\leq 10^3$ (?)	$\leq 2 \times 10^4$
Characteristic time (ns)	<u>250</u>	900	<u>0.6; 300</u>	<u>350</u>	~ 70	~ 20	~ 20	≥ 100
Rad. damage at appr. dose ^{b)} (Gy)	≤ 10	≤ 10	$\sim 10^3$	~ 10	$\sim 10^4$	$\sim 10^3$	$\geq 10^4$	Not measured; expected to be very large
Mechanical stability	Hygroscopic, fragile	Very good	Good	Good	Very good	Very good	Toxic liquid	Cryogenic liquid
References	[54, 55, 57]	[72]	[73, 74]	[61-63]	[64, 65]	[66-69]	[75]	[76-79]

a) Values are approximate, and depend on spectral matching between light source and photon detector.

b) Values are guidelines only and very substantially depending on experiment and measuring conditions.

COST/CC \$1 \$6 \$10
 COST/(RAD. LEN.)³ \$17 \$56 \$14

implied by luminosities $\sim 10^{33} \text{ cm}^{-2} \text{ sec}^{-1}$. The problem is that the fast component emits in the uv, and so efficient detection raises problems similar to those discussed in conjunction with the development of ring-imaging Cerenkov counters¹⁹.

For hadrons, sampling calorimeters are required. Figure 10 shows the measured energy resolution for various configurations, compared with the values derived from Table III. The nature of the hadronic interaction is such that a considerable fraction of the incident particle's energy is spent in reactions which do not result in an observable signal in the detector. The intrinsic shower fluctuations which limit the resolution of hadronic calorimeters are due in large measure to fluctuations in this unseen energy. Some of this energy is carried by muons and neutrinos which leave the detector without interacting. The most consequential component of "invisible energy" is that which goes into excitation and breakup of the nuclei in the absorber material. Experimentally this is most clearly demonstrated by comparing the signal from incident electrons, where these nuclear effects are absent, to that from hadrons of the same energy. Typically the electron signal is on the average 30 to 50% greater than for hadrons. Since the hadronic cascade has a large and widely fluctuating electromagnetic component (due to π^0 production), this difference in response is the dominant contribution to the resolution width unless it can be corrected for event-by-event.

A direct method for improving the energy resolution of hadronic calorimeters by correcting the nuclear fluctuations is based on the use of ^{238}U as absorber material. In the uranium absorber some of the normally invisible energy expended in nuclear breakup leads to neutron-induced fission, which in turn produces detectable energy in the calorimeter, mostly in the form of soft photons. This method of fission compensation leads to the smaller of the two values shown in Table IV for the intrinsic fluctuations in hadronic showers.

Spatial resolution and segmentation are of increasing importance as experimental goals move toward such challenges as

jet-jet effective mass reconstruction and the mapping out of energy flow in events with thousands of final-state particles. Figures 11 and 12 show an example of how a large-area, large solid-angle uranium/scintillator detector²⁰ was segmented into many hundreds of individually read out "towers" 20 cm × 20 cm. Recently, uranium/liquid argon calorimeters have been designed to achieve even higher degrees of energy- and spatial-resolution^{21,22}.

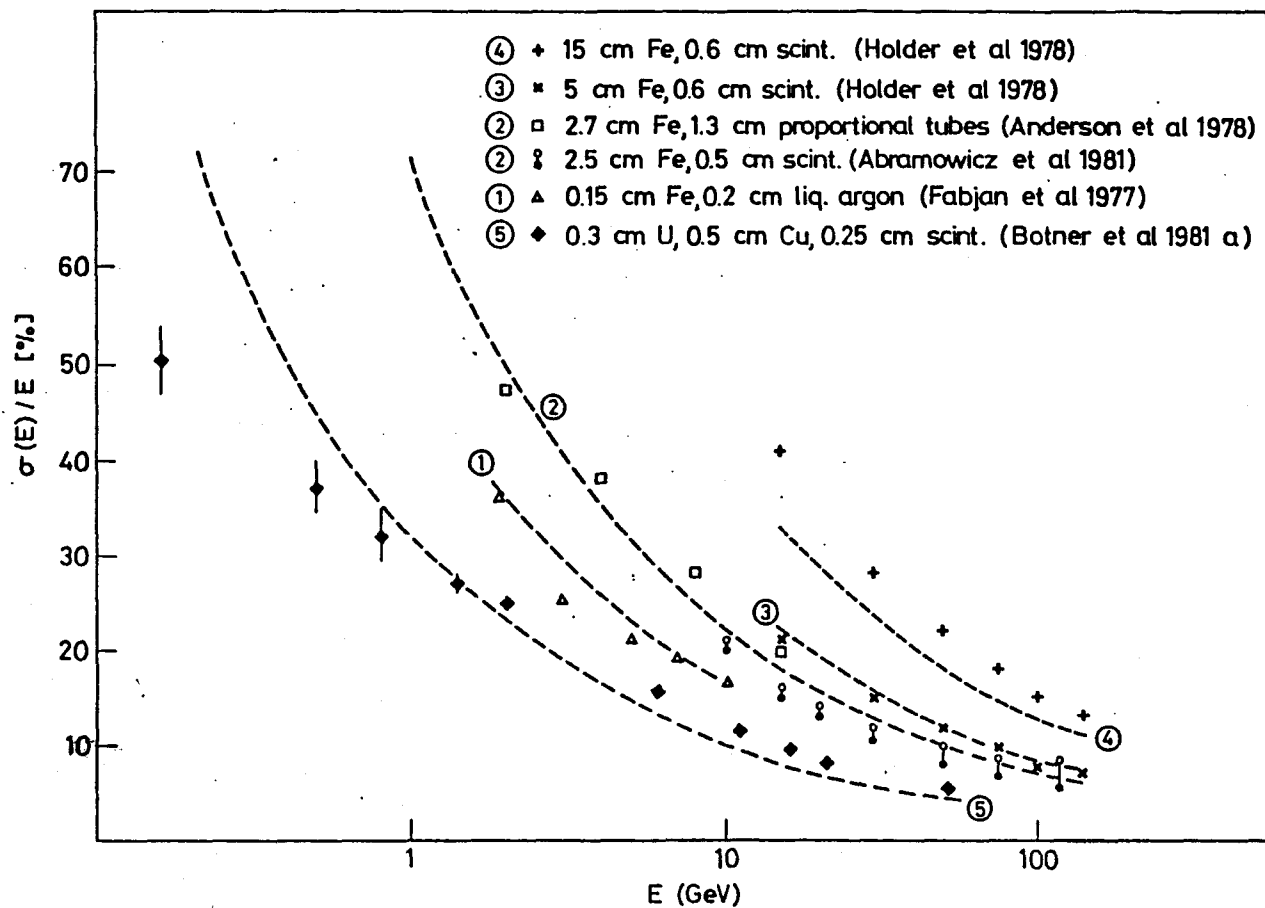


Fig. 10 Measured and predicted energy resolution for hadronic sampling calorimeters. (Ref. 17)

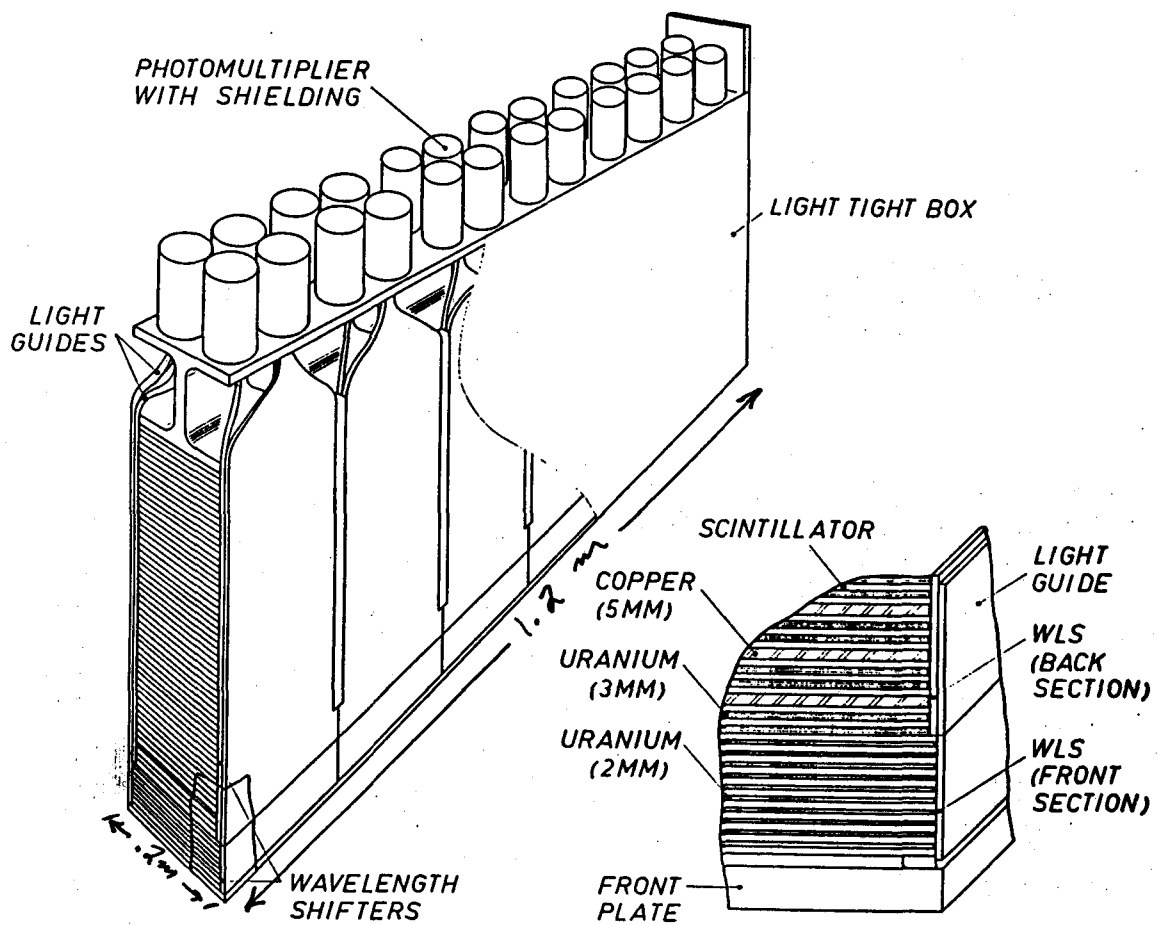


Fig. 11 A calorimeter module of the Axial Field Spectrometer.
(Ref. 20)

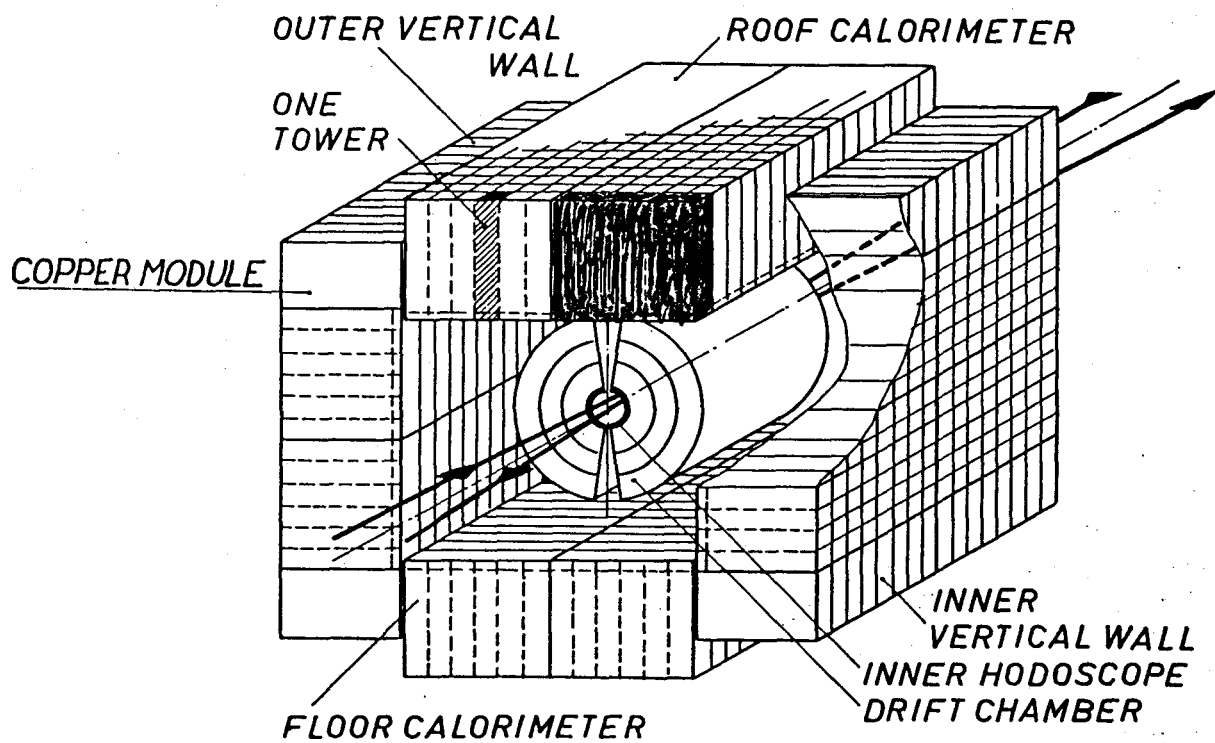


Fig. 12 The calorimeter array of the Axial Field Spectrometer at the CERN ISR. The shaded area shows the size of a single module. (Ref. 20)

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NEW PHYSICS FROM THE $(e,e'N)$ REACTION *

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ABSTRACT

New aspects of the direct nuclear knockout by the electron are studied in the general case of both polarized beams and polarized targets. The general formalism provides nine structure functions, four of which appear only when the target nucleus is polarized. In principle, a measurement of these structure functions gives complete information on the nuclear current operator. In practice, out-of-plane experiments are necessary and some structure functions, although sensitive to final state interactions, are small. On the other side asymmetries and electron polarization ratios already in coplanar kinematics have measurable size and add complementary information to the standard quasi-free nucleon knockout. In order to extend the kinematic range of energy and momentum transfers including neutron knockout, mesonic exchanges and isobar excitation in the intermediate states have to be considered. The nuclear current is therefore constructed as a nonrelativistic operator starting from the chirally invariant effective lagrangian of Peccei. The effect of the corresponding two-body contribution to the transverse nuclear current is investigated neglecting final state interactions. This contribution, already appreciable in standard $(e,e'p)$ kinematics, is important for $(e,e'n)$ reactions and in general with increasing four-momentum transfer when the nuclear response becomes more and more transverse and "parallel", i.e. concentrated along the direction of the momentum transfer.

*Invited talk presented by S. Boffi

1. INTRODUCTION

With the improved beam quality of the planned electron accelerators and associated detecting facilities new interest is found in one-nucleon emission by an electromagnetic probe in order to disentangle the different aspects of the process, i.e. reaction mechanisms and available nuclear information. In particular, it appears possible to separate the different structure functions in the nuclear response corresponding to longitudinal and transverse contributions of the hadronic current. Some work in this direction is already under way with current facilities at NIKHEF and Saclay with results indicating that probably one has to go beyond the usual impulse approximation (IA) even in the quasi-free nucleon knockout ¹ (fig. 1). The modification of the electromagnetic properties of the

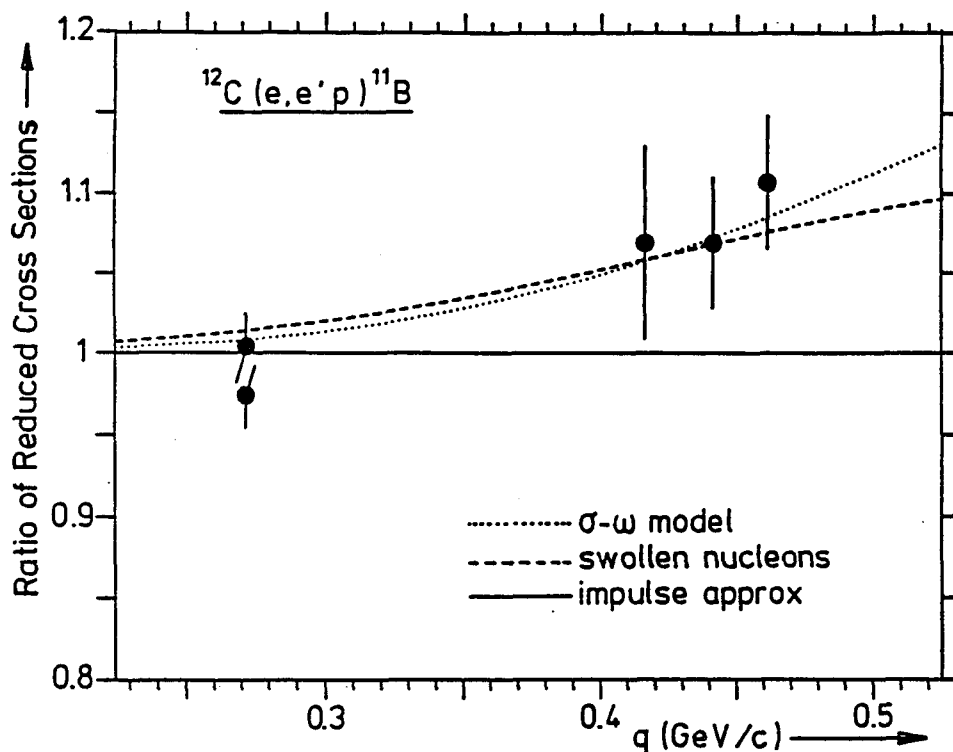


Fig. 1. Ratio of the backward to forward reduced cross section as a function of the three-momentum transfer q for the $^{12}\text{C}(e, e' p)^{11}\text{B}$ reaction leading to the ^{11}B ground state.

nucleon due to the nuclear environment is correlated to possible partial deconfinement and/or the pion cloud surrounding the nucleon. A complete experimental determination of the nuclear current is therefore desirable to gain the necessary information. This however requires taking into account spin degrees of freedom too. Therefore in the next section the general formalism is sketched to obtain the coincidence cross section for nucleon emission when both the electron beam and the target nucleus are polarized. The nine structure functions resulting in the cross section completely specify the hadronic current. Their measurement requires out-of-plane kinematics. Some examples are discussed in sect. 3 to set the scale of the relevant measurable quantities such as asymmetries and electron polarization ratios under standard quasi-free kinematic conditions. In sect. 4 meson exchange contributions (MEC) and isobar configurations (IC) in the intermediate states of the virtual photon absorption process are taken into account. Their influence on the nuclear response both in (e,e'p) and in (e,e'n) reactions is studied in sect. 5 under several kinematic conditions. Some conclusions are drawn in the final section.

2. GENERAL FORMALISM

The six-fold differential coincidence cross section is derived in this section for quasi-free nucleon knockout from a polarized target nucleus by an incident polarized electron.

An ultrarelativistic electron with initial helicity h and momentum $\vec{p}_0$ is scattered to a final momentum $\vec{p}'_0$ while a nucleon is knocked out with final momentum $\vec{p}'$. Then, the four-momentum transfer is $q_\nu^2 = \omega^2 - \vec{q}^2$, with $\omega = p_0 - p'_0$ and $\vec{q} = \vec{p}_0 - \vec{p}'_0$. In the one-photon exchange approximation and neglecting the electron mass and the nuclear recoil factor, the coincidence cross section may be written as the contraction of a leptonic tensor $L_h^{\mu\nu}$ with a hadronic tensor $W_{\mu\nu}$ along the lines of ref. 2, i.e.

$$\frac{d\sigma^h}{d\vec{p}'_0 d\vec{p}'} = \alpha^2 \frac{1}{q_\nu^2 p_0 p'_0} 2 L_h^{\mu\nu} W_{\mu\nu} \quad , \quad (1)$$

where $\alpha = e^2/4\pi = 1/137$ is the fine structure constant. Due to the initial electron helicity the lepton tensor

$$L_h^{\mu\nu} = (2q_\nu)^{-1} (1/2) \text{Tr}[\not{p}_0 \gamma^\nu \not{p}'_0 \gamma^\mu (1 + h \gamma^5)] \quad , \quad (2)$$

contains both a symmetric and an antisymmetric part under exchange of its indices ³. From the Lorentz condition for the electron current it satisfies the relations

$$q_\mu L_h^{\mu\nu} = q_\nu L_h^{\mu\nu} = 0 \quad . \quad (3)$$

The hadronic tensor

$$W_{\mu\nu} = \bar{\Sigma}_i \sum_f J_\mu J_\nu^* \quad (4)$$

is an average over initial states ($\bar{\Sigma}_i$) and a sum over final undetected states ($\sum_f$) of products of matrix elements J_μ which are written as

$$J_\mu = \int d\vec{r} e^{i\vec{q}\vec{r}} \langle \Psi_f | j_\mu(\vec{r}) | \Psi_i \rangle \quad , \quad (5)$$

i.e. matrix elements of the nuclear charge-current operator j_μ taken between the initial ($|\Psi_i\rangle$) and final ($|\Psi_f\rangle$) nuclear states. Conservation of the hadronic current requires

$$q^\mu W_{\mu\nu} = q^\nu W_{\mu\nu} = 0 \quad . \quad (6)$$

Following refs. 2,5 it is convenient to calculate the product $L_h^{\mu\nu} W_{\mu\nu}$ in the hadron Breit frame and then apply a Lorentz transformation back to the laboratory frame depicted in fig. 2. Eqs. (3) and (6) allow to express the longitudinal ($\mu, \nu = z$, i.e. along $\vec{q}$) components of the currents in terms of the time components ($\mu, \nu = 0$). Therefore, it is sufficient to consider $\mu, \nu = 0, x, y$ only in $L_h^{\mu\nu}$ and $W_{\mu\nu}$. In addition

$$\epsilon = [1 - 2 (q^2/q_V^2) \tan^2 \theta/2]^{-1} \quad (10)$$

measures the transverse linear polarization of the virtual photon exchanged by the electron scattered at an angle θ , and

$$\epsilon_L = (-q_V^2/q^2) \epsilon . \quad (11)$$

In the laboratory frame of fig. 2, where the initial nucleus is assumed to be at rest and with a polarization direction specified by the zenithal angle θ^* and the azimuthal angle ϕ^* , the six structure functions in eqs. (8) and (9) depend on the seven variables (ω , q , p' , γ , ϕ , θ^* , ϕ^*). The ϕ -dependence is rather simple; according to the experimental situation where one deals with the angle $\alpha = -\phi$ between the $\vec{q}$, $\vec{p}'$ plane and the scattering plane of the electron, one has

$$\begin{aligned} F_{00} &= f_{00} , \\ F_{11} &= f_{11} , \\ F_{01} &= f_{01} \cos \alpha + \bar{f}_{01} \sin \alpha , \\ F_{1-1} &= f_{1-1} \cos 2\alpha + \bar{f}_{1-1} \sin 2\alpha , \\ F'_{11} &= f'_{11} , \\ F'_{01} &= f'_{01} \sin \alpha + \bar{f}'_{01} \cos \alpha . \end{aligned} \quad (12)$$

With the present choice of the frame of reference (which is different from that of previous work ^{6,8}) where the x-z plane is the hadron plane, in terms of the covariant hadronic tensor components and the spherical components of the hadronic current, the nine structure functions are defined as follows:

$$\begin{aligned} f_{00} &= W_{00} = |J_0|^2 , \\ f_{11} &= W_{xx} + W_{yy} = |J_1|^2 + |J_{-1}|^2 , \\ f_{01} &= \sqrt{2} (W_{0x} + W_{x0}) = 2 \operatorname{Re} (J_1 J_0^* - J_0 J_{-1}^*) , \\ f_{1-1} &= W_{yy} - W_{xx} = 2 \operatorname{Re} (J_1 J_{-1}^*) , \\ f'_{11} &= -i (W_{xy} - W_{yx}) = |J_1|^2 - |J_{-1}|^2 , \\ f'_{01} &= -i \sqrt{2} (W_{0x} - W_{x0}) = -2 \operatorname{Im} (J_1 J_0^* + J_0 J_{-1}^*) , \end{aligned}$$

$$\begin{aligned}
\bar{f}_{01} &= -\sqrt{2} (W_{0y} + W_{y0}) &= -2 \operatorname{Im} (J_1 J_0^* - J_0 J_{-1}^*) , \\
\bar{f}_{1-1} &= W_{xy} + W_{yx} &= -2 \operatorname{Im} (J_1 J_{-1}^*) , \\
\bar{f}'_{01} &= -i \sqrt{2} (W_{0y} - W_{y0}) &= 2 \operatorname{Re} (J_1 J_0^* + J_0 J_{-1}^*) ,
\end{aligned} \tag{13}$$

where $\bar{f}_i$ is understood.

For an unpolarized target nucleus the barred structure functions $\bar{f}_{01}$, $\bar{f}_{1-1}$ and $\bar{f}'_{01}$ vanish identically. In addition, $|J_1|^2 = |J_{-1}|^2$, i.e. $f'_{11} = 0$. Therefore, in this case one recovers the expression of the cross section in terms of five structure functions already studied in ref. 6 with incident polarized electrons. The same situation occurs when the target polarization is along y, because in this case $J_{-\lambda} J_{-\lambda}^* = (-)^{\lambda+\lambda'} J_{\lambda} J_{\lambda}^*$.

The primed structure functions f'_{11} , f'_{01} and $\bar{f}'_{01}$ contribute only in the presence of polarized electrons. If in addition to an unpolarized electron beam also an unpolarized target is used, the familiar form in terms of only four structure functions, $f_{\lambda\lambda'}$, is obtained 7,8.

The relevant quantity containing all the necessary ingredients for describing the knockout process is then the matrix element of eq. (5). It is common practice in the study of quasi-free reactions to assume that the nuclear charge-current operator is a one-body operator 9. One then confines to the specific channel subspace of states $|\Psi_f\rangle$ which correspond asymptotically to a core $|E\rangle$ and a knocked out nucleon. As a result of space truncation the transition amplitude in eq. (5) is written in a one-body representation in terms of an effective transition operator $\tilde{j}_\mu(\vec{r})$ as follows 10

$$\langle \Psi_f | j_\mu(\vec{r}) | \Psi_i \rangle = \sum_{\alpha\alpha'} \chi_{E\alpha}^{(-)*} \tilde{j}_\mu \phi_{E\alpha} [S(E)]^{1/2} , \tag{14}$$

where

$$[S(E)]^{1/2} \phi_{E\alpha} = \langle E | a_\alpha | \Psi_i \rangle , \tag{15}$$

$S(E)$ being the spectral strength associated with the spectroscopic factor of the residual nucleus in state $|E\rangle$ and

$$\chi_{E\alpha}^{(-)} = (1-n_\alpha)^{-1} \langle E | a_\alpha | \Psi_f \rangle , \quad (16)$$

n_α being the occupation number ($\neq 1$) of the natural orbital α of $|E\rangle$.

Both $\phi_{E\alpha}$ and $\chi_{E\alpha}^{(-)}$ are single particle wave functions which are eigenfunctions of the same (energy dependent) Feshbach optical hamiltonian; as such, they are not orthogonal. However, the use of an effective transition operator in eq. (14) removes the orthogonality defect of the model wave functions, besides taking into account space truncation effects. Actually, the orthogonality defect in standard kinematic conditions for $(e,e'p)$ reactions is negligible ¹⁰. Therefore, one can safely approximate $\tilde{j}_\mu(r)$ by $j_\mu(r)$.

Adopting the nonrelativistic approximation of ref. 11, the charge-current operator j_μ can be expanded in powers of the inverse nucleon mass ¹². Through second order it reads

$$j_0 = F_1 + (8m^2)^{-1}(F_1 + 2 \kappa F_2) (q_\nu^2 - i \vec{\sigma} \cdot \vec{P} \times \vec{q}) , \quad (17)$$

$$\begin{aligned} \vec{j} = & (2m)^{-1}[F_1 \vec{P} + (F_1 + \kappa F_2) i \vec{\sigma} \times \vec{q}] \\ & - (8m^2)^{-1}(F_1 + 2 \kappa F_2) i \vec{\omega} \times \vec{P} \end{aligned} , \quad (18)$$

where $\vec{P} = \vec{p} + \vec{p}'$; $F_1 = F_1(q_\nu^2)$ and $F_2 = F_2(q_\nu^2)$ are the Pauli-Dirac electromagnetic nucleon form factors ($F_1(0) = F_2(0) = 1$ for protons, $F_1(0) = 0$ and $F_2(0) = 1$ for neutrons) and κ is the anomalous magnetic moment of the nucleon in nuclear magnetons ($\kappa_p = 1.79$, $\kappa_n = -1.91$).

The results presented in this paper are obtained with the above charge-current operator for the one-body part of the transition operator along the lines of previous work ^{12-15,8,6}. Other choices are possible based on a semirelativistic approach in terms of Feynman diagrams ^{16,17} and/or fully relativistic approaches ^{18,19}.

3. THE $(\vec{e}, e'p)$ REACTION ON POLARIZED TARGETS

The general formalism of the preceding section is here applied to the case of proton knockout by an electron of positive helicity from a polarized target. For each target polarization there are nine structure functions, which are functions of ω , q , p' and γ . Specific cases can be studied with target polarization oriented longitudinally (L in fig. 2), normally (N) and sideways (S) with respect to the incident electron direction ²⁰. For symmetry reasons the N-polarization reduces the number of nonvanishing structure functions to five: only f_{00} , f_{11} , f_{1-1} , f_{01} and f'_{01} survive. Without final state interactions also $f'_{01} = \bar{f}_{01} = \bar{f}_{1-1} = 0$ for L- and S-polarization of the target.

Selecting the energy of the emitted nucleon and its initial orbit, ω and p' are no more independent variables. At fixed q one can then plot the structure functions for each hole as a function of γ .

In fig. 3 and 4 the nonvanishing structure functions are given for a $p_{3/2}$ hole in ${}^7\text{Li}$ for N-polarization of the target. The final proton energy is 144 MeV, the initial electron energy is 700 MeV and $q = 2.6 \text{ fm}^{-1}$. Such conditions allow the use of two available phenomenological optical potentials, indicated by KMT ²¹ and JWS ²², to describe final state distortion. The bound state of the $p_{3/2}$ proton is taken from ref. 23 (ES). The spectral strength $S(E)$ is assumed in agreement with the extreme shell model, i.e. $S(E) = 1$. The pure longitudinal f_{00} and the pure transverse f_{11} follow the typical p-wave distribution with a maximum height comparable with the results obtained in similar cases without target polarization ⁸. However, the PWIA result is lower than the DWIA one.

As expected, the interference terms f_{01} and f_{1-1} are much smaller as in the case of unpolarized targets. On the contrary the fifth structure function f'_{01} is very large and of opposite sign with respect to f_{01} . As a consequence the cross section for protons emitted above and below the electron scattering plane will exhibit a pronounced asymmetry when dealing with incident polarized electrons and the N-polarization of the target. This asymmetry, larger than in

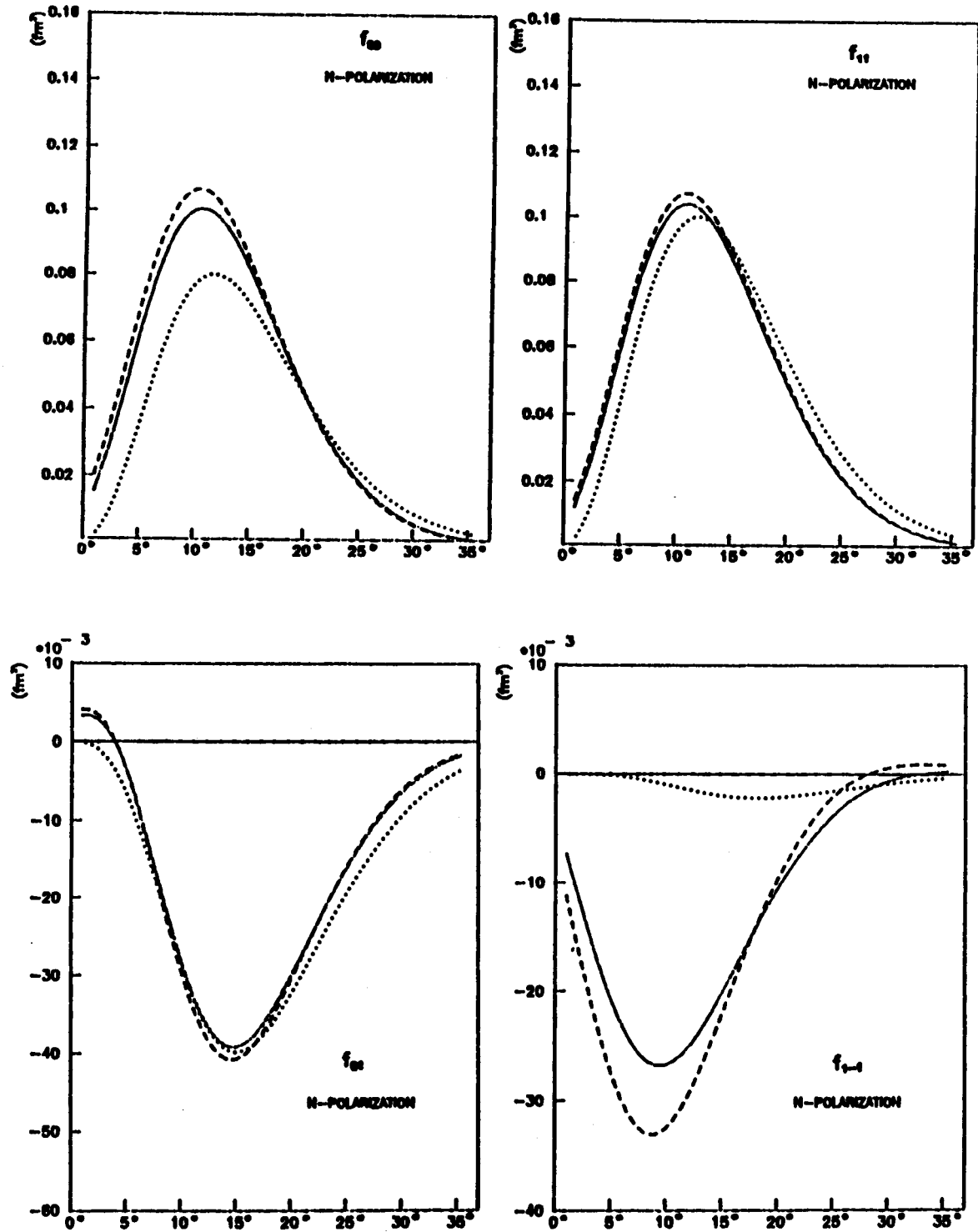


Fig. 3. Structure functions for the $p_{3/2}$ hole in N-polarized ${}^7\text{Li}$ vs. γ . The final proton energy is 144 MeV and $q=2.6 \text{ fm}^{-1}$. Bound state from ref. 23, optical model from ref. 21 (solid line) and from ref. 22 (dashed line); dotted line for PWIA.

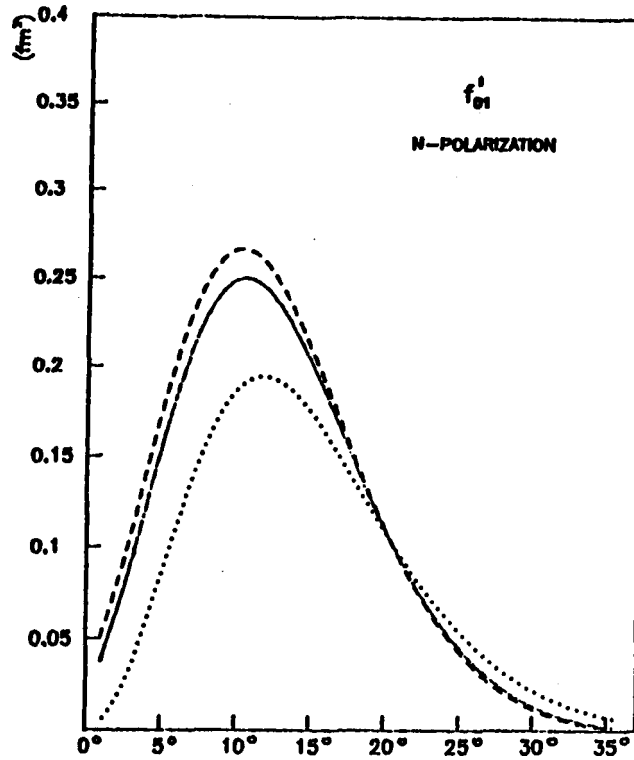


Fig. 4. The fifth structure function f'_{01} for $p_{3/2}$ hole in ${}^7\text{Li}$ under the same conditions as in fig. 3.

the ${}^{16}\text{O}$ case studied in ref. 6, is also dependent on final state interactions.

In fig. 5 two interference structure functions, $\bar{f}_{01}$ and $\bar{f}_{1-1}$, are given for $p_{3/2}$ hole in ${}^7\text{Li}$ polarized along the incident electron direction. In PWIA these two structure functions vanish. In DWIA their size is much smaller than the corresponding pure longitudinal and the pure transverse structure functions, which dominate the cross section. Therefore, their determination by means of out-of-plane experiments requires a much higher precision than the one available in the near future. However, it is interesting that also for L-polarization of the target the PWIA cross section is lower than the DWIA one for this choice of optical potentials, whereas for S-polarization the reverse is true.

With coplanar kinematics one can measure asymmetries, i.e.

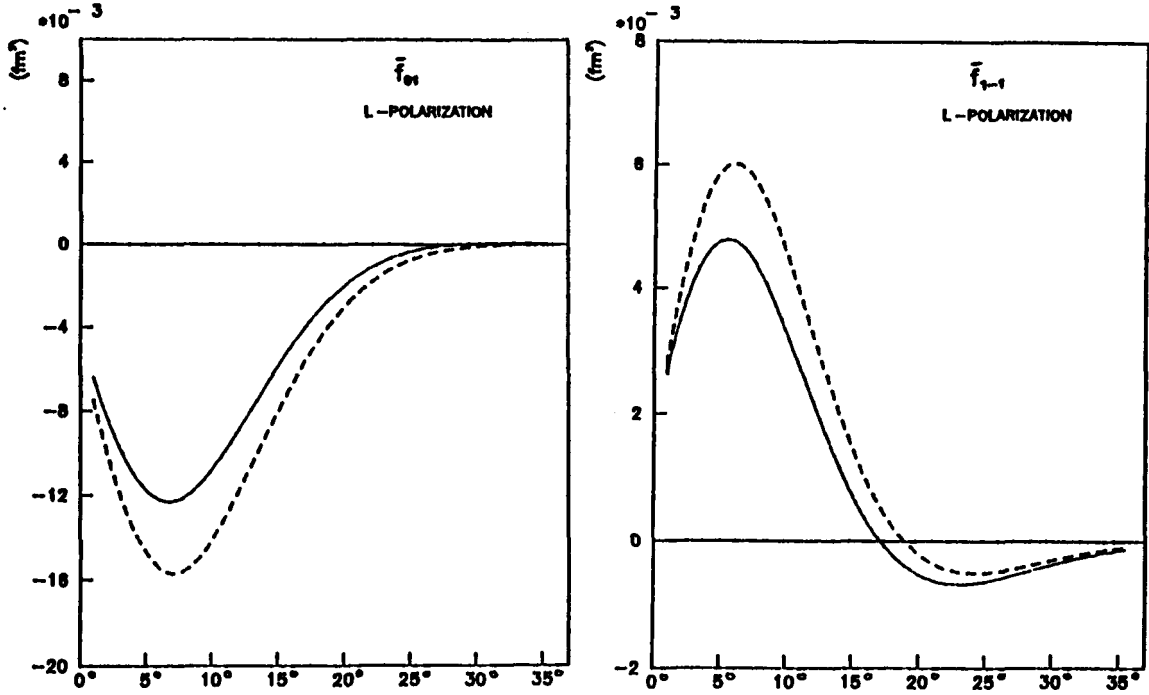


Fig.5. The interference structure functions $\bar{f}_{01}$ and $\bar{f}_{1-1}$ for $p_{3/2}$ hole in L-polarized ${}^7\text{Li}$. Other conditions are the same as in fig. 3.

$$(\Sigma^N - \Sigma^S)/\Sigma_0, (\Sigma^L - \Sigma^N)/\Sigma_0, (\Sigma^S - \Sigma^L)/\Sigma_0, \quad (19)$$

where Σ_0 is the cross section with unpolarized target and unpolarized electron. These asymmetries are determined by the behaviour of the usual four structure functions f_{00} , f_{11} , f_{01} and f_{1-1} with different orientations of the target polarization. Only two of them are independent. As an example, in figs. 6 and 7 the N-S and the L-N asymmetries are given for $p_{3/2}$ hole in ${}^7\text{Li}$ as a function of the recoil momentum ($\vec{p}_R = -\vec{p}$) under the same kinematic conditions of fig. 3. These asymmetries are not vanishing as an indication that different results are expected for different orientation of the target polarization. Moreover, the size of the asymmetry is measurable and is increased by final state interactions.

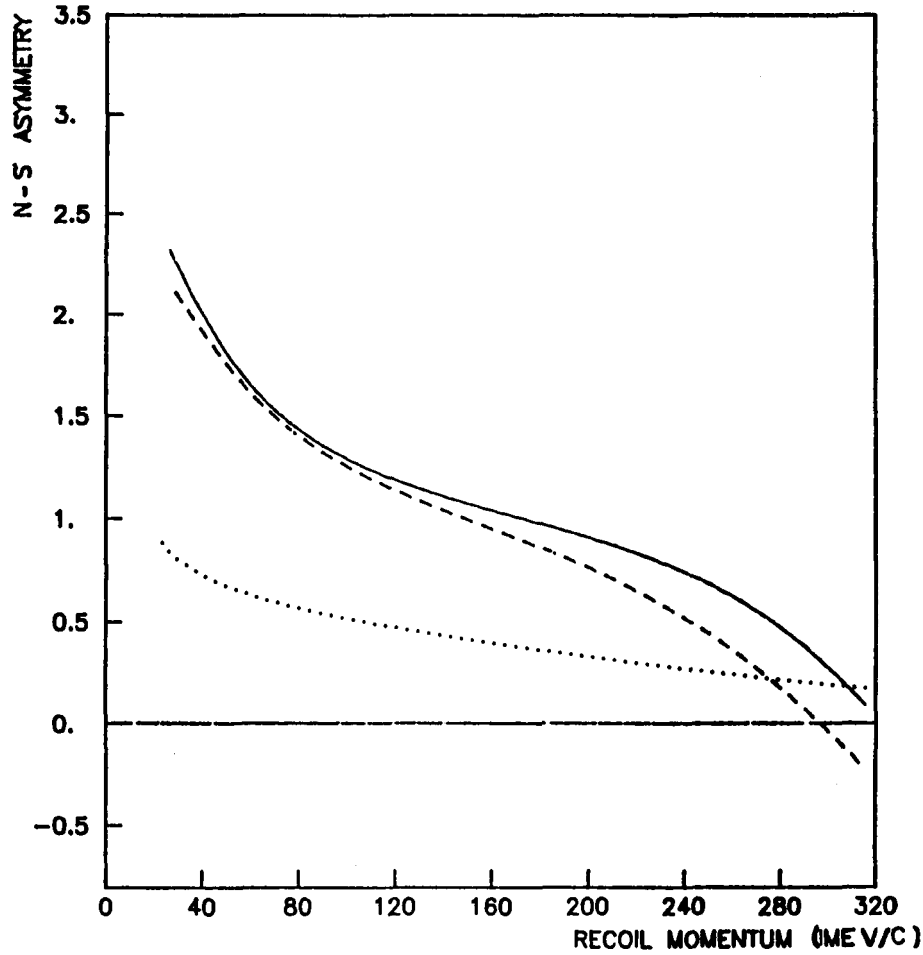


Fig. 6. Asymmetry $(\Sigma^N - \Sigma^S)/\Sigma_0$ for $p_{3/2}$ hole in ${}^7\text{Li}$ under the same conditions as in fig. 3.

Another measurable quantity with coplanar kinematics is the electron polarization ratio Δ/Σ with L- and S-polarization of the target. For N-polarization, $\Delta/\Sigma = 0$ as already noticed in ref. 20. In figs. 8 and 9 $(\Delta/\Sigma)_S$ and $(\Delta/\Sigma)_L$ are plotted for the $d_{3/2}$ hole in ${}^{39}\text{K}$ obtained with 500 MeV incident electrons and 100 MeV outgoing protons; $q = 2.2 \text{ fm}^{-1}$. The PWIA result is compared with two different choices of bound and scattering states in DWIA, i.e. bound ES states 23 with optical potential (SCH) given in ref. 24 and both bound and scattering states obtained with the hamiltonian (GR) given in ref. 25. These results show that the electron polarization ratio is large and

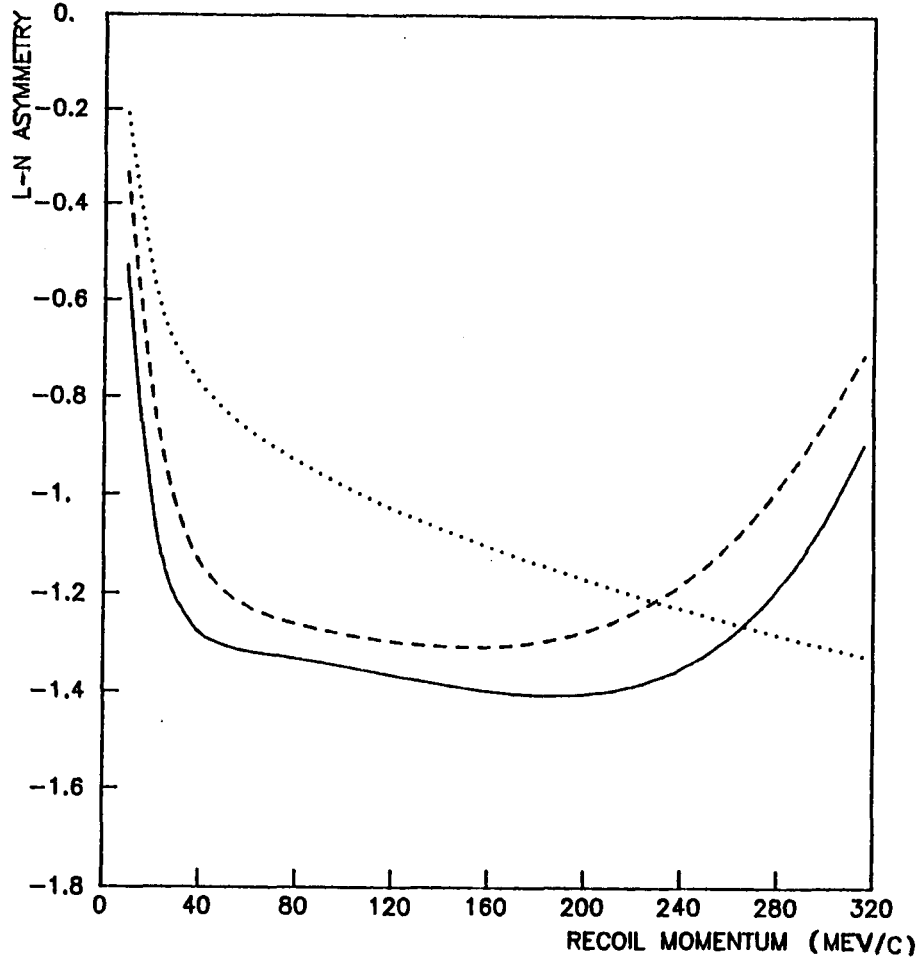


Fig. 7. Asymmetry $(\Sigma^L - \Sigma^N)/\Sigma_0$ for $p_{3/2}$ hole in ${}^7\text{Li}$ under the same conditions as in fig. 3.

can be measured, e.g., on top of the $d_{3/2}$ momentum distribution ($p_R \approx 125$ MeV/c) with a weak dependence on final state distortion.

With an aligned target the tensor analyzing power T_{20} is also measurable. According to the Madison convention ^{26,27} the cross section (1) averaged over electron helicity can be rewritten in terms of irreducible statistical tensors t_{IM} describing the polarization state of the target as

$$\frac{d\sigma}{d\vec{p}'_0 d\vec{p}'} = \Sigma_0 \Sigma_{IM} t_{IM} T_{IM} . \quad (20)$$

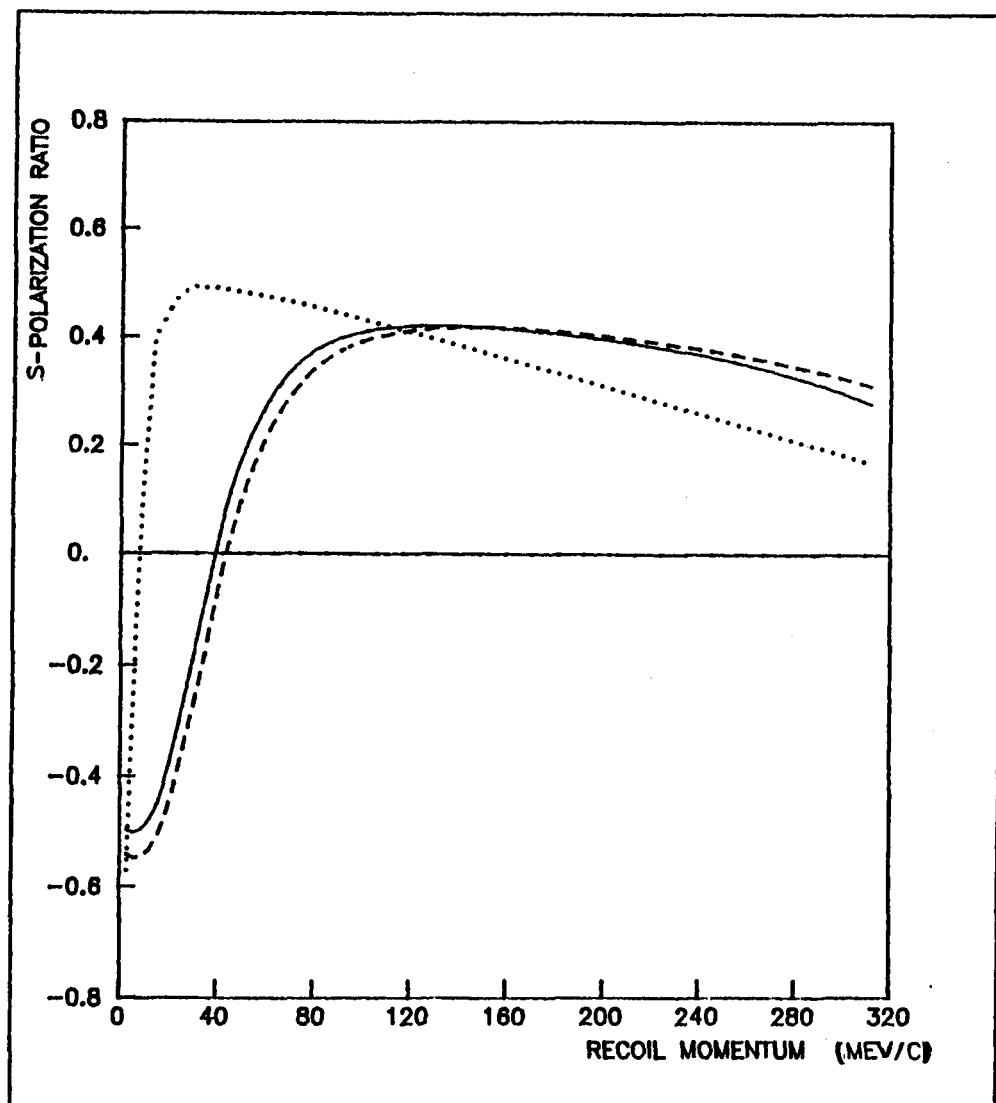


Fig. 8. The electron polarization ratio Δ/E with S-polarization of the ^{39}K target nucleus for $d_{3/2}$ hole with 100 MeV outgoing proton energy and $q=2.2 \text{ fm}^{-1}$. Bound state from ref. 23 and optical model from ref. 24 for the solid line, bound and scattering states with the hamiltonian of ref. 25 for the dashed line; PWIA for the dotted line.

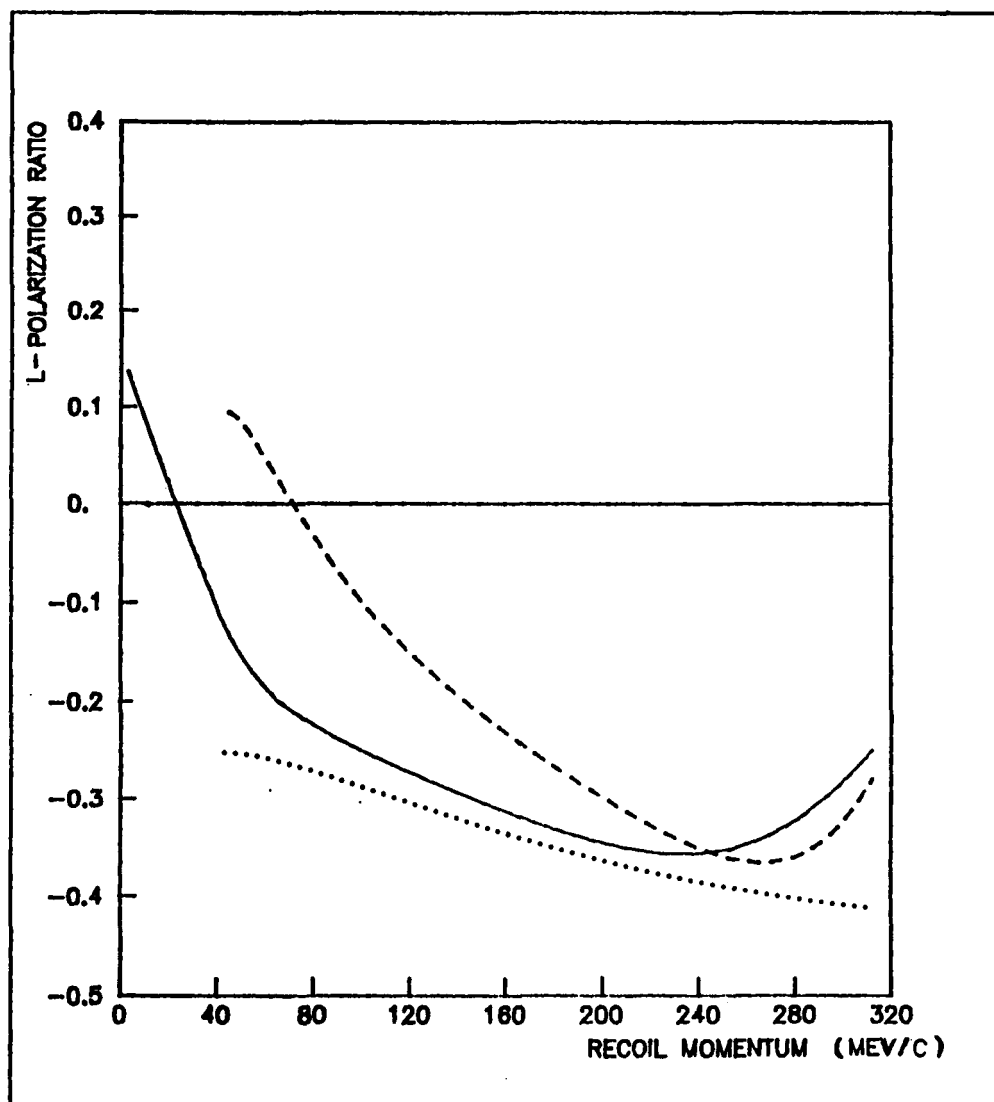


Fig. 9. The same as in fig. 8 but with L-polarization of the target.

Selecting, e.g., the target polarization aligned along the incident electron direction, T_{20} is obtained as

$$T_{20} = (\Sigma^L - \Sigma_0)/\Sigma_0 \quad . \quad (21)$$

In fig. 10 this quantity is plotted for the $d_{3/2}$ hole in ^{39}K under the same kinematic conditions as in figs. 8 and 9. As for the polarization ratio, also T_{20} (being the ratio between two cross sections) is not very much dependent on final state interactions. Its size also shows that it would be possible to measure it.

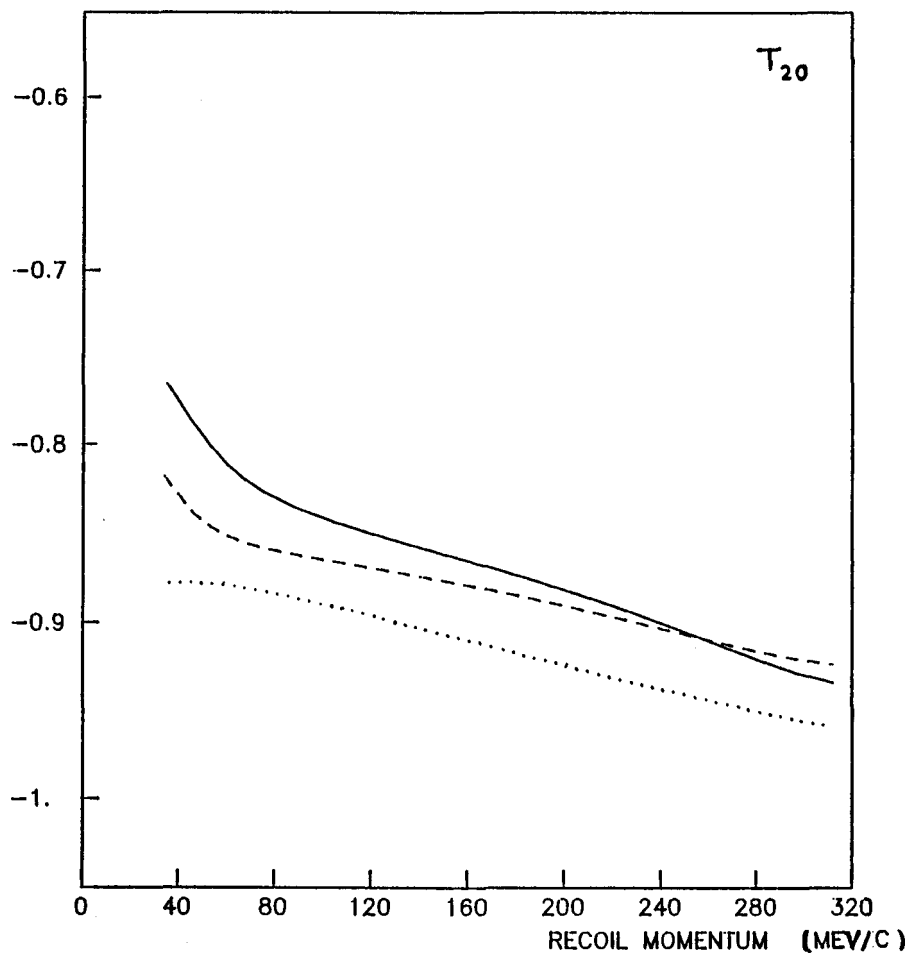


Fig. 10. The tensor analyzing power T_{20} for $d_{3/2}$ hole in ^{39}K with the same conditions as in fig. 8.

4. EXCHANGE CURRENTS AND ISOBAR CONFIGURATIONS IN NUCLEON KNOCKOUT

It is common belief that in quasi-free nucleon knockout meson exchange currents and isobar configurations in the intermediate states of the virtual photon absorption process play a minor role. This is certainly the case for deuteron electrodisintegration ²⁸ and for the inclusive (e,e') reaction ²⁹. However, for the higher four-momentum transfer allowed by future electron accelerators and for neutron knockout where the one-body current gives a small contribution MEC and IC can not be neglected. In this section the corresponding two-body currents are given and a simple model for nucleon knockout is described.

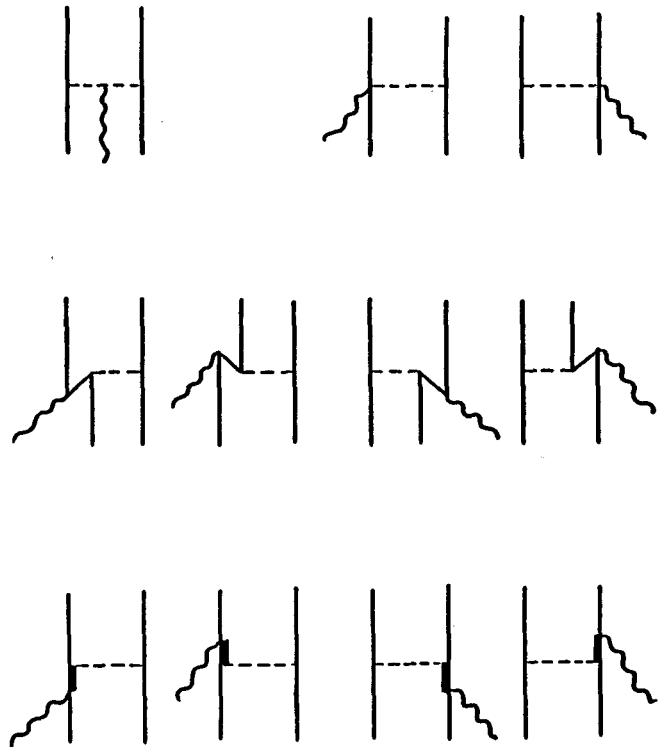


Fig. 11 Lowest order Feynman graphs used in this paper.

The chirally invariant effective lagrangian of ref. 30 is here adopted when considering the lowest order Feynman graphs of fig. 11. Expressions for the two-body current matrix elements in momentum space representation ²⁹ can be reduced nonrelativistically. The current $\vec{j}^P$ due to the pair diagrams vanishes identically, the currents due to the pion-in-flight diagrams ($\vec{j}^\pi$) and to the contact diagrams ($\vec{j}^C$) coincide in this approximation with the result obtained in refs. 31,32 with pseudoscalar coupling, i.e.

$$\vec{j}^\pi(\vec{k}_1, \vec{k}_2) = -i \frac{f^2}{\mu^2} F_1^V(q_V^2) [\vec{\tau}(1) \times \vec{\tau}(2)]_3 \frac{\vec{\sigma}(1) \cdot \vec{k}_1}{\mu^2 + k_1^2} \frac{\vec{\sigma}(2) \cdot \vec{k}_2}{\mu^2 + k_2^2} (\vec{k}_2 - \vec{k}_1) , \quad (22)$$

$$\vec{j}^C(\vec{k}_1, \vec{k}_2) = -i \frac{f^2}{\mu^2} F_1^V(q_V^2) [\vec{\tau}(1) \times \vec{\tau}(2)]_3 \left\{ \vec{\sigma}(1) \frac{\vec{\sigma}(2) \cdot \vec{k}_2}{\mu^2 + k_2^2} - \frac{\vec{\sigma}(1) \cdot \vec{k}_1}{\mu^2 + k_1^2} \vec{\sigma}(2) \right\} , \quad (23)$$

$$\vec{j}^P(\vec{k}_1, \vec{k}_2) = 0. \quad (24)$$

In eqs. (22) and (23), μ is the pion mass, $\vec{k}_1$ and $\vec{k}_2$ are the momenta transferred to the nucleons, and $f^2/4\pi=0.08$. The isovector electromagnetic form factor F_1^V has been inserted according to the suggestion of ref. 33 to ensure current conservation.

The current $\vec{j}^\Delta$ due to IC diagrams is

$$\begin{aligned} \vec{j}^\Delta(\vec{k}_1, \vec{k}_2) = & -i \frac{\kappa^* h f}{3m\mu^2} \vec{q} \times \{ 2M_1 \tau_3^{(1)} \vec{k}_1 \frac{\vec{\sigma}(1) \cdot \vec{k}_1}{\mu^2 + k_1^2} + 2M_1 \tau_3^{(2)} \vec{k}_2 \frac{\vec{\sigma}(2) \cdot \vec{k}_2}{\mu^2 + k_2^2} + \\ & + M_2 [\vec{\tau}(1) \times \vec{\tau}(2)]_3 [(\vec{k}_2 \times \vec{\sigma}(1)) \frac{\vec{\sigma}(2) \cdot \vec{k}_2}{\mu^2 + k_2^2} - \frac{\vec{\sigma}(1) \cdot \vec{k}_1}{\mu^2 + k_1^2} (\vec{k}_1 \times \vec{\sigma}(2))] \} , \end{aligned} \quad (25)$$

where

$$M_1 = \frac{2M + 3m}{M^2 - m^2} , \quad M_2 = \frac{2M + m}{M^2 - m^2} , \quad (26)$$

and $M=1232$ MeV is the isobar mass. The constants $h^2=0.290$ and $\kappa^*=5.0$ are taken from ref. 30. Consequently, at $q_V^2=0$ the strength of $\vec{j}^\Delta$ is larger by $\approx 50\%$ than the corresponding strength in the static quark model 31,32.

A dipole form factor

$$F(q_V^2) = [1 - q_V^2 / (855 \text{ MeV})^2]^{-2} \quad (27)$$

has to be placed in front of $\vec{j}^A$ as in ref. 29 to take into account the electromagnetic form factor of the isobar which corresponds to the isovector Sachs form factor G_M^V used in the static quark model 31,32.

In the above approximation only the spatial nonrelativistic two-body $\vec{j}^{(2)}$ currents (22), (23) and (25) contribute. They add to the one-body charge-current operator of eqs. (17) and (18) and are effective only in the transverse part of the hadronic current when calculating the product $L^{\mu\nu}W_{\mu\nu}$ in eq. (1).

In order to set the scale of MEC and IC contributions, preliminary results are presented in the following section according to a simple model where the emitted nucleon is described by a plane wave and the matrix elements of the hadronic current (5) are evaluated within the Fermi gas model. Also, zero spin targets and unpolarized incident electrons are considered.

Then, one assumes that the final state $|\Psi_f\rangle = |\vec{p}\vec{p}'\rangle$ is a 1p-1h state with respect to the initial Fermi gas ground state $|\Psi_i\rangle = |0\rangle$, i.e.

$$|\vec{p}\vec{p}'\rangle = \Theta(|\vec{p}'| - \vec{k}_F) \Theta(\vec{k}_F - |\vec{p}|) a_{\vec{p}}^\dagger a_{\vec{p}} |0\rangle, \quad (28)$$

where $\vec{k}_F$ is the Fermi momentum. In this approximation $\vec{p}_R = -\vec{p}$ is the recoil momentum of the residual nucleus.

Separating the hadronic current in a one-body ($j_\mu^{(1)}$) and a two-body part ($j_\mu^{(2)}$) one has for the matrix elements (5) in momentum space:

$$J_\mu = \{ j_\mu^{(1)}(\vec{p}, \vec{p}') - \Gamma_k j_\mu^{(2)}(\vec{p}', -\vec{k}, \vec{k} - \vec{p}) \delta_{\mu 1} \Theta(\vec{k}_F - |\vec{k}|) \} \\ \cdot \Theta(|\vec{p}'| - \vec{k}_F) [S_\alpha(E)]^{1/2} \Phi_\alpha(\vec{p}) \quad (29)$$

In eq. (29), $\Theta(\vec{k}_F - |\vec{p}|)$ has been substituted by a more realistic

momentum density describing the bound state α from which the nucleon has to be removed (see eq. (15)).

In this PWIA model only the usual four structure functions f_{00} , f_{11} , f_{01} and f_{1-1} survive and factorize, i.e.

$$f_{\lambda\lambda'} = S(\vec{p}_R; E) g_{\lambda\lambda'} \quad , \quad (30)$$

where $S(\vec{p}, E)$ is the diagonal part of the spectral density ³⁴,

$$S(\vec{p}, \vec{p}'; E) = \Gamma_\alpha \Phi_\alpha(\vec{p}) S_\alpha(E) \Phi_\alpha^*(\vec{p}') \quad , \quad (31)$$

containing the relevant nuclear information to be extracted from the data.

The quantities $g_{\lambda\lambda'}$ in eq. (30) are determined by the nuclear current. The contribution of $j_\mu^{(1)}$ is rather simple ⁸:

$$\begin{aligned} g_{00} &= F_1^2(q_V^2) + (4m^2)^{-1} q_V^2 F_1(q_V^2) [F_1(q_V^2) + 2\kappa F_2(q_V^2)] \quad , \\ g_{11} &= m^{-2} F_1^2(q_V^2) p'^2 \sin^2\gamma + (2m^2)^{-1} [F_1(q_V^2) + \kappa F_2(q_V^2)]^2 q^2 \quad , \\ g_{01} &= -2 \sqrt{2} m^{-1} F_1^2(q_V^2) p' \sin\gamma \quad , \\ g_{1-1} &= -m^{-2} F_1^2(q_V^2) p'^2 \sin^2\gamma \quad . \end{aligned} \quad (32)$$

MEC and IC terms of $\vec{j}^{(2)}$ affect g_{11} and g_{1-1} ; g_{01} is modified by MEC only. Their expressions are rather cumbersome and will not be given here.

5. NUCLEAR RESPONSE WITH MEC AND IC

We consider first (e,e'p) reactions with standard kinematics, i.e. emitted protons of 100 MeV by a momentum transfer of ≈ 450 MeV/c. The ^{16}O case has received considerable attention in previous work ⁸ and will be therefore adopted as a case study. In fig. 12 the four structure functions (30) are plotted vs. the emission angle γ (with

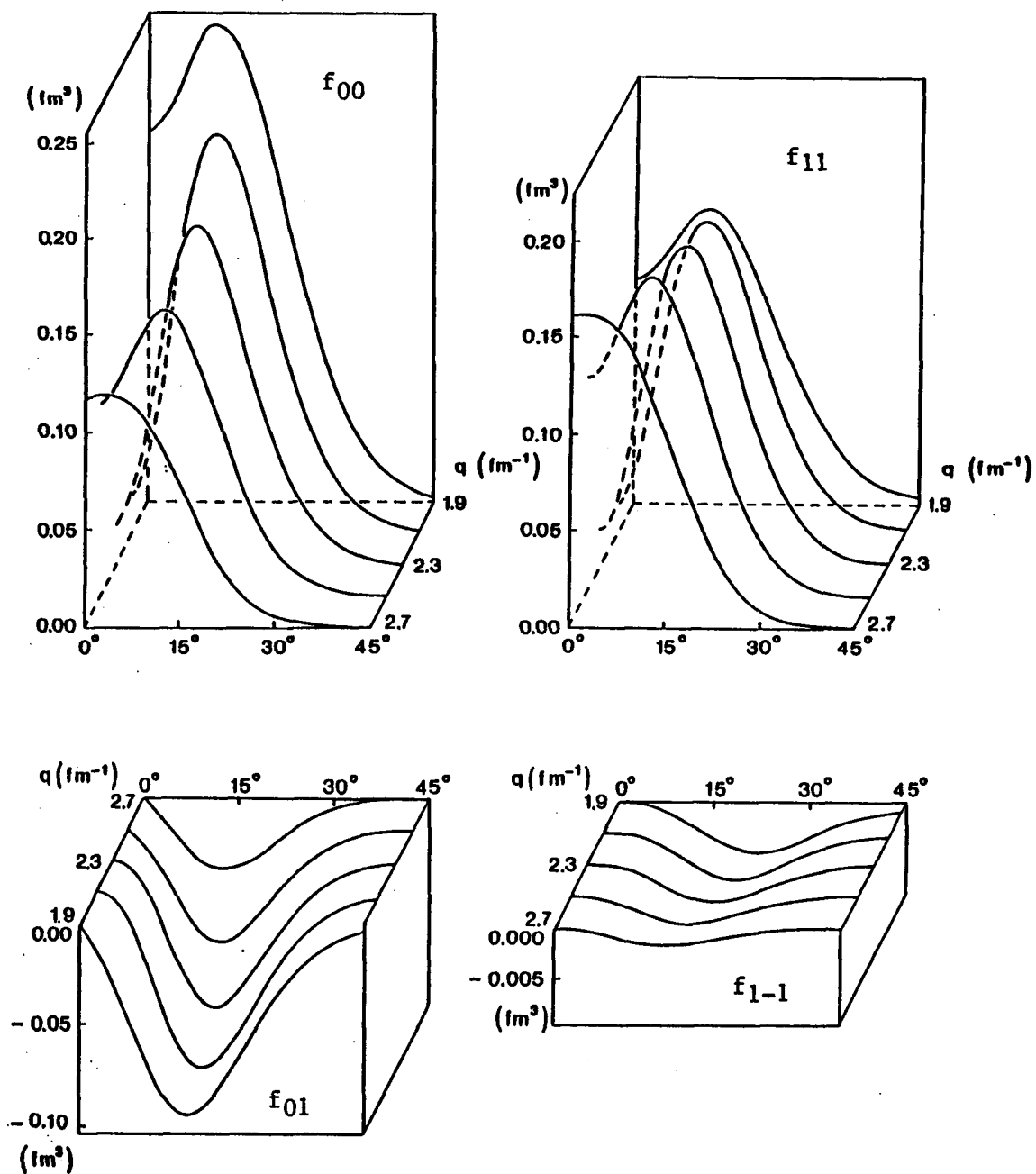


Fig. 12. PWIA structure functions of the $p_{1/2}$ hole in ^{16}O with an emitted proton of 100 MeV. Bound state from ES 23.

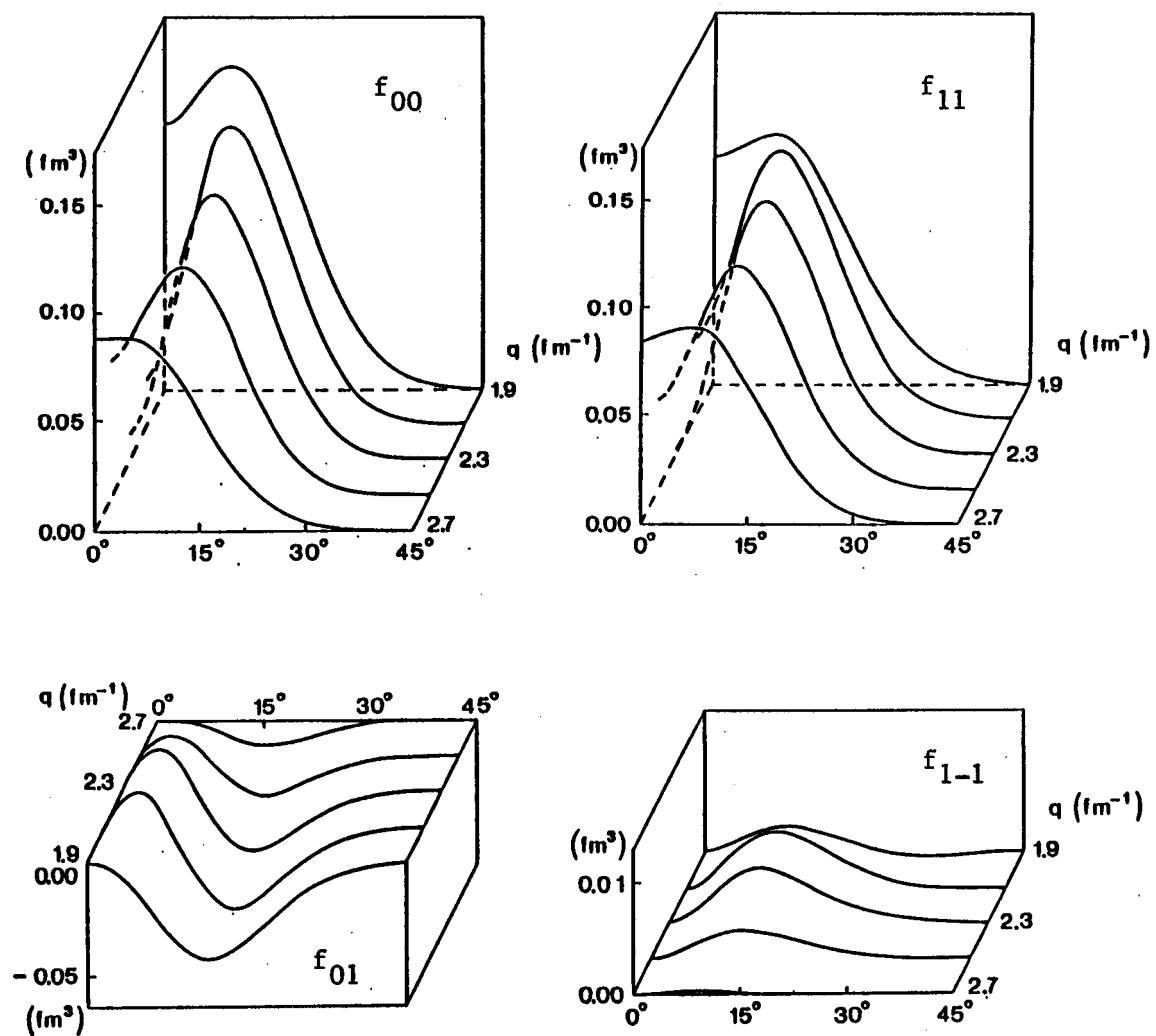


Fig. 13. The same structure functions of fig. 12 but in DWIA. Optical model potential from JA 35.

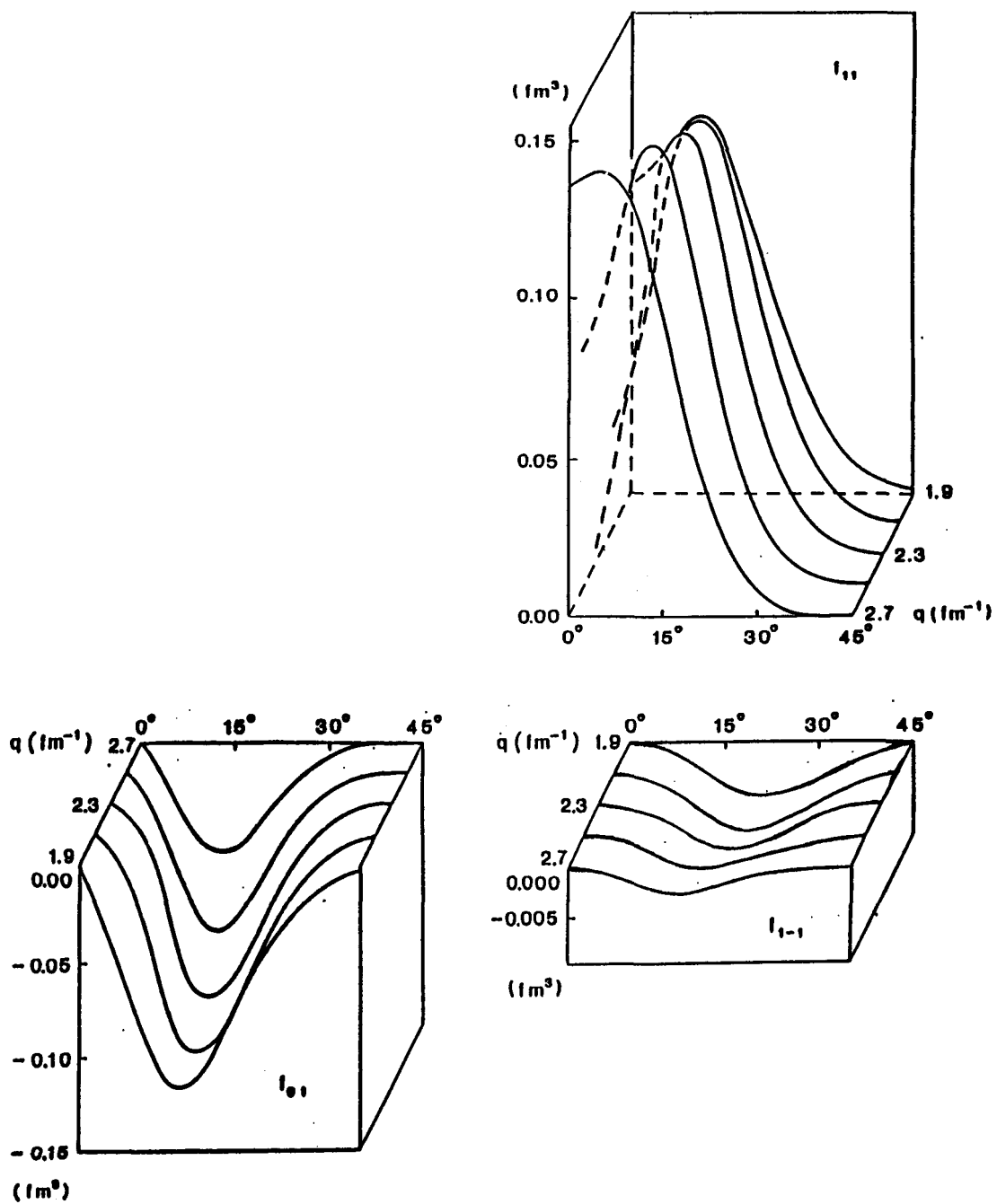


Fig. 14. PWIA structure functions of fig. 12 modified by MEC and IC.

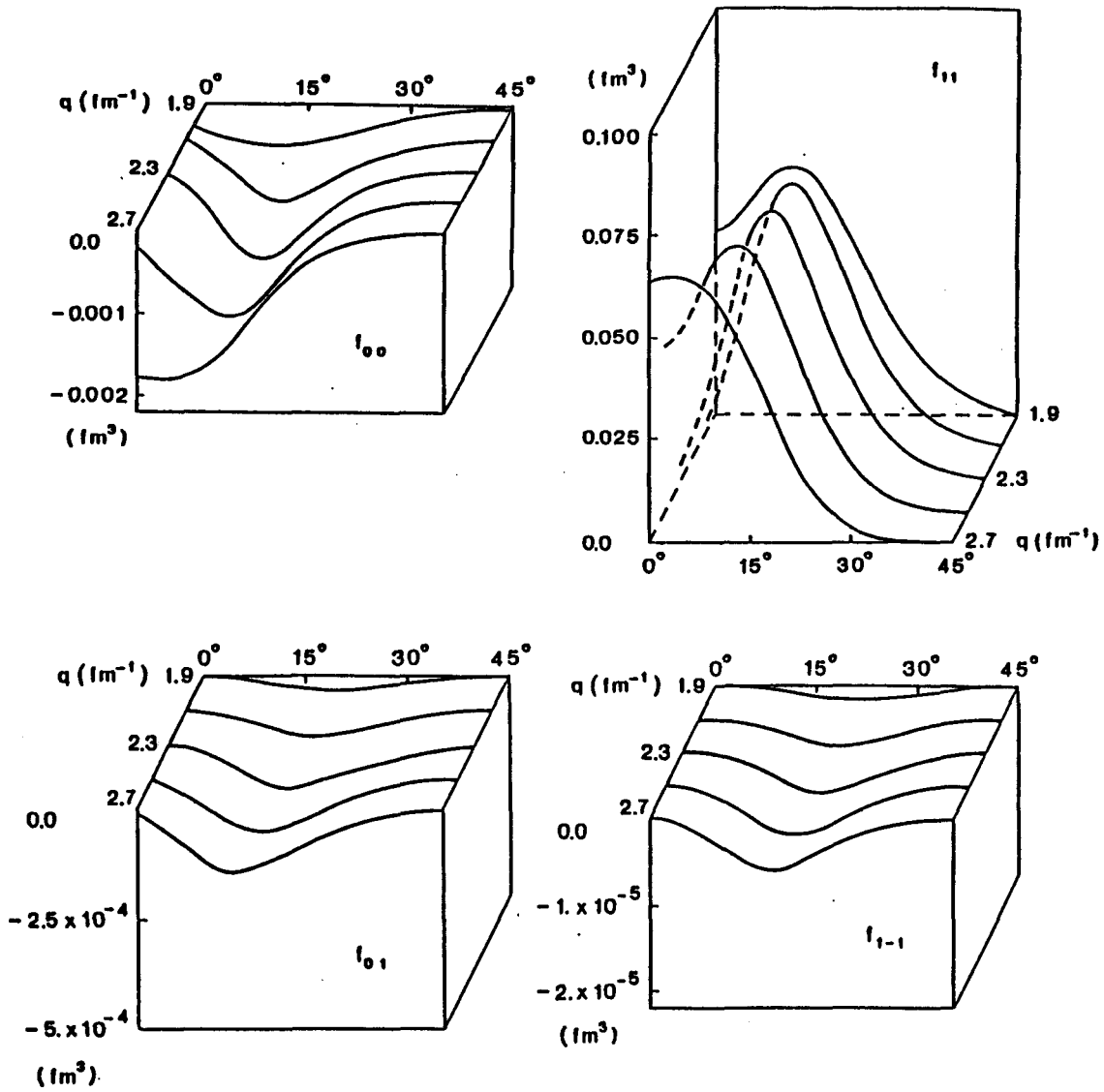


Fig. 15. PWIA structure functions of the $p_{1/2}$ hole in ^{16}O with an emitted neutron of 100 MeV. Bound state from ES 23.

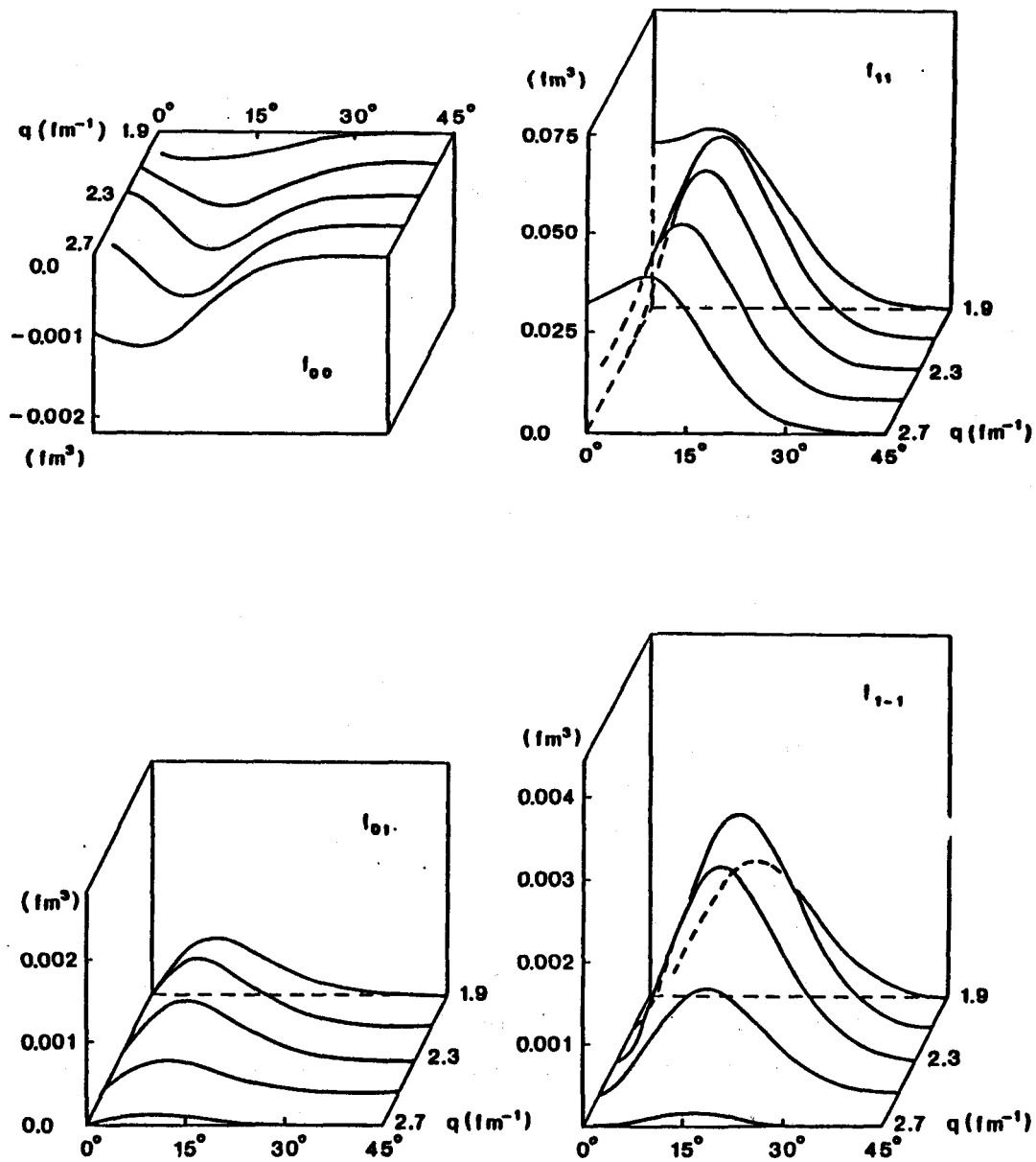


Fig. 16. The same structure functions of fig. 15. but in DWIA. Optical model potential from JA 35.

respect to $\vec{q}$) at different values of q . The bound state $p_{1/2}$ wave function is taken from ES 23 here and in all the results presented in this section. The effect of final state interactions is seen in fig. 13 where the same structure functions are calculated with the DWIA procedure of sect. 2 and ref. 8. The optical potential is the phenomenological potential JA 35. The effect of the distortion is mainly a reduction factor for the pure longitudinal and pure transverse structure functions. For the (smaller) interference terms a change of sign and a modification of the shape can occur as a consequence of their sensitivity to the coupling between the nuclear one-body spin (magnetic) current and the spin-orbit term in the optical potential 8.

MEC and IC do not affect f_{00} in the model of sect. 4. Therefore in fig. 14 only f_{11} , f_{1-1} and f_{01} appear. The role of MEC is seen from f_{01} , where IC do not contribute. In fact the contact term enhances the size of the response with respect to the one-body current by an amount which is almost cancelled by the pion-in flight contribution. This occurs also in f_{11} and f_{1-1} . The IC contribution then enhances f_{1-1} and reduces f_{11} . The effect is appreciable with the parameters of Peccei 30. With the parameters of the static quark model 31,32 agreement with the results of ref. 36 is achieved where a 25% reduction at $q=2 \text{ fm}^{-1}$ is found in the transverse response of the inclusive quasi-free electron scattering 37. The reductive effect of IC has been found also in deuteron electrodisintegration 38,39. However, the correction due to IC is of less importance if higher order terms due to isobar degrees of freedom are taken into account in the deuteron wave function 39. Here, these corrections are beyond the scope of this preliminary analysis.

The same standard kinematics is used for the $^{16}\text{O}(e,e'n)$ reaction. One can compare the corresponding PWIA structure functions given in fig. 15 with the DWIA results in fig. 16 to learn the distortion effects in the frame of direct knockout mechanism in the impulse approximation. Final state interactions have a significant role in this case, especially on the interference structure functions. Fig. 17 gives the PWIA results when MEC and IC are added. As for protons the

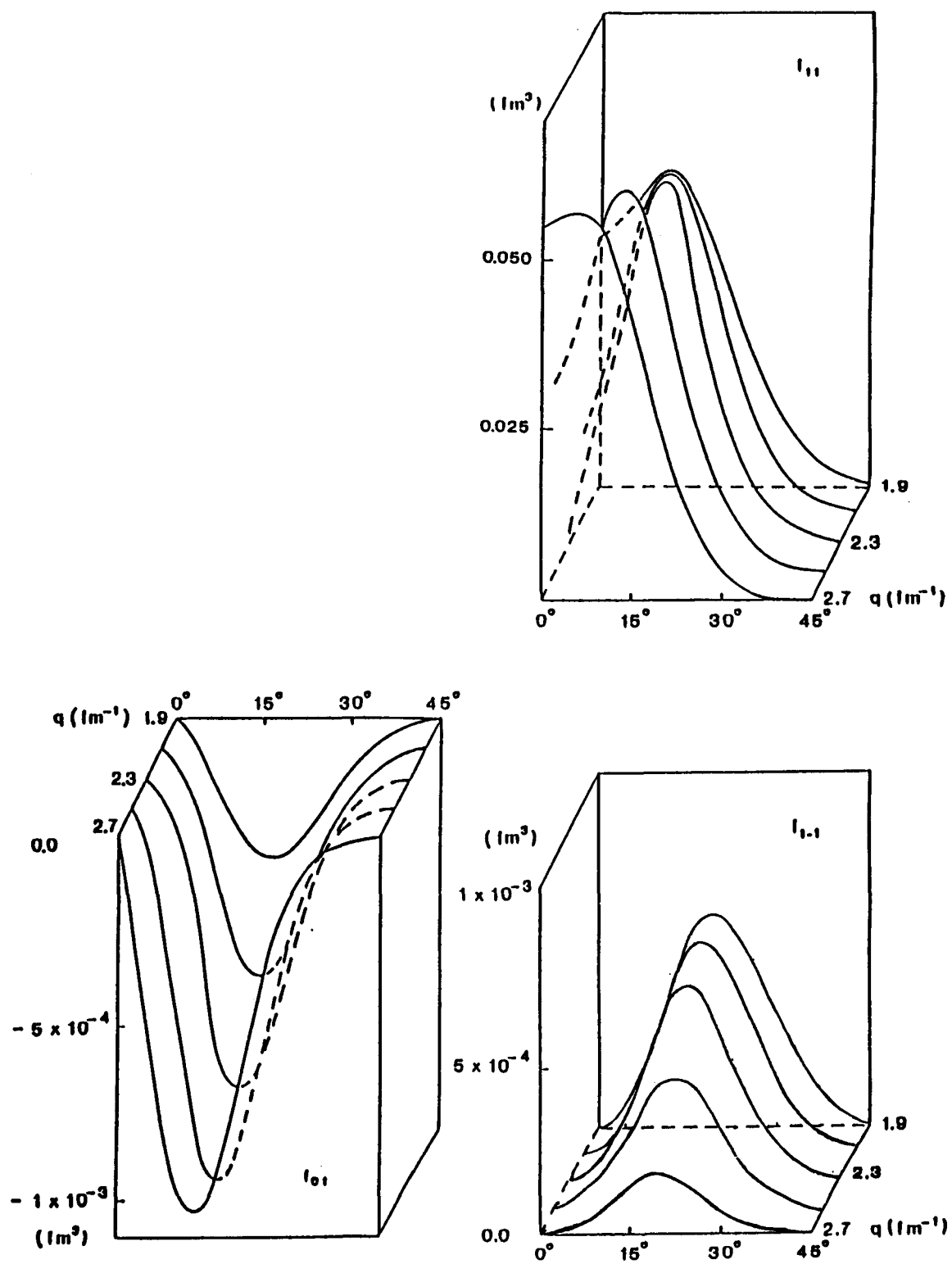


Fig. 17. PWIA structure functions of fig. 15 modified by MEC and IC.

contact and the pion-in-flight contributions are of opposite sign, but the net effect is now an order of magnitude enhancement of the size of the PWIA f_{01} . MEC and IC completely determine f_{1-1} whose sign is the opposite one of the proton f_{1-1} . These effects stem from the fact that the electromagnetic neutron form factor F_1 is very small in this conditions. Therefore the relative IC reduction on f_{11} is also much smaller than in the proton case.

At higher values of the three-momentum transfer and of the ejectile momentum the nuclear response becomes more and more peaked around $\gamma=0$, i.e. "parallel". This indicates that to study high energy and momentum transfers the parallel kinematic conditions are appropriate. This also means that only the pure longitudinal and the pure transverse structure functions survive. In the model considered here f_{00} is entirely related to the electromagnetic nucleon form factors (see eq. (32)), whereas MEC and IC can modify f_{11} . Therefore a longitudinal/transverse separation will give information on a possible modification of the nucleon in the nuclear medium.

In fig. 18 the ^{12}C transverse response is plotted for $s_{1/2}$ hole with kinematic conditions under study at BATES ⁴⁰, i.e. $q=1000$ MeV/c and $\omega = 320$ MeV and 450 MeV. This corresponds to $|p|\approx 200$ MeV/c and 0 MeV/c, respectively, with the consequence that a factor of 20 between the two kinematics comes from the different size of the involved value of the momentum distribution.

In fig. 19 the ^{40}Ca response for $d_{3/2}$ hole is given with kinematic conditions similar to those employed at Saclay to achieve the longitudinal/transverse separation ⁴¹. In fig. 20 the f_{11} response is plotted for ejectile energy of 300 MeV. The parallel kinematics at fixed ejectile momentum emphasizes the q -dependence of the response. The range of q in fig. 19 is from about 150 MeV/c to 740 MeV/c, whereas in fig. 20 it is between 510 MeV/c and 1100 MeV/c. Correspondingly, the size of the one-body response is determined by the size of the electromagnetic nucleon form factors. In particular, f_{00} is appreciable only where also F_1 is, i.e. for $(e,e'p)$ in fig. 19. In contrast, the transverse response f_{11} is highly dominated by the one-body spin (magnetic) current and the IC contribution. Therefore,

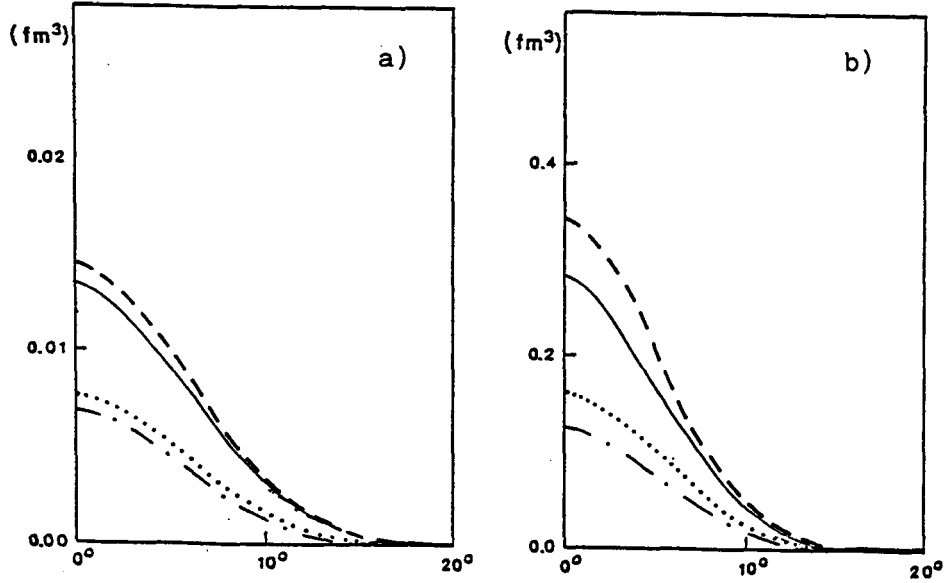


Fig 18. f_{11} of ^{12}C for $s_{1/2}$ hole at $q=1000$ MeV/c and a) $\omega=320$ MeV, b) $\omega=450$ MeV. Solid line for protons and dot-dashed line for neutrons; the one-body contribution in the two cases is represented by the dashed and the dotted line, respectively.

although small, it is not negligible under the kinematic conditions of fig. 20. Also, the $(e,e'p)$ and $(e,e'n)$ responses show a tendency to become comparable with increasing momentum and energy transfers as a consequence of their transverse character.

6. CONCLUSIONS

In this talk two new aspects in nucleon knockout by an electromagnetic probe have been considered. The first one concerns spin degrees of freedom and the possibility of performing experiments with both incident polarized electrons and polarized targets. The accuracy required to separate all the nine structure functions which characterize the nuclear response completely is probably too high for the available instrumentation in the near future. On the other side already interesting information could be gained by measuring asymmetries and electron polarization ratios. These quantities do not

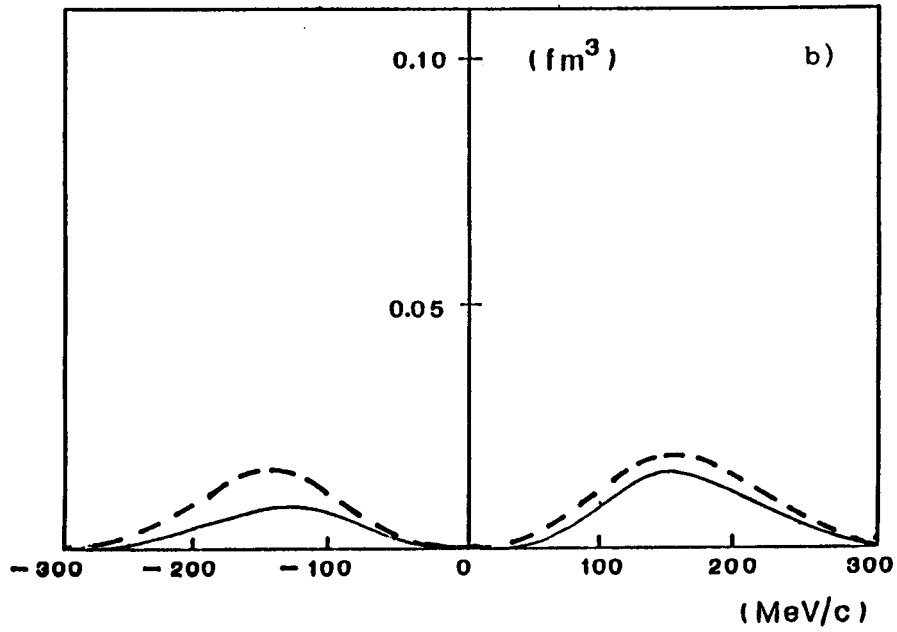
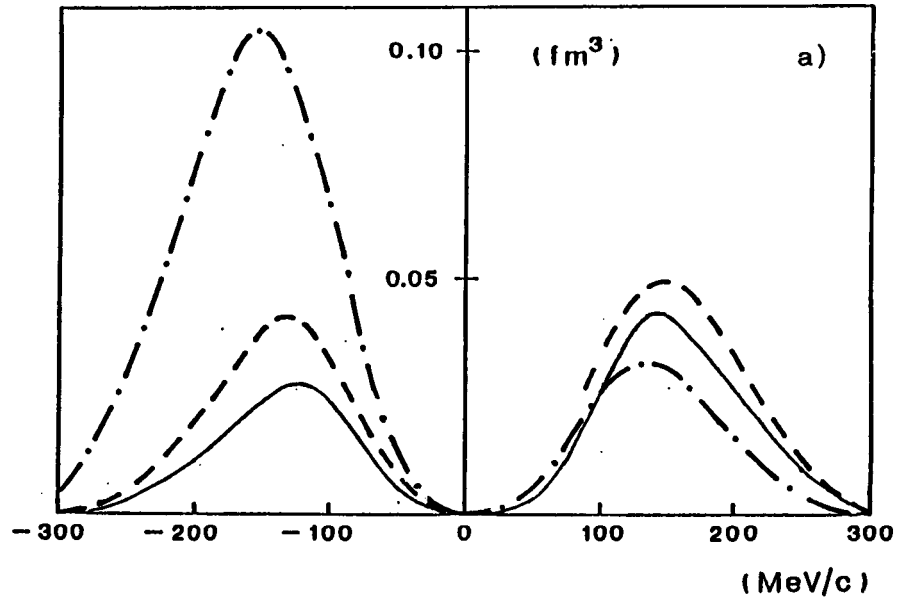


Fig. 19. Longitudinal (f_{00} , dot-dashed line) and transverse (f_{11} , solid line) response of ^{40}Ca for $d_{3/2}$ hole with ejectile energy of 100 MeV: a) protons, b) neutrons. The dashed line is the one-body contribution to f_{11} .

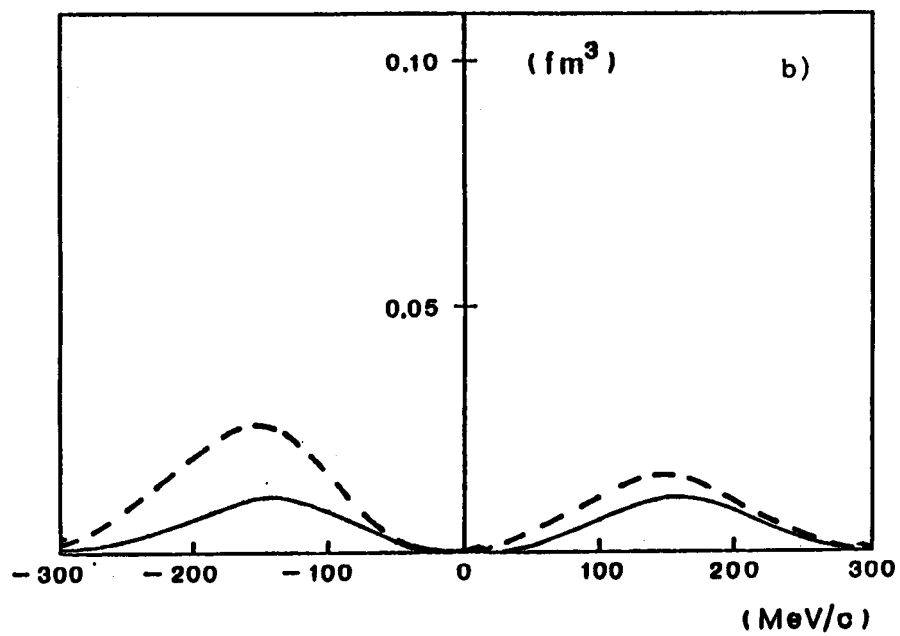
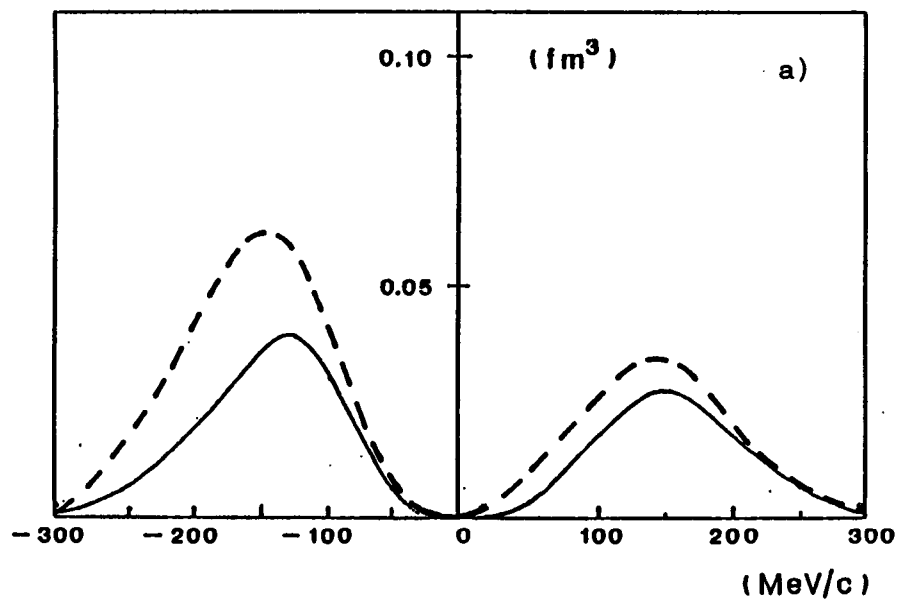


Fig. 20. The same as in fig. 19 but for 300 MeV ejectile energy.

require out-of-plane kinematics necessarily and appear to have measurable size and a fair sensitivity to final state interactions.

The second aspect is related to the advent of new facilities, like CEBAF, which will allow higher momentum and energy transfers. In the preliminary analysis presented here MEC and IC terms modifying the nuclear current have been considered to lowest order. Therefore, only the transverse part of the current is affected by two-body contributions. The contact term and the pion-in-flight terms, which are the only surviving MEC contributions in the nonrelativistic approximation used here, have compensating effects. The IC terms increase with q and reduce the response of both $(e,e'p)$ and $(e,e'n)$ reactions. With increasing (q,ω) the response becomes more and more transverse and concentrated along the direction of $\vec{q}$. Therefore, the combination of the one-body spin (magnetic) current with IC terms smoothes the difference between the $(e,e'p)$ and the $(e,e'n)$ response.

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Dynamic Structure Functions in Nuclear Medium

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A recently developed theory for the dynamic structure function of infinite Fermi systems is applied to the calculation of the longitudinal response function in nuclear matter from realistic interactions. The one particle-one hole correlated states are exactly included in the calculation, and the two particle-two hole states are also taken into account via the real and imaginary part of the optical potential. The calculated dynamic structure function is then used to estimate, at large values of the momentum transfer k , the longitudinal response function observed in e-scattering experiments. There is agreement between theory and experiments at $k = 300$ and 400 MeV/c for ^{40}Ca , ^{48}Ca and ^{238}U . At higher momentum transfer values, as $k = 550$ MeV/c, the calculated response agrees with the data on ^{40}Ca and ^{56}Fe at small values of the loss energy ω , whereas it is larger than experiment at large ω .

1. - Introduction.

It has recently been suspected that the intrinsic properties of the nucleon are substantially modified by the nuclear medium. Recent accurate measurements on scattering experiments of leptons from nuclei have raised this suspicion. The main effects observed are: (i) the quenching of the single particle strength in both elastic and inelastic e-scattering^{1,2}; (ii) the quenching of the longitudinal response function in (e,e') reactions³⁻⁵; (iii) The EMC-SLAC effect^{6,7}.

The claimed discrepancies between experimental data and theoretical estimates heavily relies on theoretical treatments of the nucleus based on

the independent particle model picture. The question we address ourselves in this contribution is to what extent the quenching observed in e-scattering experiments can be explained by means of a more realistic treatment of the nucleus, described by a conventional non-relativistic hamiltonian of the type:

$$H = \hbar^2/2m \sum \nabla_i^2 + \sum v_{ij} + \sum v_{ijk} \quad (1.1)$$

The two-nucleon potential $\sum v_{ij}$ may be any of the known realistic interactions like *Paris* ⁸⁾, *Urbana* v_{14} ⁹⁾, *Argonne* v_{14} ¹⁰⁾ etc.. which reproduces the N-N data up to ≈ 400 MeV, and the three nucleon interaction $\sum v_{ijk}$ is included to give a reasonable description of ^3H , ^3He , ^4He and nuclear matter ¹¹⁾.

Realistic calculations performed both in nuclear matter and few body nuclei with the hamiltonian (1.1) have shown large correlation effects

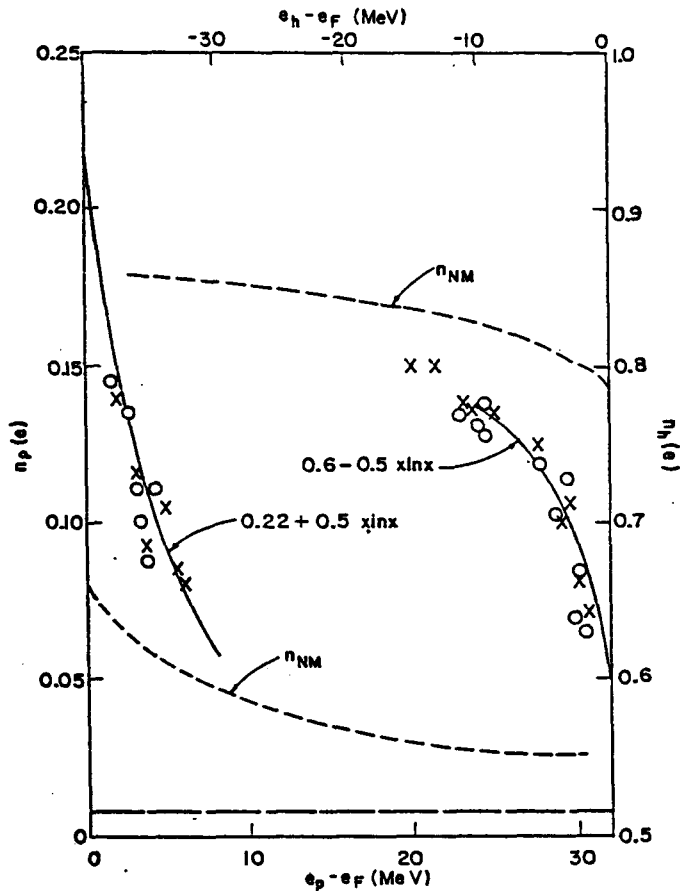


Fig.1 - Taken from ref.(14). Occupation probabilities $n(e)$ of single particle states. The dashed line gives the nuclear matter results. The crosses and circles denote the occupation numbers for proton and neutron states obtained by adding to n_{NM} the contributions due to the coupling between the single particle states and the surface oscillations.

which are strongly state dependent and whose inclusion in a mean field type of description of the system is not at all clear. It has been found¹¹⁻¹³⁾ that the ground state wave function contains an high percentage of excited two-nucleon configurations which strongly affect the nucleon momentum distribution $n(k)$. Fig.1 shows the occupation probabilities $n(e)$ extracted¹⁴⁾ from a microscopic calculation of the momentum distribution of nucleons in correlated nuclear matter. A comparison with the step function which is the independent particle model estimate of $n(e)$ evidentiates the correlation effects. The average value n^- of $n(e)$ for states in the shell just below the Fermi energy is ≈ 0.7 and n^+ is ≈ 0.1 . The quenching for elastic e-scattering may be approximated by $n^- - n^+ \approx 0.6$ and that for inelastic e-scattering by $n^-(1 - n^+) \approx 0.6$, which are in qualitative agreement with experiments.

In the present contribution we discuss the results obtained¹⁵⁾ in a microscopic calculation of the dynamic structure function $S(k, \omega)$ of nuclear matter, defined as

$$S(k, \omega) = \sum_n |\langle 0 | \rho_k^+ | n \rangle|^2 \delta(\omega - \omega_n) \quad , \quad (1.2)$$

$$\rho_k = \sum_i \exp(ik \cdot r_i) \quad , \quad (1.3)$$

where $|n\rangle$ are eigenstates of the hamiltonian (1.1), and $\omega_n = E_n - E_0$ are the excitation energies. Correlated basis function theory¹⁶⁻¹⁸⁾ is used to calculate $S(k, \omega)$. The method consists in using **orthonormalized** correlated states (ONCS) $|p_1 \dots p_n h_1 \dots h_n\rangle$ in a perturbation theory calculation of $S(k, \omega)$. The ONCS are obtained from the not orthogonal correlated states, given by

$$|p_1 \dots p_n h_1 \dots h_n\rangle = (1/\sqrt{\text{Norm.}}) (S \prod F_{ij}) |p_1 \dots p_n h_1 \dots h_n] \quad , \quad (1.4)$$

where $|p_1 \dots p_n h_1 \dots h_n]$ is a n particle-n hole Fermi gas states, and F_{ij} are two-particle correlation operators given by

$$F_{ij} = \sum f^p(r_{ij}) \text{ OP}(i, j) \quad , \quad (1.5)$$

$$O^{p=1,8}(ij) = 1; \sigma_i \sigma_j; \tau_i \tau_j; \sigma_i \sigma_j \tau_i \tau_j; S_{ij}; S_{ij} \tau_i \tau_j; L S; L S \tau_i \tau_j, \quad (1.6)$$

and S is the operator which symmetrizes the product $\prod F_{ij}$. The orthogonalization of the states (1.4) is done by a proper use of the Schmidt and Löwdin orthogonalization procedures^{15,19}. The diagonal and not diagonal matrix elements of the hamiltonian and the identity operators are calculated by using FHNC/SOC techniques^{20,21}. The correlation operator F_{ij} which fixes the ONCS basis is obtained by requiring that $|0\rangle$ gives the lowest variational binding energy per nucleon.

The main advantage of this theory, with respect to other approaches based upon mean field theory^{22,23}, is that realistic interactions can be used in the calculation, so that the effect of correlations can be studied explicitly. The perturbative series is expected to converge rapidly. For instance, the second order correction to the binding energy of nuclear matter is only $\approx 1.3 \text{ MeV}$ ²⁴ and second order perturbation theory gives good results for the optical potential of a nucleon in nuclear matter^{24,25}.

The calculated dynamic structure function is quenched with respect to the Fermi gas evaluation and its quasi elastic peak is shifted towards higher values of ω . At momentum transfer $k \geq 300 \text{ MeV/c}$ the particle-hole effective interaction has little effect on $S(k, \omega)$. In this region the main ingredients of the calculation result to be : (i) the vertex factors $\langle 0 | \rho_k^\dagger | ph \rangle$ which are of the form $\delta(k+h-p) X_0(p, h, k)$, where the function X_0 can be quenched of $\approx 15\%$ with respect to 1; (ii) the single particle excitation energies which are taken from the calculated real part of the optical potential²⁴; (iii) the spreadings of the single particle energies which give the main part of 2p-2h states contributions and have been taken from the calculated imaginary part of the optical potential²⁵. Note that the calculated and the experimental optical potential do not differ appreciably. Complex optical potentials have also been used by *DellaFiore, Lenz and Brieva*²² to study the response of ^{12}C .

The longitudinal response function, observed in electron scattering

experiments, apart from form factors and relativistic corrections, is given by the response of the nuclear matter to the proton density fluctuation operator:

$$\rho_{L,k} = \sum_i \exp(ik \cdot r_i) (1 + \tau_3(i))/2 \quad . \quad (1.7)$$

At large values of k the longitudinal response function $R_L(k, \omega)$ can be extracted from the $S(k, \omega)$ calculated in this work and compared with the available experimental data on ^{40}Ca , ^{48}Ca , ^{56}Fe and ^{238}U nuclei.

The theory is fully derived in refs. (15,19) and briefly summarized in Section II. The results are presented and discussed in Section III.

2. - Correlated basis function theory of $S(k, \omega)$.

It is convenient¹⁸⁾ to calculate the dynamic structure function $S(k, \omega)$ from the density-density correlation function:

$$S(k, \omega) = \text{Im } D(k, \omega) / \pi \quad , \quad (2.1)$$

$$D(k, \omega) = \langle 0 | \rho_k^\dagger (H - E_0 - \omega - i\eta)^{-1} \rho_k | 0 \rangle \quad , \quad (2.2)$$

where $|0\rangle$ is exact ground state of the hamiltonian (1.1) and E_0 the corresponding energy eigenvalue.

The perturbation theory is built up by separating H into an unperturbed part H_0 and an interaction term H_I defined by the following two equations:

$$\langle m | H_0 | n \rangle = \langle m | H | m \rangle \delta_{mn} = H_{mm} \delta_{mn} \quad , \quad (2.3)$$

$$\langle m | H_I | n \rangle = \langle m | H | n \rangle (1 - \delta_{mn}) = \bar{H}_{mn} \quad , \quad (2.4)$$

where $|m\rangle$ denotes a generic ONCS. Note that one should use the exact ground state rather than the ONCS $|0\rangle$. Short range correlations are already included into the ONCS, therefore $|0\rangle$ is very close to the true ground state and it can be corrected perturbatively^{12,24,25).}

The $D(k, \omega)$ is given by the following perturbative series:

$$\begin{aligned}
D(k, \omega) &= \sum_n (-)^n \langle 0 | \rho_k^\dagger (H_0 - E_0 - \omega - i\eta)^{-1} \{ H_I (H_0 - E_0 - \omega - i\eta)^{-1} \}^n \rho_k | 0 \rangle \\
&= \sum_i X_0^*(i) G_0(i) X_0(i) - \sum_{ij} X_0^*(i) G_0(i) \bar{H}_{ij} G_0(j) X_0(j) + \dots, \quad (2.5)
\end{aligned}$$

where

$$X_0^*(i) = \langle 0 | \rho_k^\dagger | i \rangle, \quad (2.6)$$

$$G_0(i) = (H_{ii} - E_0 - \omega - i\eta)^{-1}. \quad (2.7)$$

The series (2.5) is formally summed to all orders by the following Dyson-like integral equation:

$$X(i) = X_0(i) - \sum_j \bar{H}_{ij} G_0(j) X(j), \quad (2.8)$$

$$D(k, \omega) = \sum_i X_0^*(i) G_0(i) X(i). \quad (2.9)$$

Successive approximation to $D(k, \omega)$ are obtained by including in the sum 1p-1h ONCS, 2p-2h ONCS etc.. Diagrams (a) and (b) of Fig.2 represent two of the perturbative terms containing 1p-1h ONCS only. Diagrams (c-g) are examples of diagrams having 2p-2h ONCS insertions also.

In this contribution we present results on the dynamic structure function calculated in the three following approximations: (i) $S_{ph}^\circ(k, \omega)$ is obtained when only diagram (a) of Fig.2 is included in the calculation of $D(k, \omega)$ and therefore $X(i)$ is taken to be equal to $X_0(i)$ in eq.(2.9) with i being restricted to the 1p-1h ONCS only; (ii) $S_{ph}(k, \omega)$ is obtained when the integral equation (2.8) is fully solved within the the space of the 1p-1h ONCS; (iii) $\bar{S}(k, \omega)$ includes the main contributions due to 2p-2h ONCS, namely the coupling of particle (hole) states to 2p-1h (1p-2h) ONCS as in diagrams (c) and (d) of Fig.2. This can be done by substituting to the variational single particle energy $e_v(i)$, appearing in the Green function $G_0(i)$ (note that $H_{ii} - E_0 = e_v(p_i) - e_v(h_i)$), those obtained in second order CBF perturbation theory:

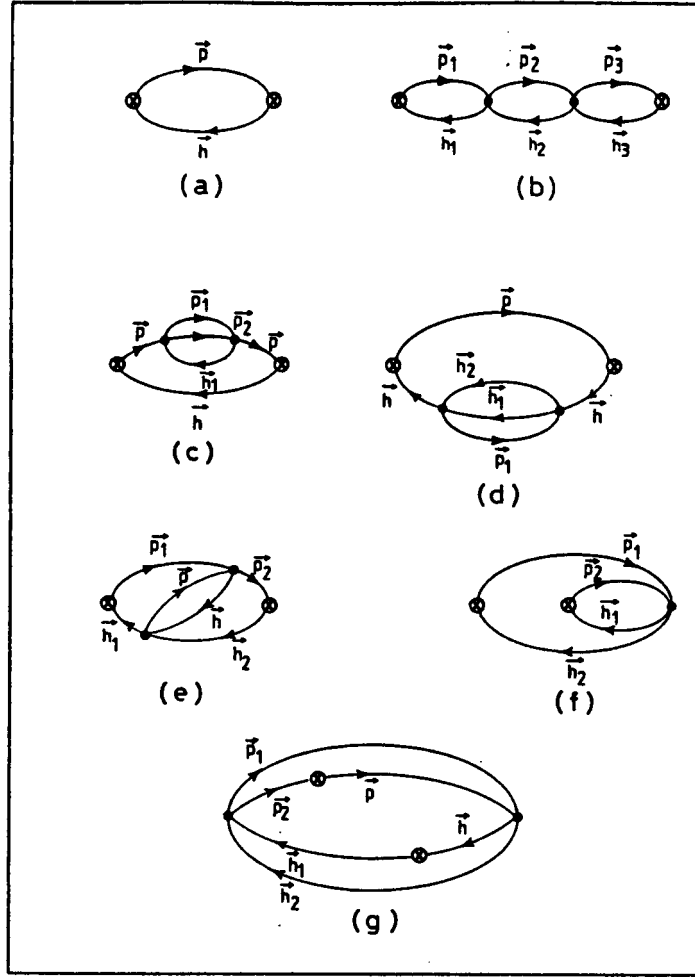


Fig.2 - Correlated basis diagrams contributing to $S(k, \omega)$.
The circled crosses and the dots denote the vertices X_0 and the non diagonal matrix elements $\bar{H}_{ij}$ respectively.

$$e(p_i, k, \omega) = e_v(p_i) + \Delta U(p_i, k, \omega) + iW(p_i, k, \omega) \quad , \quad (2.10)$$

$$e_v(q) = (\hbar q)^2 / 2m + U_v(e_v) \quad , \quad (2.11)$$

with $U_v(e)$ taken from Fig.3, and there is a similar equation for the single hole energy $e_v(h_i, k, \omega)$. A reasonable estimation of these effects on the dynamic structure function at large k is obtained by using ΔU and W from the optical potential in nuclear matter^{24,25}). The quantity $U_2 - U_v$ of Fig.3 is used for ΔU and W_0 of Fig.4, divided by the effective mass m^* ²⁶), is

is used for ΔU and W_0 of Fig.4, divided by the effective mass m^* ²⁶), is taken to approximate W in the above equation. The main effect of the imaginary part W is to spread the response S_{ph} at energy ω over a width

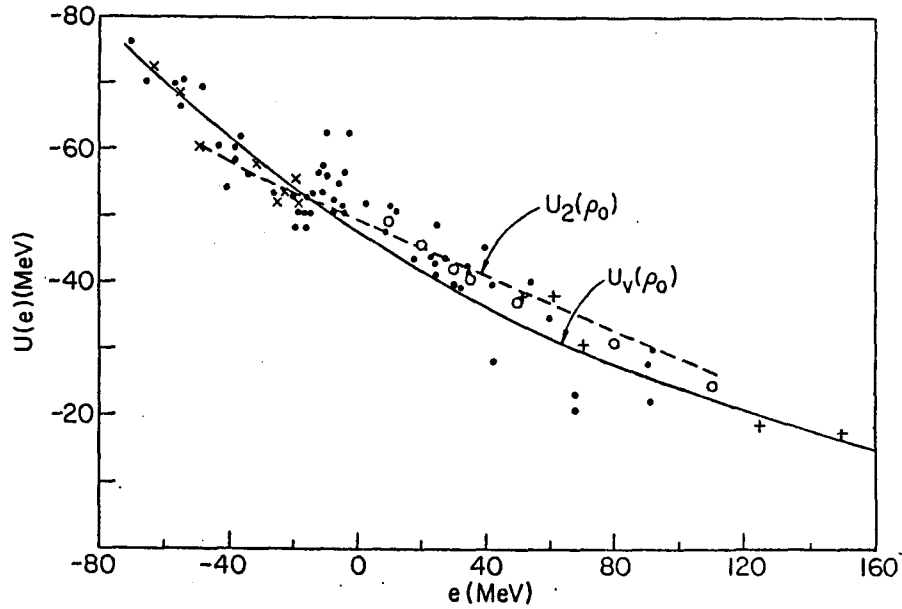


Fig.3 - Taken from ref.(24). The calculated real part of the optical potential is compared with experimental data.

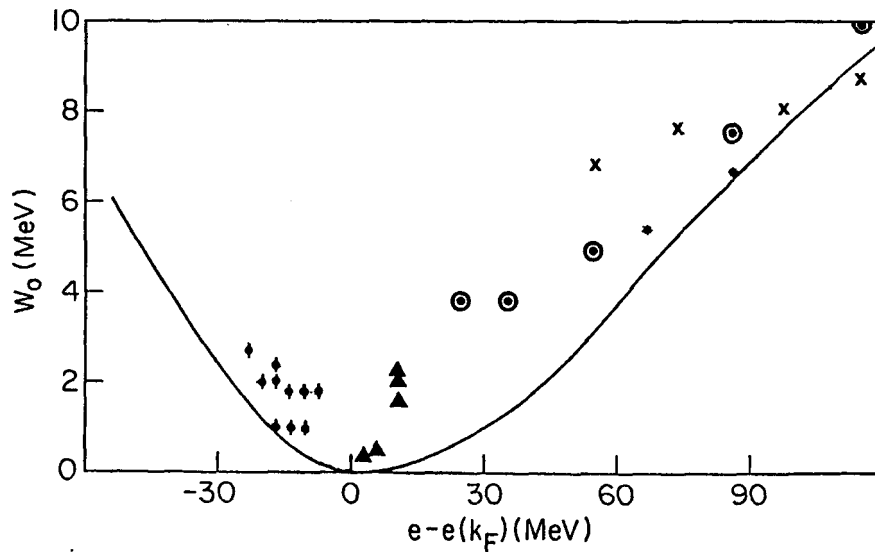


Fig.4 - Taken from ref.(25). The calculated imaginary part of the optical potential is compared with experimental data.

W. In applying the optical potential approximation we implicitly assume the width to depend on the energy ω only and obtain:

$$\overline{S}(k, \omega) = (1/\pi) \int d\omega' S_{ph}(k, \omega') W(\omega) / [(\omega' - \omega)^2 + W^2(\omega)] \quad . \quad (2.12)$$

2p-2h ONCS contributions coming from diagrams of the type (e-g) of Fig.2 have been disregarded in the present calculation. Diagrams of the type (e) modify the particle-hole effective interaction and diagrams (f,g) are generated when the $|0\rangle$ is corrected for 2p-2h admixtures^{24,25}.

The procedure to calculate the ONCS $|i\rangle$ is fully described in refs. (15,19) and will not be repeated in this contribution. Here we limit ourselves to stress that the orthogonalization amongst the states is necessary in order that both the small and large k limits of the vertex factor $X(i)$ are correctly described. This has important consequences on the sum rules: if the pair correlations F_{ij} have no momentum dependence (as in the case that the $L \cdot S$ components in eq.(1.5) are neglected) $S \prod F_{ij}$ commutes with ρ_k and therefore the state $\rho_k |0\rangle$ is a superposition of 1p-1h ONCS only. It follows that $\rho_k |0\rangle$ is not coupled to other types of excited states, as for instance 2p-2h ONCS, and the sum rule:

$$\int_0^\infty d\omega S(k, \omega) = S(k) = \langle 0 | \rho_k^\dagger \rho_k | 0 \rangle \quad , \quad (2.13)$$

is exhausted by the 1p-1h ONCS, therefore one has

$$\int_0^\infty d\omega S_{ph}^\circ(k, \omega) = \int_0^\infty d\omega S_{ph}(k, \omega) = S(k) \quad . \quad (2.14)$$

If the L^2 and the $(L \cdot S)^2$ are neglected in the interaction the following equation is satisfied:

$$[\rho_k^\dagger [H, \rho_k]] = A(\hbar k)^2 / 2m \quad , \quad (2.15)$$

which leads to the following energy weighted sum rule:

$$\int_0^\infty d\omega \omega S_{ph}(k, \omega) = (\hbar k)^2 / 2m \quad . \quad (2.16)$$

3. - Results.

The results presented in this section have been obtained for the Urbana v_{14} plus three-nucleon interaction-model^{11,23)} of nuclear matter at $k_F=1.33\text{fm}^{-1}$, under the following approximations: (i) the 1p-1h ONCS are exactly taken into account in the calculation, whereas the main 2p-2h ONCS (expressed by diagrams of the type (c,d) of Fig.2) are included via a complex optical potential (see eq.(2.11)); (ii) diagrams of the type (e-g) of Fig.2 are omitted, which is equivalent to say that the ONCS $|0\rangle$ is considered to be the exact ground state; (iii) $L\cdot S$ correlations are neglected in the calculation of $X(i)$; this implies that the sum rule (2.13) should be satisfied; (iv) the contribution of the three-nucleon interaction to $\overline{H}_{ij}$ has been approximated as in ref. (25).

The dynamic structure functions $S(k,\omega)$ at $k=0.665, 2.03$ and 3.99 fm^{-1} are given in Figs.5-7. The curves S_{FG} show the Fermi-gas estimates. The curves S_{MF} are obtained from $S_{ph}^{\circ}(k,\omega)$ by setting the vertex factors

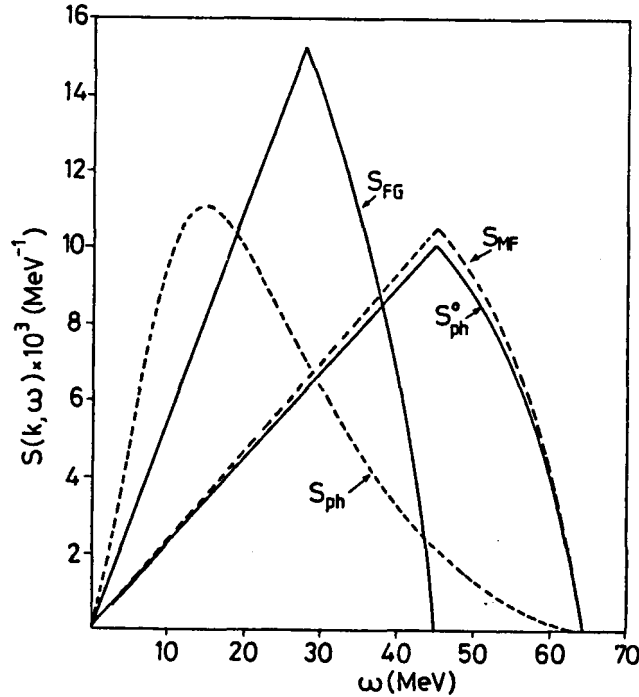


Fig.5 - Taken from ref.(15). Dynamic structure function of nuclear matter at $k = k_F/2=0.665\text{ fm}^{-1}$ in various approximations.

$X_0(i)$ equal to unity, namely

$$S_{MF}(k, \omega) = \sum \theta(|\mathbf{h}_i + \mathbf{k}| - k_F) \delta(e_v|\mathbf{h}_i + \mathbf{k}| - e_v(\mathbf{h}_i) - \omega) \quad (3.1)$$

The curves S_{MF} correspond to mean field calculations done with single particle (hole) energies taken from U_v of Fig.3, and differ from S_{ph}^o only by $|X_0(i)|^2$. In ref. (15) it is proved that X_0 has to be unity in both the limits $k \rightarrow 0$ and $k \rightarrow \infty$. In fact the two curves are quite close at $k = 0.665$ and 3.99 fm^{-1} ; at $k = 2.03 \text{ fm}^{-1}$ the vertex factors cause up to $\approx 12\%$ quenching.

A comparison between S_{ph}^o and S_{ph} shows that at low momentum transfers the particle-hole interaction H_{ij} has large effects in the dynamic structure function, but, already at $k = 2.03 \text{ fm}^{-1}$ it is negligible. This also implies that at $k > 2.03 \text{ fm}^{-1}$ diagrams of the type (e-g) of Fig.2 give a small contribution, therefore the present calculation may be adequate in this region. Drözd et al.²⁷⁾ have studied the effects of 2p-2h states on the longitudinal response function of ^{12}C and find that at $k = 300 \text{ MeV}/c$ the main effects come from diagrams of the type (c,d) of Fig.2.

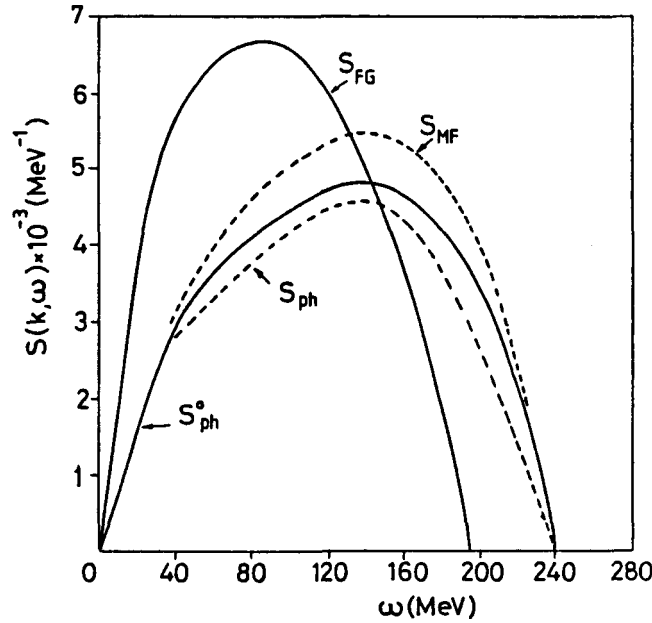


Fig.6 - Same as Fig.5 but at $k = 2.03 \text{ fm}^{-1}$.

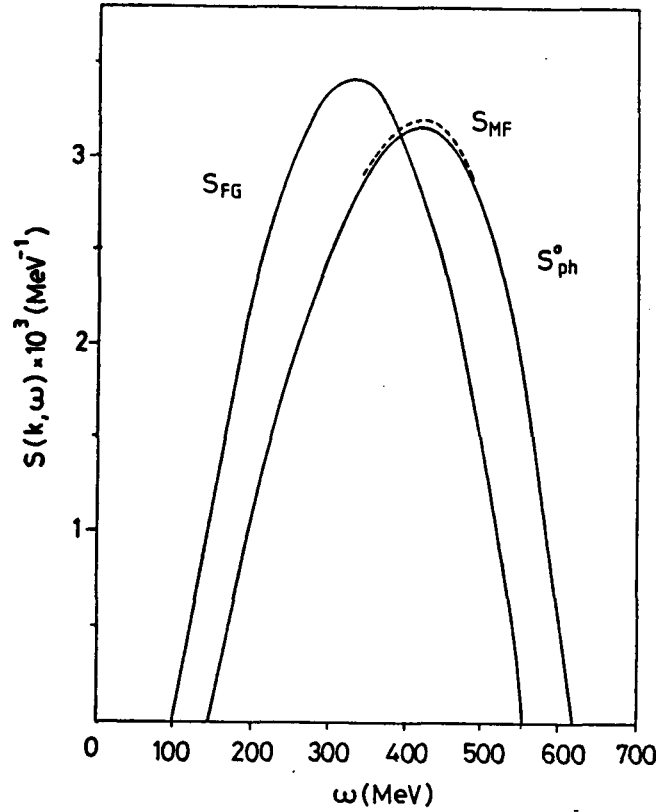


Fig.7 - Same as Fig.5 but at $k=3k_F = 3.99\text{fm}^{-1}$.

Table I displays the results for the sum rules at various values of the momentum transfer. The small differences amongst $S(k)$, $S_{ph}^o(k)$ and $S_{ph}(k)$ are partly due to numerical inaccuracy and partly to the FHNC/SOC approximation used to calculate $X_0(i)$. The ratio R given in the Table is defined by the following equation:

$$\int_0^\infty d\omega \omega S(k, \omega) = R (\hbar k)^2 / 2m \quad . \quad (3.2)$$

Since the hamiltonian used in this calculation has L^2 and $(L \cdot S)^2$ terms then eq.(2.14) is not satisfied. However these terms are weak and repulsive, therefore the ratio R is expected to be slightly larger than one.

$k \text{ (fm}^{-1}\text{)}$	$S_{FG}(k)$	$S(k)$	$S_{ph}^{\circ}(k)$	$S_{ph}(k)$	R
0.665	0.37	0.32	0.35	0.35	0.89
1.33	0.69	0.61	0.62	0.57	0.99
1.995	0.91	0.80	0.80	0.75	1.09
2.66	1.	0.88	0.91	0.84	1.12
3.99	1.	0.95	0.98	0.92	1.08

Table I - Sum rules of the dynamic structure function.

Fig.8 shows $S(k, \omega)$ calculated by using various approximations at $k=550$ Mev/c. The curve $\bar{S}_{IA}$ represents the estimate obtained from microscopic momentum distribution¹²⁾ $n(q)$ by using the impulse approximation:

$$\bar{S}_{IA}(k, \omega) = \sum n(q) \delta(e_v(|\mathbf{q}+\mathbf{k}|) - e_v(q) - \omega) \quad . \quad (3.3)$$

We have used $e_v(q)$ rather than the free single particle (hole) energies $(\hbar q)^2/2m$, since the difference $U_v(|\mathbf{q}+\mathbf{k}|) - U_v(q)$ is not always negligible in nuclear matter. At small and large ω one finds $\bar{S} \approx \bar{S}_{IA}$; in the region of

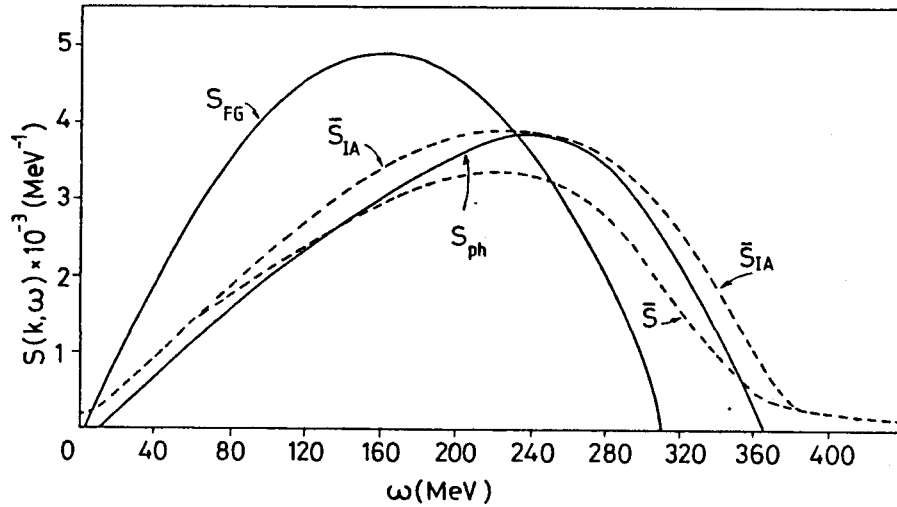


Fig.8 - Taken from ref.(15). Dynamic structure function of nuclear matter at $k = 550 \text{ MeV/c}$ in various approximations.

the quasi-elastic peak however, $\bar{S}$ is smaller than $\bar{S}_{IA}$ by $\approx 10\%$. That is consistent with the fact that the integrated structure function $\bar{S}_{IA}(k)$ is ≈ 1 , whereas $S(k) = 0.88$.

The comparison between $\bar{S}(k, \omega)$ and $S_{ph}(k, \omega)$ shows that 2p-2h ONCS reduce of $\approx 10\%$ the strength of the peak and scatter it to higher energies.

The calculation of the longitudinal response function $R_L(k, \omega)$ requires the use of the vertex factor $X_{L,0}(i) = \langle i | \rho_{L,k} | 0 \rangle$ rather than $X_0(i)$. The total integral $S_L(k)$ of the corresponding structure function $S_L(k, \omega)$ over ω , in the region $k=300$ to 550 MeV/c results to be²⁸⁾ larger than $S(k)$ not more than 9%, therefore it is reasonable to take $S_L(k, \omega) \approx \bar{S}(k, \omega)$ in this region.

The longitudinal response function $R_L(k, \omega)$ is obtained by multiplying $S_L(k, \omega)$ with the squared proton form factor. In Figs.9-11 the experimental data are compared with the responses calculated with (i) the

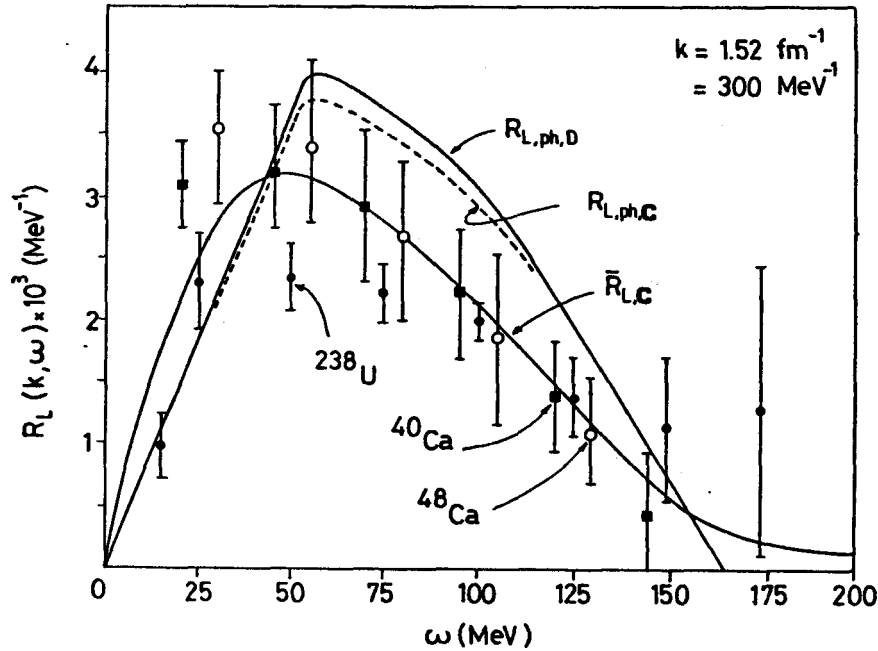


Fig.9 - Taken from ref.(15). Longitudinal response function of nuclear matter at $k = 300$ MeV/c compared with experimental data on ^{40}Ca and ^{48}Ca at $k=290\text{MeV/c}$ and ^{238}U at 300MeV/c . The error bars are large enough to contain both the *Saclay* and the *Bates* results.

dipole form factor as given in ref.(29):

$$R_{L,D}(k,\omega) = S(k,\omega) F_D^2(k,\omega) \quad , \quad (3.4)$$

$$F_D(k,\omega) = [1+(k^2-\omega^2)/(855 \text{ MeV})^2]^{-2} \quad , \quad (3.5)$$

and (ii) the *covariant* form factor as reported in ref.(30):

$$R_{L,C}(k,\omega) = S(k,\omega) F_C^2(k,\omega) \quad , \quad (3.6)$$

$$F_C(k,\omega) = [1+(k^2-\omega^2)/(832.5 \text{ MeV})^2]^{-2} [1+(k^2-\omega^2)/(1876.4 \text{ MeV})]^{-1}. \quad (3.7)$$

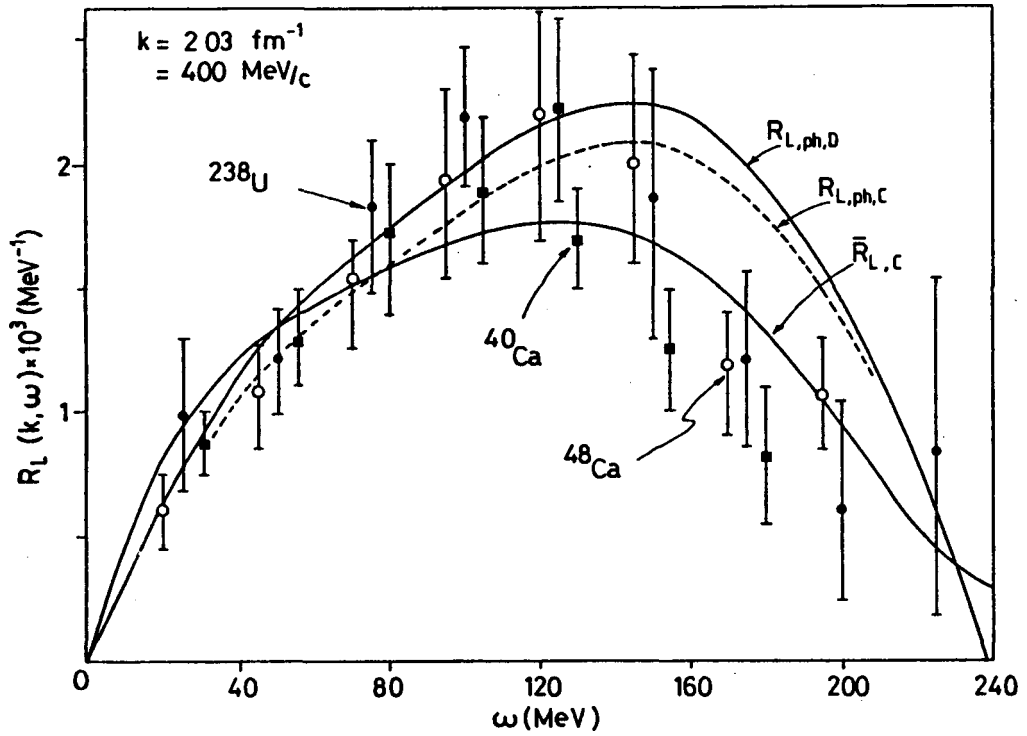


Fig.10 - Taken from ref.(15). Longitudinal response function of nuclear matter at $k = 400 \text{ MeV}/c$ compared with experimental data on ^{40}Ca and ^{48}Ca at $k = 410 \text{ MeV}/c$ and ^{238}U at $400 \text{ MeV}/c$.

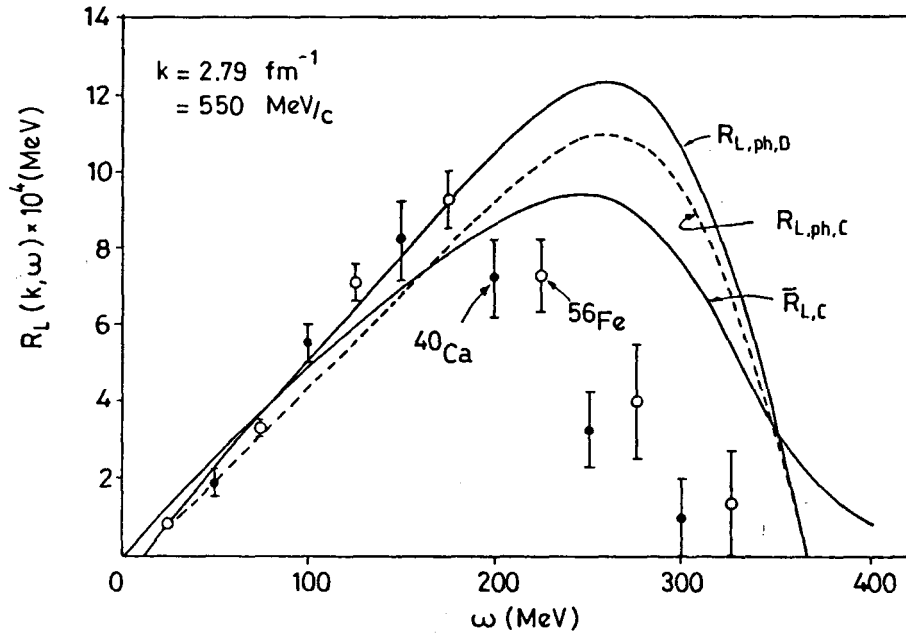


Fig.11 - Taken from ref.(15). Longitudinal response function of nuclear matter at $k = 550 \text{ MeV}/c$ compared with experimental data on ^{40}Ca and ^{56}Fe .

Note that $R_{L,C}$ is smaller than $R_{L,D}$ of $\approx 5\%$ at $k=300 \text{ MeV}/c$ and 13% at $k=550 \text{ MeV}/c$. The results displayed in Figs.9,10 show that at $k=300$ and $400 \text{ MeV}/c$ calcium and uranium nuclei give similar responses; at these values of the momentum transfer there is good agreement between our theoretical estimates and the experimental data. At $k=550 \text{ MeV}/c$ the calcium and iron data agrees with theory at lower value of ω and are too small at larger ω .

The results presented do not seem to show any particular evidence for a strong dependence of the nucleon form factor on the nuclear medium^{31,32}). In order to study this effect in some more detail both more accurate experiments and theoretical treatment of 2p-2h contributions are necessary.

Acknowledgements

S.Fantoni is grateful to the Department of Physics of the University of Illinois at Urbana-Champaign for the kind hospitality. This work has been supported by U.S. National Science Foundation under grant PHY84-15064.

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QUARK MODELS OF NUCLEONS AND NUCLEI

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ABSTRACT

We review some recent developments in our understanding of hadron structure based on the cloudy bag model and the colour dielectric soliton model. The significance of such models for a deeper understanding of nuclear structure is also discussed. Particular attention is paid to the use of deep-inelastic scattering as a tool for understanding both nucleon and nuclear structure in terms of quarks.

1. INTRODUCTION

For those of us whose deepest desire is a better understanding of the way hadronic matter is put together there is good news and bad news. Unlike most good jokes, let's take the bad news first. The bad news, which is by now pretty old, is that we have a beautiful theory, QCD, which probably explains all of the hadronic phenomena we can ever observe. Since we not only want to have fun doing physics, but also earn a living, the advent of a complete theory is threatening. To make it worse the observation of scaling, followed by logarithmic scaling violations, has confirmed both the existence of quarks and the fact that their interactions are asymptotically free.

Next the good news - QCD is so tough to deal with that it has given very little in the way of unambiguous guidance concerning the structure of nucleons and nuclei. Stated another way, there is a plethora of "QCD-motivated" models, some of which command

almost religious reverence from their adherents, and yet are mutually contradictory. In such a situation there is a need for clear insight, for imagination and most important for good experimental work to keep the theorists honest - and perhaps give us a few hints!

In this work I do not pretend to be impartial. I shall review some of the theoretical work under way at the University of Adelaide in this general area. Let me at this point acknowledge the work of S. Bass, R.P. Bickerstaff, L.R. Dodd, E. Duff, G.V. Dunne, M.A. Lohe, A. Signal, A.G. Williams (at the U. of A.) and R. Barrett, G. Chanfray, P. Guichon, L. Heller, B. Jennings, and Y.-S. Zhong (who have each spent a significant period here). Much of what follows could not have been done without their cooperation.

The structure of this article is as follows. In section 2 we define the key issues involved in relating hadronic structure to nuclear structure and reactions. Section 3 deals with recent work on the baryon number one system within the cloudy bag model and some non-topological soliton extensions of the bag. This discussion includes the colour-dielectric soliton model, the pion-nucleon σ -term and the $\Lambda^*(1405)$ resonance. One particular attempt to apply the colour dielectric model to a nuclear structure problem is described in section 4. In section 5 we examine just how disappointingly little insight deep-inelastic scattering has given us so far - particularly in view of its tremendous promise. Some concluding remarks are reserved for section 6.

2. DOES NUCLEAR STRUCTURE NEED QUARKS?

Perhaps the best illustration of the paucity of information that QCD has given us concerning hadron structure is the variety of "QCD-motivated" models presently available. Possibly the most mathematically interesting are the topological solitons,¹⁻³ where explicit mention of quarks is replaced completely by colourless meson fields with non-linear couplings capable of yielding stable,

soliton solutions. Such models are motivated by the idea that three colours might be large enough that the corrections to QCD with an infinite number of colours could be small.¹ The phenomenological achievements of this sort of model are not very impressive - typically agreeing with data at the 30% level. (However g_A is much worse.) Furthermore one must be content to believe that the structure of the non-strange and strange baryons is completely different.⁴ Nucleons falling in this general class will be labelled Skyrmions.⁵

At the other end of the spectrum the most pedantic models use familiar shell-model technology with non-relativistic kinematics (!) and pairwise effective interactions diagonalised in a limited space. Because of its simplicity this approach is very widespread. It can also be readily extended to the nucleon-nucleon (N-N) system. With regard to hadron spectroscopy it has been extremely useful because of the technical ease of dealing with centre of mass motion.⁶ Any model of this kind will be labelled a NRQM, for non-relativistic quark model.

Somewhere in the middle are the chiral bag models, and in particular the cloudy bag model (CBM),^{7,8} which has the decency to treat quarks as relativistic particles, while also respecting a fundamental symmetry of QCD - namely $SU(2) \times SU(2)$. On the debit side of the ledger all bag models have trouble with c.m. corrections, and have not been so useful for dealing with highly excited baryons. Nevertheless the phenomenological achievements for low-lying hadrons have been impressive. The strange and non-strange hadrons are not radically different beasts. In addition, the observed value of g_A is not a miracle but a simple consequence of relativity.

In the CBM the coloured quarks occupy a volume of space of radius between 0.8 and 1.1 fm - just a little smaller than the MIT bag model.⁹ Since the average separation of nearest neighbour nucleons in nuclear matter is just over 2 fm, it does not take much insight to wonder whether a complete description of the atomic nucleus should not include overlapping bags - i.e. exotic

states of six or more quarks. Baym began work on the percolation of quarks through the nucleus as much as seven years ago on just these grounds.¹⁰

The fundamental question which many of us would like to answer is whether or not a theory which does not explicitly involve colour can adequately describe nuclear matter at typical nuclear densities. By beginning with the assumption that $N_c = \infty$ is a good starting point, models of the Skyrmin type really beg the question. Of course the same is true of quantum hadrodynamics.¹¹

Colour plays a crucial role in the NRQM. There have been suggestions that not just the short-distance repulsion, but even the intermediate attraction in the N-N force could be a residue of the colour force.¹² To our taste there are two key issues (apart from the non-relativistic treatment) which make the NRQM less attractive than bags (including soliton bags). Firstly, the NRQM still suffers from spurious, long-range van der Waal's forces. Secondly, the inter-quark interactions are known (from studies of heavy quark systems) to include a significant many-body component.¹³

Very briefly this is the case, based on a process of elimination, for preferring to deal with bag models here. Having come this far the remaining question - introduced earlier when we talked about percolation - is the size of the confinement region (i.e. the bag radius). Some of the more recent, indirect evidence on this will be presented in the next section. Earlier work on the axial form-factor, baryon magnetic moments and so on, has been reviewed elsewhere^{7,8} (see also ref. (14)). The greatest shame in this regard is how little attention has ever been paid in using deep-inelastic scattering (DIS) data, to give this information unambiguously.¹⁵ After all, DIS is the only tool available which directly probes quark distributions inside nucleons. This will be mentioned again in section 5.

3. RECENT WORK ON NUCLEON STRUCTURE

3.1 The Cloudy Bag Model

For a detailed review of the CBM, including its application to πN scattering, baryon magnetic moments, spectroscopy and so on, we refer to earlier papers.^{7,8} Briefly stated the CBM is built on the success of the MIT bag model, but also guarantees chiral symmetry through the inclusion of a rather low-density cloud of virtual pions.

Our recent work has been aimed at probing more deeply the connections between the CBM and QCD. In particular, because the σ -term is one of our best clues as to the mechanism of chiral symmetry breaking in QCD, we have calculated it in the CBM.¹⁶ Before presenting the essential result some background information may be helpful.

If the full QCD Hamiltonian is $H = H_0 + H_{\text{CSB}}$, where H_0 is chiral invariant (i.e. no quark masses) and H_{CSB} is not, the σ -term is:

$$\begin{aligned}\sigma &\equiv \frac{1}{3} \sum_{i=1}^3 \langle N | [Q_i^5, [Q_i^5, H]] | N \rangle \\ &= \frac{1}{3} \sum_{i=1}^3 \langle N | [Q_i^5, [Q_i^5, H_{\text{CSB}}]] | N \rangle.\end{aligned}\tag{1}$$

Here $|N\rangle$ is the physical nucleon and Q_i^5 are the axial charges. Experimentally the value of σ is obtained by using fixed- t dispersion relations to extrapolate the on-mass-shell, isoscalar πN scattering amplitude to the unphysical point where $s = m_N^2$ and $t = +2m_\pi^2$. After correcting¹⁷ for the small t -variation between $2m_\pi^2$ and 0, one finds the most recent estimate for $\sigma = 60 \pm 8 \text{ MeV}^2$.¹⁸

By far the most popular theoretical approach for calculating σ involves chiral perturbation theory, based on picking out the leading non-analytic contributions (LNAC) to m_N . To see what this means we recall the discussion of Gasser¹⁹ which

gives

$$\sigma = \hat{m} \frac{\partial}{\partial m} m_N, \quad (2)$$

with $\hat{m} = (m_u + m_d)/2$ the average mass of the light quarks at an appropriate mass scale. In the absence of the Goldstone mechanism for breaking chiral symmetry eq. (2) would be simply

$$\sigma = \hat{m}(\Sigma_u + \Sigma_d), \quad (3)$$

with

$$\Sigma_u = \langle N | \bar{u}u | N \rangle, \quad \Sigma_d = \langle N | \bar{d}d | N \rangle. \quad (4)$$

However, in the real world the nucleon mass is corrected by a self-energy graph involving a virtual πN loop, the value of which depends on $\hat{m}$ through m_π . If the $N\pi$ contribution to m_N is δm_N we find instead of eq. (3):

$$\sigma = \hat{m}(\hat{\Sigma}_u + \hat{\Sigma}_d) + \hat{m} \frac{\partial}{\partial m} \delta m_N, \quad (5)$$

where now $\hat{\Sigma}_{u(d)}$ are matrix elements of $\bar{u}u$ ($\bar{d}d$) in the nucleon in the chiral limit ($\hat{m} = 0$). Since¹⁹ m_π^2 is proportional to $\hat{m}$, the last term in eq. (5) may be written as $m_\pi^2 \partial(\delta m_N)/\partial m_\pi^2$. A careful study of the appropriate integral reveals that to lowest order in m_π , $\partial(\delta m_N)/\partial m_\pi^2$ is proportional to m_π . Thus formally (recall $m_\pi^2 \propto \hat{m}$) we find,

$$\sigma = \hat{m} (\hat{\Sigma}_u + \hat{\Sigma}_d) + C \hat{m}^{3/2} + O(\hat{m}^2), \quad (6)$$

and $C \hat{m}^{3/2}$ is the LNAC.

A numerical study by Gasser used dipole form-factors with a mass parameter $\Lambda \in (0.5, 1.0) \text{ GeV}$. Even with Λ equal to 1 GeV, which Gasser regarded as unacceptably large, the σ -term was only equal to 53 MeV - rather low compared with experiment. With his preferred value, $\Lambda \sim 0.7 \text{ GeV}$, σ was back to 44 MeV. Thus while

the LNAC, being between 8-20 MeV, did improve the agreement with experiment there was still a significant gap.

Before we address this problem, let us consider the question any self-respecting nuclear theorist should be asking at this stage. It is actually the question at the heart of the last section. Every nuclear theorist knows that the $NN\pi$ form-factor, if represented by a dipole, has an associated mass parameter $\Lambda_\pi \sim 1300-1500$ MeV. We need only cite the Bonn potential,²⁰ or recent studies of the asymptotic D/S ratio of the deuteron.²¹ On the other hand, Gasser rejects 1 GeV as outrageously large from the QCD point of view. More than 10 years ago Primakoff²² made the natural connection between $g_A(q^2)$ and $g_{\pi NN}(q^2)$ and suggested $\Lambda_\pi \sim 0.9$ GeV. Modern studies of these two form-factors in the CBM have confirmed this result for a free nucleon. In fact, based on the theoretical work of Guichon et al.,²³ and the data for $g_A(q^2)$, we prefer $\Lambda_\pi \sim 0.9 + 0.14$ GeV. We are convinced that this value is appropriate to a free nucleon, while the much larger values coming from studies of the N-N system must include short distance corrections (e.g. quark exchange), which arise naturally when extended sources overlap!

This is the rub. A value of $\Lambda_\pi \sim 0.9 + 0.14$ GeV corresponds to a bag radius of about $1.16 + 0.20$ fm.²³ With this sort of size, nucleon overlap becomes an essential part of nuclear physics at densities near that of nuclear matter.

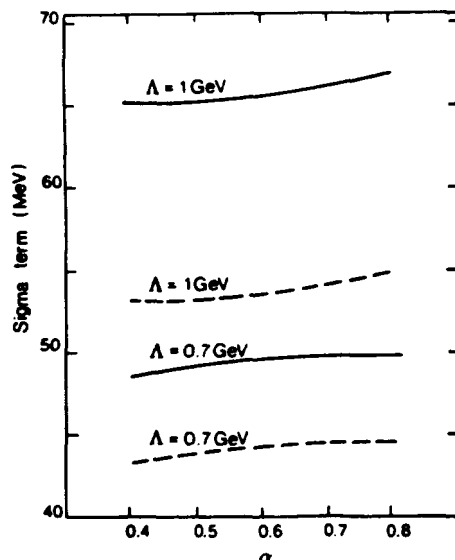


Figure 1: The σ -term plotted against $\alpha = \sqrt{3}/4$

$G_{\Lambda\Sigma^0\pi^0}/G_{pp\pi^0}$ for $\Lambda_\pi = 0.7$ and 1.0 GeV.

The dashed curves are from Gasser¹⁹, while the full curves include the $\Delta\pi$ contribution of Chanfray and Thomas.¹⁶

This crucial point made, let us finish the discussion of the σ -term. In the CBM chiral symmetry (in the Goldstone mode) is broken both by an explicit quark mass term and by a mass term for the Goldstone boson. (Of course, it is understood that these two parameters can not be independent.) The meson contribution to the σ -term can be shown, by explicit evaluation of eq. (1), to be just the expectation value of $[-\frac{1}{2} m_\pi^2 \phi^2]$ in the physical nucleon. For the fraction of time that the nucleon looks like a pion plus a nucleon bag this coincides with $m_\pi^2 \partial(\delta m_N)/\partial m_\pi^2$ of Gasser.¹⁹ However in the CBM there is another contribution of the same order, involving a virtual pion and a delta bag in the intermediate state. The role of the $|\Delta\pi\rangle$ component of $|N\rangle$ in the calculation of charge radii, magnetic moments, g_A and so on, has been thoroughly explained in the CBM.^{7,8} It is in fact crucial not only to explaining these static properties but to establishing the sense of the whole approach - that is its excellent convergence properties.²⁴

In Fig. 1 we show the result of including the $\Delta\pi$ contribution to σ , in addition to the terms calculated by Gasser. For Λ_π between 0.7 to 1.0 GeV this additional contribution varies from 5.4 to 12 MeV. The total theoretical value of σ is then between 50 and 65 MeV, which agrees rather well with the data cited earlier.

Two points remain to be made. First, the $\Delta\pi$ term is of order $\hat{m}^2$ in a strict chiral expansion, which is the reason Gasser did not include it. On the other hand, in the CBM the $\Delta\pi$ component is as important as the $N\pi$ component. Numerically it gives a significant contribution to the σ -term, which certainly should be included.

The second point is simply that this additional contribution does not alter our earlier conclusion on the size of the nucleon.

3.2 The Colour Dielectric Model:

While the bag is a beautiful mathematical toy, and can produce surprisingly good results for a number of observables, it

is clear that its sharp boundary is an idealisation. The obvious way to improve this approximation is to replace the bag by a non-topological soliton. One of the best known examples is that of Friedberg and Lee,^{25,26} where a scalar field σ is coupled to itself through an effective potential, $U(\sigma)$, and to the quarks through a Yukawa interaction, $g\sigma\bar{\psi}\psi$. For certain parameter values this model can produce quark wave functions very like those in the bag model.²⁶ Unfortunately from our point of view, the model has one unattractive feature. The potential $U(\sigma)$ is chosen to have an absolute minimum at $\sigma = \sigma_v$. Thus, in the outer limits of the soliton, where $\bar{\psi}\psi$ is small, the quarks behave as if they had a mass $g\sigma_v$. For typical parameter sets this is often less than 1 GeV, so the quarks are hardly confined. If one were only interested in an isolated nucleon this would not matter a great deal. However, in considering (say) the transition to quark matter, such a model could be misleading.

The colour dielectric model of Chanfray, Nachtman and Pirner (CNP)²⁷ is a phenomenological extension of the work of Nielsen and Patkos on QCD.²⁸ Neglecting for the present the effective gauge fields, it is described by a Lagrangian density

$$L = \bar{\psi} i \not{\partial} \psi - \frac{m}{\chi} \bar{\psi}\psi + \frac{1}{2} \sigma_v^2 (\partial_\mu \chi)^2 - U(\chi). \quad (7)$$

Here χ is the dimensionless scalar field for which σ_v sets the scale, and $U(\chi)$ is the self-interaction of χ which has an absolute minimum at $\chi = 0$ and a local minimum at $\chi = 1$. In a region of space where χ is unity, m is the quark mass. On the other hand, quark confinement occurs because $m/\chi \rightarrow \infty$ as $\chi \rightarrow 0$. For simplicity $U(\chi)$ is parameterized as a quartic with two parameters ($U(0)$ is chosen to be zero). The bag energy density B is $U(1)$, while for small χ , $U(\chi) \simeq \frac{1}{2} m_{GB}^2 (\sigma_v^2 \chi^2)$ and m_{GB} is often referred to (probably incorrectly) as the "glueball mass".

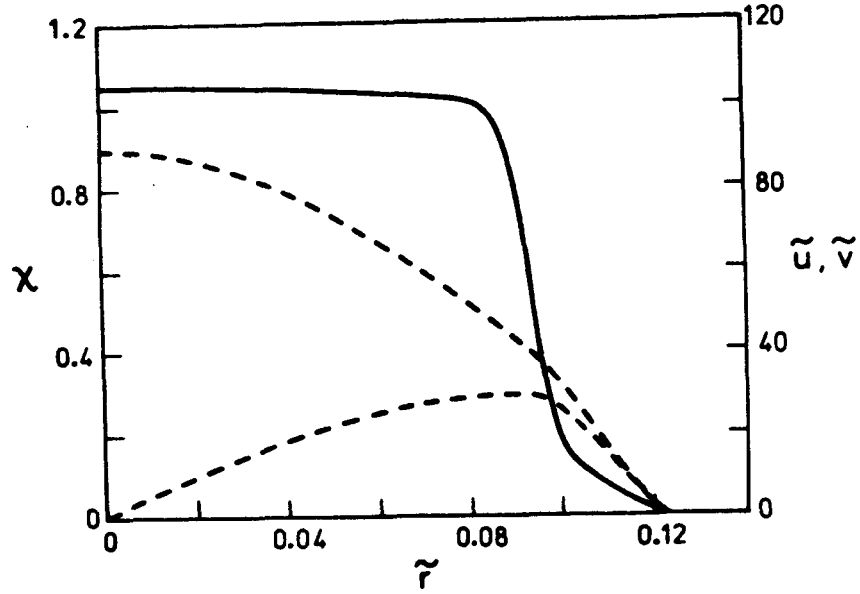


Figure 2: Bag-like solution for the colour-dielectric soliton model, showing the scalar field χ (solid) and the scaled quark wave functions ($\tilde{u}$ and $\tilde{v}$) against a scaled radius parameter ($\tilde{r} = r/r_v$) - from ref. [30]

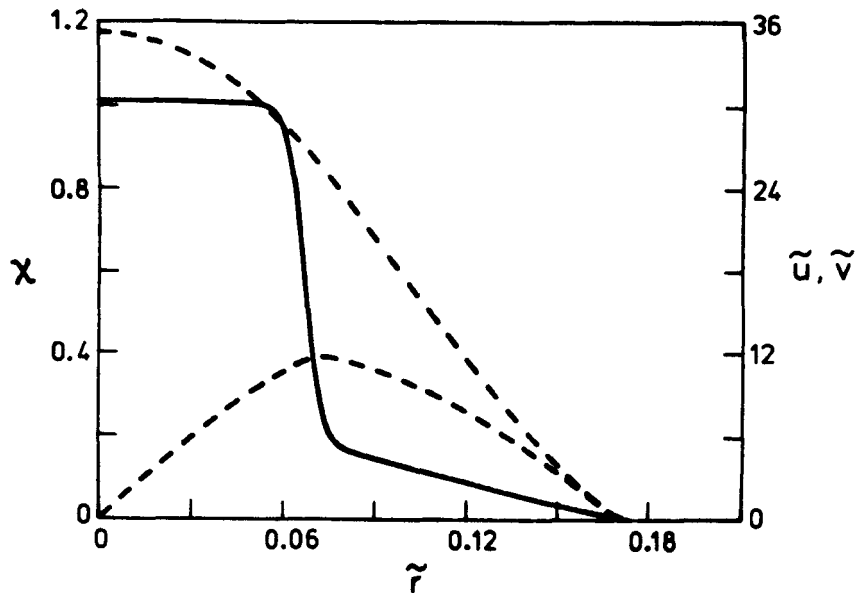


Figure 3: As for Figure 2, but for a lower quark mass - note the large region where χ is neither 0 nor 1. (From ref. [30].)

The equations resulting from a mean-field approximation to eq. (7) were solved using the numerical method of Dodd and Lohe.²⁹ In Figs. 2 and 3 we show two examples of the solutions with baryon number one obtained by Williams, Dodd and Thomas (WDT).³⁰ The dimensionless radius parameter, $\tilde{r}$, is just $\sigma_v r$, while $\tilde{u}$ and $\tilde{v}$ are the upper and lower radial pieces of the 1s wave function ψ_{1s} . Usually we choose σ_v so that the proton rms radius is 0.83 fm (i.e. $\sigma_v = 15.1$ MeV for Fig. 2 and 18.4 MeV for Fig. 3). The other static properties are then quite acceptable - e.g. $m_N = 1.04$ GeV, $g_A = 1.20$ and $\mu_p = 2.4 \mu_N$ for Fig. 2.

Perhaps the most interesting feature of the solutions is the behaviour of χ in the exterior of the soliton. Following CNP, we had assumed that for a suitable choice of $U(\chi)$, one would be able to obtain bag-like solutions in this model. Figure 2 is as close to being a step function for $\chi(r)$ as we could get without making the quark mass too large (here $m \sim 90$ MeV). Clearly χ drops rapidly in the surface, but then goes to zero almost linearly. This trend is even more obvious in Fig. 3, where $m \sim 28$ MeV.

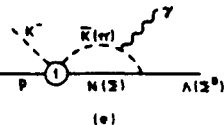
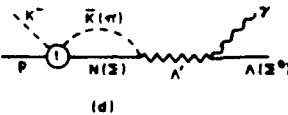
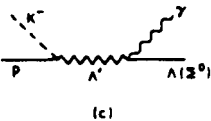
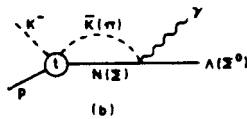
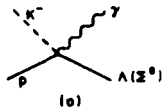


Figure 4 : Processes which contribute to the radiative decay of the K^-p atom - from Zhong et al.²¹

Especially for Fig.3 there is a large probability of finding the quarks in a region where $\chi \neq 1$, and such a solution would not be well approximated by a bag. We shall recall this caution in section 4 when we mention applications of this model to nuclear structure.

To conclude this discussion we observe that because the quarks spend much of their time in a region where χ is less than unity, there is a considerable sensitivity to the quark mass. For example, for the soliton shown in Fig. 3 a mass difference ($m_d - m_u$) of just over 1 MeV may suffice to explain the observed n-p mass difference (after electromagnetic corrections).³⁰

3.3 The $\Lambda^*(1405)$:

While this does not belong in a discussion of quark models of nucleons and nuclei, recent work on the $SU(3) \times SU(3)$ extension of the CBM has some relevance to the program of hadron spectroscopy at CEBAF. Briefly, the $SU(3) \times SU(3)$ CBM has been applied to $\bar{K}N$ and $\bar{K}N$ scattering up to about 500 MeV/c kaon laboratory momentum.³¹ The relevant result here is that the $\Lambda^*(1405)$ resonance, with $J^\pi = 1/2^-$ $S = -1$, was found to be a $\bar{K}N$ bound state rather than a 3-quark bag state in our work. Indeed we found that the lowest mass bag-state, coupled to $\bar{K}N$ and $\Sigma\pi$, would have to have a mass of at least 1650 MeV. If confirmed this result has important consequences for quark models. The challenge is how to test it.

One possible experimental test now under way at Brookhaven is to measure the radiative decays of kaonic hydrogen,

$$K^- + p \rightarrow \Lambda + \gamma, \quad (8a)$$

$$K^- + p \rightarrow \Sigma^0 + \gamma. \quad (8b)$$

By calculating the processes shown in Fig. 4, Zhong et al. found branching ratios of 1.9×10^{-3} and 1.5×10^{-3} for (8a) and (8b) respectively.³² The former is in agreement with existing data, while the latter prediction is well within the anticipated precision of the present experiment.³³ By way of comparison, the MIT bag model³⁴ predicts a ratio of $\Sigma^0\gamma$ to $\Lambda\gamma$ of between 3.3 and 6.3.

Instructive though these measurements will be, the K^-p atom is in the tail of the rather narrow $\Lambda^*(1405)$. On the other hand

the kinematic freedom of the (e,e') process would allow one to produce the $\Lambda^*(1405)$ right on resonance. Clearly some more theoretical effort should be made to see what extra information could be obtained this way.

4. MORE THAN ONE NUCLEON.

The past decade has seen an enormous amount of work on quark-level calculations of the N-N force. Since the discovery of the EMC effect (see section 5) in particular, there have also been many attempts to construct quark models of nuclei. For the Friedberg-Lee model, for example, it has been shown³⁵ that the transition from nucleon-matter to quark matter occurs at nuclear matter density! However, we argued earlier that because of the low effective quark mass between nucleons in this model it is not satisfactory for such studies. The analogous study for the model of CNP²⁷ has not yet been carried out. Nevertheless there has been a study of nuclear charge symmetry breaking in the colour dielectric model³⁶, which we now review briefly.

Chanfray et al. (CNP) assumed that for a suitable $U(\chi)$ in equ. (7) one could find a bag-like solution for a free nucleon - i.e. $\chi = 1$ for $r < R$ and $\chi = 0$ for $r > R$. (The effect of relaxing this probably unrealistic assumption (recall section 3.2) remains to be seen.) If one then considers a model nucleus consisting of colour dielectric bags arranged on a regular lattice, it is relatively easy to show that the minimum energy configuration involves a non-zero value of χ between the nucleons. In Fig. 5 we show a simple approximation in which $\chi = 1$ inside the nucleons, $\chi = \chi_N$ between the nucleons, and $\chi = 0$ outside the nucleus.

A non-zero value of χ_N has two effects. The kinetic energy of each quark is lowered by a fraction $\chi_N/2mR$, while it costs an amount $U(\chi_N)$ per unit volume between the nucleons. Both χ_N and the radius of the nucleon within the nucleus, R , are found by minimising the total energy of the system $E(\chi_N, R)$. As shown by

CNP the change in R is at most a few percent at or below nuclear matter density. Then one finds

$$\chi_N = \frac{3}{4} m_N \frac{\hat{m}}{m_u m_d} \left[\frac{8}{3} \pi R (r_0^3 - R^3) m_{GB}^2 \sigma_V^2 \right]^{-1}, \quad (9)$$

where $\frac{4}{3} \pi r_0^3$ is the volume of the Wigner-Seitz cell occupied by the nucleon (i.e. $r_0 \propto \rho^{-1/3}$).

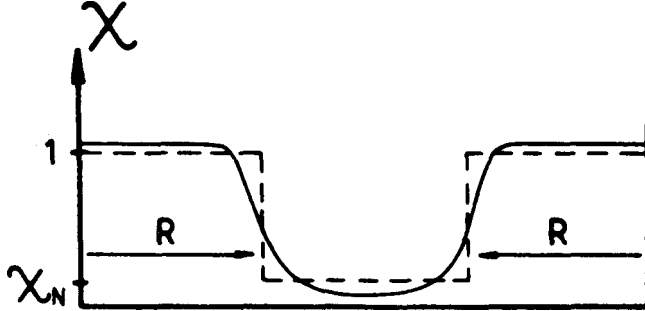


Figure 5: A schematic illustration of how the χ -field might vary along a line joining the centres of two nucleons in the nucleus (solid), and a step function approximation to it.³⁶

Before we go any further we should clarify the relationship between this model and the conventional picture of meson-mediated nuclear forces. In deriving eq. (9) and sketching Fig. 5 we have made a mean-field approximation, neglected $q\bar{q}$ creation and quark and gluon exchange (except in so far as the latter confines the quarks). We can not hope to reproduce serious nuclear structure. However we may hope to get some insight into aspects of nuclear structure related to quark confinement which are totally missing in quantum hadrodynamics. With regard to charge symmetry breaking (CSB) we can also study the large effects of an up-down quark mass difference. If the colour dielectric model is at all realistic, such effects would be expected to survive even if meson exchange or some quark exchange equivalent were to be included.

With this caution we write the total energy per nucleon as $m_N + W$ and find, following Williams and Thomas,³⁶ that

$$W = - \frac{m_N}{8R} \chi_N \sum_{i=1}^3 m_i^{-1} + \frac{4\pi}{3} (r_0^3 - R^3)^{1/2} m_{GB}^2 \sigma_V^2 \chi_N^2. \quad (10)$$

Equation (10) clearly implies a neutron-proton binding-energy difference ΔW :

$$\Delta W = \frac{m_N}{8R} \chi_N (m_u^{-1} - m_d^{-1}). \quad (11)$$

For the parameters used by CNP in connection with the EMC effect (i.e. an effective quark mass of only 800 MeV between the nucleons) this would imply a proton-neutron binding energy difference of up to 10 MeV - in the absence of Coulomb forces!³⁷ A relatively minor change in parameters can cure this problem, but we are then left with the interesting question of whether the residual CSB effects can help explain any anomaly in conventional nuclear physics - e.g. the Nolen-Schiffer anomaly.

A detailed investigation of the consequences of this model for mirror nuclei has indeed shown that there is a range of parameters for which the Nolen Schiffer anomaly can be understood.³⁶ While we can give no reasons to prefer these particular parameters, notably m_u and m_d , they do lie in the range commonly used in bag models. Even if one were to prefer other values, this CSB effect, which is effectively a new many-body force, would need to be included in any accurate study of charge symmetry violation in nuclei.

5. DEEP INELASTIC SCATTERING

It is widely appreciated that DIS provides the most direct evidence for the presence of quarks inside the nucleon. Therefore, if (as we have argued) nucleon overlap, quark exchange and even percolation play an important role in nuclear structure, DIS should be a useful tool to study these effects. It is for this reason that the EMC effect,^{38,39} which reveals a difference in the electromagnetic structure functions of heavy nuclei compared with deuterium, has caused such a flurry of activity. This discrepancy, which increases roughly as $\ln A$,³⁹ has variously

been attributed to clusters of 6, 12 or more quarks, swollen nucleons, deltas in nuclei, and so on. Exciting though this may be, our first duty must be to eliminate simple (!) corrections such as fermi motion and nuclear binding before attributing the discrepancy to more exotic causes.

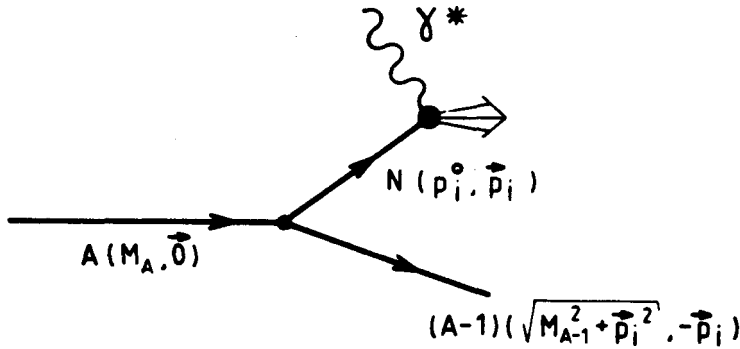


Figure 6: The simple model for binding and fermi motion corrections considered by Dunne and Thomas.⁴²

5.1 The EMC Effect With a Little Cold Water

The first estimates of the effect of fermi-motion⁴⁰ for Fe suggested a discrepancy between experiment and a theory involving only nucleons of some 30% at Bjorken- x about 0,6 - where only valence quarks contribute. However it has recently been realized that this calculation did not handle binding well enough.^{41,42} The key process, shown in Fig. 6, involves the virtual photon being absorbed by an off-mass-shell nucleon. (Here we assume that the impulse approximation is valid so that final state interactions between the $(A-1)$ nucleus and the debris of the struck nucleon can be ignored. At least for $x \gtrsim 0.2$ this is easily justified on the grounds of typical interaction times. It may be reasonable at smaller x , but this is harder to establish.)

A careful calculation of the mechanism shown in Fig. 6 implies that

$$F_2^A(x, Q^2) = A^{-1} \sum_i n_i \int_x^A dy f_i(y) F_2^N(\frac{x}{y}, Q^2). \quad (12)$$

Here n_i is the occupation number of the single particle state i ,

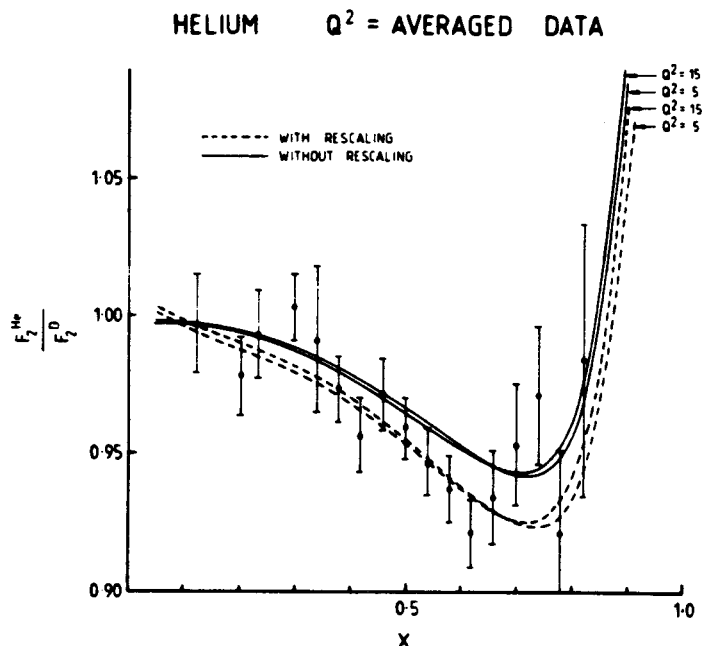


Figure 7; Comparison of the predictions of a calculation of binding and fermi motion corrections⁴² for ${}^4\text{He}$ with Q^2 - averaged data from SLAC.³⁹ The calculation uses harmonic oscillator wave functions and measured separation energies.

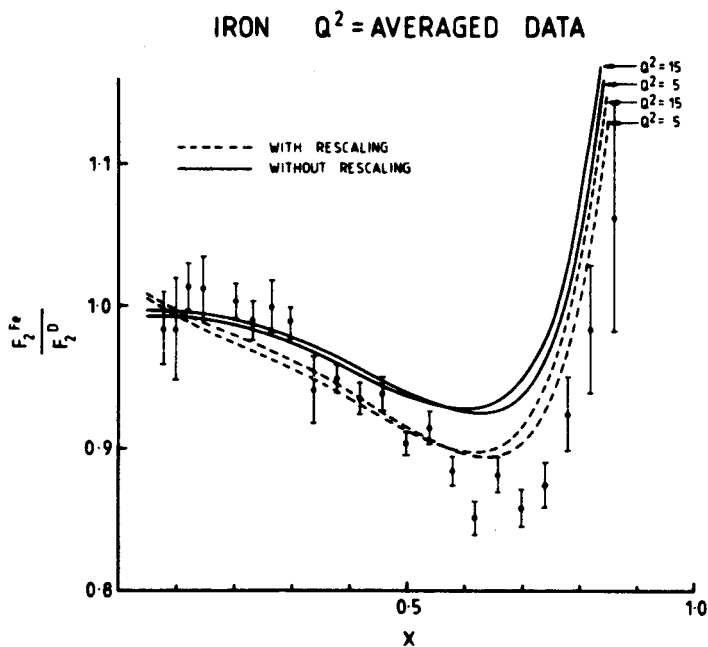


Figure 8: As for Figure 7, but for ${}^{56}\text{Fe}$.

binding energy ϵ_i , and $f_i(y)$ is determined by ϵ_i and the single-particle density. For example, for a harmonic oscillator Os state one finds⁴²

$$f_{Os}(y) = \left(\frac{m}{\pi M \omega}\right)^{1/2} \exp(-(my - m_0)^2 / m M \omega), \quad (13)$$

where $m_0 = m_N - |\epsilon|$. Clearly this peaks at $y = m_0/m_N < 1$, so that $F_2^A(x, Q^2)$ involves F_2^N at a larger value of x . In Figs. 7 and 8 we show the results of such a calculation for ${}^4\text{He}$ and ${}^{56}\text{Fe}$ - solid curves. Clearly this simple kinematic correction, which arises because the bound nucleon has (on average) less than one unit of momentum, can explain a good deal of the EMC effect.

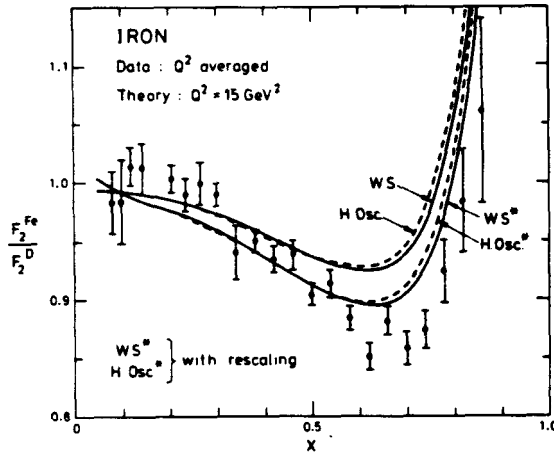


Figure 9: A comparison of the results of binding and fermi motion corrections for ${}^{56}\text{Fe}$ using a Woods-Saxon potential with the results for a harmonic oscillator.

That the residual discrepancy is not caused by using the harmonic oscillator model is shown in Fig. 9, which is a recent calculation using a Woods-Saxon single-particle potential.⁴³ Of course one could use ϵ_i as adjustable parameters⁴¹ to fit the EMC data, but then one would learn nothing fundamental.

It was pointed out by Dunne and Thomas⁴² that in writing eq. (12) it was assumed that the nucleon structure function does not depend on the nucleon invariant mass. Instead they used perturbative QCD to suggest an ansatz for the dependence of F_2^N on $p^2 \neq m_N^2$. Indeed, if $M(n, Q^2)$ is the n 'th moment of a structure function, at a photon invariant mass $-Q^2$:

$$M(n, Q^2) = \int dx x^{n-2} F_2(x, Q^2), \quad (14)$$

they suggested that

$$M^A(n, Q^2)/M^N(n, Q^2) = [\alpha(m_N^2)/\alpha(p_A^2)]^{d_n}. \quad (15)$$

Here d_n is the appropriate anomalous dimension, α the strong coupling constant and p_A^2 the invariant mass of the nucleon in target A. Equation (15) resembles the rescaling formula of Close and collaborators⁴⁴, and can be shown to imply

$$M^A(n, Q^2) = M^N(n, \xi(Q^2) Q^2), \quad (16)$$

with the rescaling factor, $\xi(Q^2)$ given by:

$$\xi(Q^2) = \left(\frac{m_N^2}{p_A^2} \right)^{\alpha(p_A^2)/\alpha(Q^2)}. \quad (17)$$

By including this correction as well as the effect of fermi motion described earlier we obtain the other curves shown in Figs. 7-9. In this case the agreement with experiment is rather good, except perhaps right in the minimum.

So where are we exactly in our analysis of the EMC data? Clearly a simple change of mass scale* associated with binding explains most of the effect. However the necessity for a guess at the off-shell behaviour of the nucleon structure function is unsatisfactory. Either we must test this ansatz in a specific model, or alternately work with a theory of on-mass-shell nucleons which would remove the need for it. (Some promising steps in this direction have been taken recently.⁴⁶) Certainly the connection between the binding connection and the work of Llewellyn Smith,⁴⁶ Ericson and Thomas⁴⁸ and Berger and Coester⁴⁹ on a pion excess being responsible for the effect, is theoretically satisfying. In

*Rather than confinement scale - as discussed at greater length by Bickerstaff and Thomas.⁴⁵

a nut-shell the nucleons are bound by extra pions, which necessarily deplete their share of the nuclear momentum. Nevertheless, until we can say unambiguously that binding and fermi-motion has been accurately calculated, without free parameters or guesses, we can conclude nothing about the more exotic speculations on nuclear structure.

While this is disappointing, it is also a powerful spur to more research effort - both theoretical and experimental. We need precise DIS data (using weak and electromagnetic probes) as a function of A , Q^2 and x in a style which has hitherto been foreign to high-energy physics. The Drell-Yan process, in extreme kinematic conditons, may also provide important information.⁵⁰

5.2 Constraining Nucleon Models:

Sad to observe, the power of DIS has not yet been exploited to choose between phenomenological models of nucleon structure. In principle this should be the best technique to resolve arguments over the size of the quark confinement region. Several years ago we used the known asymmetry of the nucleon sea to put a lower limit on the bag radius in the CBM⁵¹ - i.e. $R \gtrsim 0.87 \pm 0.10$ fm. A very similar limit has just been obtained by Signal and Thomas, using the strange sea of the nucleon.⁵² However the most direct test would involve the valence quarks alone. Jaffe pioneered the treatment of DIS from a bag over ten years ago,⁵³ but the lack of momentum conservation caused important problems - which have only been overcome in 1-space 1-time dimension. We have recently made some progress on the 3D case and hope to be able to report on the results soon.⁵⁵ There is a very definite chance that DIS will soon fulfil its promise and eliminate a number of hitherto indistinguishable nucleon models - thereby clarifying a basic issue in this whole discussion.

6. OUTLOOK

This is a very exciting time for our field. The atmosphere is full of exciting ideas and speculation on the nature of hadron and nuclear structure resulting from QCD. New machines like CEBAF have lots of questions to answer. In this review we have mentioned specifically the need for electroproduction data for states like the $\Lambda^*(1405)$. We also pointed out the necessity for a systematic study of the x , Q^2 and A -dependence of nuclear structure functions. While the most exciting work possible at CEBAF will probably be none of these, it is highly likely that it will involve tests of the general ideas we have presented - albeit refined and sharpened in the usual way. I for one look forward to these developments.

This work is supported by the Australian Research Grants Scheme and the University of Adelaide.

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HADRON PRODUCTION IN CHARGED LEPTON SCATTERING*

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ABSTRACT. A review is given on the properties of complete hadronic final states as produced in deep inelastic electron scattering in the energy domain of CEBAF. It is discussed to what extent and how CEBAF can help to improve our present understanding of electroproduction which is essentially based on the quark-parton model.

INTRODUCTION

CEBAF which will provide electron beams of energies up to 4 (6) GeV is, in the first place, designed for experiments in nuclear physics. However, about 15 to 20 years ago the energies of electron beams available for particle physicists were of the same order of magnitude. In order to test the basic concepts of the today's understanding of deep inelastic lepton scattering, namely the quark-parton model and quantum chromodynamics, higher and higher energies were aimed for. Despite this it may still be worthwhile not to forget the unsolved problems left behind in the early seventies. Due to its high duty cycle CEBAF offers the opportunity to investigate this energy domain with much better precision than previously possible so that results of a new quality can be expected.

There is a side-effect which I want to mention here as well and which in my opinion should not be considered unimportant: the co-operation of nuclear and particle physicists at CEBAF can lead to mutual inspiration, similar to the impact the EMC-effect which was found by particle physicists had on nuclear physics.

The aims of this report are the following: firstly the present knowledge about the complete hadronic final states in lepton scattering will be reviewed. Doing this the central rôle the quark-parton model plays in interpreting hadron production will be emphasized. Based on the current status of our knowledge it will be investigated how CEBAF can contribute to an improved understanding of deep inelastic scattering. Finally it will be discussed what consequences this has to the layout of the planned experiments.

This article concentrates on complete hadronic final states in lepton scattering. An excellent and very complete review of exclusive channels was given

*Invited talk given at the 1986 CEBAF Summer Workshop

by D.G. Cassel at the 1985 CEBAF workshop¹. Therefore the discussion of specific reactions will be kept short here and the interested reader is referred to the report just mentioned. The bibliography given at the end of this article is supposed to be as comprehensive as possible with respect to the experimental study of complete hadronic final states albeit it is not claimed to be complete.

KINEMATIC VARIABLES AND BASIC FORMULAE

In first order, hadron production in deep inelastic electron scattering, namely

$$e + N \rightarrow e' + \text{hadrons}, \quad (1)$$

proceeds via the exchange of a virtual photon between the lepton and the target nucleon as schematically depicted in fig. 1. Here e represents the beam particle, N denotes the target proton or neutron and e' stands for the scattered electron.

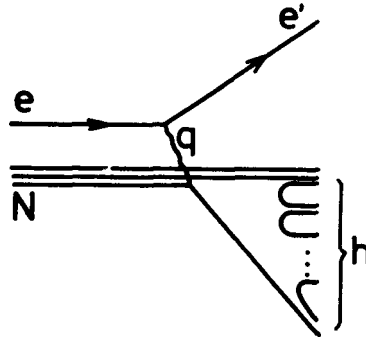


Fig. 1: Diagram of electroproduction of hadrons in the quark-parton model.

The kinematic variables commonly used to describe the lepton part of process (1) defined in terms of the four-momenta of the initial and final particles (fig. 1) are given in the following way:

$$Q^2 = -q^2 = -(e - e')^2 \approx 4EE' \sin^2 \frac{\Theta}{2} \quad \text{mass of the virtual photon squared}$$

$$W^2 = s = (q + N)^2 = M^2 - Q^2 + 2M(E - E') \quad \text{hadronic cms energy squared}$$

where E (E') is the energy of the incident (scattered) electron and Θ is the electron scattering angle, all given in the laboratory system. M denotes the nucleon mass, masses of leptons are neglected. At fixed beam energies and if beam and target are not polarized two variables are sufficient to describe inclusive lepton scattering completely. However, instead of Q^2 or W^2 other variables are also frequently used:

$$\nu = \frac{N \cdot q}{M} = E - E' \quad \text{energy of virtual photon in the laboratory system}$$

$$x = \frac{Q^2}{2N \cdot q} = \frac{Q^2}{2M\nu} \quad \text{Bjorken scaling variable (fraction of the nucleon's momentum carried by the scattered quark)}$$

Since the virtual photon is spacelike, i.e. has non-zero mass, it can have also helicity 0 (longitudinal photon) besides helicities ± 1 (transverse photon). The ratio between the fluxes of longitudinal and transverse photons is given by

$$\varepsilon = \frac{1}{1 + \frac{2(Q^2 + \nu^2)}{4EE' - Q^2}}. \quad (2)$$

For the hadronic part of process (1) the following three quantities are basically used:

$p_{||}$ momentum parallel to the virtual boson direction,
 $p_{\perp}$ momentum perpendicular to the virtual boson direction,
 φ azimuthal angle with respect to the virtual boson direction.

Instead of $p_{||}$ usually one of the three following variables are preferred:

$$x_F = \frac{p_{||}}{p_{max}^*} \approx \frac{p_{||}}{W/2} \quad \text{Feynman } x \text{ (in the centre-of-mass system)}$$

$$y = \frac{1}{2} \ln \frac{E_h + p_{||}}{E_h - p_{||}} \quad \text{rapidity}$$

$$z = \frac{N \cdot h}{N \cdot q} = \frac{E_h}{\nu} \quad \text{energy fraction (in the laboratory system)}$$

with the hadron's four-momentum h , energy E_h and maximum cms momentum p_{max}^* .

In the case of exclusive single hadron production it is customary to use

$$t = (q - h)^2 \quad (3)$$

the four-momentum transfer squared to the nucleon. t and x_F or z are correlated such that t approaching its kinematic limit t_{min} corresponds to x_F and z going to 1.

Electroproduction is considered as virtual photoproduction. The contributions from higher order QED graphs are calculated² and the results are used to extract the contribution due to one photon exchange from the experimental measurements (radiative corrections). In this picture the differential cross section for the production of a single hadron is given by³

$$\frac{d^4\sigma}{d\Omega dE' dt d\varphi} = \Gamma_T(E, E', \Theta) \cdot \left(\frac{d\sigma_T}{dt} + \varepsilon \frac{d\sigma_L}{dt} + \varepsilon \frac{d\sigma_P}{dt} \cos 2\varphi + \sqrt{2\varepsilon(\varepsilon + 1)} \frac{d\sigma_I}{dt} \cos \varphi \right) \quad (4)$$

with

Ω solid angle of scattered electron

Γ_T flux of transverse virtual photons
 σ_T contribution from transverse unpolarized virtual photons
 σ_L contribution from longitudinal virtual photons
 σ_P contribution from transverse polarized virtual photons
 σ_I contribution from transverse-longitudinal interference

For the total cross section (sum over all channels, integrals over t and φ) one gets

$$\frac{d^2\sigma_{tot}}{d\Omega dE'} = \Gamma_T (\sigma_T + \epsilon\sigma_L) \quad (5)$$

or in terms of structure functions

$$\frac{d^2\sigma_{tot}}{d\Omega dE'} = \left(\frac{d\sigma}{d\Omega} \right)_{Mott} \cdot [2W_1(Q^2, \nu) \tan^2 \frac{\Theta}{2} + W_2(Q^2, \nu)]. \quad (6)$$

It is common to introduce instead of W_1 and W_2

$$F_1 = MW_1 \quad \text{and} \quad F_2 = \nu W_2. \quad (7)$$

In the quark-parton model at very high values of Q^2 and W^2 F_1 and F_2 are predicted to depend only on x . F_2 is then given by the quark charges e_i and the distribution functions q_i of the quarks inside the target nucleon in the following way:

$$F_2(x) = x \sum_i e_i^2 q_i(x) \quad (i=u, d, \dots, \bar{u}, \bar{d}, \dots). \quad (8)$$

If the scattering occurs only from spin $\frac{1}{2}$ partons the Callen-Gross relation

$$F_2(x) = 2xF_1(x) \quad (9)$$

is valid and from equation (6) one can derive

$$\frac{d^2\sigma_{tot}}{dQ^2 dW^2} \propto \frac{1}{Q^4} \frac{1 + (1 - \frac{\nu}{E})^2}{M\nu} F_2(x). \quad (10)$$

The strong fall-off with Q^2 makes it difficult to obtain experimental results at large Q^2 (large x) values. The high duty cycle of CEBAF will help to improve the situation (see discussion further down).

The quark-parton model formulae can be extended to describe hadron production which in this model is interpreted as a two-step process: firstly, on a short time scale, the virtual photon is absorbed by a quark inside the nucleon and secondly, on a longer time scale, this struck quark and the spectator 'diquark' (remnant target system) fragment into hadrons. The differential cross

section to produce a specific hadron h per lepton scattering event at a given value of x and per energy interval dz can be expressed by

$$\frac{1}{\sigma_{tot}(x)} \frac{d\sigma^h}{dz} = \sum_i \epsilon_i(x) \cdot D_i^h(z) \quad (11)$$

where the summation runs over all quarks and antiquarks. The structure of this formula reflects the picture of a process in two independent steps by the factorization into $\epsilon_i(x)$ which describes the virtual photon quark scattering and $D_i^h(z)$ which describes the fragmentation of the quark of flavour i into the hadron h . In the simple quark-parton model the fragmentation function $D_i^h(z)$ does not depend on any event variable, i.e. W^2 or Q^2 . This is called scaling. Thus, factorisation and scaling represent two distinct predictions of the quark-parton model and their observation in the experimental data gives strong support to the validity of the quark-parton picture. In formula (11) $\epsilon_i(x)$ represents the probability for scattering off a quark (antiquark) of flavour i with the energy fraction x inside the target nucleon. Using equations (8) and (10) one gets

$$\epsilon_i(x) = \frac{e_i^2 q_i(x)}{\sum_j e_j^2 q_j(x)}. \quad (12)$$

Instead of $\frac{1}{\sigma_{tot}(x)} \frac{d\sigma^h}{dz}$ its Lorentz invariant form $\frac{E_h}{\sigma_{tot}(x)} \frac{d^3\sigma^h}{dp^3}$ is also frequently used and denoted as 'structure function' $F(z)$ or $f(z)$.

GENERAL PROPERTIES OF HADRONIC FINAL STATES

In this section the general properties of complete hadronic final states are discussed. The data presented stem mainly from experiments with 4π detectors. To obtain full coverage of the phase space bubble chambers (muon beams!) and later on streamer chambers have been used at Cornell, DESY and SLAC. This implies that no results are available which comprise large acceptance together with γ detection and charged particle identification. The results shown below refer to charged particles which in general are assumed to be pions. If necessary contributions due to protons are subtracted. This is done on the basis of measurements from electronic detectors.

In fig. 2 the dependences of the mean charged multiplicities on Q^2 and W^2 are shown^{4,5,6,7}. The dominant effect is the approximately linear rise of $\langle n \rangle$ with $\ln W^2$. This behaviour is also observed at higher energies⁸ and can be easily understood in the following picture^{9,10}: it is well known that the height of the rapidity distribution of hadrons does not depend on energy near $y=0$. On the other hand the width of the y distribution (central plateau) increases with $\ln W^2$. Since the multiplicity can be obtained by integration of the rapidity distribution over the complete y range the observed $\ln W^2$ behaviour follows

naturally. However, it is surprising that this picture works already at very low W where resonance production and the influence of single exclusive channels are important.

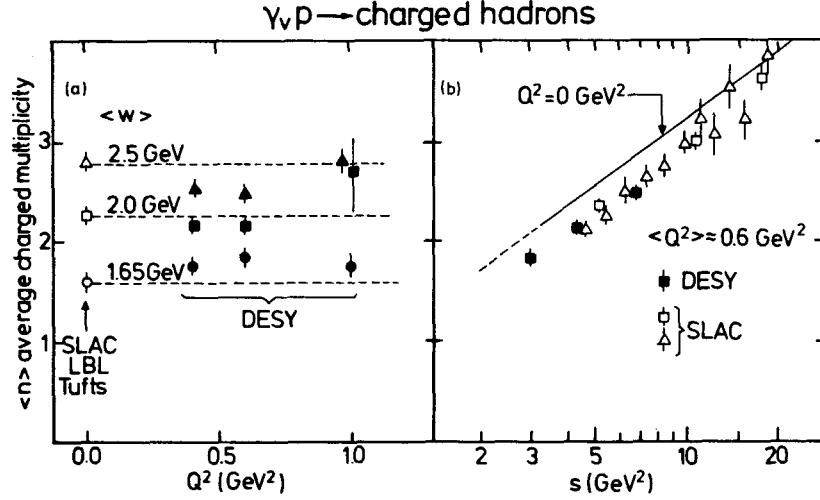


Fig. 2: Average charged multiplicity (a) as a function of Q^2 and (b) as a function of $W^2 = s$. The full symbols refer to data from ref. 5. The measurements at $Q^2 = 0$ GeV² (in (b) drawn as solid line) were extracted from ref. 4. The open squares and triangles represent results from ref. 6 and 7 respectively.

The situation with respect to the Q^2 dependence is much less clear. In fig. 2a, starting from $Q^2 = 0$ GeV², no obvious trend can be observed. At higher energies a slight increase of $\langle n \rangle$ with Q^2 at fixed W^2 is found¹¹. In both cases the interpretation is not straight forward. Especially at low W one would have to understand the contributions from the various exclusive channels which exhibit very different Q^2 dependences in order to be able to interpret the behaviour of $\langle n \rangle$. One has also to take into account that varying Q^2 at fixed W implies a variation of x as well and therefore the contributions from the different flavours of the struck quark vary with Q^2 .

To summarize: only a very weak Q^2 dependence of $\langle n \rangle$ is observed. To obtain a better understanding very high statistical precision would be necessary over a large Q^2 range. CEBAF can certainly help to clarify the situation. On the average only about two charged particles are produced per event so that no major difficulties are to be expected from high multiplicity events but the electromagnetic background may cause problems in determining the event topologies.

In the quark-parton model it is expected¹² that the hadrons produced in the fragmentation of a quark will retain its quantum numbers, e.g. the charge. The net charge of the forward going hadrons should therefore reflect the charge of the scattered quark and thus allow to determine the quark charges via

$$\langle Q_{forw} \rangle = \sum \epsilon_i(x)(e_i \mp \eta). \quad (13)$$

The sum runs over all quarks ($-\eta$) and antiquarks ($+\eta$). ϵ_i is defined in equation (12). η , the mean charge of quarks created in the fragmentation process, takes into account that, since hadrons and not individual quarks are experimentally observed, not all quark-antiquark pairs formed during fragmentation are completely contained in the sample of the forward going hadrons ($x_F > 0$). The limiting values of η are 0 (u, d and s quarks are equally probable) or $\frac{1}{6}$ (s quarks are suppressed).

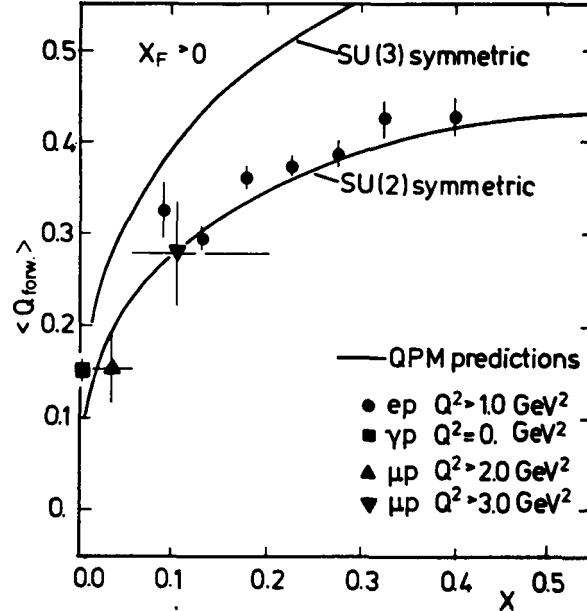


Fig. 3: Mean net charge of the hadrons going forward ($x_F > 0$) as a function of Bjorken x . The data are taken from the following references: circles - ref. 13, squares - references 14 and 15, triangles - ref. 16.

Near $x = 0$ where scattering occurs from sea quarks and antiquarks ($\langle Q_{forw} \rangle$) is expected to be zero. With increasing x scattering from u quarks (proton target) becomes more and more dominant and $\langle Q_{forw} \rangle$ should increase up to $(\frac{2}{3} - \eta)$ at $x = 1$. This behaviour is indeed experimentally observed^{13,14,15,16} (fig. 3). The data points lie slightly above the quark-parton model prediction¹⁷ with complete s quark suppression.

These results were confirmed at higher energies¹⁸ where good separation between forward and backward going hadrons could be demonstrated and the validity of the quark-parton model is out of question. One may therefore conclude that, since the quark charge retention can be observed at rather low W , the quark-parton model is applicable even at rather low energy.

After the presentation of integral quantities like total multiplicity and charge flow now the differential momentum distributions of final state hadrons will be discussed. In fig. 4 the Lorentz invariant normalized cross section (for the definition see end of previous section, here: integrated over $p_\perp^2$, averaged over φ) for the production of negative pions is shown as a function of Feynman

x in various W and Q^2 ranges⁵. The distributions have maxima at $x_F = 0$ and fall steeply towards $x_F = 1$ and $x_F = -1$. The stronger decrease backward can be understood because due to baryon number conservation there must be always a baryon in the final state which is preferentially among the target fragments thus leaving less energy for meson production.

Within the limited precision of the data no significant Q^2 dependence can be observed. There is also agreement with photoproduction¹⁵ of negative pions when the contributions from elastic ρ^0 production (diffractive process) are removed. Such a behaviour (Feynman scaling) is expected from the quark-parton model and we will come back to this point further down.

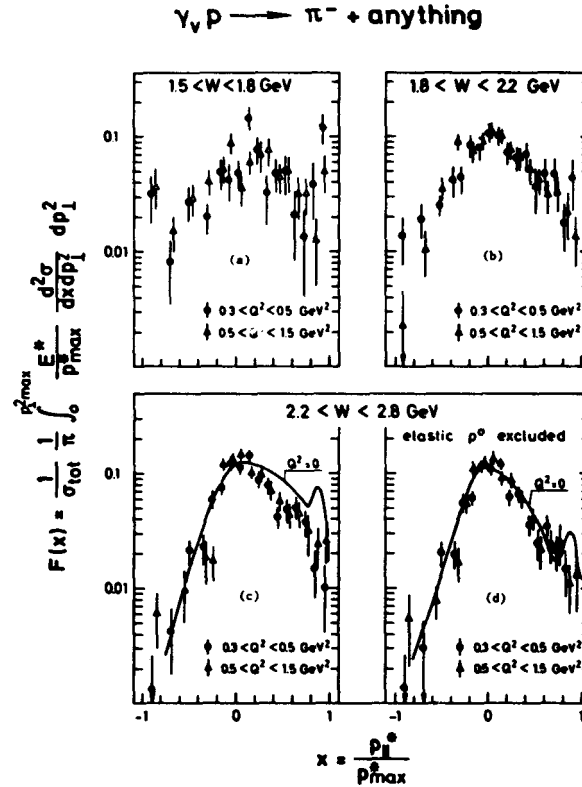


Fig. 4: Inclusive π^- structure function⁵ as a function of x_F in different W and Q^2 intervals. In (d) contributions from the reaction $ep \rightarrow ep\rho^0$ are removed. The solid lines in (c) and (d) represent data from ref. 15.

The production of baryons exhibits a completely different pattern as can be seen in fig. 5. There the same quantity as in fig. 4 is depicted as a function of x_F for the production of Λ hyperons¹⁹ and it is compared with a smoothed representation of various results^{9,20,21} on proton production. A production is about ten times less frequent than proton production. This further supports the observation made in connection with quark charge retention (fig. 3) that in the fragmentation process the creation of $s\bar{s}$ pairs is less probable than that

of $u\bar{u}$ and $d\bar{d}$ pairs.

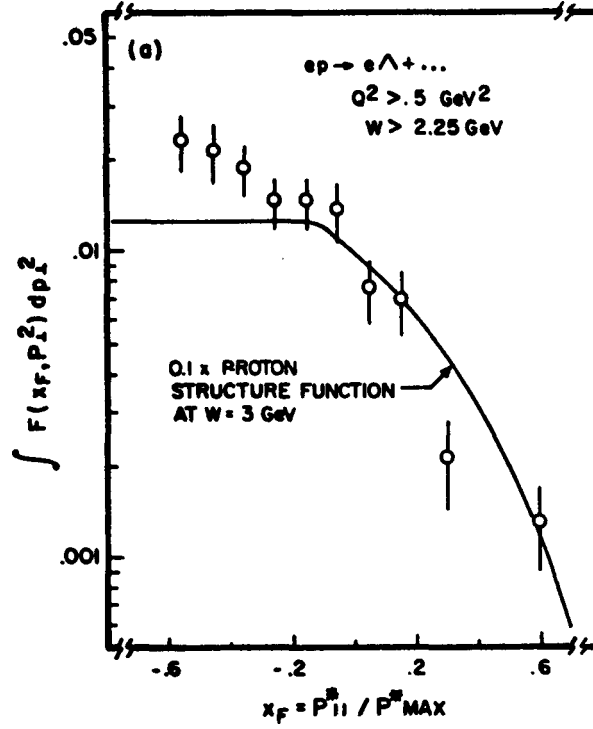


Fig. 5: Structure function for inclusive electroproduction of Λ hyperons¹⁹. For comparison proton production^{9,20,21} scaled by 0.1 is shown as well (solid line).

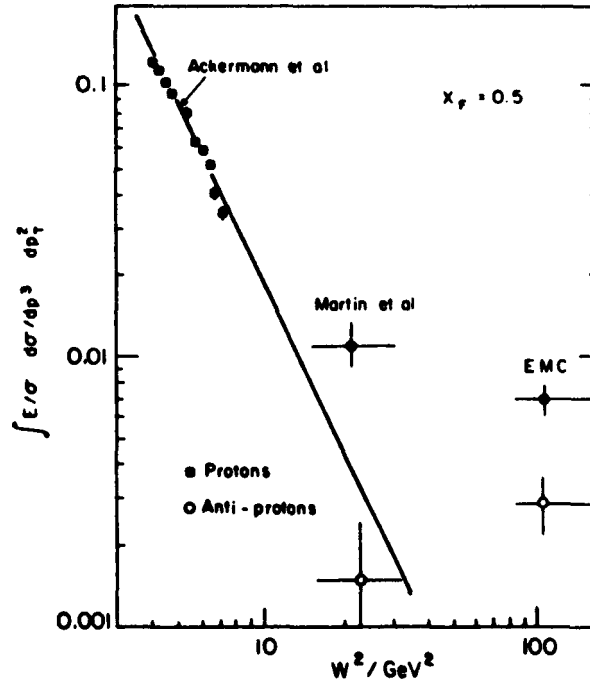


Fig. 6: Relative proton yield at $x_F = 0.5$ as a function of W . The data are taken from ref. 20 (squares) and references 23 and 24 (circles).

Baryon production occurs mainly at negative x_F and shows a rather flat

behaviour there. In the forward hemisphere the distributions fall off very steeply with increasing x_F , more steeply than in the case of pions (fig. 4). Both facts can be understood by interpreting baryons as predominantly being target fragments. At low energies those target remnants can also spill into the forward direction. With increasing W the spill-over becomes less important²² as can be seen in fig. 6 where the relative proton yield is shown as a function of W^2 . The results at high energies^{23,24} do not follow the trend extrapolated from the low energy points²⁰. This observation together with further results on proton and antiproton production in μp scattering²⁵ and e^+e^- annihilation²⁶ proves that baryons are also produced in quark fragmentation. However at CEBAF energies in both the forward and backward hemispheres baryons must basically be regarded as target remnants. In order to study quark fragmentation into positive mesons it is indispensable to be able to identify protons.

In the quark-parton model hadron production is expected to be independent of the event variables, i.e. the fragmentation functions D_i^h should only depend on x_F (or z). This prediction is usually tested by comparing measurements of $\frac{1}{\sigma_{tot}} \frac{d\sigma^h}{dz}$ (equation (11)) at different Q^2 or W (see also discussion in connection with fig. 4). This procedure is not completely correct since varying Q^2 at fixed W or vice versa results in a variation of x . If the hadron spectra from u and d quarks do not agree hadron production cannot be expected to be independent of Q^2 and W . These complications can be avoided if one considers the sum of π^+ and π^- production. Starting from equations (10), (11) and (12) one can show that

$$\frac{1}{\sigma_{tot}} \left(\frac{d\sigma^{\pi^+}}{dz} + \frac{d\sigma^{\pi^-}}{dz} \right) = D_u^{\pi^+}(z) + D_u^{\pi^-}(z) \quad (14)$$

and

$$\frac{1}{\sigma_{tot}} \frac{d\sigma^{\pi^0}}{dz} = D_u^{\pi^0}(z). \quad (15)$$

These results are obtained using the following charge conjugation and isospin relations among the fragmentation functions:

$$D_u^{\pi^+} = D_d^{\pi^-} = D_{\bar{d}}^{\pi^+} = D_{\bar{u}}^{\pi^-} \quad (16)$$

$$D_u^{\pi^-} = D_d^{\pi^+} = D_{\bar{d}}^{\pi^-} = D_{\bar{u}}^{\pi^+} \quad (17)$$

$$D_u^{\pi^0} = D_d^{\pi^0} = D_{\bar{d}}^{\pi^0} = D_{\bar{u}}^{\pi^0} \quad (18)$$

The fragmentation functions for π^0 's are correlated to those for π^+ and π^- mesons via

$$D_q^{\pi^0} = \frac{1}{2} (D_q^{\pi^+} + D_q^{\pi^-}). \quad (19)$$

So the quark-parton model predicts that

$$\frac{1}{\sigma_{tot}} \frac{d\sigma^{\pi^0}}{dz} = \frac{1}{2\sigma_{tot}} \left(\frac{d\sigma^{\pi^+}}{dz} + \frac{d\sigma^{\pi^-}}{dz} \right) \quad (20)$$

and both π^0 production and the sum of π^+ and π^- production do not depend on x .

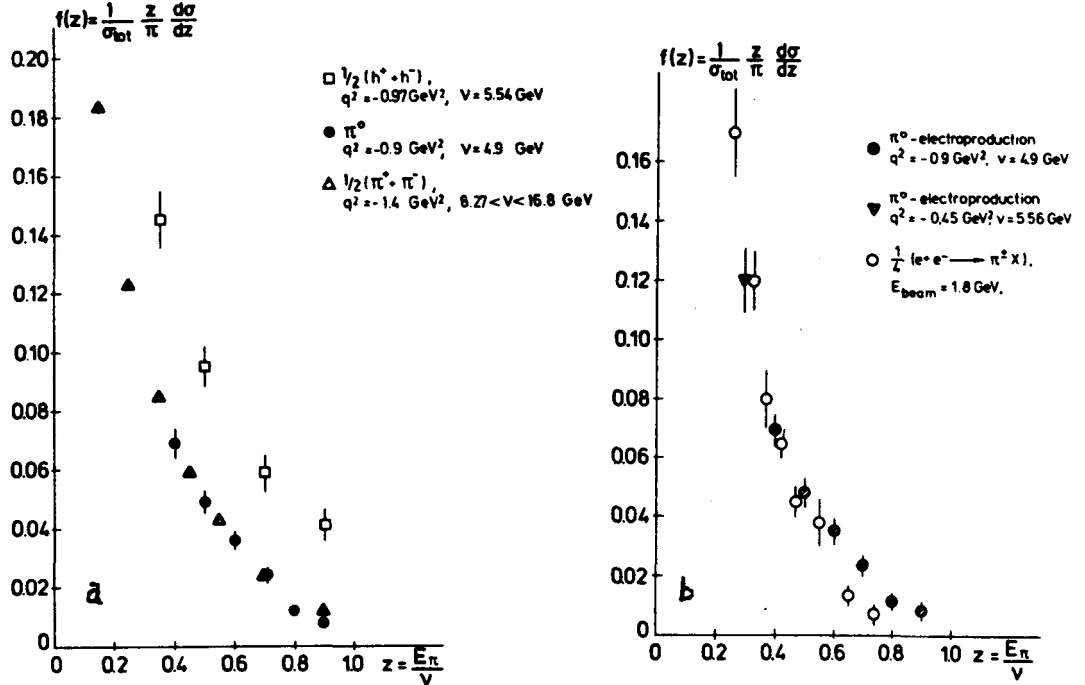


Fig. 7: (a) Comparison of π^0 structure functions²⁷ with charged pion²³ and hadron²⁸ electroproduction. (b) Comparison of π^0 electroproduction²⁷ with the production of charged pions in e^+e^- annihilation²⁹.

The validity of relation (20) is experimentally confirmed as can be seen in fig. 7a: π^0 production²⁷ is compared with the mean of π^+ and π^- production²³ and the distributions agree very well although the x values are slightly different. The results containing all charged hadrons²⁸ lie well above the pion data. This shows again the importance of particle identification.

Since the fragmentation process should be the same whether the fragmenting quarks are produced in electron scattering or in e^+e^- annihilation the fragmentation functions measured in these reactions are expected to agree. In e^+e^- annihilation equations similar to (14) and (15) can be derived but with an additional factor 2 on the right hand side since q and $\bar{q}$ are always created in pairs so that each event contains two quark jets. Thus π^0 electroproduction has to be compared with the quarter of $\pi^+ + \pi^-$ production in e^+e^- annihilation²⁹. This comparison is shown in fig. 7b. Rather good agreement is observed but

the e^+e^- distribution is somewhat softer.

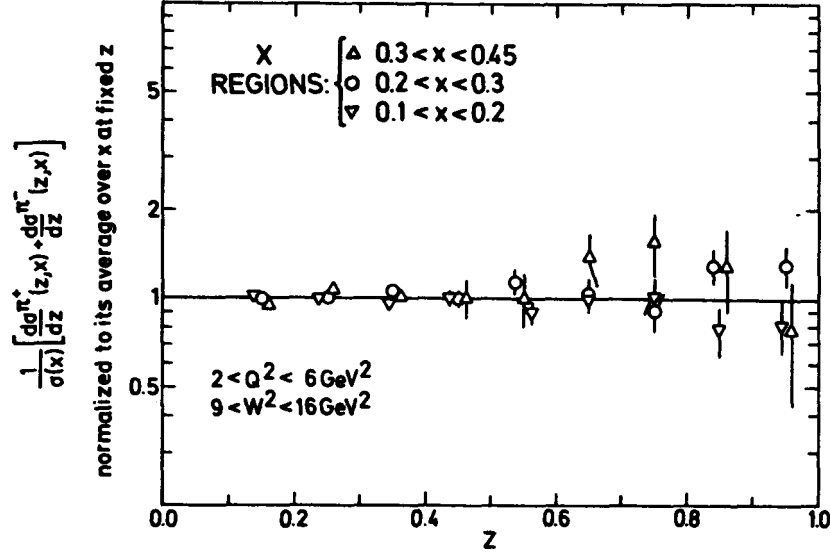


Fig. 8: Sum of the z distributions for π^+ and π^- electroproduction³⁰ in three different x intervals divided by their average over x .

Equation (14) was obtained based on the fundamental quark-parton model assumption of factorisation. A test of this basic property was performed by the DECO collaboration³⁰ (fig. 8). Within the statistical and systematic uncertainties $\frac{1}{\sigma_{tot}} \left(\frac{d\sigma^{\pi^+}}{dz} + \frac{d\sigma^{\pi^-}}{dz} \right)$ does not show any significant dependence on x . However the data are also consistent with next to leading order QCD calculations³¹.

To summarize the discussion so far one must say that it still remains open whether deviations from the quark-parton model predictions have been observed at low energies. Furthermore, even if such deviations will be detected in the future it will stay unclear whether they can be unambiguously attributed to perturbative quantum chromodynamics. The contributions from diffraction, mass effects and final state interactions will have to be understood first.

In the quark-parton model the final state hadrons are interpreted as either quark (forward hemisphere) or diquark fragments (backward hemisphere). In the centre-of-mass system a two jet structure is therefore expected, the two jets being collinear, similar to the event structure found in e^+e^- annihilation. In ep scattering the directions of the jet axis and of the virtual photon should be strongly correlated, deviations being due to the intrinsic transverse momentum of the quark in the target proton. To investigate the validity of this expectation in ep scattering events the jet axis is determined by sphericity

$$S = \frac{2}{3} \min \left(\frac{\sum_i p_{\perp i}^2}{\sum_i p_i^2} \right) \quad (21)$$

or thrust

$$T = \max \left(\frac{\sum_i |p_{||i}|}{\sum_i |p_i|} \right) \quad (22)$$

where the sums run over all charged hadrons and $p_{\perp i}$ and $p_{\parallel i}$ are determined in the centre-of-mass system with respect to the sphericity and thrust axis respectively. The distributions³² of $|\cos\Theta|$ (Θ : angle between the sphericity or thrust axis and the direction of the virtual photon) shown in fig. 9 exhibits a strong correlation between the jet axis and the virtual photon direction as expected. Events due to diffractive processes, e.g. $ep \rightarrow ep\rho^0$, are suppressed.

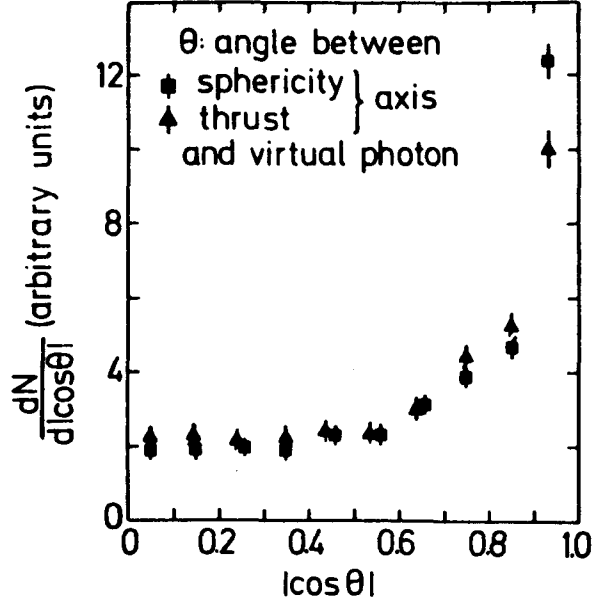


Fig. 9: Distribution of $|\cos\Theta|$ for events from ep scattering with more than three final state hadrons at $9\text{GeV}^2 < W^2 < 16\text{GeV}^2$ and $Q^2 > 1\text{GeV}^2$.

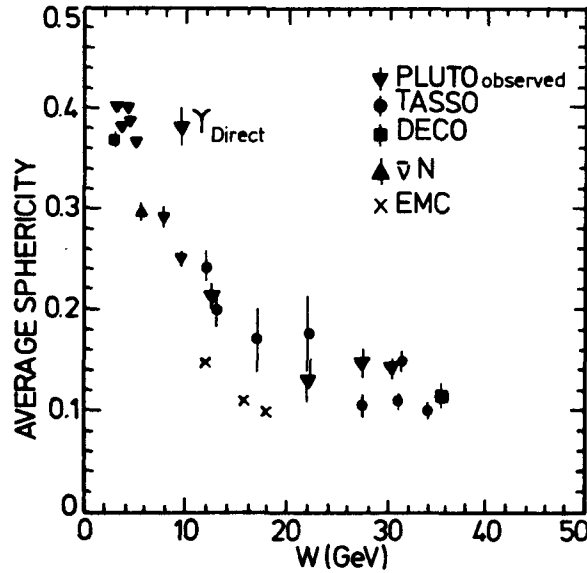


Fig. 10: Average sphericity as a function of W . The references are given in the text.

From the width of the $|\cos\Theta|$ distribution the mean value of the intrinsic trans-

verse momentum of the quark inside the target proton was estimated to be 0.7 ± 0.2 GeV.

The values of the sphericity (equation (21)) can vary between 0 and 1. Two jet like events have sphericities near 0 whereas sphericity values near 1 represent spherical events. In fig. 10 the mean sphericity is shown as a function of W for e^+e^- annihilation (PLUTO³³, TASSO³⁴) and lepton nucleon scattering (DECO³⁵, $\bar{\nu}N$ ³⁶, EMC³⁷). $\langle S \rangle$ decreases with increasing W for both types of reactions thus giving evidence for stronger collimation of the two back-to-back jets with rising energy. The results from lepton scattering lie systematically below those from e^+e^- annihilation. This can be due to the contribution of heavy quark fragmentation in e^+e^- annihilation which is absent in lepton scattering and also to the slightly lower charged hadron multiplicities observed in lepton scattering⁸ compared with e^+e^- annihilation.

BEYOND THE QUARK-PARTON MODEL

In the previous section it became clear that our present understanding of electron scattering is almost completely based on the quark-parton model. As could be shown many very different aspects of the experimental results can be successfully interpreted in this common framework. However, the simple quark-parton model cannot be the end of the story since it lacks a dynamical basis to be regarded as a theory. Quantum chromodynamics is commonly considered as the correct theory of strong interactions but up to now quantitative predictions are very difficult to achieve and are usually only obtained for high momentum transfer where perturbative methods can be applied. The deviations from the quark-parton model which are observed at low (CEBAF) energies are due to the fact that low Q^2 photons do not resolve the parton structure of the target. So more collective processes are expected to occur. These will be discussed in this section.

At low W near the elastic limit resonance production is important. This results in a distinct pattern of the total electroproduction cross section^{38,39} as a function of W showing several clear maxima which are correlated to baryon resonances. However as one can see in fig. 11 this structure gets somewhat less pronounced when Q^2 increases. It was also found⁴⁰ that scaling of νW_2 can be extended into the resonance region if one uses $x' = \frac{Q^2}{Q^2 + W^2}$ as scaling variable instead of $x = \frac{Q^2}{Q^2 + W^2 - M^2}$ and if the resonance peaks are averaged out. Furthermore, the contributions from longitudinal photons (σ_L) to the total cross section are only of the order of 10 - 20 % in the resonance region³⁹. These observations lead to the conclusion that in general the quark-parton picture can be used even at rather low Q^2 and W . Only very specific reactions need different explanations.

In the simple quark-parton model for $Q^2 \rightarrow \infty$ σ_L is expected to be zero

since only transverse photons (helicity ± 1) can be absorbed by the quarks which have spin $\frac{1}{2}$. The determination of σ_L is therefore important if mechanisms others than those described in the quark-parton model should be investigated.

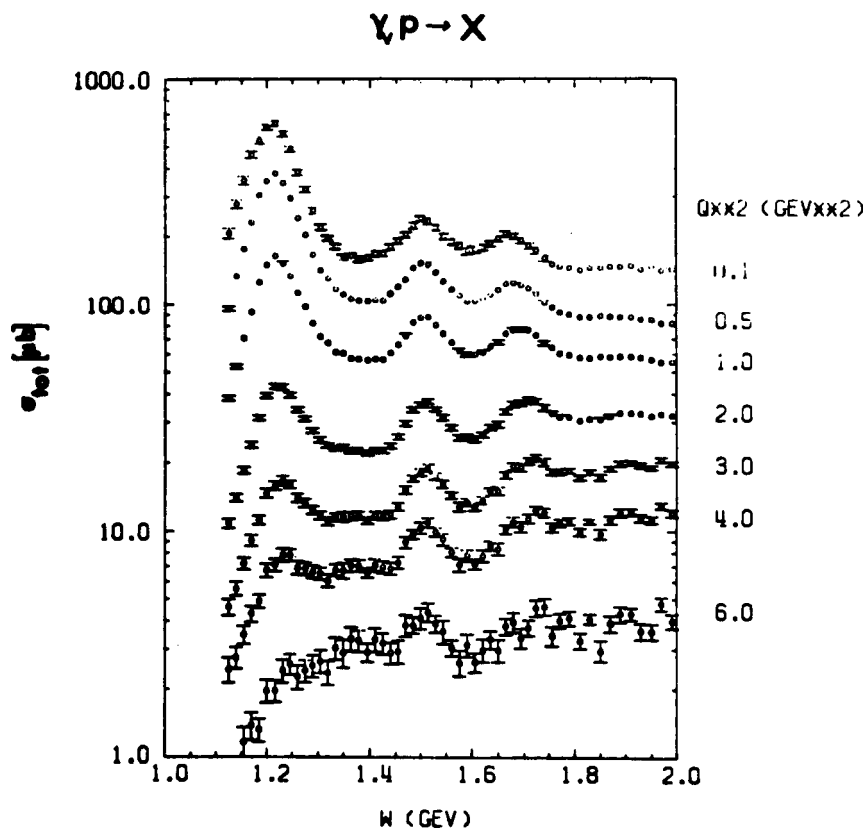


Fig. 11: Total electroproduction cross section^{38,39} as a function of W at different Q^2 values.

It follows from equations (4) and (5) that in order to extract σ_L one has to measure at different values of ϵ keeping Q^2 and W fixed. This can only be achieved by changing the beam energy. Since σ_L usually is very small^{39,41} compared to σ_T a rather long lever arm in ϵ is needed. The consequences of this for experiments at CEBAF will be discussed in the next section.

An example of such a measurement⁴² is given in fig. 12. It becomes obvious from this figure that besides high statistics small systematic uncertainties must be aimed for between the measurements at different beam energies (i.e. ϵ).

A summary of results on $R = \sigma_L/\sigma_T$ is shown in fig. 13. In fig. 13a measurements⁴³ of R at relatively low Q^2 ($\leq 10\text{GeV}^2$) are given as a function of x . All data points are significantly different from zero and lie systematically above the QCD expectation. In contrast to that at high Q^2 ($\geq 10\text{GeV}^2$) no deviations of R from zero are found⁴². So one can expect that scattering from

objects other than bare spin $\frac{1}{2}$ quarks becomes important at low Q^2 .

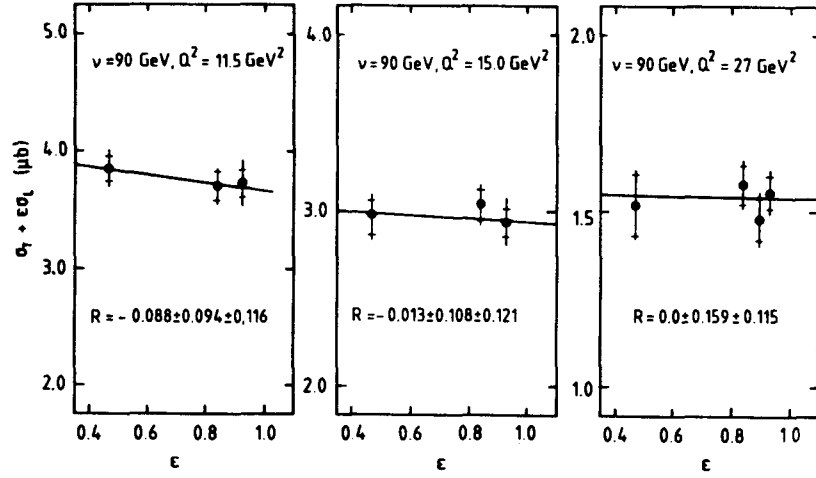


Fig. 12: Total electroproduction cross section⁴² as a function of ϵ . To the statistical errors (inner error bars) the systematic errors are added. The straight lines represent fits from which the values of R are derived.

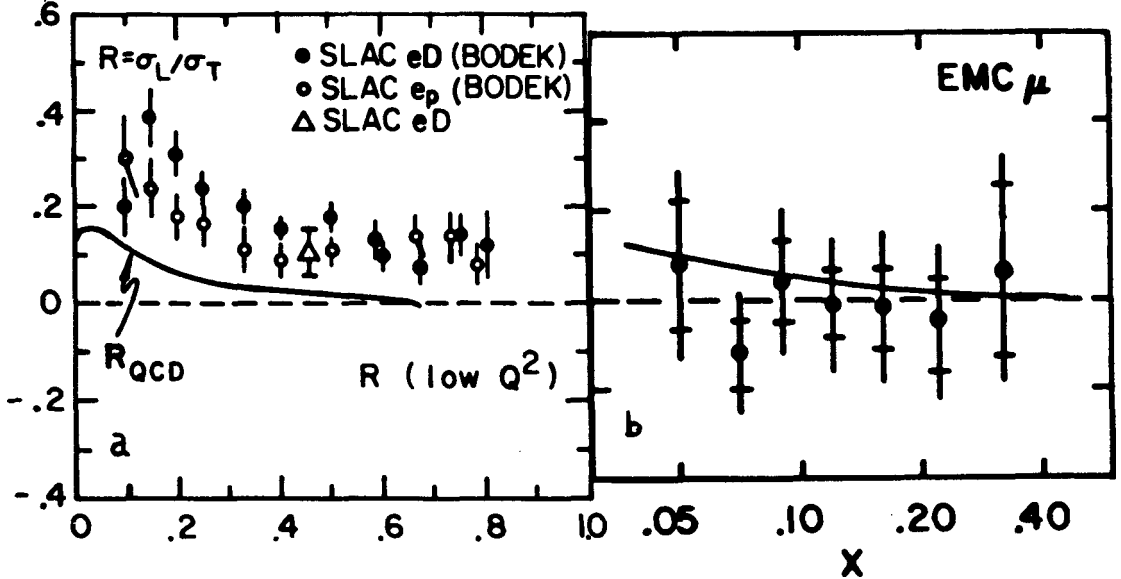


Fig. 13: $R = \sigma_L/\sigma_T$ as a function of x (a) at low Q^2 ($\leq 10 \text{ GeV}^2$)⁴³ and (b) at high Q^2 ($\geq 10 \text{ GeV}^2$)⁴². The leading order QCD predictions (solid lines) are discussed in references 42 and 43 respectively.

It was pointed out in the introduction to this article that electroproduction of exclusive channels was extensively discussed in the last year's CEBAF Workshop¹. Nevertheless I would like to add a few comments to what was said there. My remarks will refer to the production of π^+n , $K^+\Lambda(\Sigma^0)$ and ρ^0p in ep scattering and I will concentrate on alternative ways to understand these channels compared with the discussion in ref. 1.

Although the contributions from longitudinal virtual photons to the total cross section are only of the order of 10 - 20 % they can be much stronger for specific exclusive channels. At small momentum transfer the reaction $ep \rightarrow e\pi^+n$ is even dominated⁴⁴ by σ_L (fig. 14). This is easily understood by the dominance of the one pion exchange contribution in the framework of Born term models by which this process can be successfully described. At t near t_{min} (forward scattering) no spin transfer to the nucleon occurs so only helicity 0 (longitudinal) virtual photons can couple to the exchanged pion. In this picture the pion form factor $F_\pi(Q^2)$ can be extracted from measurements of π^+n production in a rather model independent way.

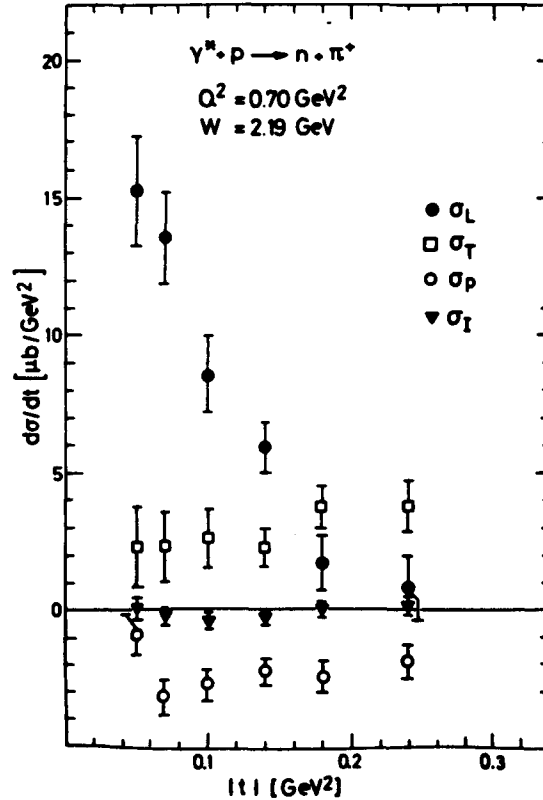


Fig. 14: $d\sigma/dt$ for electroproduction of π^+n as a function of $|t|$. The contributions due to the various helicities and polarisations of the virtual photons are separated⁴⁴.

This is the conventional picture but one can look at this reaction also somewhat differently as schematically depicted in fig. 15. Here scattering from pre-existing pions inside the target is assumed⁴⁵. The measurements of the reaction $ep \rightarrow e\pi^+n$ would then allow to determine the probability distribution $G_{\pi^+/p}(x)$ to find a pion of momentum fraction x in the proton from the following ansatz:

$$\frac{d^2\sigma}{dx dQ^2} = G_{\pi^+/p}(x) \frac{d\sigma}{dQ^2}(e\pi^+ \rightarrow e'\pi'^+) \quad (23)$$

However the pion form factor is needed as input and the cross section for elastic $e\pi$ scattering has to be calculated. So it remains open whether this alternative approach to the interpretation of π^+n electroproduction can be successfully pursued.

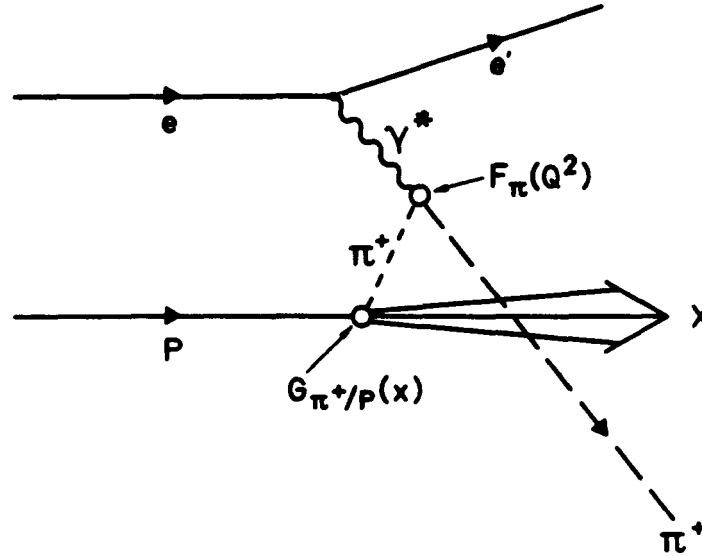


Fig. 15: Diagram of electron scattering from virtual pions pre-existing in the target proton.

In analogy to $ep \rightarrow e\pi^+n$ one would expect that the exclusive channels $ep \rightarrow eK^+\Lambda$ or $eK^+\Sigma^0$ could be used to determine the kaon form factor $F_K(Q^2)$ in the space like region. However one kaon exchange is not so clearly dominating⁴⁶ these processes as one pion exchange does in π^+n production. It was observed⁴⁷ that the Σ^0/Λ ratio drops with increasing Q^2 . In the picture of one kaon exchange this could be understood if the coupling constant of the exchanged kaon to the nucleons were greater in the case of $K^+\Lambda$ production than for $K^+\Sigma^0$ production ($g_{\Lambda NK}^2 \gg g_{\Sigma NK}^2$). This is not really supported by the values of these coupling constants derived from various experiments⁴⁸. So since the dominance of one kaon exchange is not unambiguously established in these reactions another explanation⁴⁹ of the dropping Σ^0/Λ ratio may be more profound. It is based on simple quark-parton model arguments and predicts Σ^0/Λ to decrease with increasing x where scattering off u quarks dominates. There the remaining target system is preferentially in a state of isospin $I = 0$ which combined with a strange quark will result in a Λ hyperon. Then at $x \rightarrow 1$ Σ^0 suppression naturally follows but strictly speaking this picture applies only to the behaviour of the transverse part of the cross section. Since the experiments suggest a non negligible contribution from longitudinal photons ($\sigma_L \approx \sigma_T$) to $K^+\Lambda$ production none of the two mechanisms alone may be able to explain the experimental data. Separation of σ_T and σ_L is therefore needed to improve the present status of understanding.

Exclusive ρ^0 production (i.e. $ep \rightarrow e\rho^0p$) clearly is the first place where

the vector meson dominance (VMD) model should apply. In this model the dependence of the cross section on Q^2 can be approximately parametrized as^{50,51}

$$\sigma(Q^2) = \sigma(0) \frac{1}{\left(1 + \frac{Q^2}{m_\rho^2}\right)^2} \left(1 + \varepsilon \xi^2 \frac{Q^2}{m_\rho^2}\right). \quad (24)$$

This formula implies that with increasing Q^2 the contribution from longitudinal photons gets more important (second term within the brackets on the right hand side) if ξ^2 is different from zero. This parameter ξ^2 can be determined from the Q^2 dependence of the exclusive ρ^0 production cross section by comparing the functional dependence (24) with the experimental results (fig. 16). The data^{52,53,54,55} which were obtained at relatively high beam energies are best described by the curve with $\xi^2 = 0$. This would mean that the contributions from longitudinal photons disappear, at least at high W . This may not be surprising since at low W of less than 3 GeV a decreasing importance of σ_L with increasing W was already observed⁵⁶.

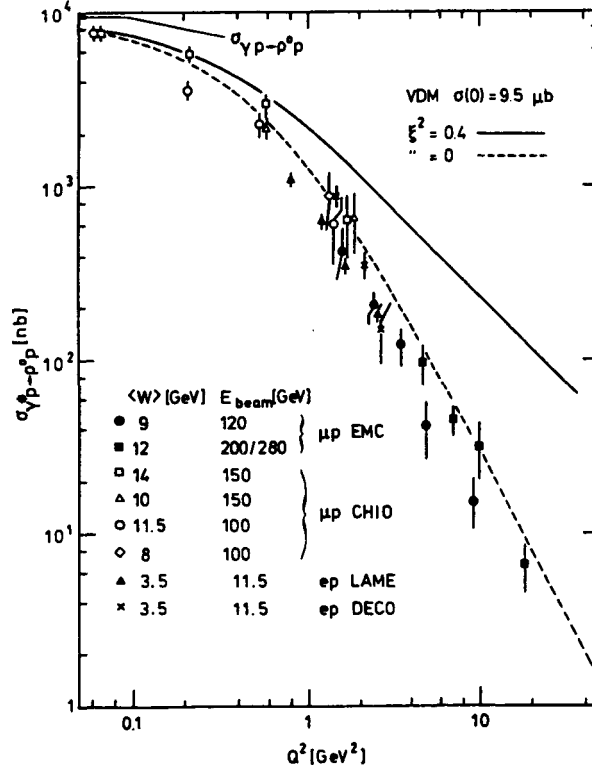


Fig. 16: Total cross section for exclusive ρ^0 electroproduction as a function of Q^2 . The data are taken from measurements carried out by the following collaborations: EMC⁵², CHIO⁵³, LAME⁵⁴, DECO⁵⁵. The lines represent VMD calculations⁵⁰ with different values of the parameter ξ^2 .

In order to determine σ_L in a model independent way measurements of

the exclusive ρ^0 production cross section as a function of ϵ have been carried out at high W . The results are shown in fig. 17 in which results from EMC⁵² and CHIO⁵³ are combined. No increase of $\sigma_T + \epsilon\sigma_L$ with ϵ can be observed thus confirming the absence of contributions from longitudinal photons.

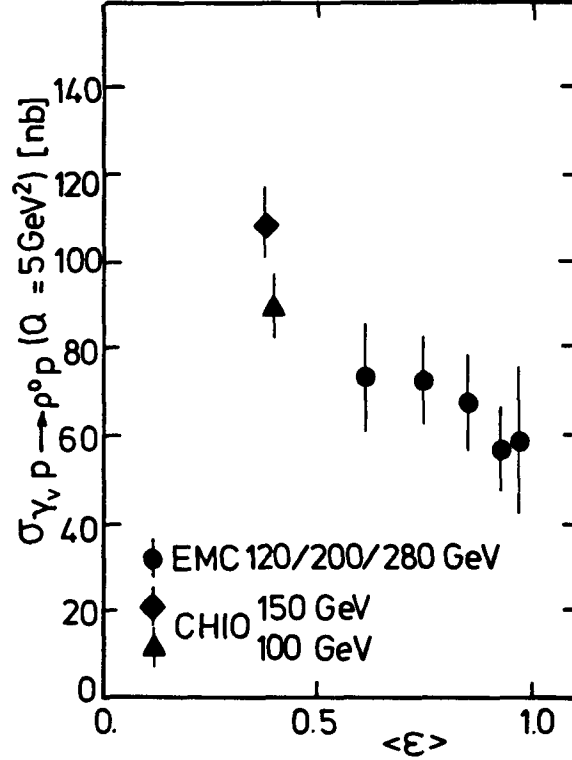


Fig. 17: Total cross section for exclusive ρ^0 electroproduction as a function of ϵ . The data stem from EMC⁵² and CHIO⁵³.

At low energies s-channel helicity conservation (SCHC) was found to work in ρ^0 production⁵⁶. This means that the helicities of the virtual photon and the ρ^0 meson agree. In this picture the angular distribution of the decay pions in the ρ^0 rest frame can be described by

$$W(\cos\theta) = \frac{3}{4}[1 - r_{00}^{04} + (3r_{00}^{04} - 1)\cos^2\theta] \quad (25)$$

where θ is the polar angle between the ρ^0 direction and the π^+ meson. Integrations were performed over the azimuthal angles of the π^+ with respect to the ρ^0 direction and of the ρ^0 production plane with respect to the lepton scattering plane.

In the VMD framework the parameter ξ^2 is related to the density matrix element r_{00}^{04} via

$$r_{00}^{04} = \frac{1}{1 + \frac{m_\rho^2}{\epsilon\xi^2 Q^2}}. \quad (26)$$

In fig. 18 r_{00}^{04} is shown as a function of Q^2 . Fitting equation (26) to these data gives $\xi^2 = 0.40 \pm 0.13$, in clear disagreement with what was concluded from the Q^2 dependence of the cross section (fig. 16). So SCHC seems to break down at high W since longitudinal photons do not contribute to ρ^0 production whereas from the analysis of ρ^0 decay a sizeable fraction of the ρ^0 mesons are found to have helicity 0 ($r_{00}^{04} \neq 0$). Exclusive ρ^0 production no longer behaves like a diffractive process at high W .

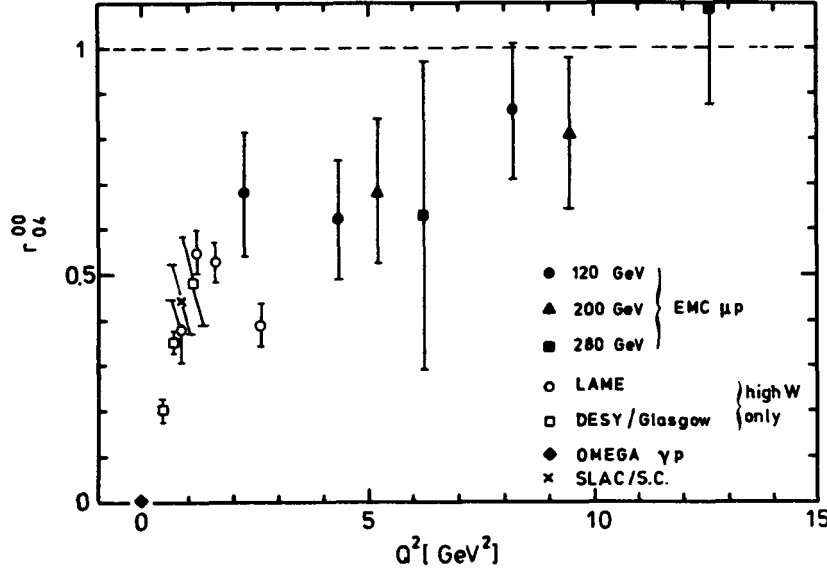


Fig. 18: ρ^0 density matrix element r_{00}^{04} . Results from various experiments are shown: EMC⁵², LAME⁵⁴, DESY-Glasgow⁵⁶, SLAC⁵⁷, OMEGA⁵⁸.

The breakdown of SCHC was found by the analysis of the Q^2 dependence of $\sigma(ep \rightarrow e\rho^0 p)$ and of r_{00}^{04} over a large Q^2 range. The high Q^2 data stem from experiments performed at high beam energies. It clearly would be interesting to see whether the same observations could be made also at low energies but high Q^2 . The high duty cycle of CEBAF which allows to reach high Q^2 may help in answering this question.

REMARKS ON EXPERIMENTS AT CEBAF

The kinematic range in which experiments at CEBAF are carried out is the transition region to the deep inelastic regime. The Q^2 variable is the most important parameter describing this transition since it is a measure of the resolving power of the probe. The aim should be to reach the highest Q^2 values possible. This is not easy since the cross section drops dramatically with increasing Q^2 as shown in fig. 19. There in the Q^2 - W^2 plane contours of constant total cross sections are drawn which are obtained from equation (10) using the parametrisation of F_2 as given in ref. 42. The beam energy is taken to be 4 GeV. The curves stop at the kinematic limits which are obtained under

the assumption that the minimum energy to measure the scattered electron is 0.5 GeV. It is clear from this picture that primary energies higher than 4 GeV would be desirable in order to reach high Q^2 also at larger W^2 values.

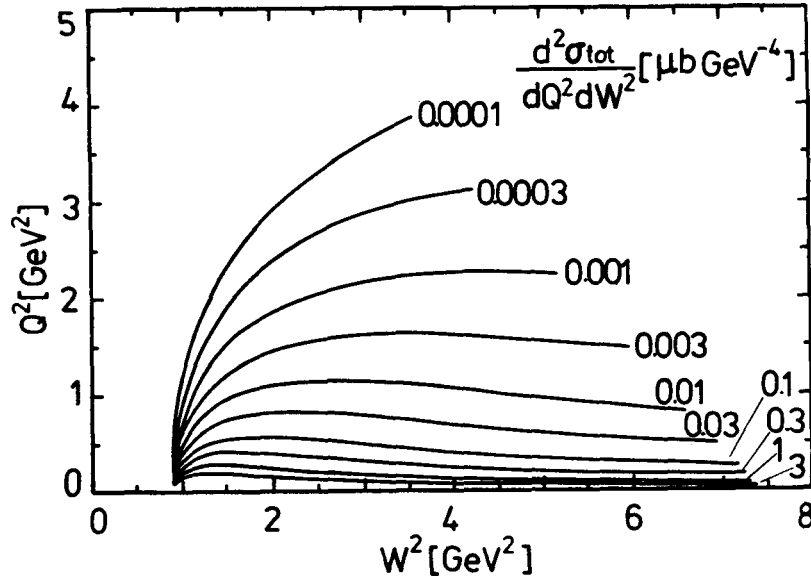


Fig. 19: Contours of $\frac{d^2\sigma}{dQ^2 dW^2}$ in the Q^2 - W^2 plane for the beam energy of 4 GeV.

Experimental results at high x are very scarce. This has essentially the same reason as why high Q^2 data are lacking too since large x implies large Q^2 . At low W high x can be reached more easily (fig. 20) so CEBAF is an ideal place to study the high x region.

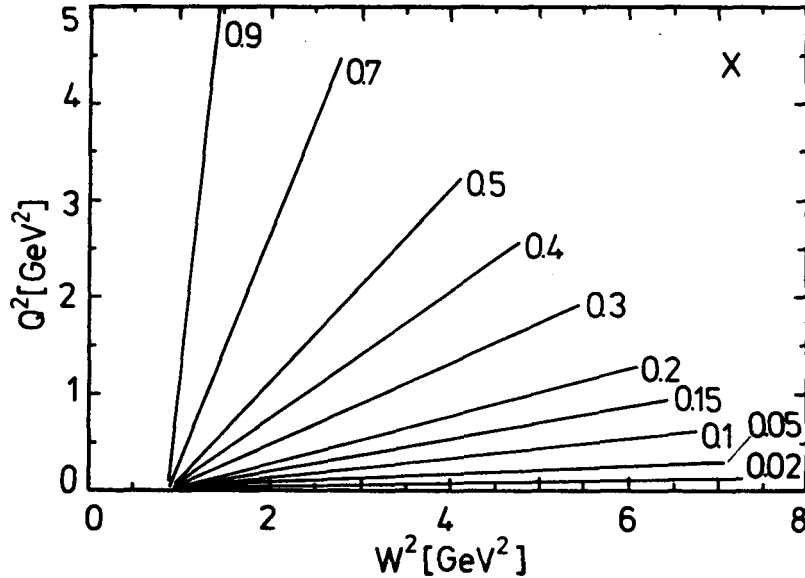


Fig. 20: Contours of x in the Q^2 - W^2 plane.

Another very important task of experiments at CEBAF is the separation between σ_T and σ_L . This needs measurements at different beam energies in order to vary ε while Q^2 and W^2 stay unchanged. The kinematic range in which such a σ_T/σ_L separation can be performed is very limited as one can see in fig. 21. Contours for the same 10 values of ε are shown for beam energies of 4 GeV (solid lines) and of 2.5 GeV (dotted lines) respectively. Also in this case the demand for higher energies is obvious.

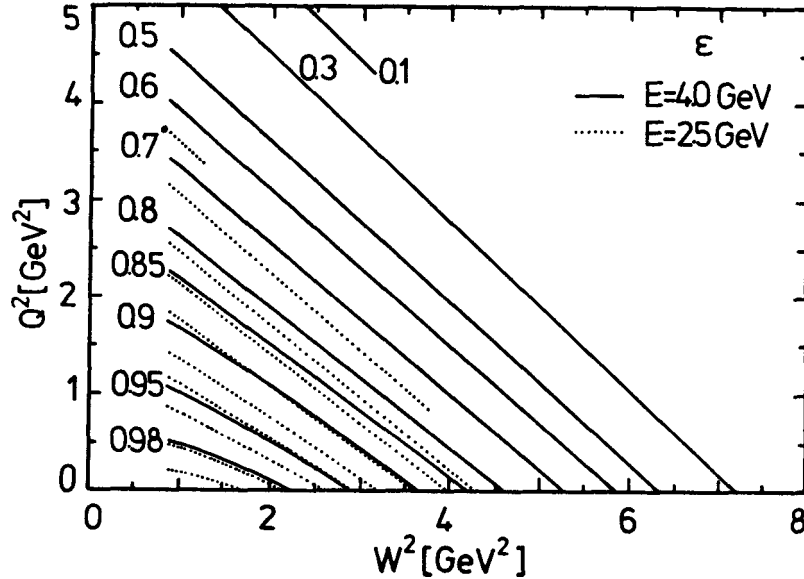


Fig. 21: Contours of ε in the Q^2 - W^2 plane. The solid lines refer to the beam energy of 4 GeV, the dotted lines to 2.5 GeV. The values of ε are the same for both energies but only the 4 GeV lines are labelled.

To summarize, the most interesting region is the high Q^2 (high x) domain. This implies that the scattered electron has to be detected and identified at rather large scattering angles and low momenta. There the π^- background is sizeable and good electron/pion separation is important. In order to obtain enough statistics the solid angle in which the scattered electron can be accepted must be large.

In general one must say that experiments at CEBAF dealing with particle physics must be rather sophisticated since a lot of experimental results are already available from Bonn, CEA, Cornell, DESY, NINA and SLAC. In order to be able to improve our understanding of this subject only experiments of advanced technologies should be envisaged.

The high duty factor of CEBAF will allow to obtain high statistics. To take full advantage of the potentially higher statistical accuracy the systematic uncertainties must be kept to a minimum. Therefore the following requirements should be fulfilled: Since hadron masses cannot be neglected at CEBAF energies complete particle identification is necessary. To be able to reconstruct the topology of the events with high reliability detectors of good space and

time resolution are needed near to the target. One should also try to avoid that the performance of the detector differs significantly in various angular and momentum ranges. This is especially important when data from different primary energies are combined as it is necessary for the σ_T/σ_L separation.

The electromagnetic background induced by the beam causes major problems in the bending plane of the magnetic field. To solve this difficulty properly is unavoidable. One should however keep the blind region of the detector around the beam to a minimum since the majority of the hadrons are created with rather small angles relative to the beam. An open dipole magnet with a superconducting tube⁵⁹ to protect the beam region against the magnetic field is in my opinion the most promising solution for a large acceptance detector.

SUMMARY

It has been demonstrated that the quark-parton model allows to understand to a great extent the properties of hadronic final states produced in electron scattering, even at rather low Q^2 and W . However, at those low Q^2 and W values additional mechanisms are needed to interpret the experimental results. The explanations developed so far have still large uncertainties so that data of much better quality are necessary in the transition region to the domain of the quark-parton model. Large acceptance detectors with good particle identification and photon detection should be aimed for. The high Q^2 (high x) region should be especially taken care of. Separation between the contributions from transverse and longitudinal photons is important. In order to fully exploit the potentially good statistical accuracy which is provided by the high beam flux of CEBAF it is of great importance to keep the systematic uncertainties as small as possible.

ACKNOWLEDGEMENTS

The invitation to give a summary report on hadron production in lepton scattering at the CEBAF Workshop offered me the opportunity to critically reconsider the main part of my scientific activities of the past fifteen years. I appreciate this very much and I hope that the physical understanding which evolved during this period has been able to help the preparation of a valuable particle physics programme at CEBAF.

I want to thank Prof. F. Gross for having invited me to this workshop which I enjoyed very much because of the friendly atmosphere and lively discussions, especially in the working groups. During the preparation of this talk I received valuable remarks from Drs. J. Gayler and T. Sloan which I gratefully acknowledge.

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**NUCLEAR PHYSICS AT PEP –
FIRST TEST AND FUTURE PLANS**

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ABSTRACT

A test run of internal target nuclear physics at the PEP storage ring is described. The Time Projection Chamber (*TPC* -2γ detector) was used to detect the inelastically scattered electron and complete hadronic final state in the interaction of 14.5 GeV electrons with D_2 , Ar and Xe gas targets. The data comprises mostly low- x low- Q^2 events, but some deep inelastic scattering as well. The future possibilities of a dedicated nuclear physics program at PEP are outlined.

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I. Introduction

Nuclear physics with internal targets at storage rings is principally motivated by experiments either requiring the use of very thin targets, or for which only very thin targets are available. An instance of the former is the measurement of slow nuclear recoils¹⁻⁴, and obviously of the latter are certain polarized targets, *e.g.* $\vec{H}e$ and $\vec{d}$ for which realizable densities at present are only on the order of 10^{15} cm^{-2} . Such experiments may be performed with reasonable luminosities $\mathcal{L}$ in virtue of the large stored currents possible at rings. Specifically regarding electron scattering experiments in the multi-GeV range, at present coincidence measurements could only be performed at a handful of e^+e^- facilities in the world until dedicated accelerators such as CEBAF are completed. This is so as the bunch structure in *PEP* for example more closely approximates a cw beam than that of the direct SLAC beam.

This report describes a brief test run performed in February, 1986 at the *PEP* storage ring at SLAC with 14.5 GeV electrons and nuclear targets. This work originated from a few chance conversations with a number of Stanford graduate students who informed us that even at a *base pressure in the interaction regions of 10^{-9}* , 20% of the events were *beam-gas interactions*, and not e^+e^- physics! Subsequently contacts were made with the directors of the TPC- 2γ group to sort existing data tapes for nuclear physics event, and after a brief study the decision was made to proceed with a proposal to SLAC and TPC- 2γ for a dedicated run with a simple gas target. We are very grateful to Prof. Burton Richter, director of the Stanford Linear Accelerator Center, and Dr. Piermaria Oddone, spokesman for *PEP-4 - PEP-9* for their extensive cooperation.

The focus of the study from the physics perspective was the measurement of exclusive hadronic final states in deep inelastic scattering. Nuclear modifications of quark structure functions have been studied in inclusive e , μ scattering and in charge-changing ν interactions. Nevertheless, much less has been learned from

hadronic final states, and the entirety of that has come from the μ work with limited statistics and resolution. A more immediate and practical aim of the run was to establish rates and backgrounds with a view to establishing a dedicated nuclear physics interaction region at *PEP* in the future.

The report is outlined as follows. Chapter II will consist of a brief discussion of the *PEP* ring and environment; chapter III concerns the *TPC* (*PEP-4*) detector, and the 2γ (*PEP-9*) spectrometers for small-angle tagging. Chapter IV describes how the nuclear physics run was carried out, and gives a very preliminary view of the data. Lastly, our plans for the future at *PEP* are summarized in Chapter V.

II. The *PEP* Storage Ring

The *PEP* storage ring became operational for high energy physics in 1981. The ring has a circumference of 2.2 km, and consists of six straight sections where the interaction regions (IR's) are situated, and six arcs. The straight sections are long, 130 meters, to assist cleaning up the synchrotron radiation from the last dipoles in the arcs and thus minimize the x-ray flux to the inner radius of the detector. Furthermore, the vacuum in the interaction region must be as good as possible for e^+e^- physics. The straight sections minimize the synchrotron radiation hitting the walls, and are equipped with more than 3000 liters/sec ion pump capacity. Typical base pressures of low- 10^{-9} torr are limited by thermal outgassing of UHV construction material (300-series baked stainless steel, etc.). (By contrast, the synchrotron radiation load on the vacuum chamber in the arcs is approximately 40 watts/cm^2 , with the base pressure consequently being in the low- 10^{-8} range.)

Fig.1 shows the layout of the ring with the most recent experiments in the even-numbered regions. *PEP* normally runs at 14.5 GeV for high energy physics. The stored current is typically 20 mA per beam in a 3-bunch structure. Thus beam crossings occur at each IR every 2.4 μs .

A detailed analysis of what limits the thicknesses of gas targets that may

be introduced into the *PEP* storage ring has been carried out by Spencer⁵. The only consideration of importance for compatibility of an internal target nuclear physics program with a high energy program going on simultaneously is the limitation of beam lifetime due to bremsstrahlung from the gas target. (The compatibility of internal target physics and synchrotron radiation research that will be carried on at *PEP* by *SSRL* has *not* been addressed yet; the low-emittance lattice being developed for *SSRL* operation very likely will raise other issues, *e.g.* growth in phase space due to multiple scattering.) The storage lifetime as a function of gas pressure is given by:

$$\frac{T_o}{\tau} \simeq \left[\frac{4}{3} \ln \left(\frac{\gamma}{\delta\gamma} \right)_{RF} - \frac{5}{6} \right] 4\sigma_o Z(Z+1) \ln \left(\frac{183}{Z^{1/3}} \right) \times \left[\frac{N_A}{A} \rho^{STP} l_t \left(\frac{P}{760} \right) \left(\frac{273}{T} \right) \right] \quad (1)$$

Some of the constants and variables that may not be obvious are: l_t , the length of the target in the beam; $\sigma_o = \alpha r_e^2$, where r_e is the classical electron radius; T_o , the time of revolution around the ring; and $\left(\frac{\delta\gamma}{\gamma} \right)_{RF}$ is the RF capture bucket width of the ring. This last quantity can be $\simeq \pm 1\%$.

Fig. 2 shows both the maximum permissible densities, and corresponding beam-gas luminosity $\mathcal{L}_{beam-gas}$ for a variety of possible target gases spanning the periodic table. These values were based on an acceptable storage lifetime of 2 hours, and a beam current of 20mA. The highest attainable luminosities, 10^{33} , are for hydrogen and deuterium, the target densities here being limited to a few tens of *nanograms/cm²*. Heavy gases imply maximum target thicknesses of a few *nanograms/cm²*; rate considerations may dictate lower thicknesses yet. How one achieves and monitors such low prescribed densities will be briefly touched upon in Chapter V.

Fig.1 also shows the immediate future home of nuclear physics at *PEP*, IR-8, and where experiments will be conducted when longitudinally polarized electrons become a reality. IR-6 is the preferred region for polarization work, as some modifications to the lattice will need to be made on either side of the

straight section (introduction of special dipoles and solenoids) and thus regions neighboring RF sections are excluded. Furthermore, introduction of a spin-precessor 180° across from the IR is required in some schemes, necessitating an open region there too⁶.

III. The Time Projection Chamber

The *TPC-2 γ* complex (cutaway perspective view, Fig. 3) is a marriage of two detectors representing different physics interests. The important features of each will be briefly described below.

1. The Time Projection Chamber (*TPC*)

The principle of the *TPC* was first conceived by Nygren in 1974^{7,8}. Fig. 4 shows the *TPC* in schematic as it exists at *PEP*. The *TPC* is a 2 meter diameter, 2 meter long volume of high-pressure gas (8.5 atm. of *Ar - Methane*) with a central high voltage membrane nominally at -60 kV. The chamber is situated in a 1.3 T solenoidal magnetic field. Ionization along the tracks of charged particles drift towards either endcap with a speed of 16 μ s per meter; passing through the gated grid, the track arrives at an array of 183 anode wires (per 60° sector) in a time sequence obviously related to the dip angle of the track. Furthermore, since there are 18 cathode pad rows per sector from the inner to the outer radius of the *TPC*, the induced ionization on the pads provides both the azimuthal angle and magnetic rigidity of the particle. The highly redundant measurement of dE/dx along the track by the anode wires coupled with the momentum resolution $(\Delta p/p)^2 = (0.06)^2 + (0.006p)^2$ results in the best particle identification in any high energy detector today.

As an integral detector, the *TPC* has inner and outer drift chambers for triggering purposes, $10X_0$ hexagonal and pole-tip calorimeters (40 layers of *Pb-3-view* stereo wire chambers), and muon identification. The principle interest of the *TPC* collaboration (*PEP-4* experiment) is the study of $e^+e^- \rightarrow q\bar{q}$ annihilation.

2. The 2γ Spectrometer⁹

The 2γ collaboration (*PEP-9* experiment) studies hadron production from the fusion of two virtual photons. As the electron and positron scatter but do not annihilate, and these processes are strongly peaked at low- Q^2 , good small-angle tracking and calorimetry are essential.

Electrons (or positrons) and hadrons pass through a Cerenkov detector, Drift Chambers 1-5 (actually 15 planes of u , v , y chambers each resolving the left-right ambiguity, the anodes being twin sense wires closely spaced), a Time-of-Flight Hodoscope, and Calorimetry. The latter consists of a segmented hexagonal NaI array (20 - 100 mr), and a Pb-Scintillator sandwich shower counter (100 - 200 mr). Electrons or positrons scattered at wider angles may be seen in the Pole-Tip Calorimeter in the *TPC* itself. Momentum analysis is performed by a Septum magnet (Fig. 5) which provides a moderately large magnetic volume, while being field-free on the beam axis. The North and South ends of the *PEP-9* Spectrometer are symmetric.

IV. Nuclear Physics with the *TPC*

Described below are some details concerning the operation of the *TPC* during the test run in February, 1986, and a first look at the data taken.

1. Experimental Details

The most straightforward (and primitive) gas target was developed for the experiment, namely a gas bleed system which would introduce gas continuously into the straight section in the vicinity of the *TPC*. The target gas was therefore uniformly distributed throughout the entire region of the *TPC-2 γ* detector, the pressure profile diminishing at each pumping station as one moves from the center of the IR up and down the tunnel. The entire pressure bump extended over tens of meters in either direction. Nevertheless, in view of the diluteness of the targets (the maximum central pressures being $(3.0)10^{-7} \text{ torr}$ for D_2 , $(5.0)10^{-8}$ for Ar , and $(2.5)10^{-8}$ for Xe), the influence on the stored beam was negligible (compare Fig. 2).

The low flow rates required ($10^{-4} - 10^{-5} \text{ torr} - \text{liter/sec}$) to bring the

equilibrium pressure up to the desired levels was achieved with the gas injection system of Figure 6, designed and constructed at LLNL. The essential features of its operation is as follows. The Intermediate Volume (between V-6 and V-7) is charged with research purity grade gas at the delivery pressure of the regulator (usually 22 p.s.i). Valve V-6 is closed, and V-7 opened, allowing the gas to expand into the Control Volume. Valve V-9 is opened, presenting the Control Valve V-L with a pressure of approximately 90 *torr*. Valve V-L is a precision conductance that allows fine variations in the leak rate into the *PEP* ring at the 4 *m* point through a 65' stainless steel line.

Figure 7 shows a side view of the *TPC*- 2γ detector in schematic for a nuclear physics event. The electron (or positron) scatters at small angles, usually into the NaI detector. The hadronic final state (usually protons and pions, with relatively fewer kaons or heavier nuclear fragments, *e.g.* deuterons) is identified either in the *TPC*, or in the 2γ spectrometer. The interaction may occur anywhere along the beam pipe; however due to the angular focussing of the tracks due to fixed target kinematics, it is most desirable for the vertex to be somewhat upstream of the beam crossing point ($z = 0$), in order for hadrons at small angles to be detected with good momentum and dE/dx resolution.

Four triggers were implemented for nuclear physics running. Three required a *single tag* in *PEP*-9, *i.e.* $E > 2\text{GeV}$ in the NaI, or $E > 10\text{GeV}$ in the Shower Counter, with either (i) a hit in the *TPC* inner drift chamber, (ii) a *ripple* (an outside-in time ordering of hits on the *TPC* pad rows), or (iii) a hit in the *PEP*-9 muon chambers which could result from a fast pion. The fourth trigger required high energy deposition ($> 2\text{GeV}$) in the Pole-Tip Calorimeter, to look for high- Q^2 deep inelastic scattering. As during the dedicated nuclear physics running, colliding beams were still provided in IR-2, the usual *double-tag* triggers designed to enhance 2γ physics were specifically disabled.

Figure 8 maps the four-momentum transfer $Q^2 = -q^2 = 4EE'\text{Sin}^2\frac{\theta}{2}$ and the Bjorken scaling variable $x_{Bj} = Q^2/(2M\nu)$ as a function of the electron

(or positron) scattering angle θ and final lab energy E' . (In the usual way, M is the nucleon mass, and $\nu = E - E'$.) The calculations are appropriate for $E = 14.5\text{GeV}$ and the kinematical regions of x and Q^2 accessed by each calorimeter is marked.

With a circulating current of approximately 15 mA per beam, the event rate in the TPC- 2γ detector was 6 - 8 sec^{-1} . The corresponding dead time was upwards of 50%. Independent data tapes are written for PEP-4 and PEP-9. All of our analysis to date has been performed with the PEP-9 code OASIS.

Figure 9 shows the excellent particle identification of the TPC bases on momentum and dE/dx . Ambiguities remain at the cross-over points resulting from the relativistic rise in energy loss, *e.g.* for $\pi - K$ separation. The distribution of the origin of tracks with respect to the coordinate along the beam direction, z is shown in Figure 10. The sharp peak ($\sigma \approx 6\text{cm}$) is due to e^+e^- physics and the underlying continuum is due to beam-gas nuclear physics. Fiducial cuts are made in software at $\pm 15\text{cm}$ and $\pm 80\text{cm}$, however events are accepted within the beam-crossing region if they contain at least one proton. An actual distribution of events in x, Q^2 (Figure 11) for $e^\pm$ single tags manifests the kinematic boundaries associated with the maximum angle of the NaI and Shower Counter calorimeters. (Note the sparse region corresponding to smaller angles in the Shower Counter due to deterioration of the scintillator there.)

2. Atlas of Events

It is worthwhile looking at a representative sample of nuclear physics events, both to see the variety of physics in $e - A$ interactions at 14.5 GeV as well as to appreciate what kind of information the TPC provides (and does not provide). The figures that will be presented are essentially those available as run-time graphics (including color-coding of the tracks according to particle type!).

a. Quasielastic Scattering

Figures 12 and 13 show a top and endcap view respectively of quasielastic $e - p$ scattering in ^{40}Ar . The scattered positron is detected in the North Shower

Counter (13.0 GeV , $\theta_{e'} = 6.5^\circ$), and the recoil proton ($p_p = 1.765 \text{ GeV}/c$) is seen in the *TPC*. While the energy resolution of the Shower Counter presently is somewhat poor, the high x , Q^2 ($0.84, 2.4(\text{GeV}/c)^2$) is corroborated by the angle and momentum of the proton. The endcap view shows that the proton has punched through not only the superconducting coil and hexagonal calorimeter, but also two layers of iron in the muon identifier.

b. *Delta Electroproduction*

Figure 14 displays the event $d(e, e' \Delta)$ with the Δ identified by the invariant mass of the subsequent decay $\Delta \rightarrow p\pi^-$. While the large momentum ($p_\Delta = 1.8 \text{ GeV}$) and backward angle of the Δ is somewhat surprising, it should be noted that its transverse momentum is roughly balanced by the scattered e^+ . The spectator proton is not observed; although there is some energy deposition in the North Pole Tip Calorimeter (which could be hit near its inner radius without leaving a track in the *TPC*), it is more likely that it ranged out in the beampipe.

c. *Associated Strangeness Production*

The excellent particle identification and missing mass resolution of the *TPC* are evident in the event of Figure 15, $^{40}\text{Ar}(e, e' \Lambda K^+)$. The $p\pi^-$ has nearly the invariant mass of the Λ . (For events with separated vertices ($c\tau_\Lambda = 7.89 \text{ cm}$) the energy losses in the beampipe wall are not treated correctly except with specialized programs which track the decay partners back to the secondary vertex.) The error perimeters (2σ in the momentum and lifetime plane) all lie on their expected loci and are well away from any crossover points. The K^+ is not seen to decay in this event ($c\tau_{K^\pm} = 370.9 \text{ cm}$); in many events however the kaon track will show a kink indicating the decays $K \rightarrow \mu\nu, K \rightarrow \pi\pi^0$.

d. *Deep Inelastic Scattering*

A deep inelastic scattering event with a ^{40}Ar nucleus appears in Figure 16. Here $Q^2 = 3.7(\text{GeV}/c)^2$ and $x = 0.44$. Nothing other than the tag (inelastically scattered electron) appears in *PEP-9*, whereas the *TPC* shows a $2\pi^+ 2\pi^- p$

hadronic final state. In terms of the fragmentation variable $x_F \equiv p_{||}/p_{max}$, where $p_{||}$ is for each observed hadron the component of momentum along the direction of the three-momentum transfer q . Only 47% of the three-momentum transfer is accounted for in this event, although $x_F = 0.23$ for one π^+ . In a *wire hit* plot two further tracks are evident, which were not picked up by the analysis routine, as they crossed the *TPC* at too shallow an angle for tracking and identification.

e. *Massive Target Fragmentation of Xe*

A spectacular example of target fragmentation is seen in Figure 17, resulting from electron scattering with an energy deposition of 10 *GeV* in a *Xe* nucleus. The event corresponds to absorption of a nearly real photon ($Q^2 = 0.07, x = 0.004$). Two noteworthy features of the event are (i) that most of the particles are emitted at very wide angles ($\approx 90^\circ$), and (ii) the large d/p ratio (see figure 18). The number of deuterons is large even regarding that of the four tracks out of range on the momentum- dE/dx plot (starred on the radial and endcap view), three are protons and one is a kaon.

Such events however are rare (only three with $M \geq 12$). Figure 19 displays the charge balance in the *TPC* only for each target. This data suggests that most of the particles observed result only from the struck nucleon. The argument is as follows: in an ideal detector that would see even particles of very low momentum, scattering from a proton(neutron) would result in a charge balance $\langle \Sigma(q_i) \rangle$ of 1(0) independent of final multiplicity, so long as there was no intranuclear cascade or large energy deposition to the target spectator. The probability of interacting with a nucleon is given in the most naive quark-parton model by the weighted sum of squares of u, d quark charges in the nucleon. Summing over the protons and neutrons in the target nuclei results in an expected charge balance of 0.6 for d and 0.51 for *Xe*. The actual values are 0.58 for d , and 0.93 for *Xe*. While this is by no means constitutes firm evidence, it at least leaves open the possibility that there is a target spectator a majority of the time. One could contemplate experiments where conventional particle-gamma spectroscopy on

the target spectator is performed in coincidence with detecting the inelastically scattered electron or other particles in the jet regime from the struck nucleon. This would be of obvious importance in measuring *e.g.* F_2 for nucleons within specific nuclear shell states, in view of the controversy surrounding the importance of nuclear binding effects¹⁰.

3. Bose-Einstein Correlations in $\pi^\pm\pi^\pm$

As an example of analysis that is important to do for leptonproduction with nucleon and nuclear targets, we discuss Bose-Einstein correlations in like-sign pion production. Briefly, quantum statistics imposes correlations (anticorrelations) between identical bosons (fermions) at small relative four-momenta. The correlation function, which thus exhibits an enhancement (depletion) at small relative momenta, reflects the space-time evolution of the emitting source.

The basic ideas are summarized in Figure 20. It is recalled that in the Young experiment, there exists one source and one detector (*i.e.* there is only single-photon detection at each point on the plane of observation). An interference pattern results because there are two indistinguishable paths, and thus two amplitudes for a photon to be emitted from the source and arrive to a particular point on the screen.

In the prototype treatment of Bose-Einstein correlations, there are two point sources and two detectors that detect identical particles in time coincidence. Now there are two amplitudes for producing a coincidence; a *direct* and a *crossed* diagram ($1 \rightarrow a, 2 \rightarrow b$, and *v.v.*). If the sources and detectors are sufficiently separated that the amplitudes are indistinguishable, they will interfere, and the coincidence probability will oscillate in the angular separation of the detectors. For a realistic source, extended in space and with a finite lifetime, what results is a correlation function:

$$R \equiv \frac{\rho^{BE}(p_1, p_2)}{\rho^{UNCORR}(p_1, p_2)} \quad (2)$$

which shows a maximum at $q \equiv p_1 - p_2 = 0$ and falls off in a smooth way

(depending on the form of the source) to an asymptotic value of 1 for large q . Here $p_{1,2}$ are the four-vectors of the identical bosons, and $\rho^{BE}(p_1, p_2)$ and $\rho^{UNCORR}(p_1, p_2)$ are the two-boson densities with and without the incorporation of quantum statistics. (How one determines an *experimental* density without the influence of Bose-Einstein statistics will be discussed later.) As an example, a space-time source of the form:

$$S(r, t) \propto e^{-r^2/2\sigma^2} e^{-t/\tau} \quad (3)$$

results in a correlation function:

$$R(p_1, p_2) = 1 + \frac{\exp(-q^2 r^2)}{(1 + (q_0 \tau)^2)} \quad (4)$$

where q , q_0 are the three-momentum and energy of q .

(Bose-Einstein correlations are referred to variously under several names. Originating in long-baseline radio astronomy around 1950, it is also known as the Hanbury-Brown/Twiss effect after the original workers. It is typically called HBT in relativistic heavy ions, and GGLP (Goldhaber-Goldhaber-Lee-Pais) in particle physics. It is frequently termed intensity interferometry, and second-order interferometry.)

Figure 21 shows Bose-Einstein correlations in like-sign pion production both from $e^+e^- \rightarrow q\bar{q}$ jets (coincidentally from the TPC Collaboration at PEP¹¹), as well as from pp and $p\bar{p}$ interactions from the ISR¹² at $\sqrt{s} = 63\text{GeV}$.

The e^+e^- data are plotted in (i)(a),(i)(b) against q_T and q_0 , the variables conjugate to the size and lifetime of the emitter, respectively. The variable q_T is the component of the three-momentum difference transverse to the sum momentum of the two pions, and for this plot a cut has been imposed of $q_0 < 0.2\text{GeV}$. The variable q_0 is the magnitude of the energy difference between the pions, and a cut here has been made of $q_T < 0.2\text{GeV}/c$. The approximately Lorentzian shape of the distribution in q_0 is noteworthy. The characteristic source sizes (Fig. 21, (ii),(iii)) derived from the work are surprisingly small –

about $0.65 fm$. Furthermore, subdividing the data in angle between the jet axis and pion pair momentum gives evidence for an elongated source with a major-to-minor axis ratio of 2 or 3 to 1. This is consistent with a string or flux-tube picture of quark hadronization, at least qualitatively. The pp , $p\bar{p}$ experiments (Fig. 21 (iv)) showed no difference between them in source size, and found sizes not greatly bigger than those in e^+e^- production, here typically $1.1 fm$.

In spite of the paucity of statistics in the present run, the Hanbury-Brown/Twiss effect is evident in our data even with just a few hundred $\pi^\pm\pi^\pm$ events per target gas, although it is unrealistic to expect an definitive quantitative analysis. (Note that typical published reports on pion interferometry work with a data sample of many tens of thousands of like-sign pions.) Nevertheless, it should be mentioned that Bose-Einstein correlations in leptonproduction from nuclear targets should be expected to reveal two additional facets not contained in e^+e^- or $p\bar{p}$ production, and therefore are important to pursue in the future. First, within the language of the string model of hadronization, it might be expected that the nuclear medium might change the value of the string constant due to color screening; this might be reflected in different source dimensions for the fast pions resulting from the fragmentation of the struck quark. Second, there will be a component of the pions, lower in energy, due to excitation of the target spectator.

A very preliminary analysis of BE correlations in like-sign pion events has been performed for the data of the February run. All events with two or more like-sign pions were sifted from the data set, and every event was manually scanned to ensure that there were indeed at least two distinct π^+ or two distinct π^- . Scanning is essential, as occasionally the pattern recognition program will break a single track either as it crosses a sector boundary or the central membrane, and thus a single pion will give rise to two tracks extremely close in momentum, *i.e.* $q \approx 0$, the worst place such a problem could occur. The uncorrelated density $\rho^{UNCORR}(p_1, p_2)$ was formed by combining pions from *different* events. This is one of two usual methods for deriving the uncorrelated density,

the other by using $\pi^+\pi^-$ pairs. There are two advantages of event mixing over $\pi^+\pi^-$ pairs. First, with unlike-sign pions, care must be exercised in the low- q data to remove residual correlations due to resonance decay, *e.g.* $\omega \rightarrow 3\pi$. Second, since the number of combinations of pions one can make from event mixing is enormous, the contribution to the statistical error in R from the denominator can be made trivially small. This is important for a small data sample as in the present case, where the statistical error in $\rho^{BE}(p_1, p_2)$ is substantial.

The ratio R is displayed for all three targets in Figure 22. The variable it is plotted against is $q \equiv \sqrt{-(p_1 - p_2)^2}$, as there were too few events to either do a fit to a two-dimensional surface, or to make cuts in q_o or q_T . (There are less than 200 like-sign pions in each sample, perhaps half of the total data for each target.) Nevertheless, the rise at low- q ($< 0.1\text{GeV}/c$) is evident in each case. The results are sufficiently encouraging that work is underway to select all the nuclear physics events collected during the high-field running, 1984-1986 due to beam-gas events from the residual vacuum, mostly CO , CO_2 . It may be estimated that the total number of events collected parasitically exceed the statistics of our dedicated gas injection run by an order of magnitude.

V. Future Plans

We want to discuss briefly our vision of what a permanent dedicated nuclear physics IR would be like at PEP, and some physics that would be ideally accomplished there.

In our discussions with the SLAC management, it seems most appropriate for the near term that further nuclear physics experiments take place in IR-8 (previous DELCO region). The first round experiments contemplated will involve both polarized and unpolarized targets, but with unpolarized electrons only. The introduction of longitudinally polarized electrons in PEP for nuclear physics experiments is straightforward⁶ but will involve some modification to the lattice on either side of the nuclear physics IR, and possibly at the antipodal point of the ring. For these reasons, IR-6 would be ideal as there are no RF

accelerating sections adjoining it in the straight section, and IR-12 (the region 180° across from it) is presently unoccupied. Nevertheless, upgrading PEP for polarized electrons will not be trivial, and should be considered some years away.

1. Conceptual Design of IR-8

■ A preliminary conceptual design of what a dedicated nuclear physics interaction region would look like is shown in Figure 23. While in the longer term one may contemplate the acquisition or construction of a comprehensive 4π detector (*e.g.* a *TPC*), we are beginning from the point of view that the first experiments will be performed with a conventional multi-arm magnetic system of moderate solid angle and resolution. This follows both from the nature of the physics (two- or three-particle coincidence measurements such as $(e, e'N)$, $(e, e'\pi)$, $(e, e'K^+K^-)$ where the fixed target kinematics focusses the scattered electron and hadrons to forward angles), and from practical considerations (the availability of existing magnets, drift chambers, etc.). The central piece of the system is most of the *PEP-9* spectrometer, which cannot be used by *TPC-2* γ for future running as the high-luminosity upgrade for PEP demands that the final quadrupoles be brought close in, right to the ends of the *TPC* proper. A septum spectrometer system (Figure 5) is attractive for storage ring physics in that it provides a field-free region for the beam axis, and a transverse magnetic field elsewhere (solenoidal magnets are notoriously poor for momentum analysis of small-angle tracks). The *PEP-9* magnets may be used back-to-back, and in fact differing from Figs. 23, 5, it will be advantageous to make the yokes actually contiguous, with a new coil winding only around the nose-piece. (At present, the $\int Bdl$ for the magnets is nowhere near the design value, both due to the magnets being only one meter deep, and to half of the amp-turns being outside the region of interest (Fig. 5).) Particle identification will be accomplished by a combination of TOF and magnetic rigidity. Separation between e , π will be done by Cerenkov detectors (not shown on two of three arms), and calorimetry (either Pb-scintillator sandwich or Pb-glass). Ultimately, two movable magnetic arms could be instrumented, one small solid angle

spectrometer bending in the horizontal plane, and a big magnetic volume with a large vertical acceptance, bending in the vertical plane. The inner shielding wall for IR-8 is 27' from the beam axis rather than 21' as is true for the other IR's and as drawn. This provides greater angular range for a pivoting arm.

For unpolarized targets, the FNAL warm gas jet target concept looks attractive from the viewpoint of simplicity and cost. In recent years, it has been used both at FNAL and the Brookhaven AGS by the Purdue group; its design and performance have been reported by Stringfellow¹³. The FNAL target is at present on sight at LLNL for profiling studies using Rutherford elastic scattering from low energy cyclotron beams. One of the principal objectives is to design de Laval nozzles (which form the supersonic jet) which can produce targets of thicknesses in the $1 - 10 \text{ ng/cm}^2$ range, or about an order of magnitude smaller than those used in the $p - A$ studies.

One issue that will be new to storage ring physics will be precise determination of integrated beam-gas luminosity. An interesting option that needs to be studied is the monitoring of Møller electrons (beam electron-atomic electron scattering).

2. Measurement of the Pionic Sea via $(e, e'\pi)$ from Nuclear Targets

An example of one of the first experiments that should be performed at PEP is virtual photoproduction of pions from nuclear targets in quasielastic kinematics. The motivation of this experiment derives from the nuclear structure function anomaly (*EMC effect*). Figure 24 summarizes the experimental situation showing for both the electron scattering data from SLAC and the new EMC muon data, the ratio of $\sigma(A)/(A\sigma(D_2))$. Shortly after the original EMC data, Llewlyn-Smith¹⁴ argued that by falling short of the energy-momentum sum rule (in the nucleon sector only) empirically by about 5%, the high- x depletion for iron was reproduced. If the missing momentum were to appear in an increase in the pionic sea for nuclei (as might be expected due to the OPE contribution to the nucleons' binding), then the low- x rise in the ratio of $\sigma(A)/(A\sigma(D_2))$

would be explained as well. (Subsequently, the effects of nuclear binding were discussed¹⁰ quantitatively from a shell-model basis.) Pursuing the question of the low- x range, Pirner and coworkers reanalyzed pion electroproduction data from a hydrogen target in terms of a pionic distribution function of the nucleon, finding a pionic component of approximately 3% in the proton, carrying 0.6% of the nucleon's momentum¹⁵. They urged that a similar study be conducted for nuclear targets.

Specifically, what they advocate is summarized in Figure 25. Coincidence electron scattering where one particle is detected along with the scattered electron is sensitive to four independent structure functions: σ_L (longitudinal), σ_T (transverse), σ_P (transverse-transverse interference), and σ_I (longitudinal-transverse interference). The interference terms may be isolated just at one beam energy, by the $e' - \pi$ angular correlation, whereas a Rosenbluth separation is necessary to isolate the longitudinal component.

This is of the essence in this problem as only longitudinal photons can couple to the pion, which is spin-0. The bottom half of Figure 25 shows in fact that only σ_L (dominated by the t-channel diagram) has an angular distribution which rises steeply towards the pion pole. Thus in the kinematic regime of that experiment, the interpretation of pion production in terms of quasielastic knockout of pions considered as partons of the nucleus seems validated. It is also significant that for the deuterium target the π^+/π^- ratio is nearly 1 at low- $|t|$, whereas for pions resulting from the coupling of virtual photons to quarks, this ratio would be 4 (These limits obviously arise from counting the squares of parton charges Σq_i^2 either at the hadron or quark level.). We are presently studying whether such an experiment could be done solely with the septum spectrometer. Here the gas jet target would be close up to the magnet such that the spectrometer would subtend a large enough solid angle to detect both the electron and pion over a large kinematical region.

VI. Acknowledgements

This test could not have been performed without the encouragement of Professor Burton Richter, Director of the Stanford Linear Accelerator Center (SLAC), and Dr. Piermaria Oddone of LBL, spokesman of the PEP-4/9 experiment. We gratefully acknowledge the participation and continuing interest on behalf of the entire PEP-4/9 collaboration, who are too numerous to cite individually here. As interfacing with the PEP vacuum system was a critical aspect of this experiment, special thanks go to Norm Dean and Irwin Levy of the SLAC vacuum group, and Darrel Erikson and Dave Appelgate of the MFE division of LLNL. The efforts of the PEP Operations Staff in an unusually difficult year should not be overlooked. The participation of John Molitoris during the run was appreciated.

This work was performed under the auspices of the U.S. Department of Energy by the Lawrence Livermore National Laboratory under contract number W-7405-ENG-48.

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FIGURE CAPTIONS

1. Overview of the PEP storage ring.
2. Maximum density and corresponding luminosity for various gases as targets in the PEP ring. A circulating current of $20mA$ is assumed, as well as a minimum storage lifetime of two hours.
3. Cutaway perspective view of the entire $TPC-2\gamma$ detector. The detector is symmetric about the beam crossing point.
4. Different views of the TPC :
 - a. The TPC proper, perspective view;
 - b. The TPC endcap. The cathode pad rows underneath the anode wires, used for track position and momentum information, are not shown.
5. View of one of the PEP-9 septum magnets along the beam axis.
6. The LLNL gas injection system. The gas composition is monitored at two points, one on the gas system (RGA1), and the other on the ring (RGA2).
7. Side view of the $TPC-2\gamma$ detector, with a nuclear physics event represented by the scattered electron being detected in the NaI calorimeter, and final state hadrons in the TPC. Most of the triggers resulted from electron (positron) tags in the NaI array ($20 - 100mr$) or shower counter ($100 - 200mr$).
8. Q^2 and x as a function of $E_{e'}$ and $\theta_{e'}$ of the scattered electron for the nuclear physics run.
9. dE/dx vs. momentum plane showing the excellent particle identification of the TPC detector. For this plot vertex cuts appropriate to nuclear physics have been made.
10. Track distribution in the TPC detector during the nuclear physics run. The beam crossing vertex peak is prominent as the run took place during colliding beam conditions.
11. Actual distribution of events in x , Q^2 plane.
12. Top view of a quasielastic scattering event.

13. Endcap view of the event of the previous figure.
14. Delta electroproduction from deuterium. The invariant mass of the $p\pi^-$ is slightly less than the peak of the Δ resonance.
15. Associated strangeness production from an ^{40}Ar target.
16. A deep inelastic scattering event from an argon nucleus.
17. Massive target fragmentation in a Xe nucleus.
18. dE/dx vs. particle momentum for the event in the previous figure. Of the four particles off the scale in dE/dx , three are protons and one is a kaon.
19. Histograms of total event charge in the *TPC* for d , Ar , Xe targets.
20. Schematic diagrams:
 - a. Young interference experiment;
 - b. Prototype Hanbury-Brown & Twiss effect;
 - c. Bose-Einstein correlation from an extended source.
21. Bose-Einstein Correlations from:
 - i. $e^+e^- \rightarrow q\bar{q}$ jets in the *TPC* ¹¹ in:
 - a. q_T (for $q_0 < 0.2 \text{ GeV}$) and in;
 - b. q_0 (for $q_T < 0.2 \text{ GeV}/c$).
 - ii. From the *TPC* work, as a function of event multiplicity:
 - a. the source radius;
 - b. the chaoticity parameter.
 - iii. The radius parameter derived from the *TPC* work, subdividing the data in various angular bins with respect to the jet axis. The dashed, solid and dotted lines correspond to the assumption of an ellipsoidal source with transverse dimension r and longitudinal dimension Kr , $K = 1, 2, 3$.
 - iv. r as a function of q_T for pp , $p\bar{p}$ collisions at the ISR.
22. Very preliminary Bose-Einstein Correlations in $e - A$ interaction (this work) for d , Ar , Xe targets. Only about half of the data is used in each sample.

23. Conceptual design of the future Nuclear Physics Interaction region IR-8.
24. Data from EMC, BCDMS and SLAC on the scaled ratio of cross sections in deep inelastic regime as a function of x . (This is not equal to $F_2^A(x)/(AF_2^N(x))$ except with specific assumptions about the ratio $R \equiv \sigma_L/\sigma_T$).
25. Two diagrams:
 - a. (Top) diagram and definition of variables for the coincidence measurement $(e, e'\pi)$;
 - b. (Bottom) the distribution in $|t|$ for the four independent cross sections for pion electroproduction from hydrogen. Only σ_L shows a steep rise towards the momentum transfer direction.

MARK II

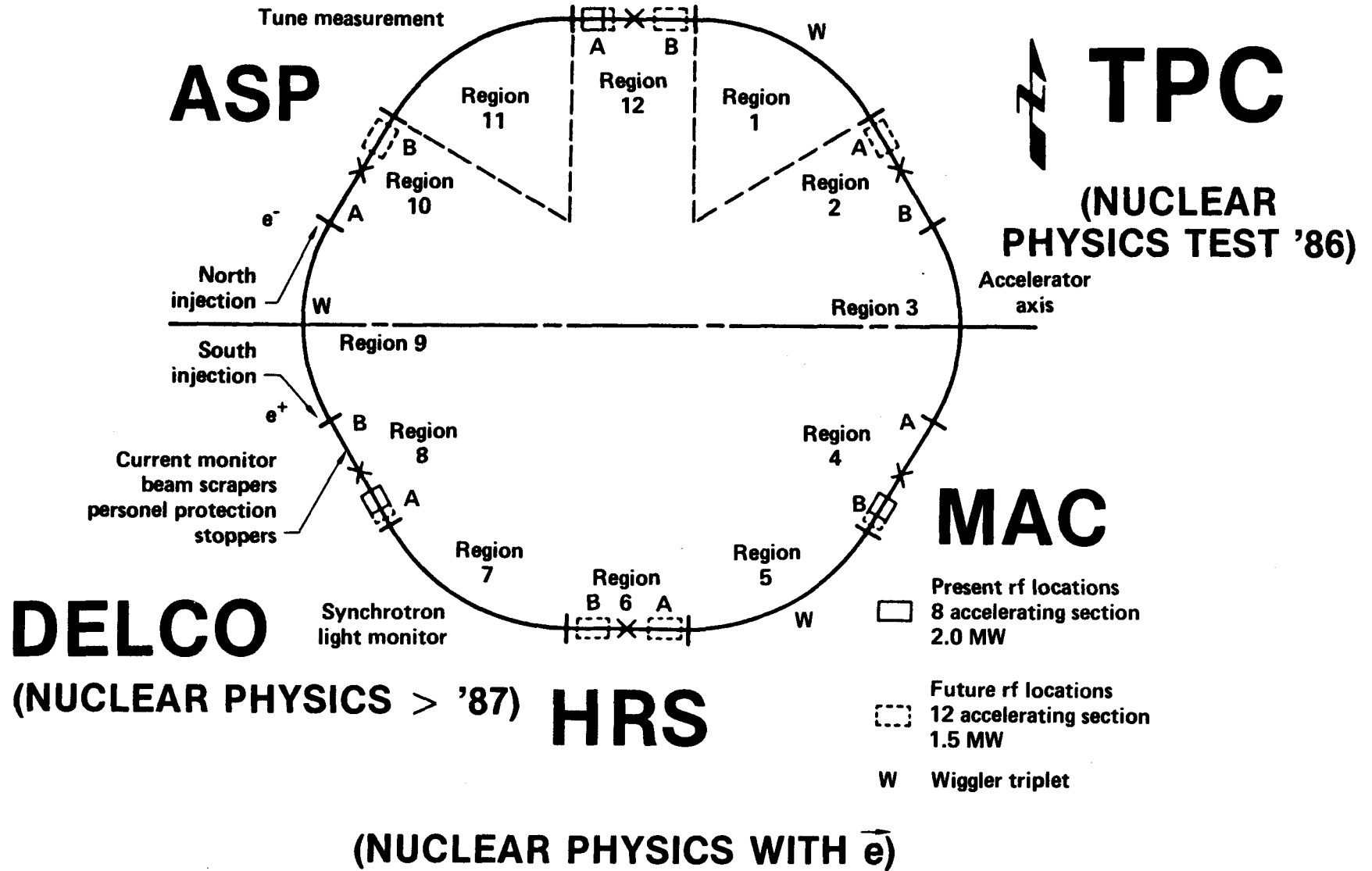


Fig. 1

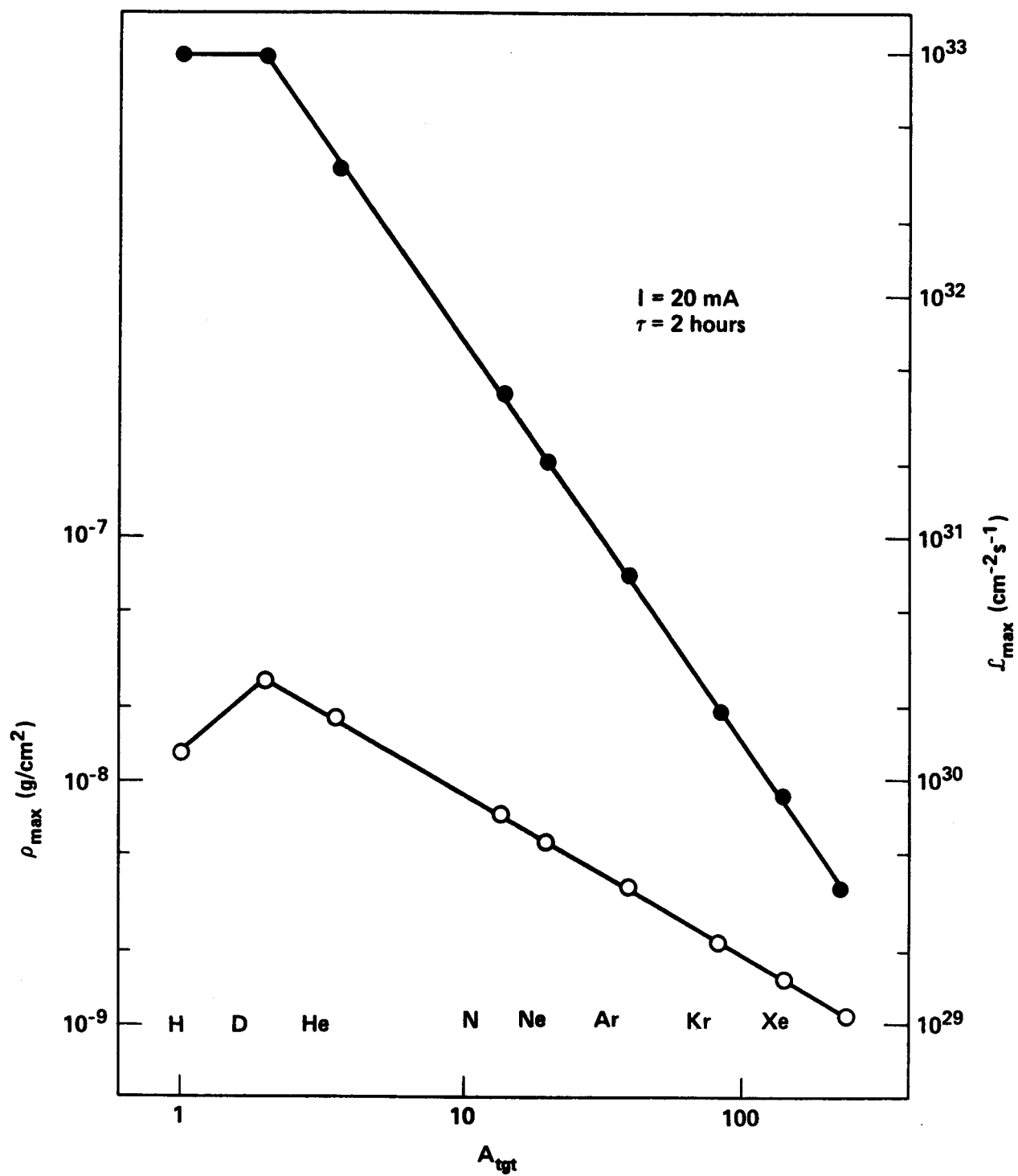


Fig. 2

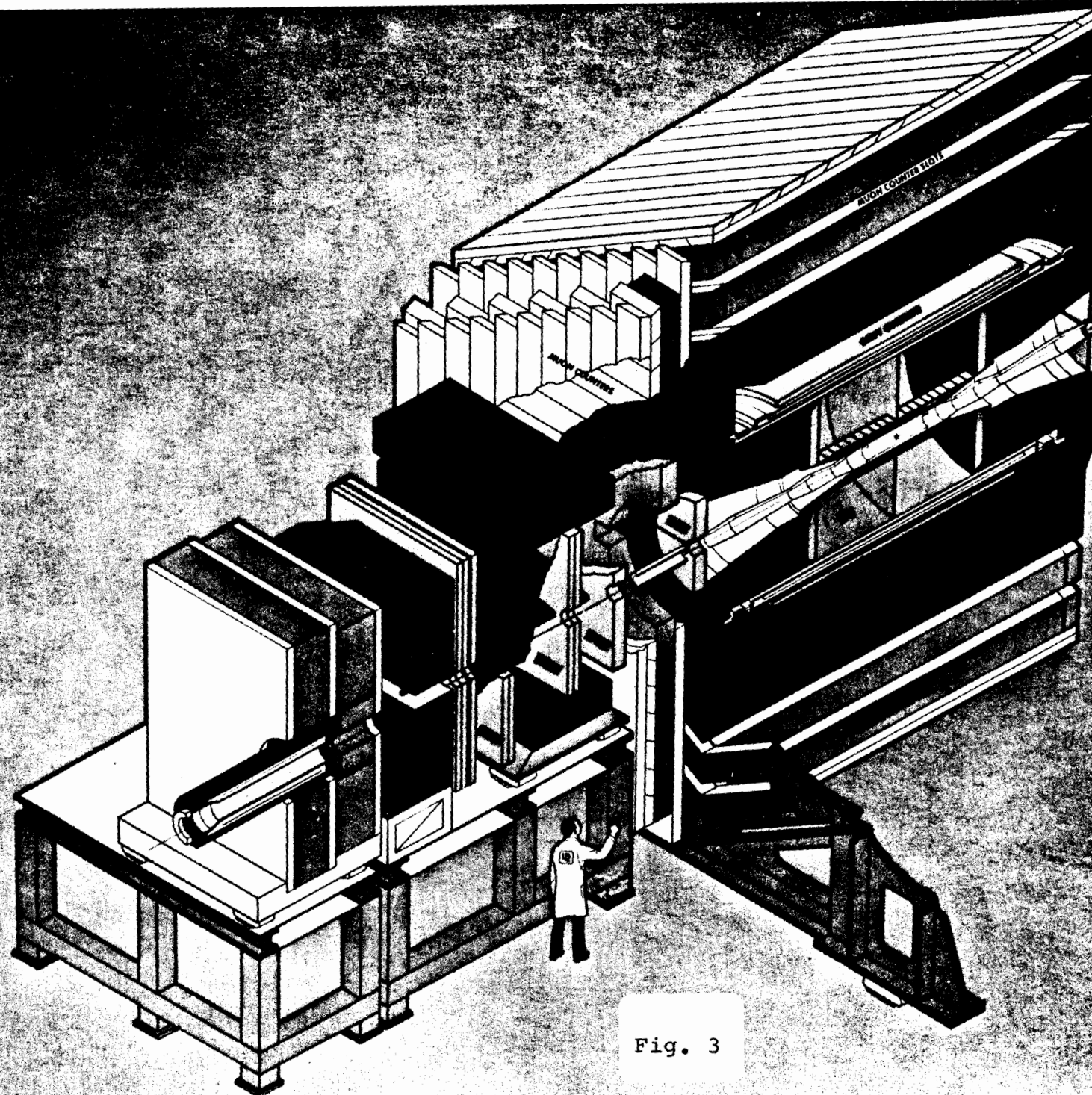


Fig. 3

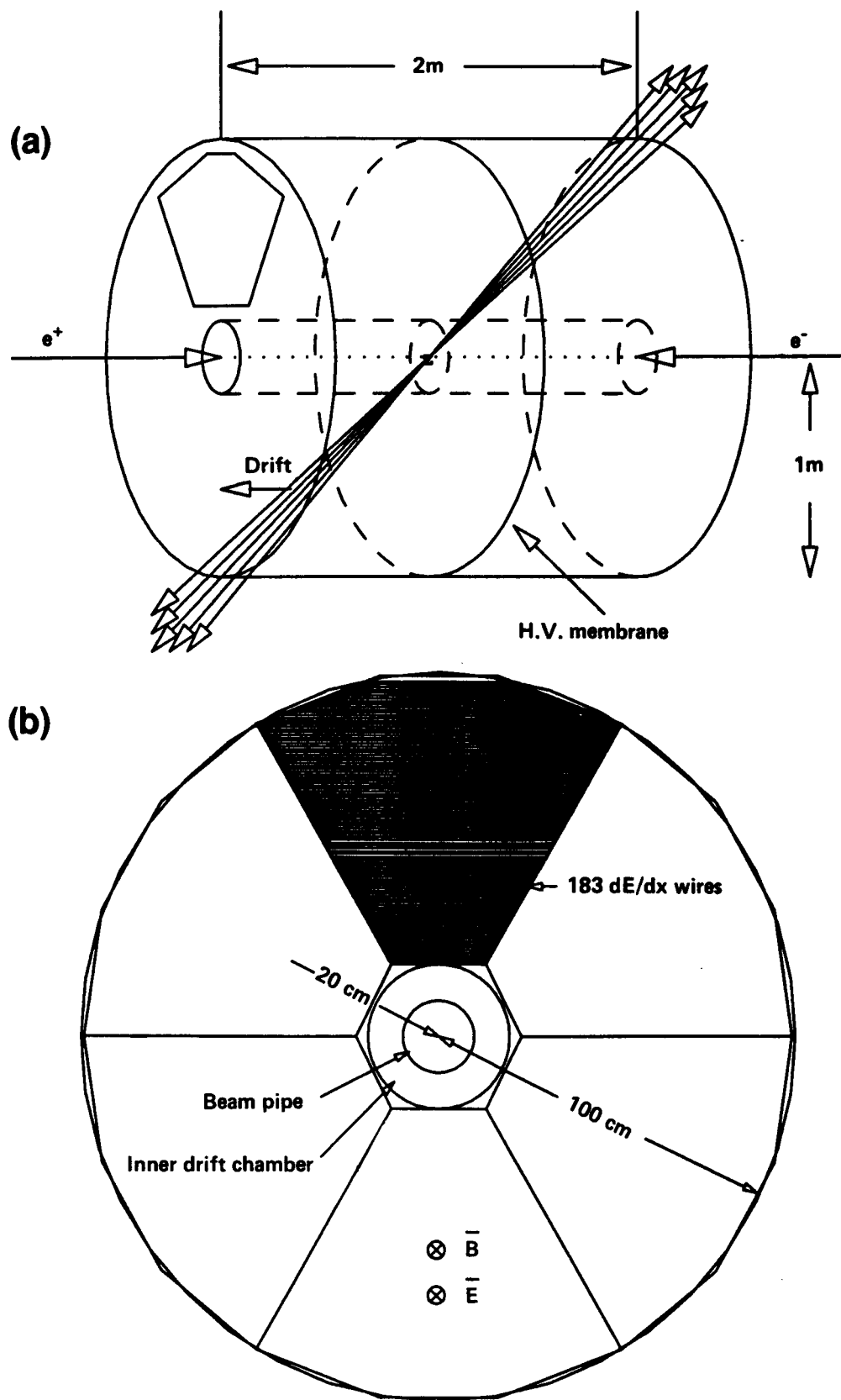


Fig. 4

PEP-9 septum magnets

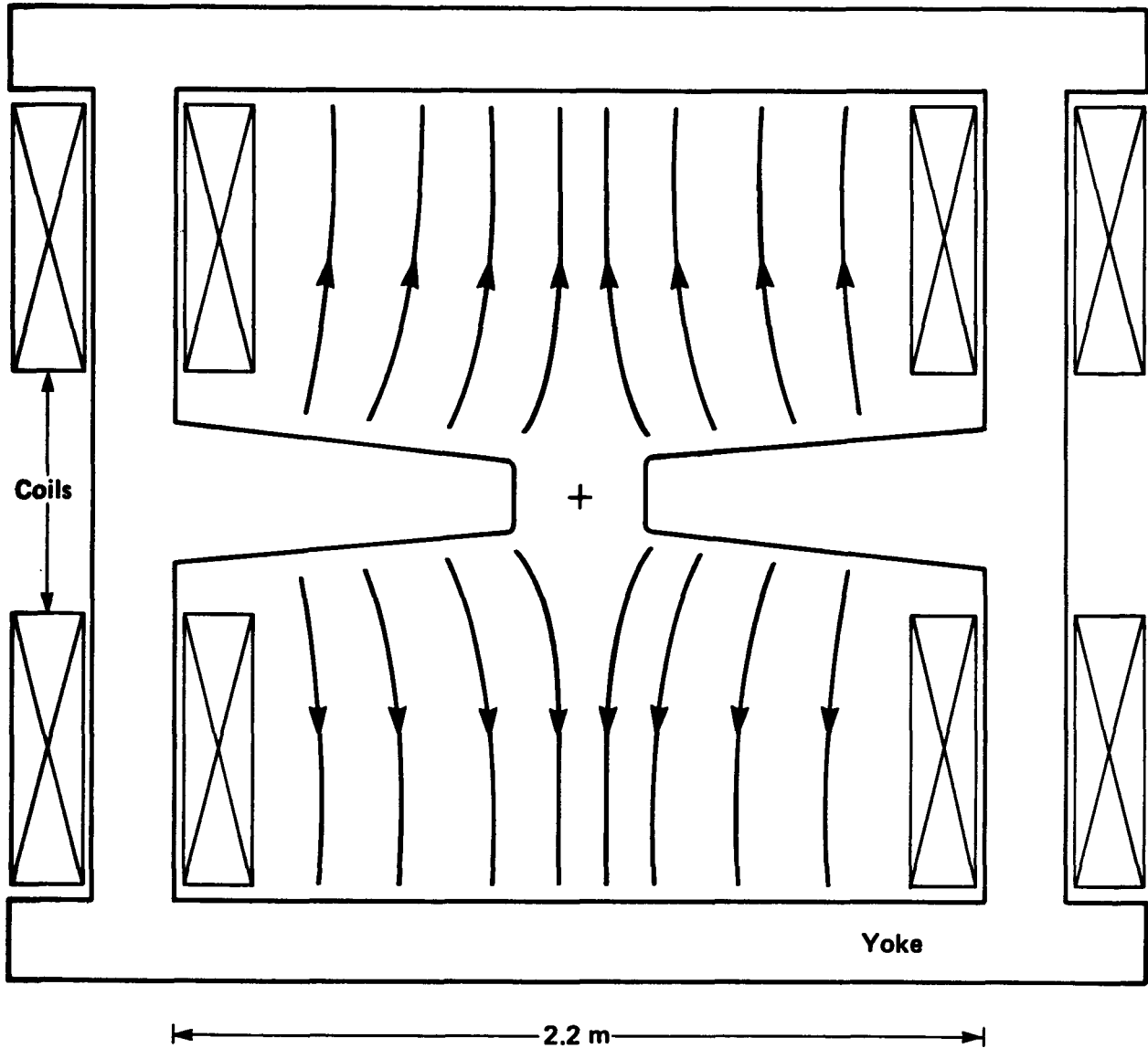


Fig. 5

LLNL GAS INJECTION SYSTEM (PEP IR-2)

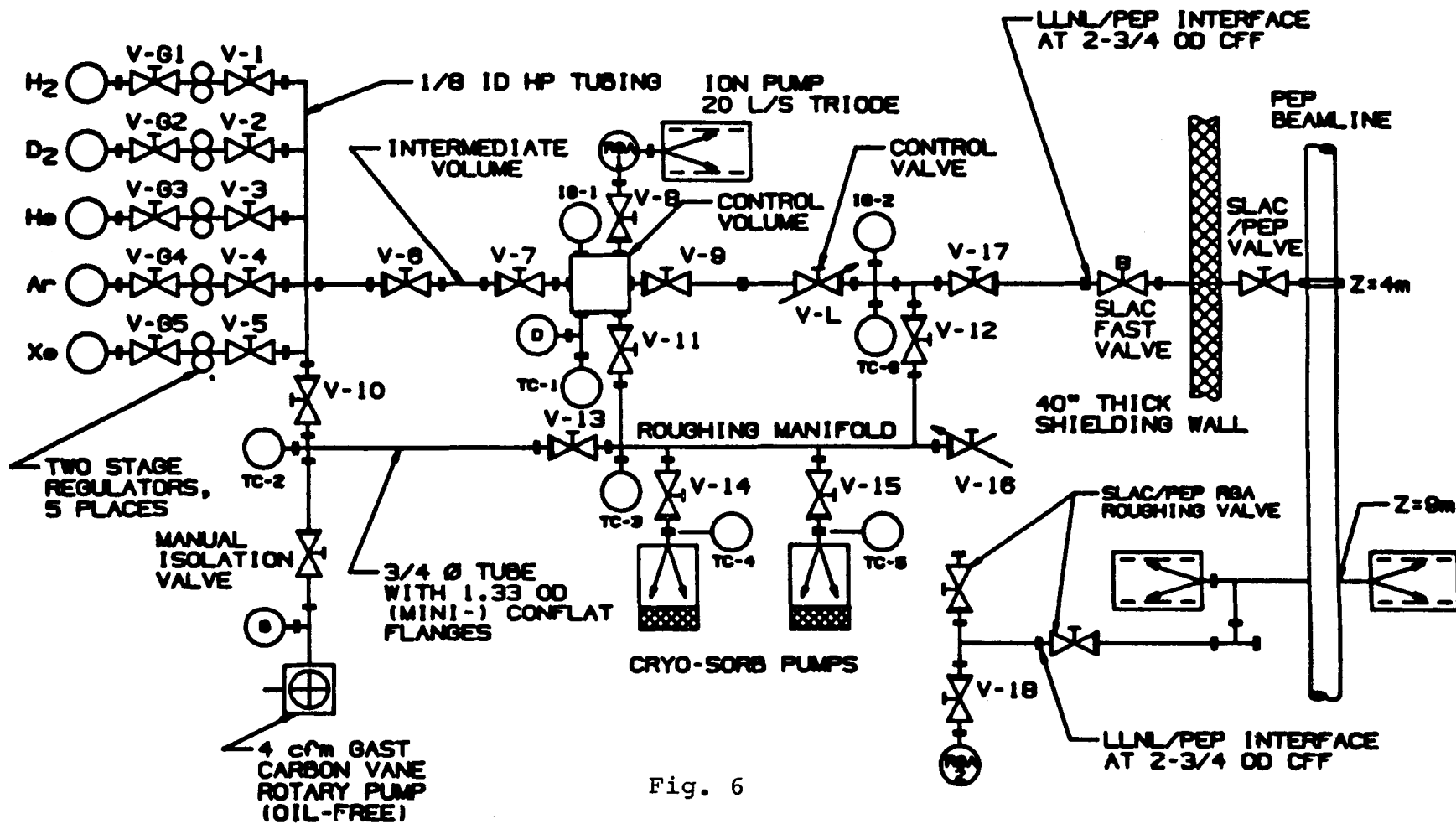


Fig. 6

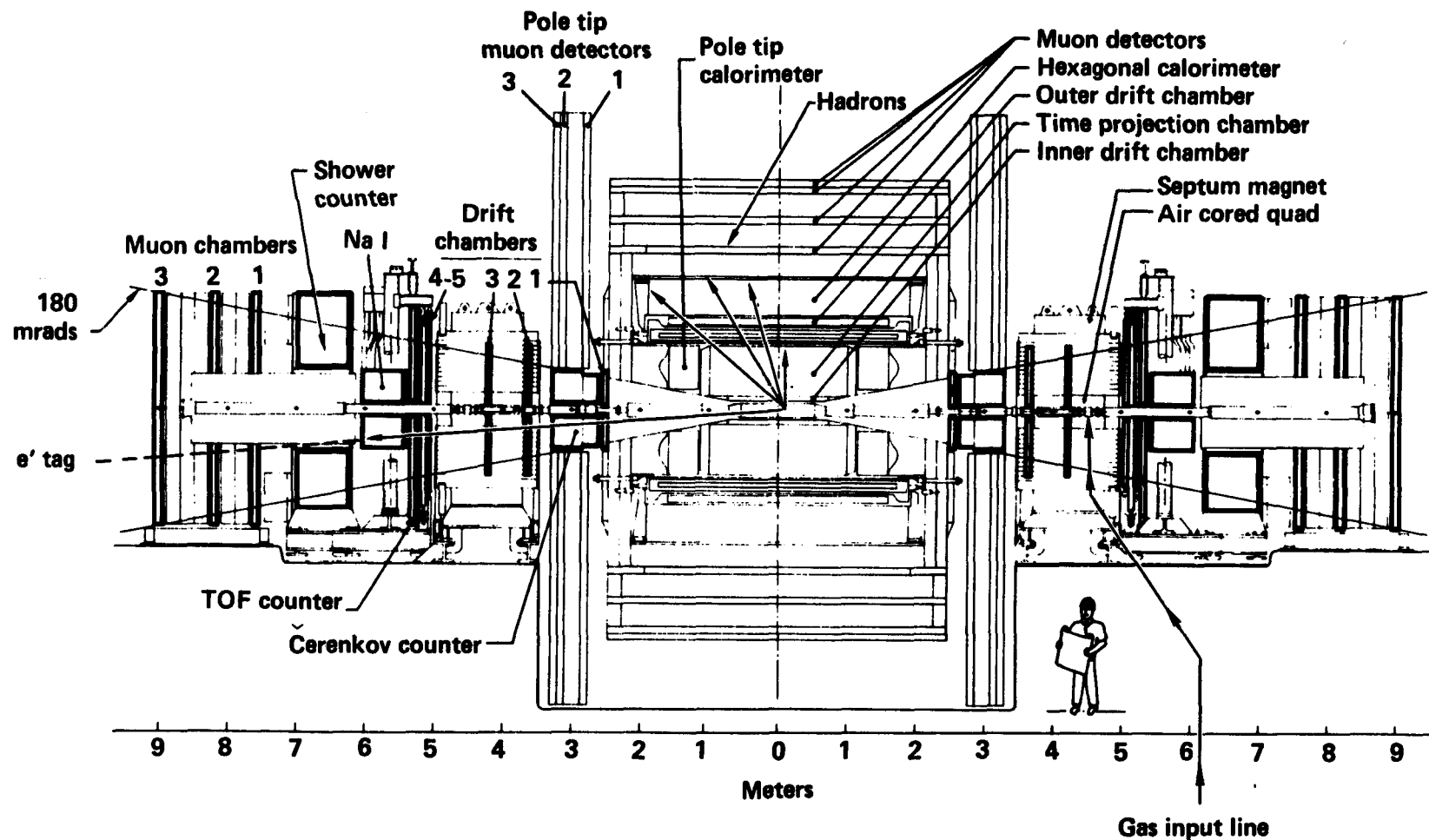


Fig. 7

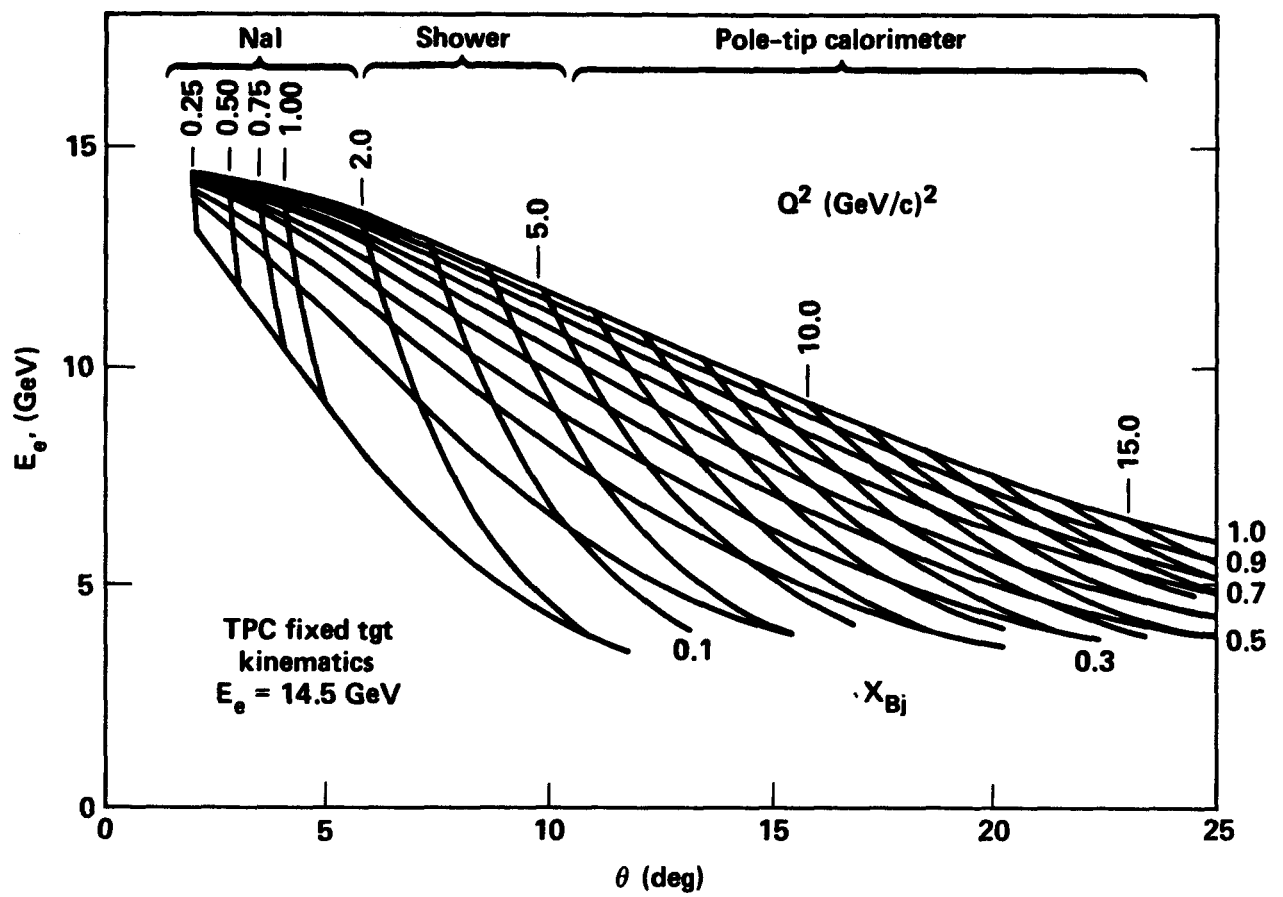


Fig. 8

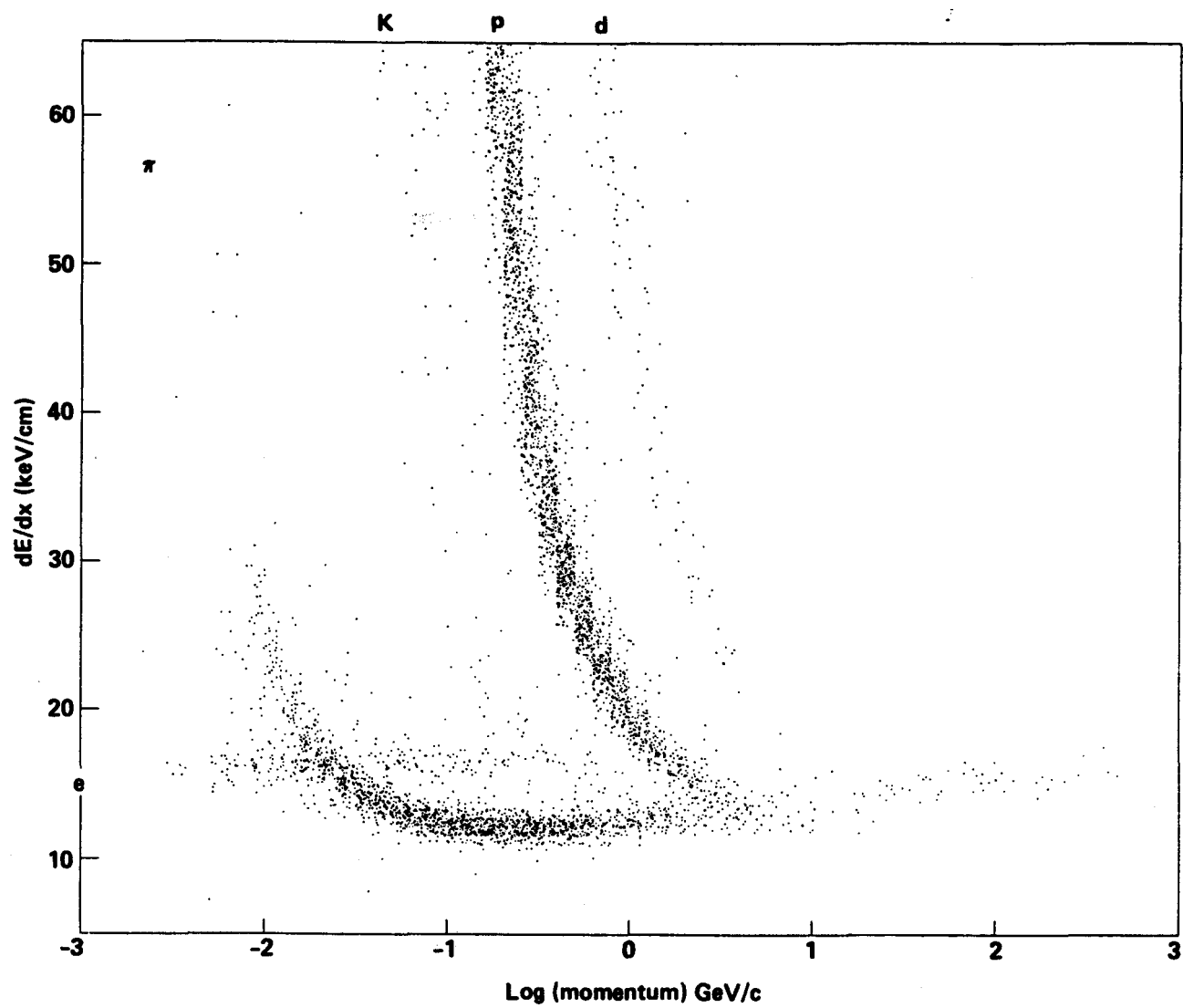


Fig. 9

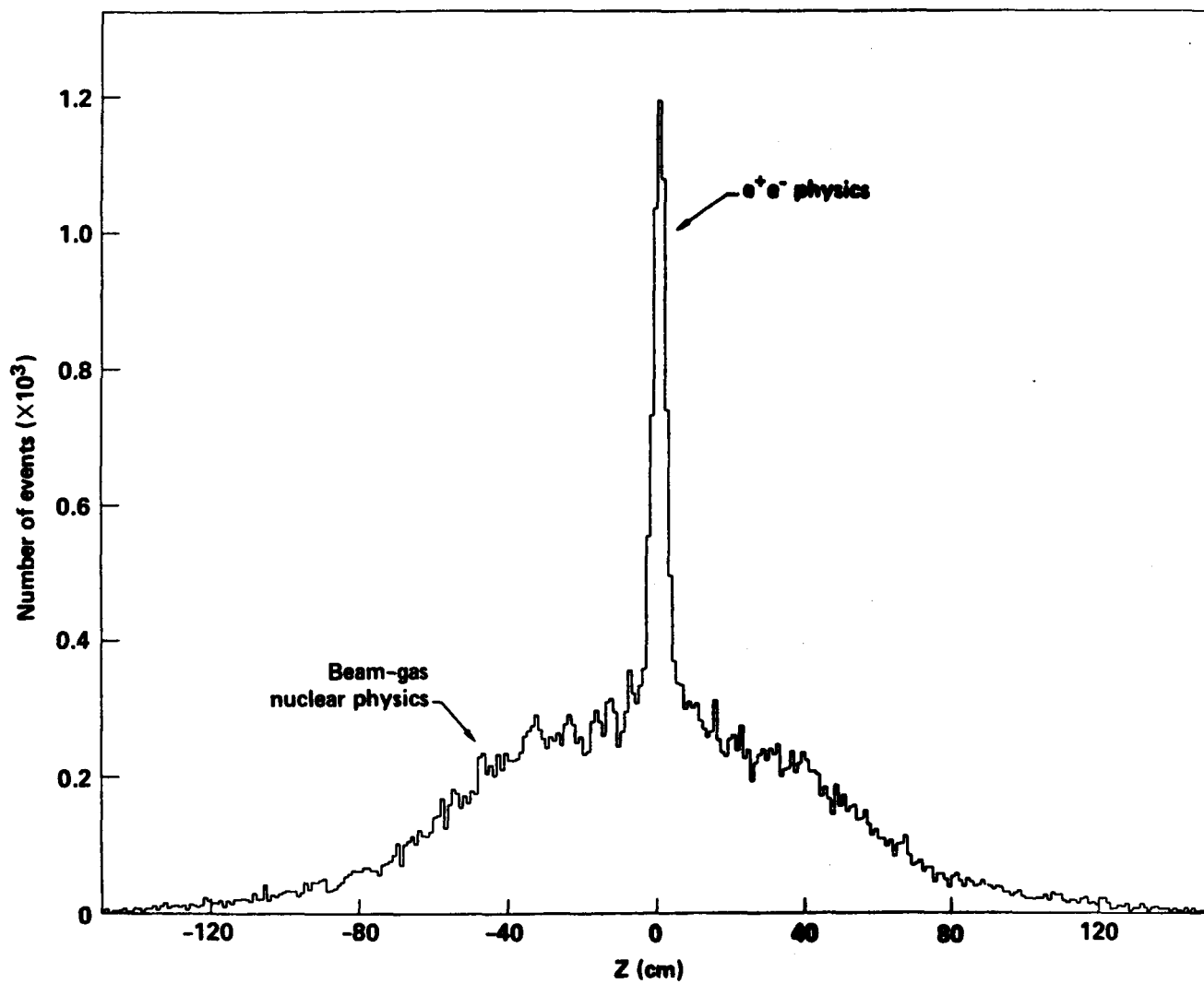


Fig. 10

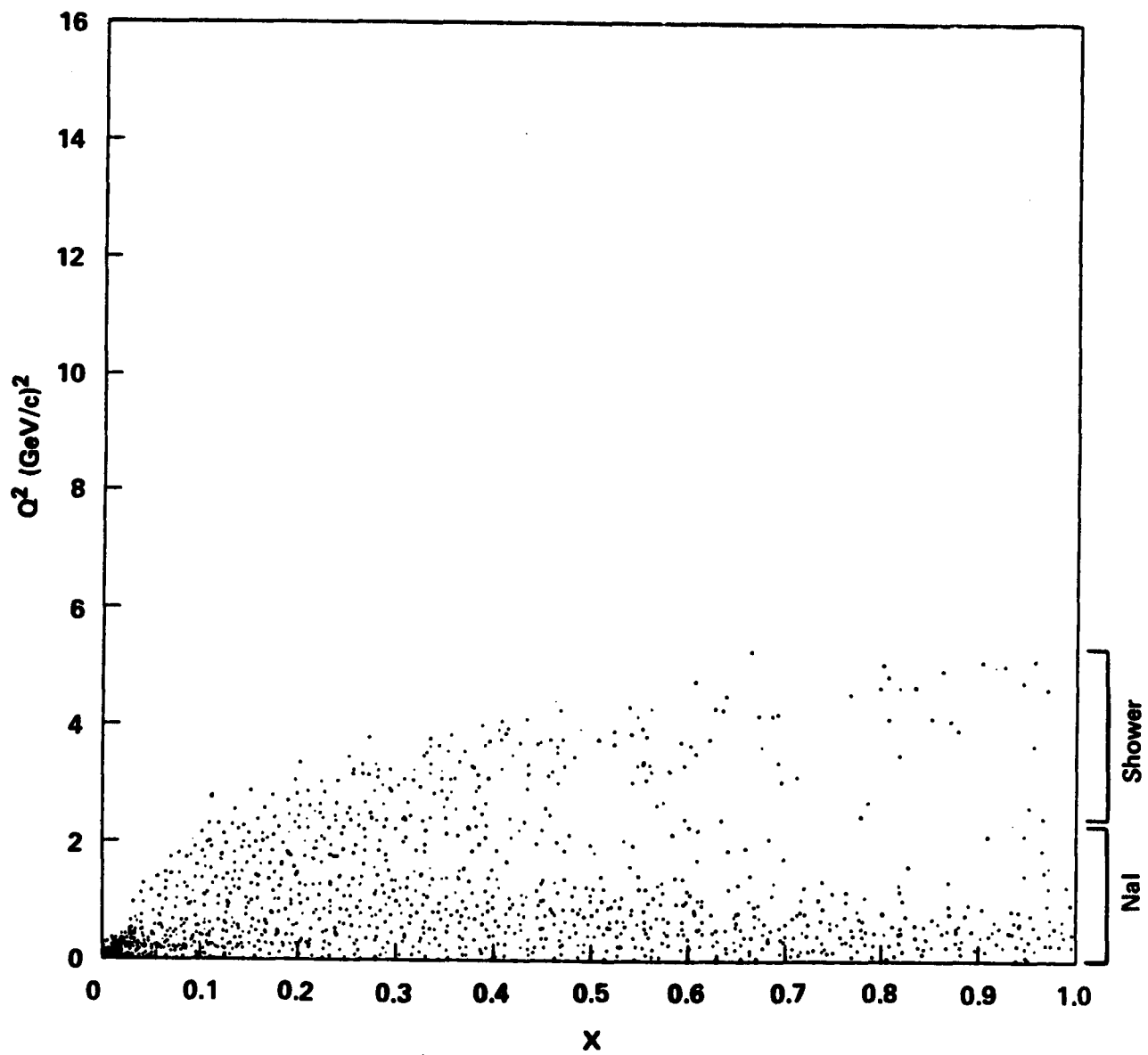


Fig. 11

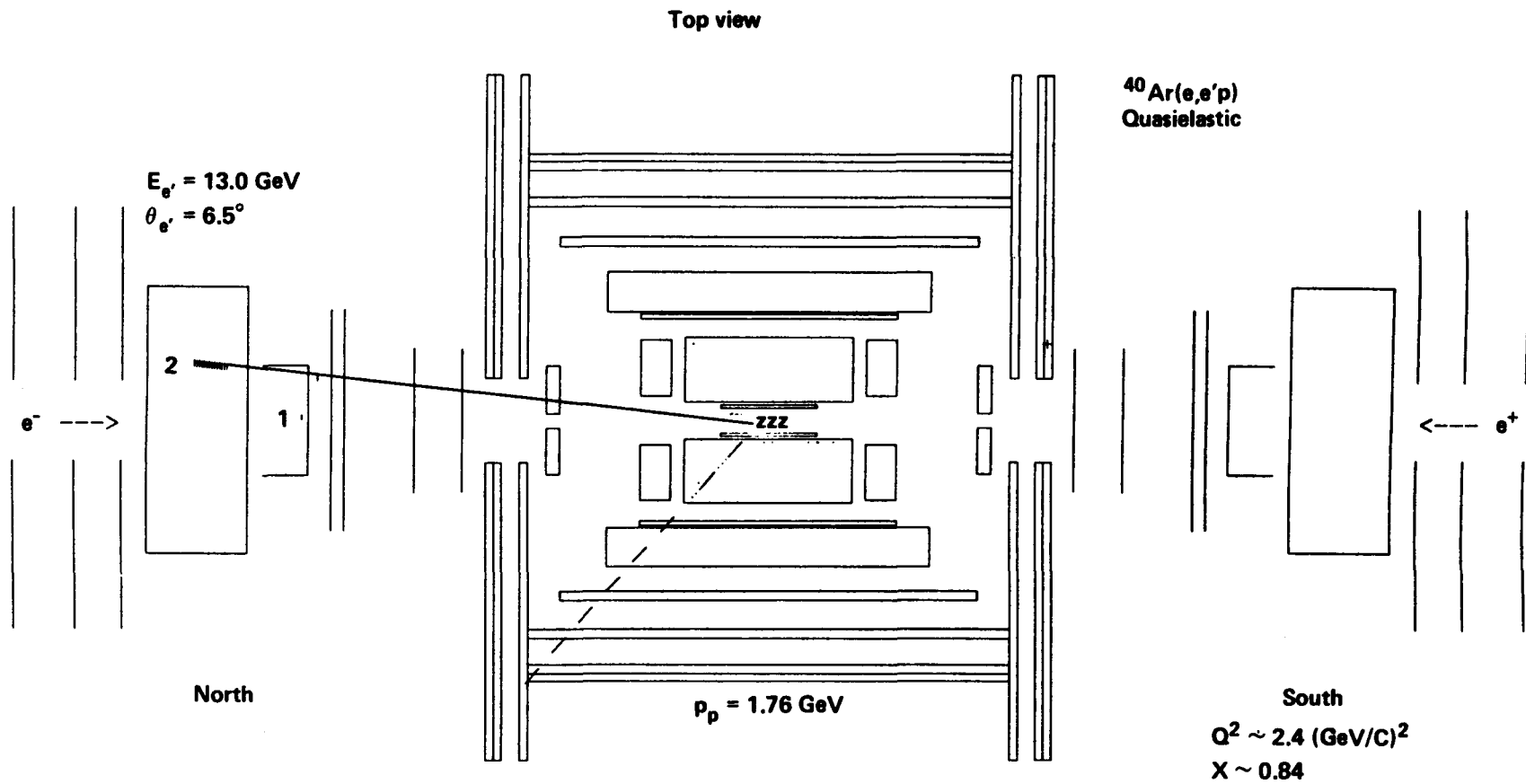


Fig. 12

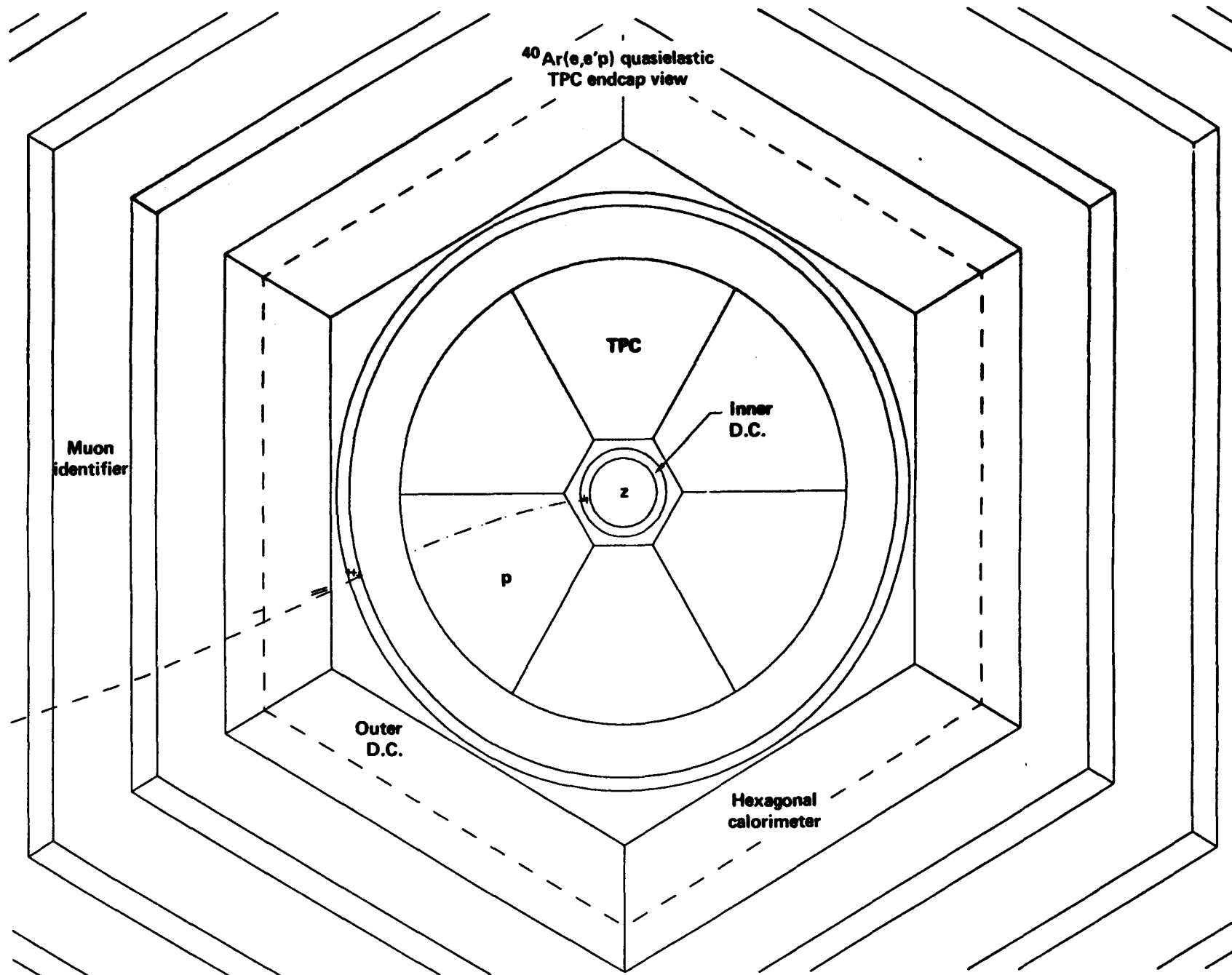


Fig. 13

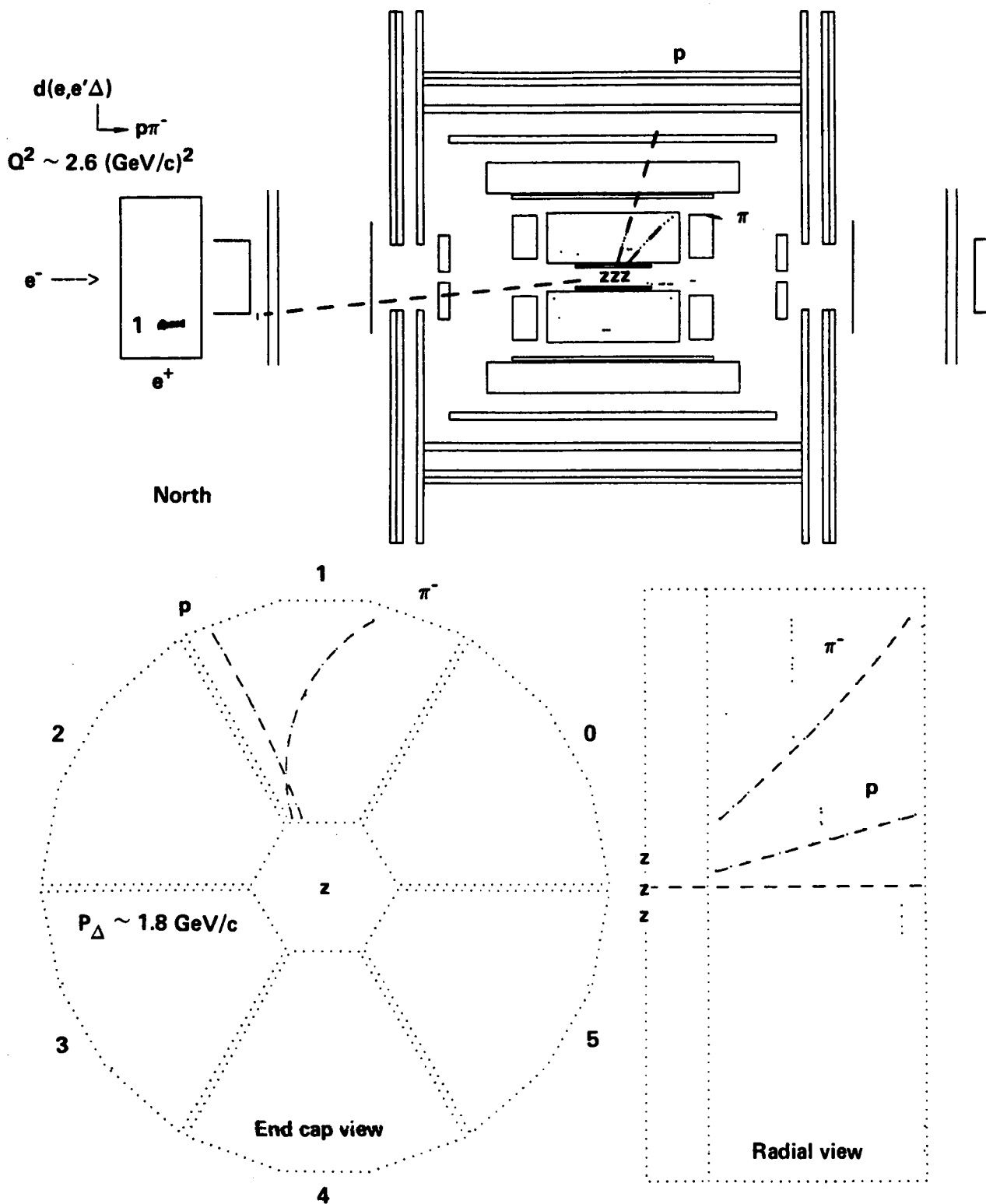


Fig. 14

Associated production
 $^{40}\text{Ar}(e, e'K^+\Lambda)$

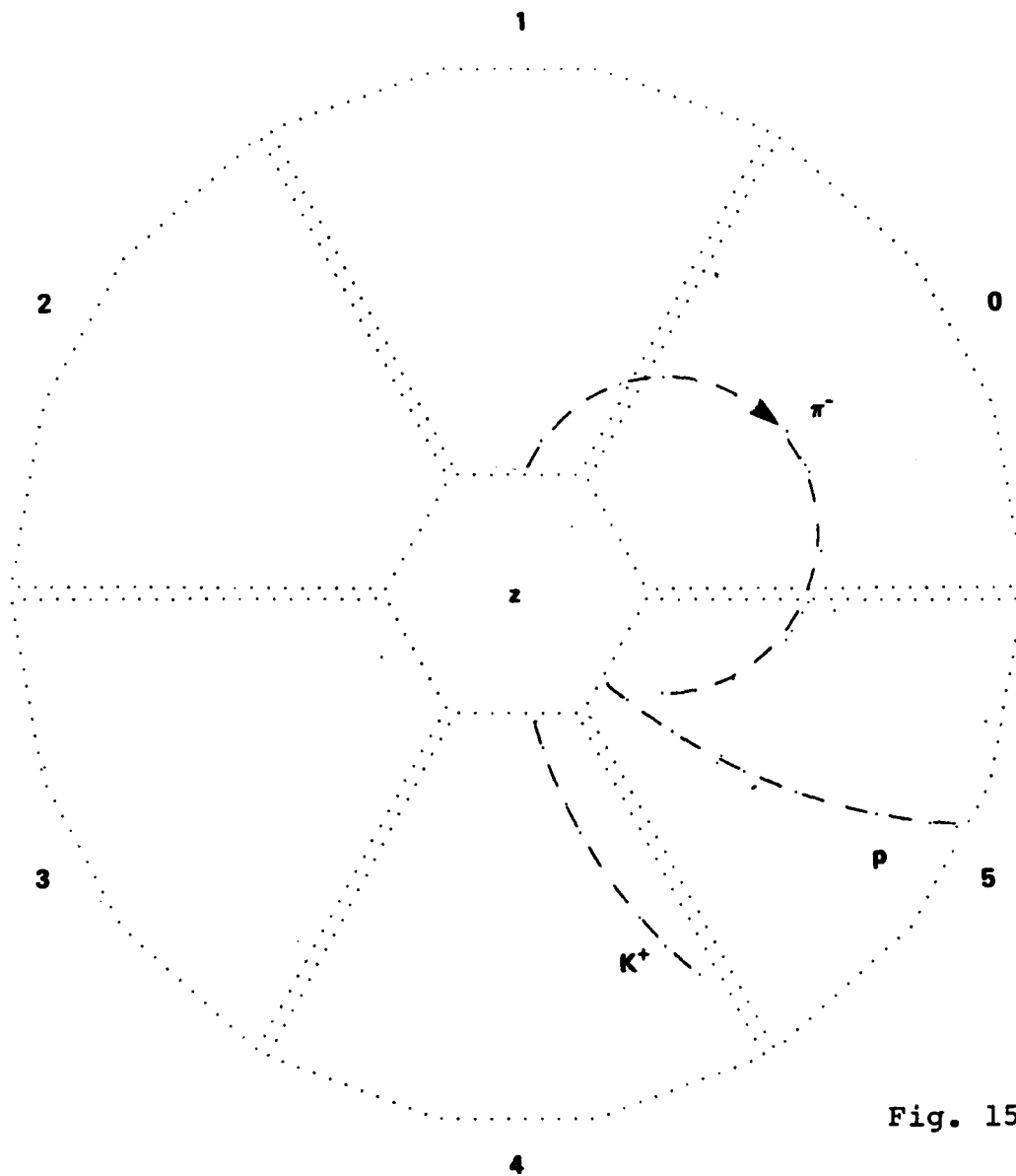
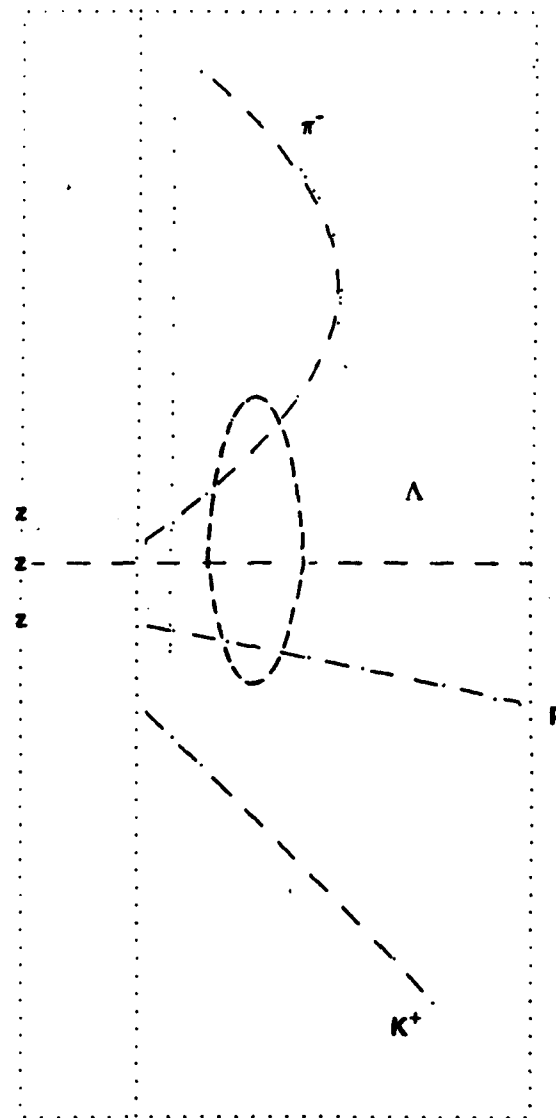


Fig. 15



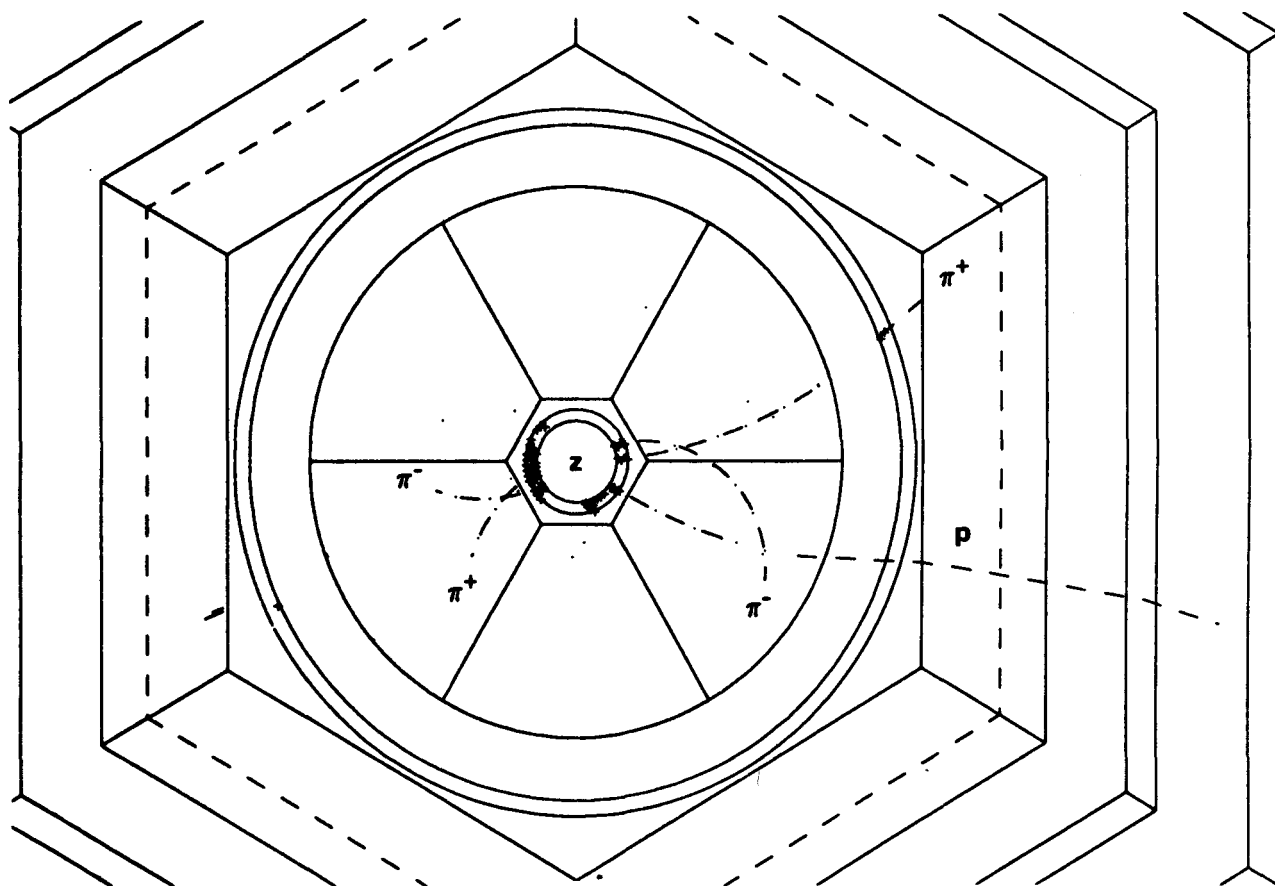
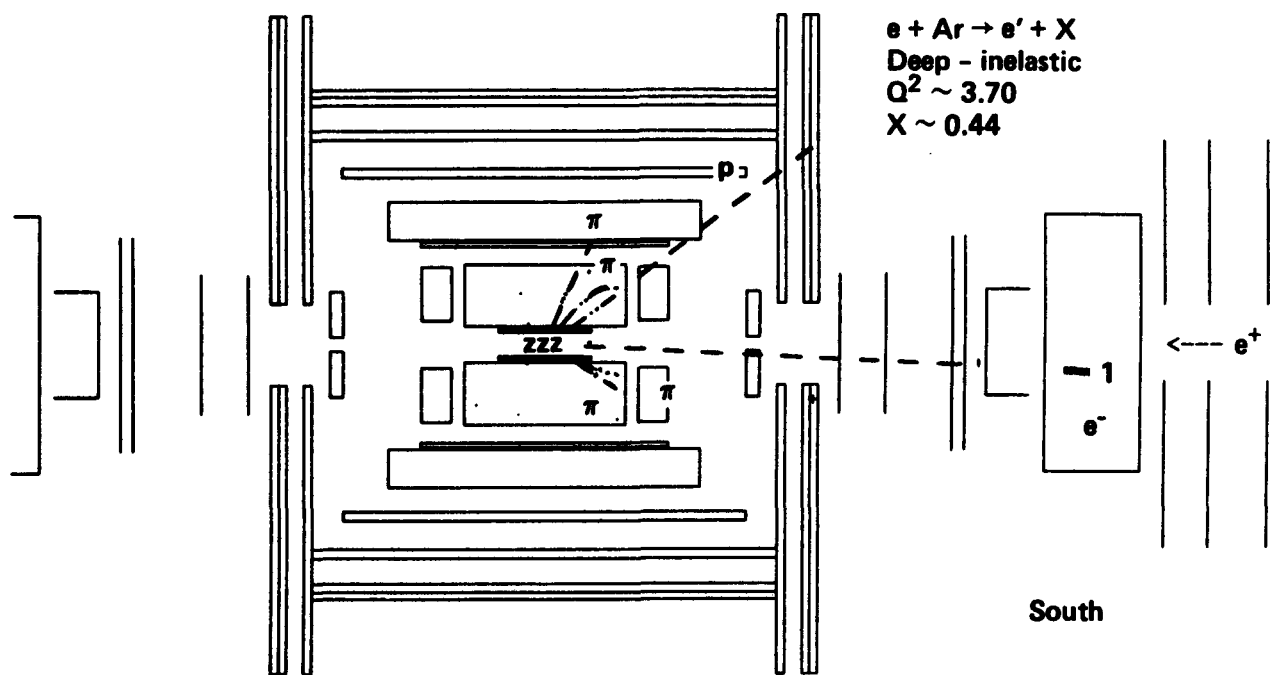


Fig. 16

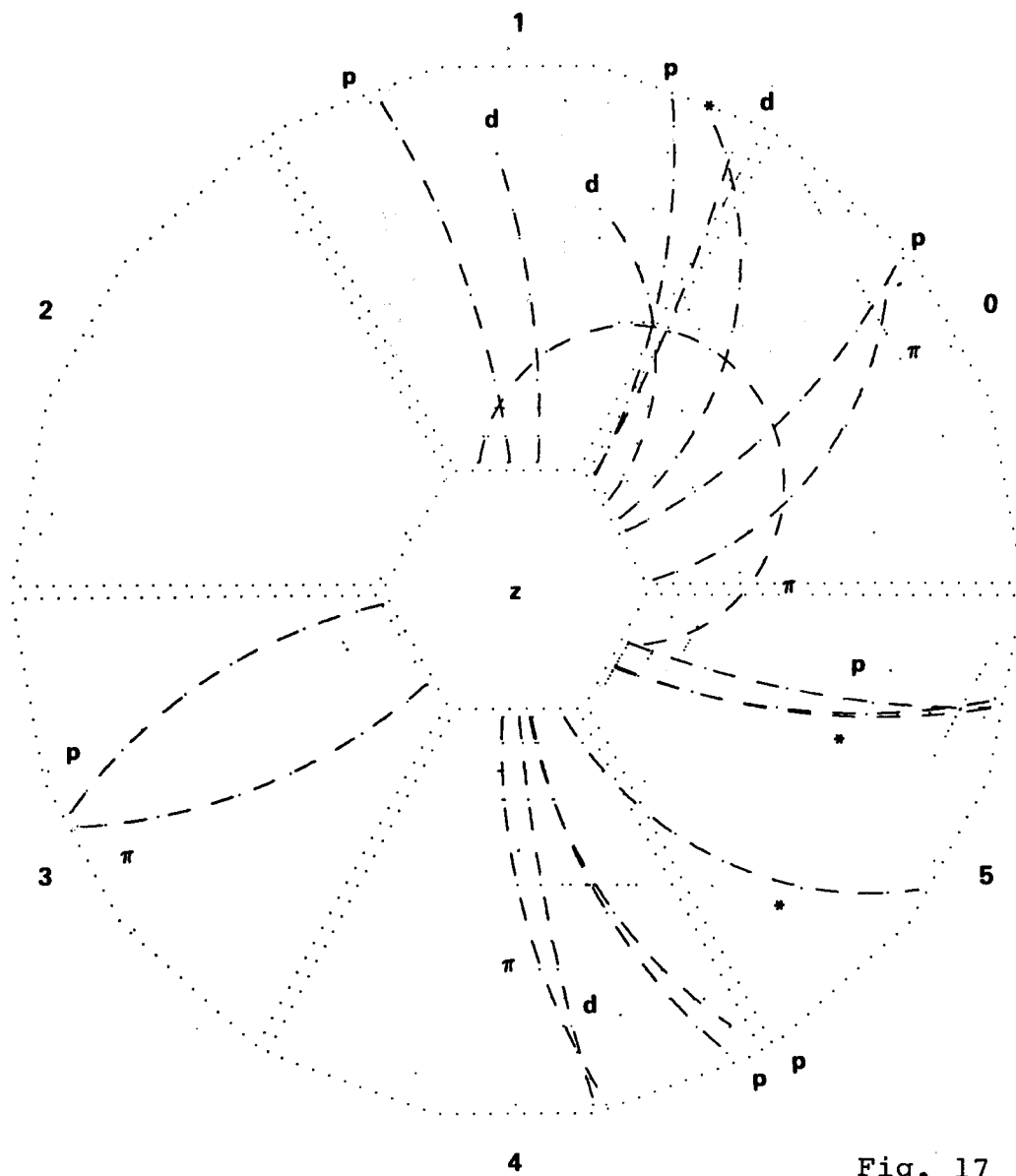
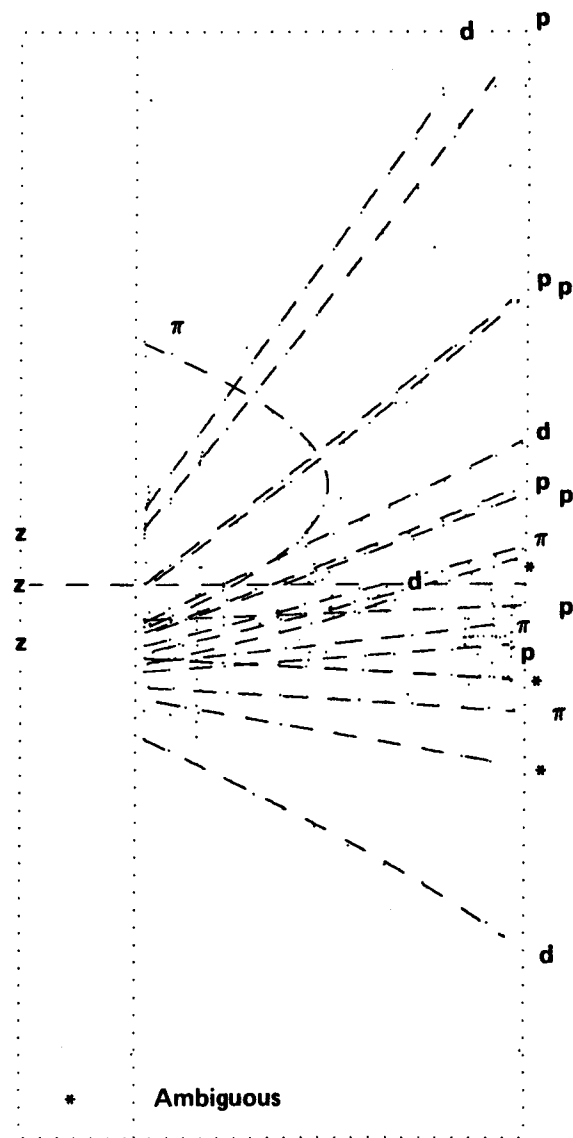


Fig. 17



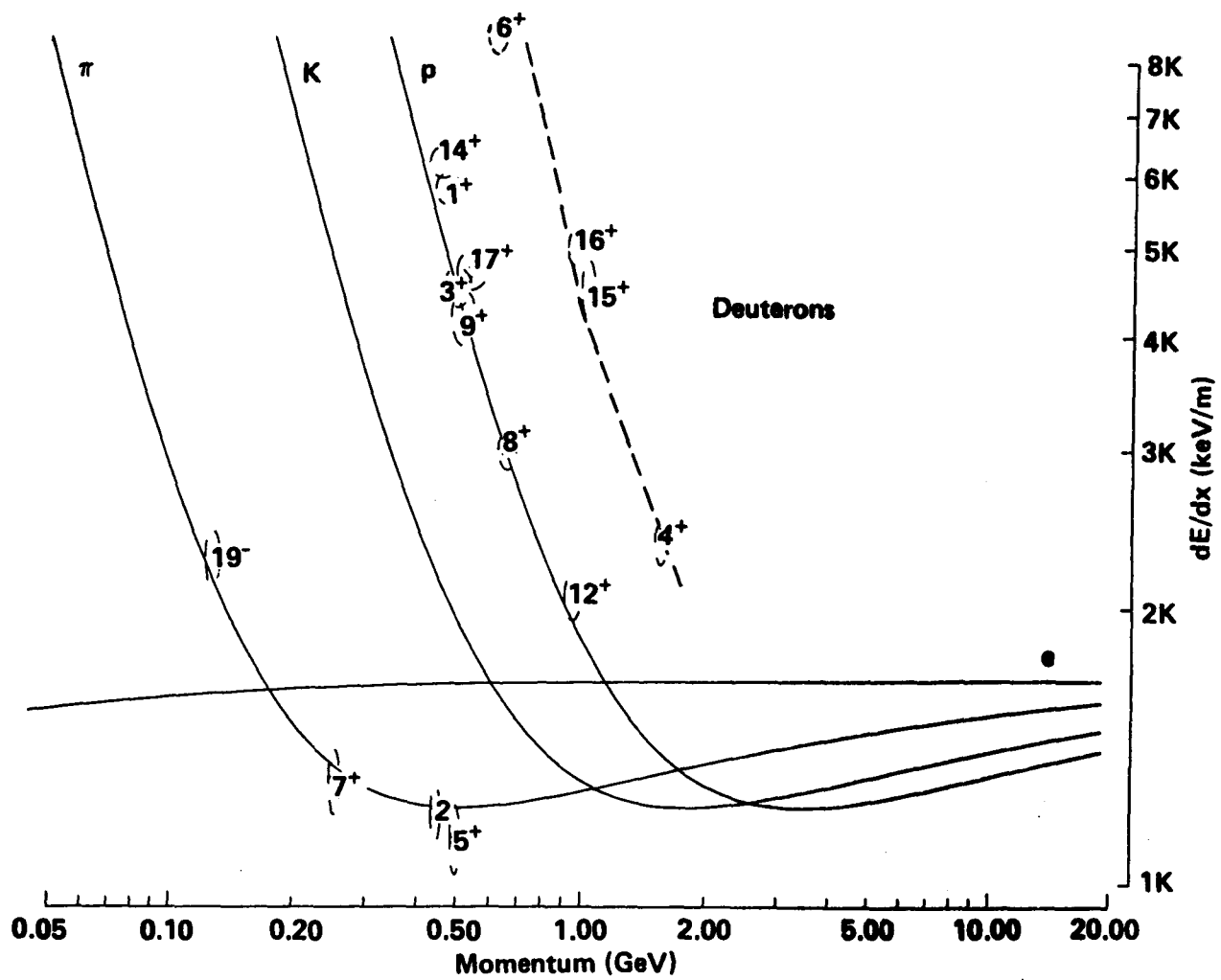


Fig. 18

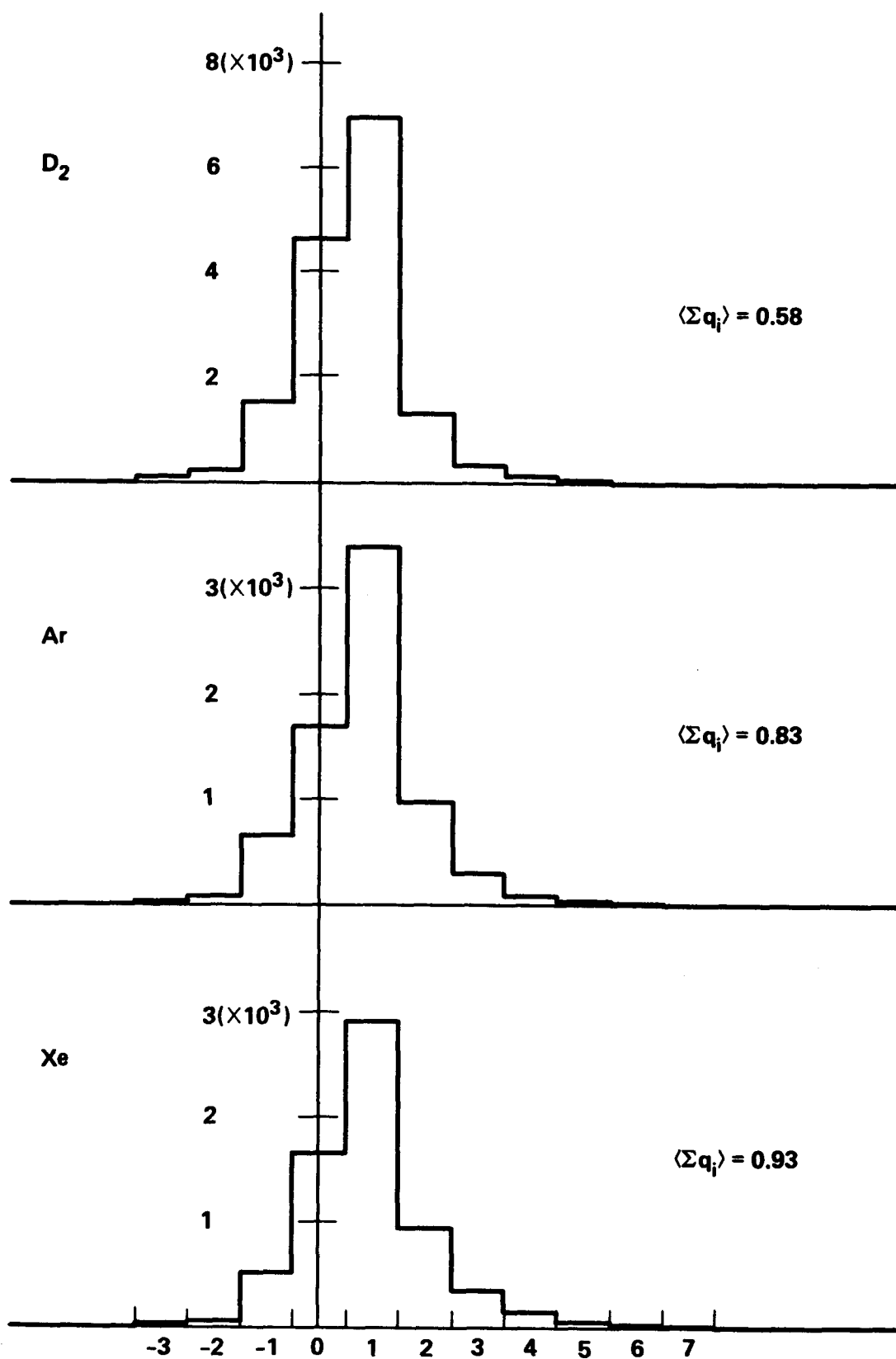
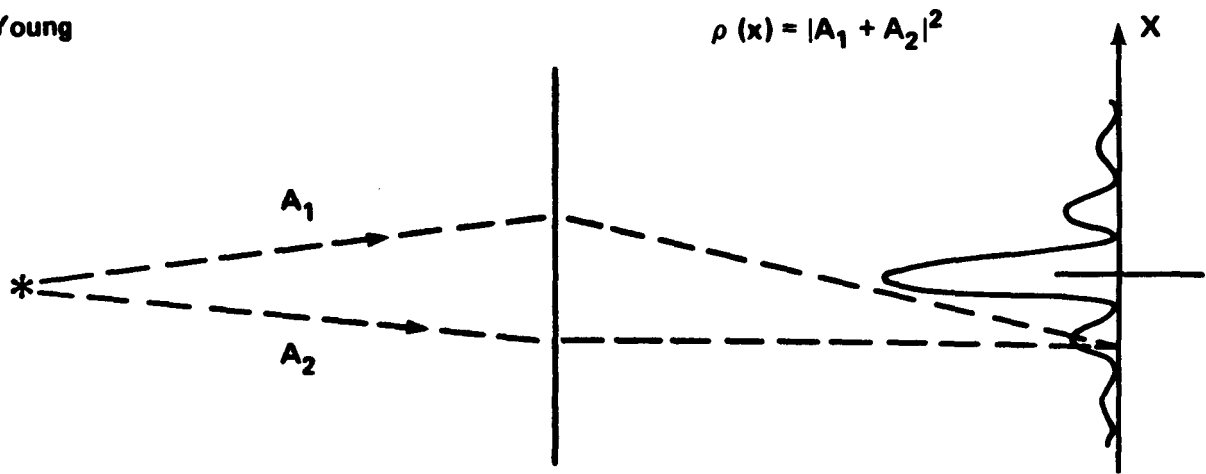
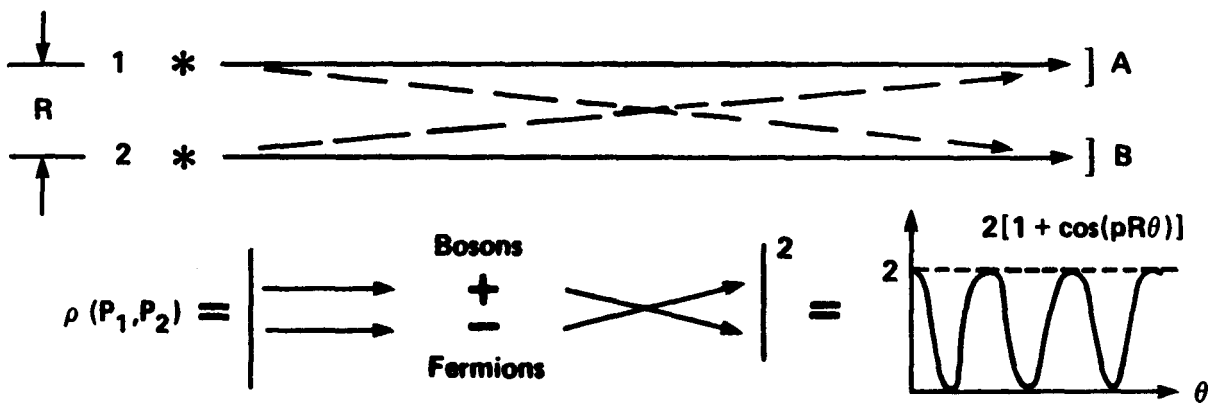


Fig. 19

(a) Young



(b) Prototype Hanbury-Brown/Twiss



(c) H-B/T extended space-time source

e.g. $S(r, t) \propto \exp(-r^2/2\sigma^2) e^{-\nu\tau}$

$$\rightarrow R \equiv \frac{\rho_{B-E}(p_1, p_2)}{\rho_{UNCORR}(p_1, p_2)} = 1 + \exp(-q^2 r^2) / (1 + (q_0 \tau)^2)$$

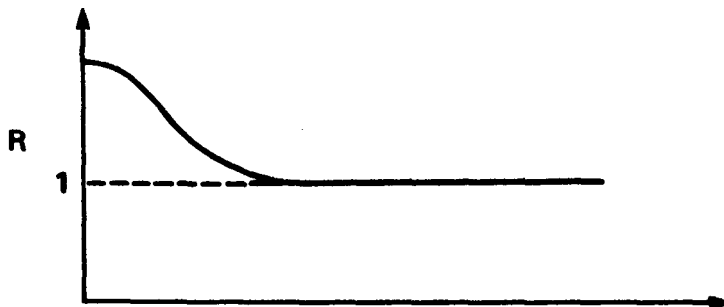


Fig. 20

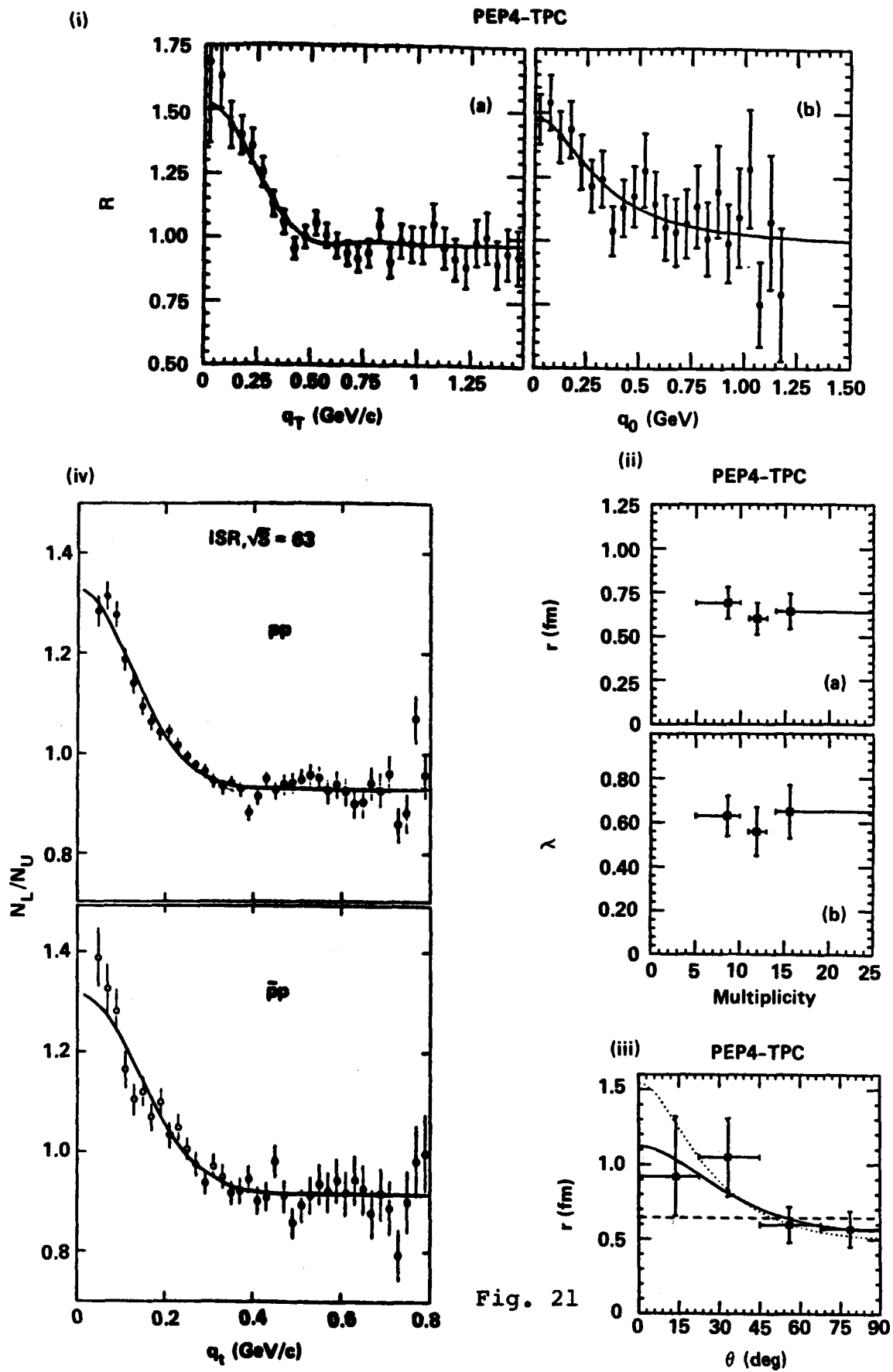


Fig. 21

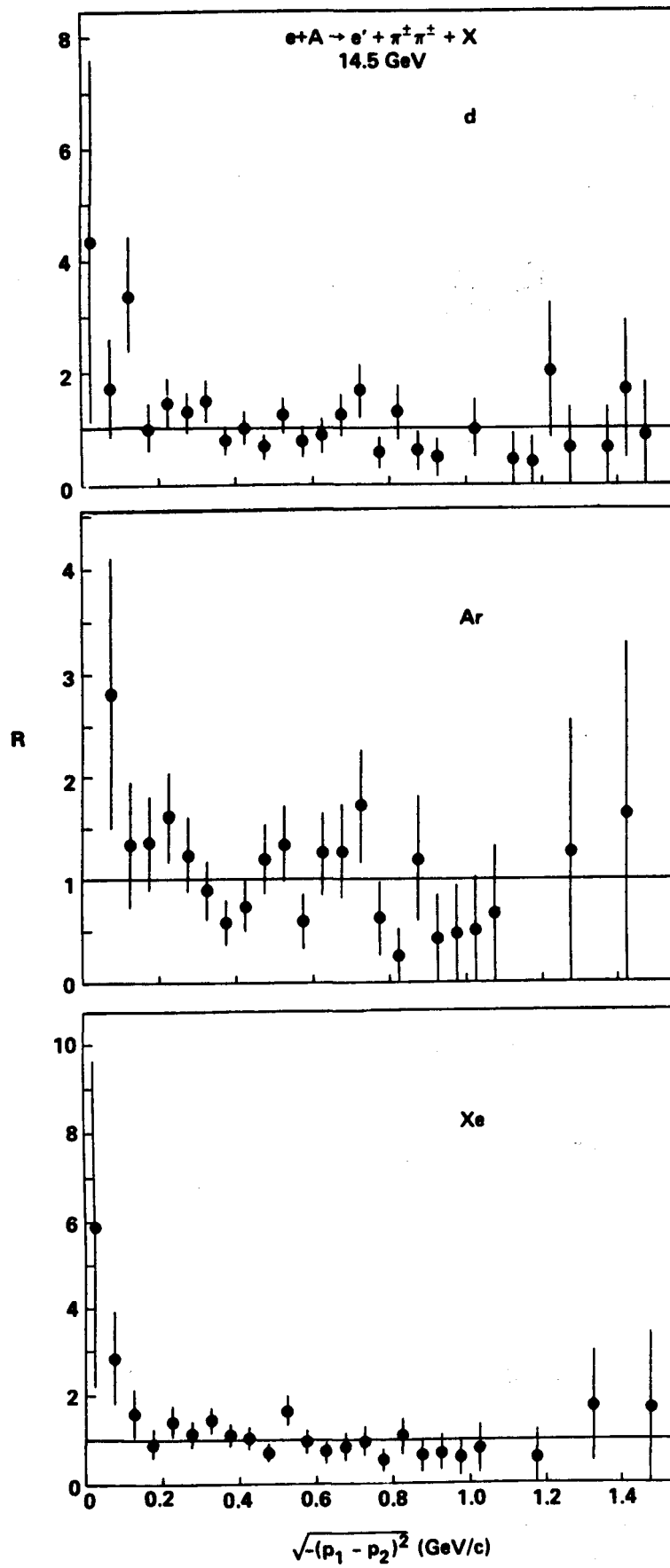


Fig. 22

Fig. 23

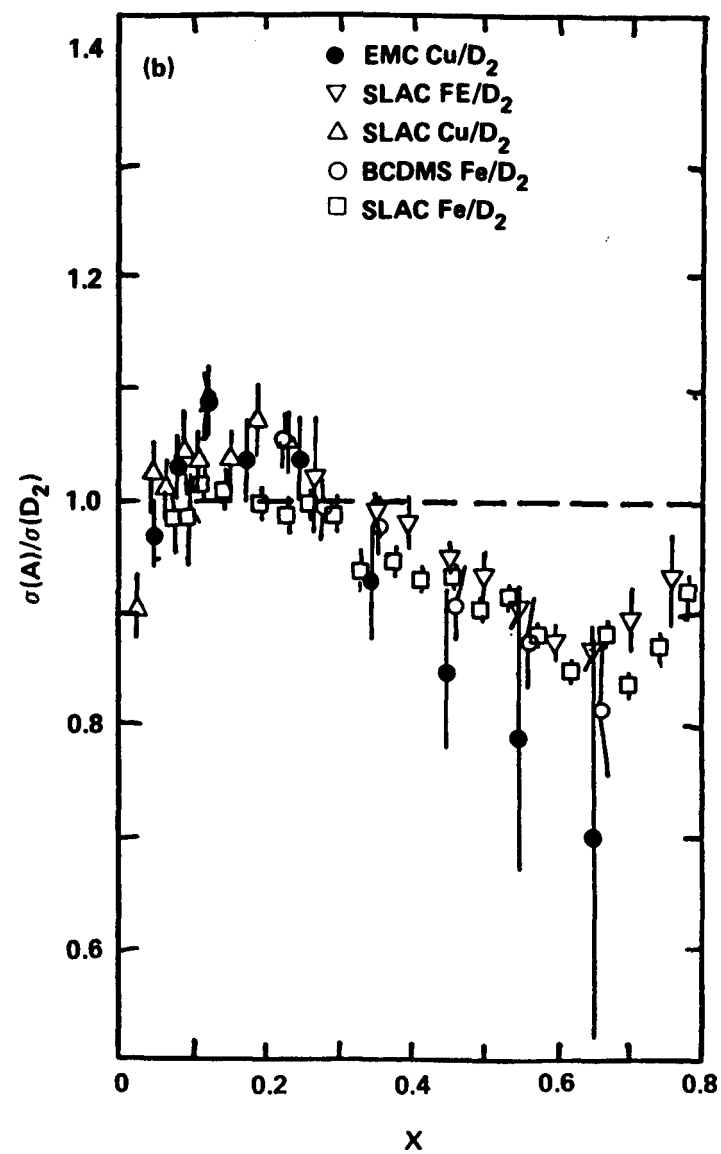
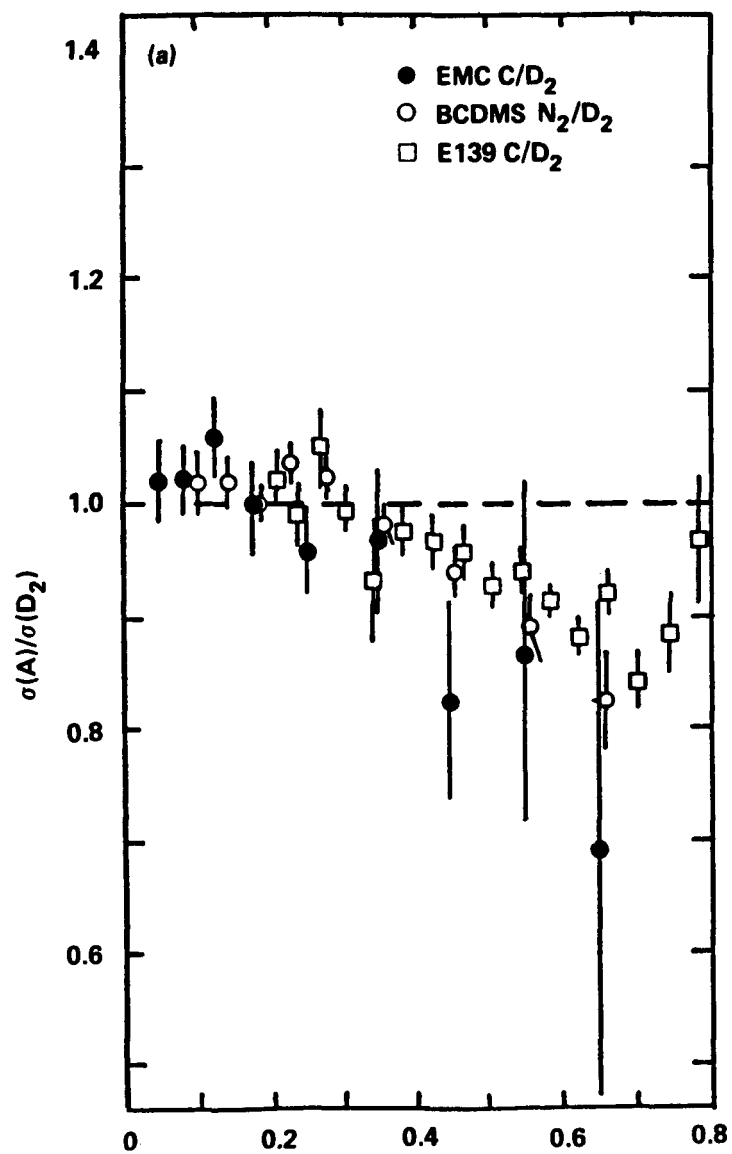
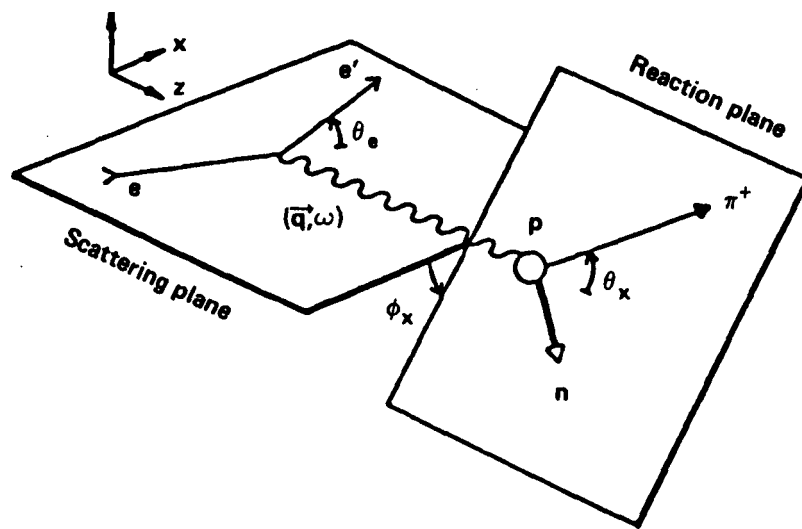


Fig. 24



$$2\pi\sigma(t, \phi) = \sigma_T + \epsilon\sigma_L + \epsilon \cos(2\phi)\sigma_p + \sqrt{2\epsilon(\epsilon+1)} \cos\phi \sigma_l$$

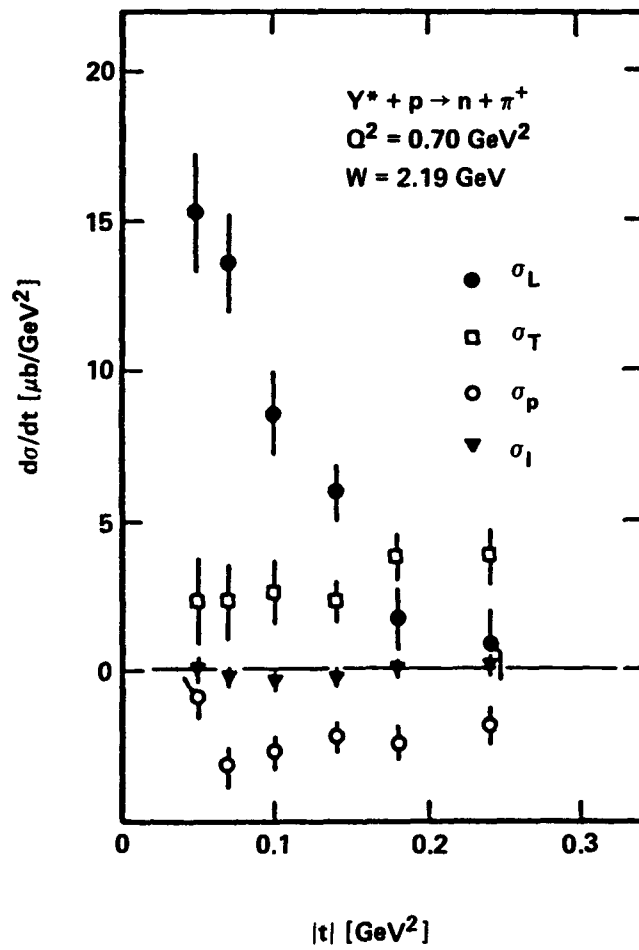


Fig. 25

NUCLEAR MEDIUM EFFECTS ON THE NUCLEON PROPERTIES

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The idea that the nuclear medium may modify the properties of the nucleon is not really a new one, at least for nuclear physicists. There has been for many years speculations on the enhancement of the pion field in nuclei as compared to the field surrounding free nucleons. More recently the idea of multiquark clusters as a possible explanation to various nuclear phenomena has been proposed by a number of authors. But until these last few years the intrinsic properties of the nucleon such as its mass, baryonic or electromagnetic size, magnetic moments, were hardly suspected to be substantially modified by the nuclear environment. The question raised essentially from two categories of experimental investigations, in quite different kinematical regimes and it remains at the present time a controversial subject.

In this paper I try to review the present status of experimental investigations and some of the recent theoretical developments relevant to this topic.

EXPERIMENTAL INVESTIGATIONS

1. Deep inelastic scattering of leptons from nuclear targets

An extensive amount of data is now available on scattering from various nuclei of:

- muons in the four-momentum range $q^2 = 9 - 200 \text{ GeV}^2$ by the EMC collaboration and, more recently by other groups (BCDMS, NMC)¹.
- electrons in the four-momentum range $q^2 = 2 - 5 \text{ GeV}^2$ at SLAC²

These measurements have shown that the structure function

$F_2(x, q^2)$ of the nucleon in a nucleus differs from that of the free nucleon: its behaviour with the Björken variable $x = q^2/2M\omega$ (ω = energy transfer) was interpreted as a shift of the quark momentum distribution towards small x (e.g. low momenta) in nuclei. Moreover the effect was found increasing with the number of nucleons A . Numerous theoretical interpretations of this observation have been proposed. One underlying idea in many of them is that the correlation length between quarks increases with the nuclear density. A more restrictive formulation of this idea was the hypothesis that the size of a nucleon embedded in a nucleus is actually increased: typically the nucleon would be "swollen" by 15% in iron to account for its modified structure function.

2. Quasi-elastic electron scattering from nuclei.

In the kinematical region $\omega \approx q^2/2M_N$ the nucleus is essentially seen as a Fermi gas of quasi-free nucleons. This description proved to be fairly successful to account for the cross-sections as a function of ω : they appear as a broad peak centered at $\omega = q^2/2M_N + E_1$ (E_1 = average binding energy of the nucleons). The width of that peak is proportionnal to the Fermi momentum k_F of the nucleus.

2.1. Inclusive low q^2 measurements with σ_L / σ_T separation.

A significant achievement in the eighties was the experimental separation of the response functions R_T and R_L to transverse and longitudinal virtual photons in inclusive (e, e') measurements up to $\approx .3 \text{ GeV}^2$. Various nuclei: ^3He , ^{12}C , ^{40}Ca , ^{48}Ca , ^{56}Fe , ^{238}U have been studied at Saclay and Bates³. Except in the case of ^3He no model is presently able to reproduce both of these response functions. In conventionnal calculations⁴ (in Schrödinger formalism) the longitudinal part is overestimated by 20 to 40 %. As a consequence the Coulomb Sum Rule (CSR) is not saturated even at the highest momenta reached in these experiments (fig 1). Collective effects are usually treated in the RPA framework: in recent semi-classical calculations⁵ the detailed shape of the response as a function of ω

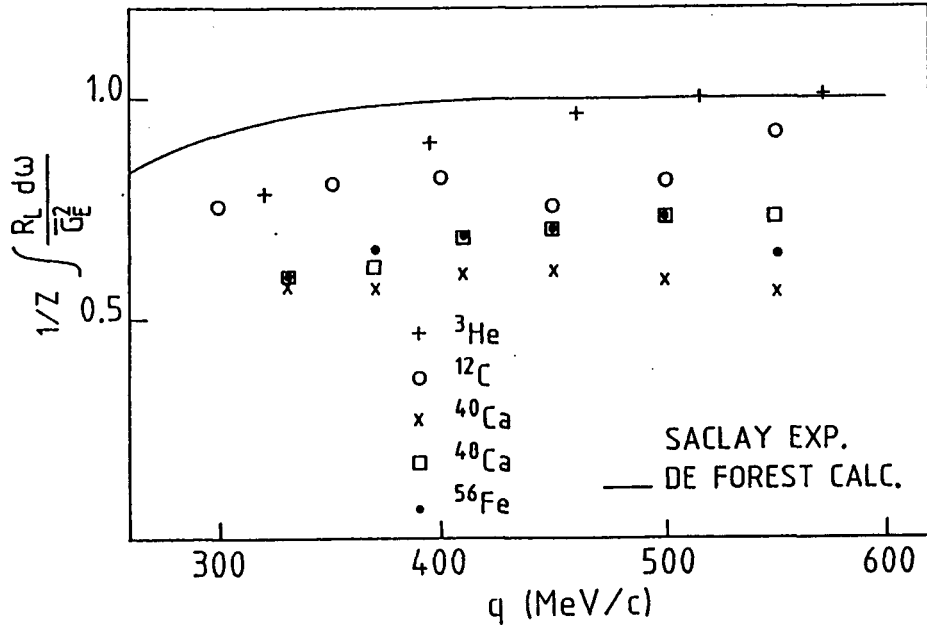


Fig. 1a. Coulomb Sum Rule normalized to unit charge. Data are from the Saclay experiments (Ref. 3). The calculation is by De Forest (Ref 50).

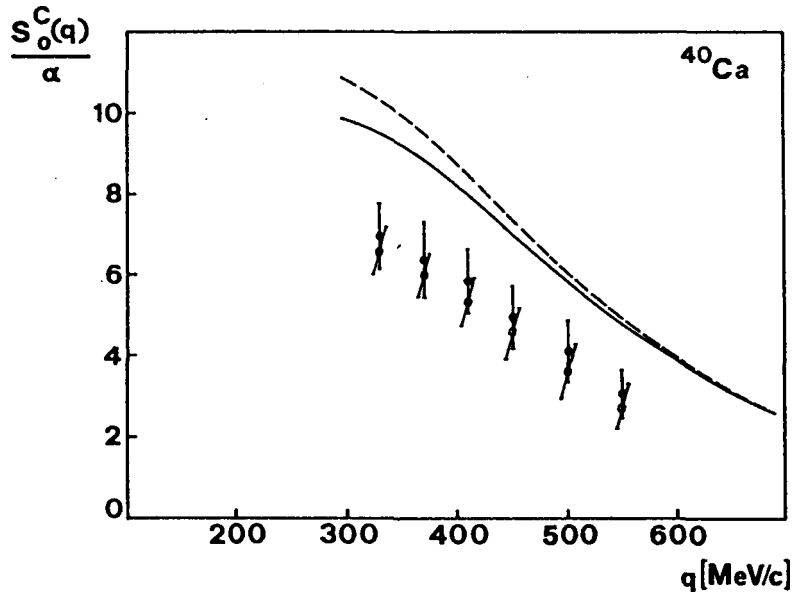


Fig. 1b. Longitudinal Sum Rule. Effect of short-range correlations: calculation by G. Orlandini and M. Traini (Ref. 51).

Dashed curve: pure H.O. wave function

Full curve: H.O. + tensor force + S.R. correlations

is quite well reproduced but the magnitude remains much too high (fig. 2).

Several explanations have been invoked to explain this quenching of the longitudinal response: short range correlations, coupling of 2p-2h states to the T=1 part of the response⁶, relativistic effects due to the Dirac spinor nature of the nucleon and again swelling of the nucleon. Indeed some empirical increase of the charge radius (typically 15 to 20 % for ⁴⁰Ca) and anomalous magnetic moment accounts for such a quenching^{5,7}.

2.2 Inclusive high q^2 measurements without σ_L / σ_T separation

At higher four-momentum transfers, extensive measurements have been performed at SLAC⁸ on ³He up to 4 GeV². Results are now available on Al and Fe from the NE3 experiment⁹. All of these measurements confirm the scaling behaviour of the inclusive cross-section in terms of one single variable y :

$$d^2\sigma(q^2, \omega) / (Z\sigma_p + N\sigma_n) d\omega = F(y)dy$$

Though there are current discussions on the best choice of this scaling variable y (linked in some way to the longitudinal component of the struck nucleon momentum¹⁰) the y -scaling is a well-established experimental fact in the low energy side of the quasi-elastic peak: it is generally interpreted as reflecting the single particle character of the quasi-free process.

Sick pointed out¹¹ that changing the scale of the q^2 dependance of the nucleon form factors G_E and G_M destroys the y -scaling. This argument tends to impose severe limitations to an increase of the nucleon electromagnetic radius in nuclei:

$$R_N^{\text{Bound}} \leq 1.03 R_N^{\text{Free}}$$

was found a realistic limit. However one should notice that in the kinematics of present day experiments the response function is dominantly transverse so that the limitation stands mainly for the magnetic radius.

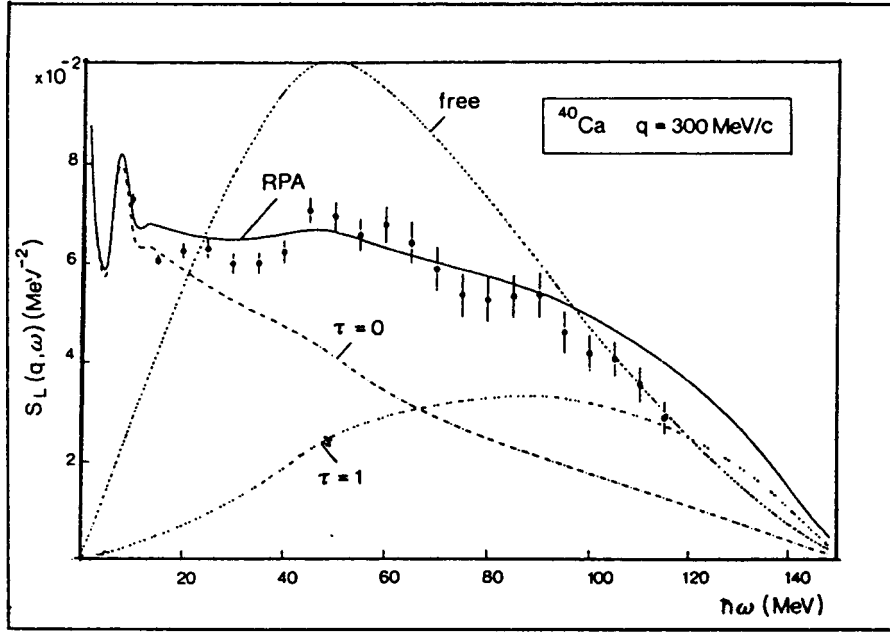


Fig. 2a . The longitudinal response function in $^{40}\text{Ca}(e,e')$ Saclay measurements (Ref. 3) compared with semi-classical R.P.A. calculations by Alberico et al.⁵ (full line). In dotted lines, ("free") the independent particle response.

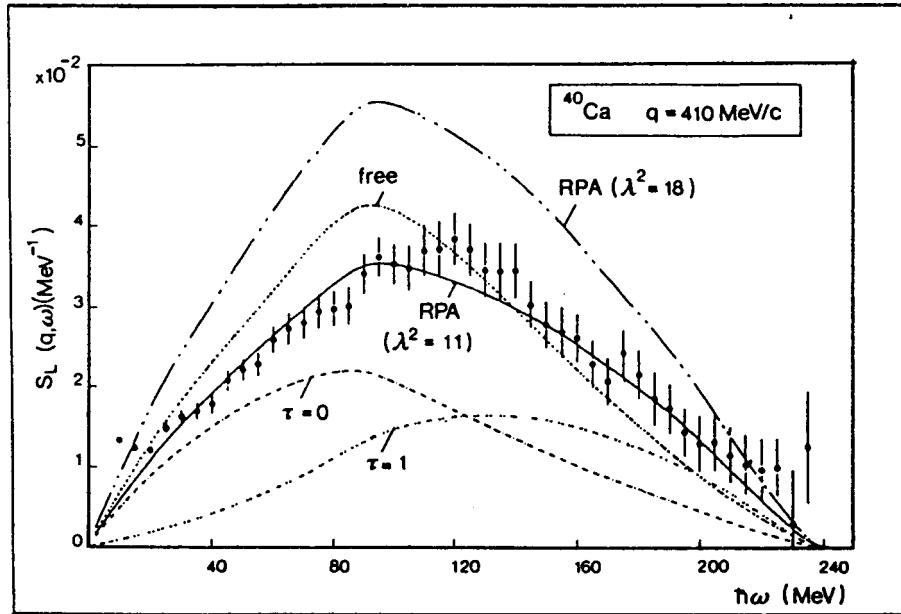


Fig. 2b . Same as Fig. 2a, for $q = 400 \text{ MeV/c}$. The double-dot-dashed line is the result of the calculation with the free proton form factor. To obtain the continuous curve, the proton charge radius has been increased by 22%.

2.3 Exclusive (e,e'p) measurements with σ_L / σ_T separation

The purpose of these experiments is to select the one nucleon knock-out process. In PWIA the cross-section would factorize in two parts:

$$d^6\sigma = K \sigma_{ep} S(E_m, p_m)$$

where σ_{ep} is the elementary electron-bound proton cross-section and $S(E_m, p_m)$ is the spectral function. Working in "parallel" kinematics (outgoing proton parallel to the virtual photon three-momentum) suppresses interference terms σ_{TT}, σ_{TL} and the σ_{ep} elementary cross-section writes:

$$\sigma_{ep} \propto \sigma_{Mott} (W_T(q^2) + \epsilon W_L(q^2))$$

ϵ is the polarization of the virtual photon.

Fixing E_m and p_m to keep the spectral function constant, and measuring a backward and a forward scattering angle allows to extract the ratio W_L/W_T for a given q^2 . Recent measurements on ^{12}C and ^6Li at NIKHEF¹² and on ^{40}Ca at Saclay¹³ have shown that this ratio is quenched by 20 to 40%, which appears consistent with the inclusive measurements (fig. 3).

One further step is, keeping E_m and p_m constant, to vary q^2 in the maximum range available and compare the behaviour of the cross-section with the free proton one. This has been done in the Saclay experiment on ^{40}Ca up to .6 GeV². The data seem to rule out a large modification in the electromagnetic form factors of the bound proton. However the information on the electric part remains too poor to draw definitive conclusions on the charge radius.

Two limitations to the PWIA factorization approximation appear in the analysis of these experiments:

1. The distortion of the incoming photon.

2. The distortion of the outgoing proton. DWIA calculations by Boffi were used for the NIKHEF and Saclay data. However this crucial problem of Final State Interaction certainly deserves further studies⁴⁶.

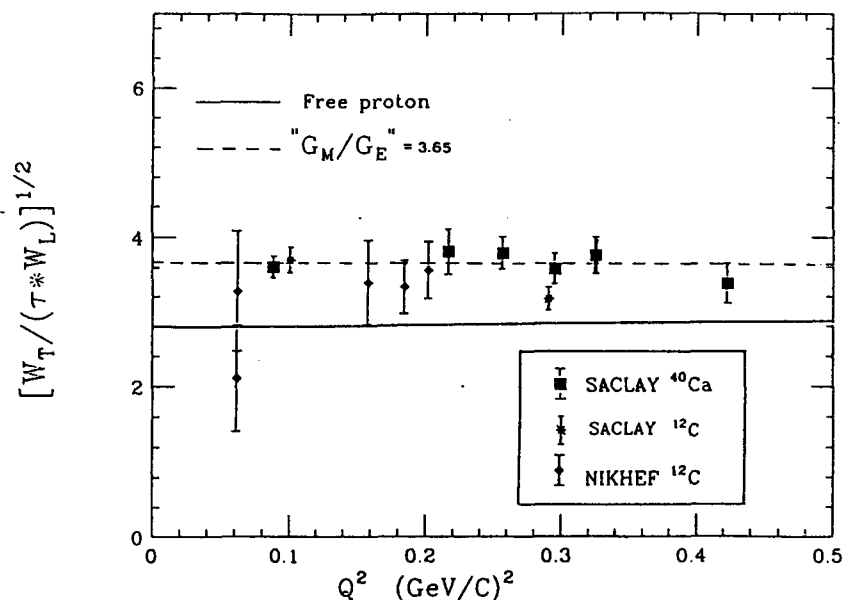


Fig. 3a. The measured ratio W_T/W_L in Saclay and NIKHEF (e,e'p) measurements. The quantity $[W_T/(\tau W_L)]^{1/2}$ would be $G_M/G_E \approx 2.79$ for a free proton. The average measured value for ^{40}Ca is 3.65.

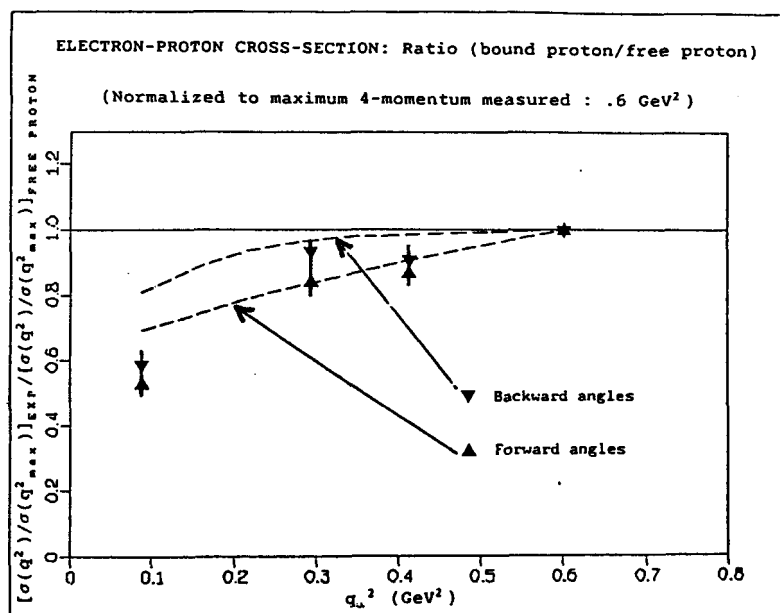


Fig. 3b. The q^2 dependance of the cross-section in the Saclay $^{40}\text{Ca}(e,e'p)$ experiment. The measured cross-section is divided by the free proton cross-section. This ratio is normalized to the $.6 \text{ GeV}^2$ point. The dashed curves show the behaviour expected if one extrapolates to this point the ratio W_T/W_L measured up to $.42 \text{ GeV}^2$, without modifying G_E or G_M .

OFF-SHELL EFFECTS AND AMBIGUITIES IN THE NUCLEONIC CURRENTS

In analyzing quasi-free electron scattering experiments one usually assumes that the nuclear current is just the sum of individual nuclear currents: this is the **Impulse Approximation**. Since nucleons are in fact off-shell one must define rules to extrapolate the free nucleon current. Prescriptions are needed for the nucleon spinor, the nucleon form factors and the current operator. The simplest choice is to keep the spinor and form factors of the free proton. An ambiguity remains in the choice of the operator: this problem has been carefully investigated by Mougey and Frullani¹⁴ and by T. De Forest¹⁵ who showed that, imposing Lorentz and gauge invariance, these ambiguities are small. This assumption does not mean that the Impulse Approximation is a good one: it only means that it is well defined.

Going further means either changing the nucleon spinor or its form factors, or both. The first alternative occurs in relativistic theories of nuclear dynamics.

RELATIVISTIC THEORIES OF NUCLEAR DYNAMICS

Since the original work by Walecka¹⁶, there has been considerable development in field theoretic descriptions of nucleonic spinor fields interacting with mesonic fields, basically a scalar field σ and a vector field ω . One further step was to incorporate in this picture the π and ρ mesons¹⁷. This Dirac phenomenology has achieved many successes in accounting for a large variety of nuclear phenomena. It has been recently formulated in terms of a renormalizable theory called Quantum Hadrodynamics (QHD)¹⁸.

Two essential features are to be retained from this relativistic approach.

1. The nucleon motion should be described by the Dirac equation.

2. The self-energy term in the Dirac equation arises from the intense attractive σ and repulsive ω fields.

Approximating the fields by their classical values, the self-energy term reduces to:

$$\Sigma = V_s + \gamma^0 V_0$$

V_s is then a constant scalar potential and V_0 a constant vector potential. Typical values for nuclear matter are:

$$V_s = -400 \text{ MeV} , V_0 = 300 \text{ MeV}$$

The effect of the σ and ω fields is to shift the mass and energy scales in the Dirac equation:

$$M \Rightarrow M^* = M + V_s$$

$$E \Rightarrow E^* = E - V_0$$

In the non-relativistic reduction V_s and V_0 appear in the combination $V_s + V_0$ which produces the conventional mean potential. On the contrary $M + E$ terms which appear in the small components of the Dirac spinor tend to double their free space value. The existence of corrective terms containing $V_s - V_0$ in the coupling to negative energy states is the source of the relativistic effects. In electron scattering processes these effects could manifest themselves at two levels:

1. The nuclear wave function, which should be calculated using relativistic Hartree or Hartree-Fock method.
2. The nucleon current, which should be modified by the shift in mass and energy scales. Work is in progress on this subject¹⁹.

Several attempts to explain inclusive data on ^{12}C and ^{40}Ca in the σ - ω framework have been developed recently^{20,21}. They do not seem to fully solve the problem of reproducing simultaneously the transverse and longitudinal response function. Further work is needed, especially in view of the new (e,e'p) experiments.

Abandoning the point-like picture of the nucleon requires incorporating some knowledge of its internal structure. This is the place to ask what can be learnt from QCD.

WHAT QCD TELLS US

The fundamental problem of understanding the QCD vacuum structure remains far from a first principles solution. To understand how it is modified in nuclear matter is certainly not an easy task. The kinds of question to be answered are: What happens with the properties of nucleons when they are not surrounded by a sufficient amount of vacuum ? Does the pressure of the vacuum (the bag constant) decrease, thus resulting in an increased size ? etc...⁴⁹

The most promising attempts to solve "exactly" non-perturbative QCD, namely lattice gauge calculations, are presently too rough to set a scale on such parameters as the mass or size of bound nucleons. Some amount of phenomenology is thus unavoidable.

QCD Sum rules and QCD condensates

There are current attempts to relate the scale of dimensional quantities such as masses and radii to the QCD condensates: chiral condensate $\langle 0 | \bar{q}q | 0 \rangle$ and gluon condensate $\langle 0 | G_{\mu\nu}^a G_a^{\mu\nu} | 0 \rangle$. The argument is based on the so-called QCD Sum rules²², and estimates of the dynamical quark m_q and gluon m_g masses are derived from values of the condensates. Currently accepted values are²³:

$$\langle \bar{q}q \rangle = -.016 \text{ GeV}^3, \quad m_q = 300-320 \text{ MeV}$$

A rough approximation of the nucleon mass gives then :

$$M_N = [-8\pi^2 \langle \bar{q}q \rangle]^{1/3} \cong 1 \text{ GeV}$$

which supports the idea that the chiral condensate is the main ingredient in the nucleon mass²⁴. Some understanding of how the condensates behave in the nuclear environment can be achieved in the framework of effective Lagrangian models.

QCD and dynamical rescaling

The idea of a dynamical change of scale associated with a growing correlation length between quarks has been proposed by Jaffe²⁵ and developed by Close, Roberts and Ross²⁶ from a careful analysis of Deep Inelastic Scattering data. By looking at the q^2 dependance, they observed that a good fit to a large fraction of the

data can be obtained by assuming:

$$F_2^A(x, q^2) = F_2^N(x, \xi q^2)$$

For large q^2 the rescaling parameter ξ increases with A in the range of 1 to 3.

Solving the evolution equation of the structure functions in first order QCD they found that the rescaling parameter ξ would be given by:

$$\xi(q^2) = (R_A^2/R_N^2)^{\alpha(\mu^2)/\alpha(q^2)}$$

where R_A and R_N are the confinement size of quarks in the nucleus and in the nucleon. The normalization point μ^2 appears as a "natural" scale of QCD; essentially $\mu^2 \simeq 1/R^2$. Typically, for $q^2 = 20 \text{ GeV}^2$, and for ^{56}Fe , $\xi \simeq 2$ leads to $R_A/R_N \simeq 1.15$.

The rescaling hypothesis tells nothing about the mechanism which affects the quark confinement. Close et al. originally used a crude bag overlap model, but they assumed their argument to be more general. However, they did not state that the change in confinement size actually corresponds to a physical swelling of the nucleon. Several estimates of the ratio R_A/R_N have been proposed in the framework of dynamical models of the nucleon.

DYNAMICAL MODELS OF THE NUCLEON

Chiral symmetry and the Skyrmion

The simplest effective lagrangean model describing the interaction of a spinor ψ with a four-component chiral field (σ, π) is the nonlinear σ -model²⁷, which serves as a basis for several models of the nucleon. In this framework, the quark condensate $\langle \bar{q}q \rangle$ is represented phenomenologically by the expectation value in the vacuum of the scalar field σ . The nucleon mass then becomes:

$$M_N = g \langle \sigma \rangle$$

The remark that $\langle \sigma \rangle$ does not reach its free value in nuclear matter was first made by Noble²⁸, in relation with the nucleon effective mass M^* of relativistic mean field theories. It would come from the neighbouring nucleons acting as additional attractive sources of σ -field²⁴. In the σ -model, the nucleon mass is exactly

$M_N = g f_\pi$ and the increase in nucleon size is associated with a decrease of f_π in nuclei. This argument is somewhat refined in the Skyrmion model²⁹.

Let us recall that the skyrmion is built on large N_c (N_c = number of colours) arguments. Following these arguments, the nucleon radius is constrained to behave like $R = C \sqrt{g_A}/f_\pi$. In a medium of baryonic density ρ , the "quasi-skyrmion" would acquire a modified size: $R(\rho) = C \sqrt{g_A(\rho)}/f_\pi(\rho)$. Neglecting higher order corrections in $1/N_c$, the constant C is unchanged.

From a study of axial currents in nuclei Rho³⁰ obtained an estimate of $R(\rho)$ in nuclear matter. He found:

$$\begin{aligned} f_\pi(\rho)/f_\pi &= M_N^*/M_N = .62 \\ g_A(\rho)/g_A &= .8 \\ R(\rho)/R &= 1.44 \end{aligned}$$

The value of M_N^* comes from a relativistic mean field approach. Note that the calculated value of g_A is consistent with experimental data.

Friedberg-Lee soliton

In this soliton bag model³¹, quarks interact only with the scalar field σ which generates the confinement. The dynamics of the soliton is governed by the self-energy term, which is a quartic potential $U(\sigma)$. The presence of neighbouring nucleons acting as additional sources of field modify this potential term.

Jädel and Peters³² studied a particular class of solutions and produced a dependance of the nucleon radius and magnetic moment with the nuclear density. They found $R(\rho)/R(0) = 1.40$ for nuclear matter density and ≈ 1.15 for iron.

Celenza-Rosenthal-Shakin soliton

Starting from a scheme analogous to the Friedberg-Lee soliton, this model³³ aims at some synthesis of the soliton approach with the relativistic models of nuclear dynamics. For that purpose, it incorporates the scalar σ and vector ω fields in addition to the π

and ρ fields to build the soliton. Celenza et al obtained a solution with static electromagnetic properties close to that of the nucleon.

In the nuclear medium, additional σ and ω fields due to the surrounding nucleons modify the quark self energy Σ_q :

$\delta\Sigma_q = a + \gamma^0 b$, where the scalar piece a appears as a shift of the dynamical quark mass. They obtained $\delta m_q = -147$ MeV, starting with $m_q = 600$ MeV.

Predictions for electromagnetic form factors lead to a nucleon swollen by 0 to $\simeq 20\%$, depending on the nuclear density. The decrease in G_E was found in agreement with the quenching of the longitudinal response function R_L observed for medium nuclei³⁴.

In recent developments of this approach the authors give a justification for the high value of their quark dynamical mass $m_q \simeq 600$ MeV. Half of it could be due to the gluon condensate $\langle G_{\mu\nu}^a G_a^{\mu\nu} \rangle$. They develop the idea that this quantity could play an essential role as an order parameter of QCD and allow to define mass and radius scale³⁵.

The colour dielectric model

Proposed by Chanfray, Nachtmann and Pirner³⁶, this soliton model presents some similarities with the Friedberg-Lee model. In analogy with a dielectric, the QCD vacuum is described as a medium with a colour dielectric constant $\epsilon = 0$, which confines coloured objects into a bag (inside of the bag $\epsilon \simeq 1$). This parameter ϵ is related to a local order parameter $\chi(x)$ ($\epsilon = \chi^4$) calculated using lattice techniques.

In the nuclear medium, the χ field is enhanced in the outside region to a non-zero value. The resulting leakage of the quark wave function generates a binding effective potential. The depth W_N of this potential (which is essentially the average binding energy) can be related to the increase in confinement radius and is calculated as a function of nuclear density. For example, in ^{56}Fe , they find $\chi_N = .016$, $W_N = -33.9$ MeV, $\Delta R/R = .15$.

PIONIC EFFECTS ON THE BOUND NUCLEON STRUCTURE

Another class of approaches to medium modifications of the nucleon emphasize the role of long range phenomena ($d \gtrsim 1\text{fm}$). Since these phenomena are essentially carried by the pion field, their basic idea is that the peripheral structure of the nucleon, i.e. its pion cloud, is distorted by the nuclear environment. Implications of an enhanced pion field for Deep Inelastic Scattering have been investigated by several authors^{37,38}.

Recently, this kind of approach has been extended by M. Ericson and Rosa-Clot³⁹ to the electromagnetic polarizability of the nuclear medium and the charge radius of the bound nucleon, two physical quantities which present some similarity. A semi-quantitative estimate leads to a large increase in mean square charge radius $\Delta \langle R_c^2 \rangle$ of .4 to .5 fm², resulting in a large quenching of the proton electric form factor.

CONCLUSIONS

A number of theoretical arguments support the hypothesis of properties of the nucleon modified by the nuclear medium. Relativistic effects, distortion of the pion field, change in the quarks confinement volume are candidates to generate these medium effects.

The idea of changing a single mass or length scale when going from the free space to the interior of a nucleus appeared an elegant way to build a bridge between short and long distance nuclear phenomena.

However, present day experimental data bring more and more evidence against the idea of a swollen nucleon.

On the one hand, it seems now likely that at least a large part of the EMC effect can be explained by a careful treatment of Fermi smearing and nuclear binding effects⁴⁰. Whether there may be some deep connection between the QCD dynamical rescaling argument and the nuclear binding explanation remains a subject of speculation⁴¹.

On the other hand, the y-scaling behaviour of the inclusive

quasi-free cross-section, the q^2 behaviour of the exclusive $^{40}\text{Ca}(e,e'p)$ cross-section have brought strong limitations to an increase of the proton magnetic radius.

There were current suggestions that the nucleon magnetic moment appears larger in nuclei than in free space^{7,42,43,44}. T.E.O. Ericson and A. Richter⁴⁵ derived an upper limit of $\simeq 2\%$ to such an increase in the $A = 3$ and $A = 12$ nuclei by using the scaling between isovector magnetic moments and the corresponding axial matrix element from beta decay.

Nevertheless the longitudinal response function in quasi-free scattering remains a puzzle and the accuracy of present data is not good enough to rule out a quenching of the proton electric form factor. More data are needed, both inclusive (e,e') and exclusive $(e,e'p)$ to separate the transverse and longitudinal response functions in a wider momentum range.

New inclusive experiments are in progress or planned in the region $q^2 < .6 \text{ GeV}^2$. An inclusive experiment (NE9) is planned at SLAC on a series of nuclei near 1 GeV^2 . Exclusive measurements on ^4He , ^6Li are in progress at Bates, NIKHEF, Saclay. Extension of these coincidence experiments to larger q^2 will have to wait for the next generation of multi-GeV CW electron machines^{47,48}.

Future studies should not be confined to the nucleon in its fundamental state. The existence and the behaviour of the nucleon resonances in the nuclear medium open an entirely new chapter of nuclear physics. Covering a large (q^2, ω) domain will bring new insights on the connection between the traditional low energy nuclear physics and the high energy domains where quarks and gluons degrees of freedom are expected to enter into play. This is clearly a major goal of future electron facilities.

ACKNOWLEDGEMENTS

I would like to thank the organizers of the 1986 CEBAF Summer Study Group and Workshop for warm hospitality during my stay at CEBAF in June 1986. I am indebted to all the participants to the working group on "Single Nucleon Densities" and to my colleagues of

the "coincidence" group at Saclay Linac for stimulating and helpful discussions.

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RECENT PROGRESS IN NON-PERTURBATIVE QCD: COLOUR DIELECTRIC LATTICE GAUGE THEORY AND NUCLEI

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1. GENERAL INTRODUCTION

Perturbative quantum chromodynamics (QCD) has been extensively studied and agrees well with experiment. As an example, let us look at the angular distribution of two-jet events measured in the UA1 experiment at CERN (see Fig. 1). The experimental data ¹ follow the $1/(1 - \cos \theta)^2$ behaviour of Rutherford scattering. Even small visible corrections from scaling violations (Q^2 -dependence) in the structure functions are calculable. The gluon exchange manifests itself in the same way as photon exchange in electrodynamics. However perturbative QCD only becomes applicable at length scales $1/Q^2 \ll (0.1 \text{ fm})^2$. The interesting physics of interacting hadrons and nuclei, is given by a scale of 1 fm to 10 fm. It therefore needs non-perturbative methods to tackle the QCD Lagrangian. Various

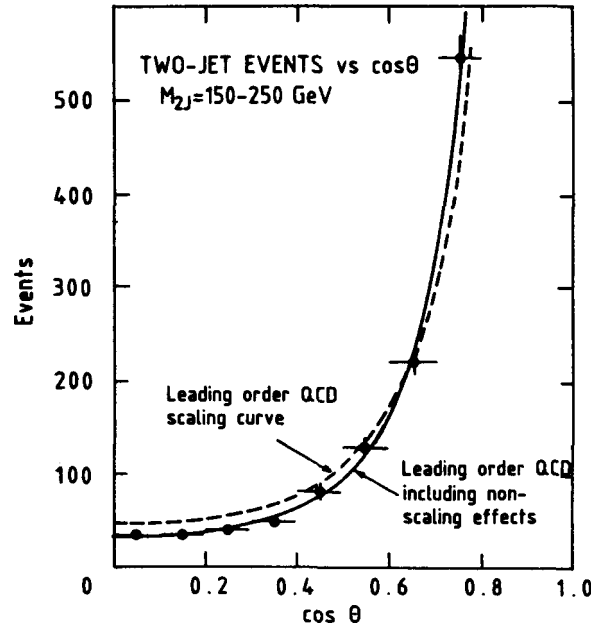


Fig. 1 Two jet events as a function of $\cos \theta$ from Ref. 1. The $qq \rightarrow qq$, $gg \rightarrow gg$ and $gq \rightarrow gq$ processes give the Rutherford scattering distribution (dashed line) $d\sigma \propto 1/(1 - \cos \theta)^2$. The full line includes the scaling violations in the structure functions.

approaches have been proposed. Since an expansion in the strong coupling constant $\alpha_s(Q^2)$ of QCD is not possible, an expansion in $1/N_c$ has been suggested, where N_c is the number of colours. ^{2,3} One obtains a theory of weakly interacting mesons with couplings $g_M = O(1/N_c)$ and masses $O(1)$. Baryonic solutions have masses $O(1/g_M) = O(N_c)$, which possibly indicates that they are soliton solutions of the meson-field equations. The existence of a gluon master field proposed by Witten ⁴ in the $N_c \rightarrow \infty$ limit would allow quantitative calculations, but up to now no results in four dimensions are available. There are calculations trying to relate the QCD Lagrangian to a Skyrme-type effective action assuming chiral symmetry breaking. ⁵ But most of the progress in this field has been made in applications of Skyrme-type models. ⁶

A different mostly numerical approach to non-perturbative QCD has been followed during the last years, namely lattice QCD. ⁷ By discretizing continuous space-time on a finite lattice one considers a finite number of gauge-field configurations, which can be simulated numerically. The finite lattice constant regularizes an otherwise diverging field theory. The lattice method has developed into a field where the computational possibilities play an important part in the discussion. However there are also serious, not yet understood, physics problems concerning the treatment of systems with finite fermion density. ⁸

The work I am going to describe here has involved a moderate numerical effort up to now. Its aim is primarily conceptual. Can we define an effective QCD theory, where the short-wavelength gluon degrees of freedom are eliminated in favour of a colour-neutral field characterizing the vacuum and giving confinement. The role of the polarizable medium is played by the high-frequency parts of the gauge fields. Since this effective theory has the same appearance as electrodynamics of dielectric media, it is called colour dielectrics.

From the 'macroscopic' Hamiltonian

$$\mathcal{H} = \int \sum_{i=1}^3 \frac{1}{\epsilon} [\vec{D}_i^2 + \vec{H}_i^2] d^3x$$

it is easily seen that colour displacement fields $\vec{D}_i$ and magnetic fields $\vec{H}_i$ cannot subsist in a medium where $\epsilon = 0$, because the energy associated with them would be infinite. The vacuum therefore has $\epsilon = 0$. The dynamics of confinement is reduced to find the value of the dielectric field in every point in space around the colour charges. The region with $\epsilon = 0$ forms a bag around the $\epsilon \neq 0$ region inside the hadron, where colour is present and not influenced by vacuum fluctuations.

My paper will consist of three parts. In Section 2 I give an introduction to the main physics of lattice gauge theory. Section 3 gives an outline of the colour dielectric model and first numerical results on the effective action after one block-spinning step. This part of the paper is based on joint work with J. Wroldsen. ⁹ Section 4 is devoted to a report about phenomenological work, done in collaboration with G. Chanfray and A. Schuh, which uses the colour dielectric model as a starting point for calculations of the free nucleon and nuclear matter or finite nuclei.

2. INTRODUCTION TO CONCEPTS OF LATTICE GAUGE THEORY

For simplicity we are going to discuss non-Abelian SU(2) chromodynamics of gluon fields. That means the gluon fields have a 2×2 matrix representation by the Pauli matrices $\vec{\sigma}^i$ ($i = 1, 3$) as

$$A_\mu(x) = \frac{1}{2} \sum_{i=1}^3 A_\mu^i \sigma^i$$

Having discretized the three space and one time dimensions we introduce as relevant degrees of freedom the link variables

$$U_\mu(x) = \exp i g \int_x^{x+e_\mu} A_\nu dx_\nu = a_0 + i \vec{\sigma} \vec{a} \quad (1)$$

which symbolize the integral of the gauge field between a site x of the lattice to the neighbouring site $x + e_\mu$, one lattice spacing away from x in the $\mu = (1, 4)$ direction. This link variable has a 2×2 matrix representation $U_\mu = a_0 + i \vec{\sigma} \vec{a}$ given by four numbers $(a_0, \vec{a})$. It is unitary, which means that the product of a link in the positive μ -direction with the same link traversed in the opposite μ -direction gives the unity operator

$$U_\mu(x) U_\mu^\dagger(x) = U_\mu(x) U_{-\mu}(x+e_\mu) = 1 \quad (2)$$

Unitarity restricts the four quantities $(a_0, \vec{a})$ to the unit circle $a_0^2 + \vec{a}^2 = 1$. The effect of a gauge transformation $G(x)$ on $U_\mu(x)$ can be seen by expanding U_μ for small gauge fields and using the standard continuum gauge transformation $A_\mu(x) \rightarrow A'_\mu = (1/i) [\partial_\mu G(x)] G^{-1}(x) + G(x) A_\mu(x) G^{-1}(x)$. We get

$$\begin{aligned} 1 + i A_\mu dx_\mu &\Rightarrow \\ 1 + \left\{ [G(x+dx) - G(x)]/dx_\mu G^{-1}(x) + i G(x) A_\mu(x) G^{-1}(x) \right\} dx_\mu &(3) \\ \approx G(x+dx) (1 + i A_\mu dx_\mu) G^{-1}(x) & \end{aligned}$$

Therefore a link $U_\mu(x)$ transforms by applying the gauge transformations G and G^{-1} on the end points $(x + e_\mu)$ and x of the link. Consequently, the natural choice for a gauge-invariant action is made from closed loops, the smallest of which are plaquettes with four links. Indeed this is the form of the Wegner-Wilson^{10,11} action $S(U)$ (see Fig. 2)

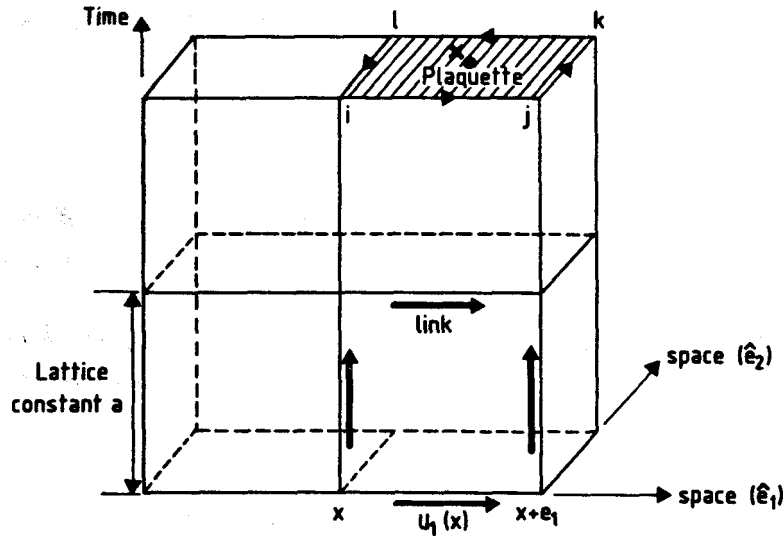


Fig. 2 Schematic picture of a lattice with the link variables and a basic plaquette.

$$S(U) = \sum_{\text{all plaquettes}} \left[1 - \frac{1}{2} \text{tr} (U_{ie} U_{ek} U_{kj} U_{ji}) \right] \quad (4)$$

$$= \sum S_{\square}$$

The partition function is given by the sum over all possible link configurations weighted with the 'Boltzmann factor'

$$Z = \int \mathcal{D}U e^{-\beta S(U)} \quad (5)$$

Here the integral can always be thought of as a multiple integration over the constrained vector variables $a_0, \vec{a}$ at all links, i.e.

$$\int \mathcal{D}U = \int \prod_{\text{links}} \pi da_0^e d\vec{a}_e \delta(a_0^2 + \vec{a}_e^2 - 1) \quad (6)$$

The QCD coupling constant g enters into the expression $\beta = 4/g^2$. In the strong coupling limit, i.e. for small β , we see that the exponential allows values of $S_{\square} \rightarrow 2$ which are related to strong fluctuations of the links, e.g.

$$(UUUU) = \begin{pmatrix} -1 & 0 \\ 0 & -1 \end{pmatrix}$$

For weak coupling or large β the 'Boltzmann' weight prefers $S_{\square} \rightarrow 0$, i.e. trivial link variables near the identity, namely

$$(UUUU) = \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}$$

Part of the normal numerical problem is associated with the necessity of working in the weak-coupling domain, where the fluctuations of the gauge fields are suppressed and the 'thermalization' of the system takes very long times.

An important point concerns the continuum limit of the action given in Eq. (4), when the lattice spacing a goes to zero. Take a plaquette in the $\hat{e}_1, \hat{e}_2$ plane and expand all link variables around the centre of the plaquette; then one gets (Fig. 2)

$$S_{\square} = 1 - \frac{1}{2} \text{tr} \left[\exp i g \bar{A}_1 \left(x - \frac{1}{2} a \hat{e}_2 \right) a \right. \\ \exp i g \bar{A}_2 \left(x + \frac{1}{2} a \hat{e}_1 \right) a \\ \exp -i g \bar{A}_1 \left(x + \frac{1}{2} a \hat{e}_2 \right) a \\ \left. \exp -i g \bar{A}_2 \left(x - \frac{1}{2} a \hat{e}_1 \right) a \right] \quad (7)$$

Let us use the Baker-Hausdorf formula $e^x e^y = e^{x+y+1/2[x,y]}$ to derive

$$S_{\square} = 1 - \frac{1}{2} \text{tr} \exp \left\{ i a^2 g (\partial_1 A_2 - \partial_2 A_1) - a^2 g^2 [A_1, A_2] \right\} \quad (8) \\ = 1 - \frac{1}{2} \text{tr} \exp i a^2 g F_{12}$$

where the commutator term does not vanish, because of the non-Abelian nature of the gauge fields. For small lattice spacing a , we can expand the exponential. The trace of the first-order term $\text{tr} a^2 g F_{12}$ vanishes when we integrate over all group elements, and the second-order term gives the desired continuum action for $\beta = 4/g^2$:

$$\beta \sum S_{\square} = \frac{1}{8} \text{tr} \beta g^2 a^4 \sum_x F_{\mu\nu}(x) F_{\mu\nu}(x) = \\ = g^2 \beta / 4 \cdot \int d^4 x \text{tr} \frac{1}{2} F_{\mu\nu} F_{\mu\nu} \quad (9)$$

Asymptotic freedom can be expressed in this language as a relation between the QCD coupling constant g and the lattice spacing a , which has the following form for SU(2)

$$g^2 / 4\pi = \frac{3\pi}{(11/2) \ln(1/\Lambda^2 a^2)} \quad (10)$$

Here Λ is the dimensional size parameter of SU(2) QCD, i.e. Λ_{QCD} . In order to obtain the continuum limit one has to make the lattice spacing a sufficiently small, which makes g also small. Then one is faced with the problem of slowing down the equilibration of the system together with a real lattice size of small dimensions (1-2 fm).⁴ Recently calculations with $(24-32)^4$ lattices have been started. Turning the above formula (10) around, a numerical simulation of a physical dimensional parameter m^2 is considered as scaling when it behaves like the physical measurable QCD constant, namely as

$$m^2 = \text{const} \cdot \Lambda^2 = \text{const} \cdot \frac{1}{a^2} \exp(-6\pi^2/11 \cdot \beta) \quad (11)$$

Taking a numerical expectation value at several values of β , a comparison with Eq. (11) establishes whether one is close enough to the continuum limit. Clearly numerical lattice simulations can only give us physical quantities expressed in units of Λ , which has to be taken from experiment.

The numerical technique we use is the heat-bath method.¹² Working our way through the lattice we generate random four vectors $a_0, \vec{a}$ for each link according to the weight $e^{-\beta S(U)}$ which is determined by the neighbouring links. Once the system is in equilibrium we can evaluate any desired correlation function.

3. COLOUR DIELECTRICS

Colour dielectrics is based on regarding the QCD vacuum as a medium characterized by a dielectric field. It is a theoretical approach which promises to construct an effective action for a bag model directly from the QCD Lagrangian. Clearly for many experiments in nuclear physics at higher energies we need a theoretical description which at the same time contains the long-range features of confinement without giving up the elementary quark and gluon degrees of freedom. After all it is the search for these QCD degrees of freedom which motivates the experimental work at new accelerators such as CEBAF. Studying the transition from asymptotically free QCD at short distances to confining hadron dynamics at long distances is the exciting task we are faced with. For this problem a transformation is ideal which maps a theory with a lattice spacing a into one with lattice spacing $2a$. Since such a mapping or block-spin transformation can be carried out repeatedly we are led to an effective model Lagrangian for long-distance hadron physics. We define a chain of field theories at $a_0, 2a_0, 4a_0, \dots, 2^n a_0$, then the continuum limit $a_0 \rightarrow 0$ is achieved by letting $n \rightarrow \infty$ such that $\lim 2^n a_0 \rightarrow \text{finite}$.

In Nielsen and Patkos¹³ a perturbative calculation of one block-spinning step towards the colour dielectric Lagrangian is made. Mack¹⁴ introduces models of effective actions and a definition of the colour dielectric field, which is used in our numerical work. For a general discussion of the renormalization group method in Monte Carlo calculations we refer the reader to the article of Wilson.¹⁵

Constructing the block variables for the effective action is a somewhat arbitrary process. We proceed in the following way. First we select every second hypercube on the

lattice. Then we make gauge transformations on all sites of a chosen hypercube in order to make its link variables $\tilde{U}$ close to the identity, i.e. $\frac{1}{2} \text{Tr } \tilde{U} \approx 1$. We iterate this gauge transformation program approximately 5 to 9 times for each hypercube. Once this is done all the information is in the links U_i which connect the individual hypercubes. In 4 dimensions there are 8 interconnecting links, the average of which defines the new link variables V times the colour dielectric field χ . This scalar field χ is necessary, otherwise the new link variables V would not be unitary.

$$\chi \cdot V = \frac{1}{8} \sum_{\substack{i=1 \\ \text{connecting}}}^8 U_i \quad (12)$$

It can easily be shown that for unitary matrices $V = b_0 + i\vec{\sigma} \cdot \vec{b}$ and $U_i = a_0^i + i\vec{\sigma} \cdot \vec{a}_i$ we get

$$\det(\chi b_0 + \chi i\vec{\sigma} \cdot \vec{b}) = \chi^2 (b_0^2 + \vec{b}^2) = \chi^2$$

and

$$\begin{aligned} \det\left(\frac{1}{8}(\sum a_0^i + i\vec{\sigma} \sum \vec{a}_i)\right) &= \\ &= \left(\frac{1}{8} \sum a_0^i\right)^2 + \left(\frac{1}{8} \sum \vec{a}_i\right)^2 \leq \max_{V_i} (a_0^i + \vec{a}_i^2) = 1 \end{aligned} \quad (13)$$

Therefore, without limiting our generality, χ will take values between 0 and 1.

For the new effective action we assume the following form:

$$\begin{aligned} e^{-S_{\text{eff}}(\chi, V)} &= \\ \exp\left[-\beta' \sum_0 \left(1 - \frac{1}{2} \text{tr}(\chi V)(\chi V)(\chi V)(\chi V)\right) - \sum_{\vec{x}} (m^2 \chi^2 + \lambda \chi^4)\right] \end{aligned} \quad (14)$$

The new plaquettes are now of length $2a$, and it is natural to decorate each new link V with the colour dielectric field χ , $0 \leq \chi \leq 1$. That is how we averaged the interconnecting gauge configurations anyway. In the continuum limit the above action will give us the desired result of a dielectric theory,

$$\begin{aligned} \mathcal{L} &= \frac{1}{2} \text{tr } B_{\mu\nu} B_{\mu\nu} \chi^4(x) + \mathcal{L}(\chi) = \\ &= \frac{1}{2} \sum_{i=1}^3 (\vec{E}_i^2 - \vec{B}_i^2) \cdot \epsilon + \mathcal{L}(\chi) . \end{aligned} \quad (15)$$

where

$$\epsilon = \chi^4$$

i.e. the colour dielectric constant $\epsilon(x) = \chi^4(x)$ modifies the electric and magnetic field strengths $\vec{E}_i^2$ and $\vec{B}_i^2$. In simulations on an 8^4 lattice using ~ 600 sweeps through the lattice, we obtained preliminary results (Fig. 3) for the dielectric constant $\langle \epsilon \rangle = \langle \chi^4 \rangle$. For strong coupling ϵ approaches zero, whereas, at $\beta = 2$, ϵ increases drastically. From simulations of the string tension it is well known that at $\beta = 2$ a cross-over transition takes place, which is reflected in the behaviour of ϵ . Using Schwinger-Dyson equations,¹⁶ we derive the new effective couplings β' , m^2 , and λ . The self interaction of the colour dielectric field is shown in Fig. 4. It documents the tendency for the effective action to assume its minimum near $\chi = 0$ when the coupling g is large. This supports the basic hypothesis of the dielectric theory of confinement. Colour sources are the source of electric inductions $\vec{D}_i$, which are related to the colour electric fields $\vec{E}_i$ by $\vec{D}_i = \chi^4 \vec{E}_i$. Therefore, the $\vec{D}_i$ fields can only be non-zero where $\epsilon = \chi^4$ is non-zero. If $\epsilon = 0$ is the unique classical vacuum then $\chi \neq 0$ costs energy and this prevents the $\vec{D}$ fields from spreading. As a result, a flux tube will form between quarks of opposite charges in the region where $\epsilon \neq 0$ only. Colour is

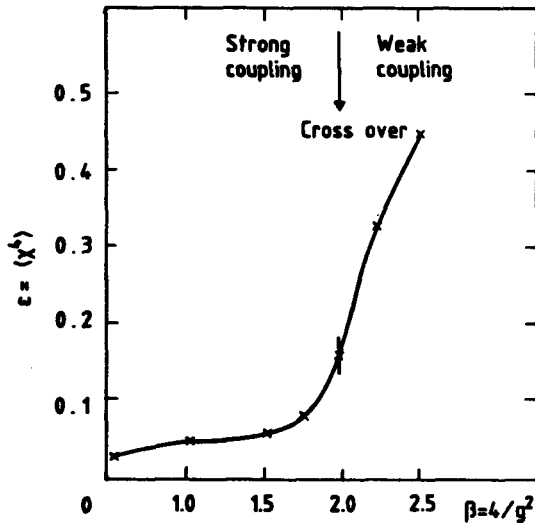


Fig. 3 Preliminary results for the dielectric constant $\langle \epsilon \rangle = \langle \chi^4 \rangle$ in an 8^4 lattice for different coupling constants $\beta = 4/g^2$.

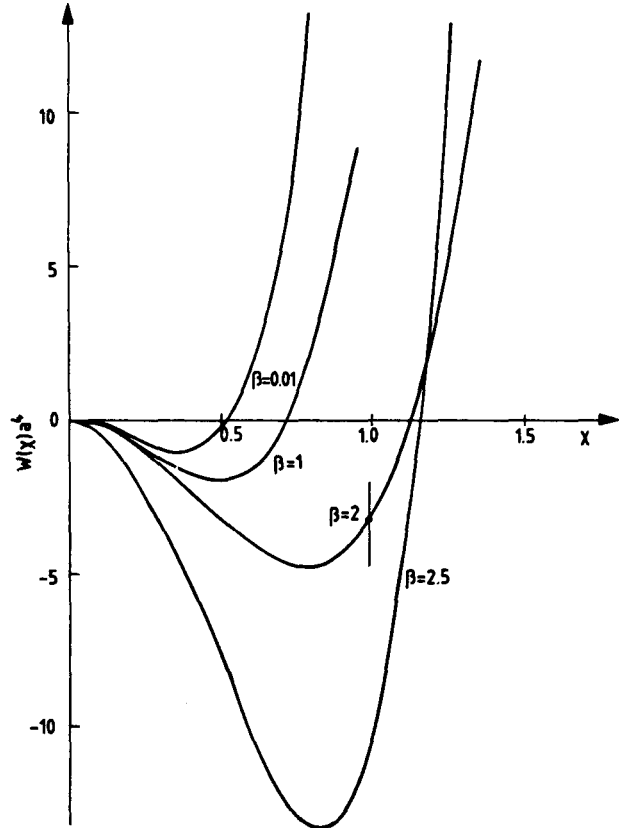


Fig. 4 Preliminary results for the χ -effective action $W(\chi)a^4$ as a function of β .

confined. In the next section we will see some phenomenological applications of this mechanism. The numerical calculations have just started and need more statistics to confirm the preliminary observed trends. Especially more simulations in the scaling region are needed.

4. PHENOMENOLOGY OF COLOUR DIELECTRICS

The original starting point for a colour dielectric quark model came from the need to understand nuclear binding and the EMC effect in the framework of a bag-like model.¹⁷ In order to understand the colour dielectric properties of nuclei, one has to introduce the coupling of the dielectric field to the quarks. In the lattice formulation the quark-gluon term is

$$S = \frac{1}{2} a^3 \sum_{i,j} \bar{\psi}_i \gamma_\mu e_\mu U_{ij} \psi_j \quad (16)$$

which destroys a quark at the lattice site j and connects it with a link U_{ij} to the site i where it is created. Averaging over different links will replace U by χV or expanding $V = 1 + i g a B_\mu e_\mu$, we get, collecting terms $\propto a^4$,

$$S_{eff} = a^3 \sum_i \bar{\psi}_i \gamma_\mu e_\mu \chi_i (1 + i g a B_\nu e_\nu) (\psi_i + e_\mu \partial_\mu \psi_i) - a^4 \sum_i \bar{\psi}_i \psi_i m_q \quad (17)$$

$$\Rightarrow \int d^4x \{ \bar{\psi}(x) \chi(x) (\gamma \partial + i g \gamma B) \psi(x) - m_q \bar{\psi} \psi \}$$

Using the covariant derivative

$$D_\mu = \partial_\mu + i g B_\mu$$

the total phenomenological Lagrangian has the form

$$\mathcal{L} = \mathcal{L}(\chi) + i \chi \bar{\psi} \gamma_\mu (\partial^\mu + i g B^\mu) \psi - \bar{\psi} \psi m_q - \frac{1}{4g^2} \chi^4 B_{\mu\nu} B_{\mu\nu} \quad (18)$$

where

$$B_{\mu\nu} = (\partial_\mu + i g B_\mu) B_\nu - (\partial_\nu + i g B_\nu) B_\mu$$

and $\mathcal{L}(\chi) = T - W$ is parametrized with

$$W(\chi) = B\chi^2 (A+6 - 2(A+4)\chi + (A+3)\chi^2) \quad (19)$$

and

$$T(\chi) = \frac{1}{2} \sigma_V^2 (\partial_\mu \chi)^2 \quad (20)$$

This phenomenological form of $W(\chi)$ is chosen in accordance with our previous work¹⁸ on the pure gluon theory, where we set $\chi = (\sigma_V - \sigma)/\sigma_V$. Note that the effective pure gluonic Lagrangian of Eq. (18) is similar to the Friedberg-Lee soliton model^{19,20} in this case. The dielectric constant is defined in the Lagrangian of Eq. (18) as $\epsilon = \chi^4$.

The fermion part of the colour dielectric model in Eq. (18), however, is totally different from the Friedberg-Lee model. In particular we note that the colour dielectric field leads to absolute quark confinement, since with $\langle \chi \rangle = 0$ in the vacuum the quarks can no longer propagate, i.e. their effective mass becomes infinite. To demonstrate this we introduce a new effective fermion field $\tilde{\psi}$ by rescaling the original field ψ

$$\tilde{\psi} = \sqrt{\chi} \psi \quad (21)$$

which shifts the χ -dependence from the kinetic term in Eq. (18) to the mass term $(mq/\chi)\tilde{\psi}\tilde{\psi}$.

As a first approximation, we consider the B_μ 's as perturbative gluon fields, dropping the non-linear parts of the Lagrangian. Phenomenologically we consider the χ -part as sufficient to describe the 0^{++} glueball and associate the χ -propagator with the glueball mass, i.e. $W(\chi)$ at $\chi = 0$ can be expanded as

$$\frac{1}{2} \left. \frac{\partial^2 W}{\partial \chi^2} \right|_{\chi=0} \chi^2 = \frac{1}{2} m_{GB}^2 \sigma_V^2 \chi^2 = B(A+6) \chi^2 \quad (22)$$

In an earlier paper,²¹ we have used the parameters $A = -2$, $\sigma_V = 160$ MeV, and $B^{1/4} = 203$ MeV, which reproduce rather well the following physical observables:¹⁸ the transition temperature $T_c = 160$ MeV to the gluon plasma, a glueball mass $m_{GB} = 720$ MeV, and a string tension $\sqrt{E}/\ell = 450$ MeV. In this calculation of the nucleon we have used the same parameters A , $B^{1/4}$, and σ_V . We note in passing that we also tried to calculate the nucleon with these parameters in the Friedberg-Lee soliton bag model²⁰ and always obtained nucleon masses $m_N \geq 2$ GeV. This is related to a general problem of the original bag model which must use different bag constants B for the heavy and light quark systems, typically $B^{1/4}(c\bar{c}) = 200$ MeV²² and $B^{1/4}(qqq) \approx 150$ MeV.²³ We will show how different colour dielectric fields in the interior describe both systems with the same effective Lagrangian.

The only remaining input parameter is the current quark mass on a hadronic size scale. We have argued ¹⁷ that the average light quark mass should lie between 20 MeV and 30 MeV, i.e. it is somewhat larger than in the estimates based at $Q^2 = 1 \text{ GeV}^2$. Different parameters have been chosen by Broniowsky et al. ²⁴ and Thomas et al. ²⁵ Especially in the last paper the whole region of parameter space has been systematically explored.

The equations of motion for the χ -field and the rescaled ψ -field (now written without $\sim$), are derived from the effective action for N quarks of equal mass m_q

$$(\vec{\alpha} \cdot \vec{p} + m_q/\chi) \psi = E \psi \quad (23)$$

and

$$-\sigma_V^2 \vec{\nabla}^2 \chi + \frac{\partial W}{\partial \chi} = N \bar{\psi} \psi \frac{m_q}{\chi^2} \quad (24)$$

Numerically, we solve the equations for the potential $\tilde{W}(\chi) = W(\chi - \delta)$, which has a minimum at $\chi = \delta = 10^{-3}$ such that $m_q/\delta \geq 20 \text{ GeV}$. Decreasing δ even more we do not find any variation in the results. Qualitatively, one understands the solution of the χ -field in the presence of the quark fields by regarding the effective potential

$$W_{\text{eff}}(\chi) = W(\chi) + N \bar{\psi} \psi m_q/\chi \quad (25)$$

which has a minimum away from $\chi = 0$.

Approximating the single-particle densities as

$$\bar{\psi} \psi \approx \frac{1}{2} \bar{\psi} \gamma_0 \psi = \frac{1}{2} / \left(\frac{4\pi}{3} R^3 \right) \quad (26)$$

we estimate the scalar density $\bar{\psi} \psi = 0.12 \text{ fm}^{-3}$ for a hadron size of $R = 1 \text{ fm}$. Then we get for $\bar{\chi}$ as the minimum of $W_{\text{eff}}(\chi)$ ($m_q = 20 \text{ MeV}$, $m_{\text{GB}}^2 \sigma_V^2 = 0.013 \text{ GeV}^4$):

$$\chi(0) \approx \bar{\chi} = 0.16 \quad (27)$$

In Fig. 5 we show the exact solutions of the coupled equations for $\chi(r)$. Indeed the value of χ inside the hadron is not very different from zero [$\chi(0) = 0.17$]. Comparing this value with $\chi(0) = 0.6$ of the $3\bar{3}$ flux tube ¹⁸ solution, we see even less change of the vacuum value of χ . The corresponding quark wave functions (Fig. 6) have a reduced small component compared to the MIT bag model, since the effective quark mass inside the bag amounts to $m_q/\bar{\chi} \approx 120 \text{ MeV}$. In effect the colour dielectric model interpolates between the constituent quark model $m_{\text{eff}} \approx 300 \text{ MeV}$ and the original MIT bag model $m_q = 0$.

The r.m.s. radius of the quark distribution is 1.2 fm, which is larger than the experimental r.m.s. radius of 0.83 fm. The magnetic moment of $2.66 \mu_N$ ($\mu_p = 2.79 \mu_N$) and the ratio of $g_A/g_V = 1.32$ are excellent results. Since $\bar{\chi}$ inside the bag is rather small it is clear that the shape of the effective potential $W(\chi)$ near $\chi = 1$ does not influence the nucleon

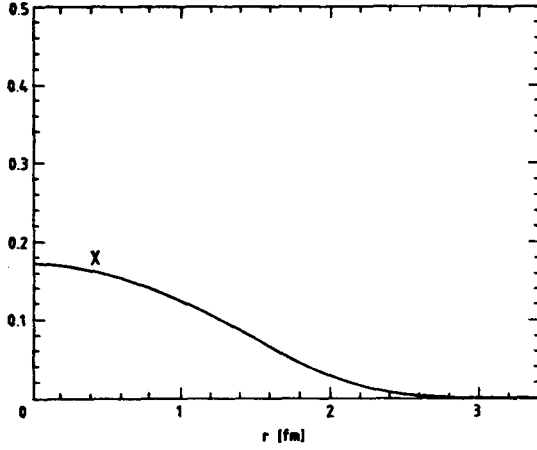


Fig. 5 The colour dielectric field $\chi(r)$ for the nucleon.

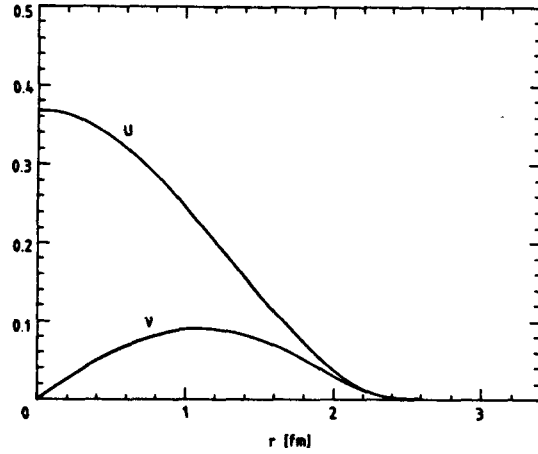


Fig. 6 The big $u(r)$ and small $u(r)$ components of the quark wave functions of the nucleon.

mass. Calculations with different sets of parameters B and A do not change the result as long as $m_{UB}^2 \sigma^2$ is kept constant. The shape of the effective potential W around $\chi = 0$ determines the dynamics. This is however different for the two phase solutions of Ref. 25, where in the interior of the bag $\chi \rightarrow 1$.

In the same model as used in Ref. 21 we have calculated ²⁶ soliton matter; i.e. we have looked for quark and χ -field solutions which are periodic at $r = r_0$, the size of a Wigner-Seitz cell. Let us recall that the nuclear matter density ρ and r_0 are related as

$$1 / \left(\frac{4\pi}{3} r_0^3 \right) = \rho$$

i.e. r_0 gives the radius of the free spherical cell available to a nucleon in nuclear matter. We take $(\partial\chi/\partial r)/r_0 = 0$ and $\partial(\psi^\dagger\psi)/\partial r = 0$ as new boundary conditions at the surface of the spherical cell. The fermion boundary condition allows two solutions, which correspond in the non-relativistic limit to the symmetric and antisymmetric solutions. They give the energies at the lower- and the upper-band edge. The shape of the quark densities is rather different for both cases. The lowest energy quarks are *strongly delocalized*; they would belong to deeply bound 'nucleons'. At the top of the band the quarks are *well localized*, even compressed in comparison to their free valence quark wave function. This model is certainly very crude, but it could tell us that the quark distributions in nuclei are much less obvious than a naïve interpretation of standard nuclear physics results can teach us. Interestingly enough most existing data cover only valence nucleons such as the famous mapping out of the $3s_{1/2}$ nucleons in Tl and Pb with electron scattering. In a schematic model ²⁷ similar to that described in our earlier paper, ¹⁷ we have also calculated the gluon cloud corrections to the nucleon properties in nuclear matter. The non-vanishing colour dielectric field allows the quarks to leak out of their bags reducing the kinetic energy of quarks, whereby nucleons get bound, and in addition it allows a gluon cloud outside the

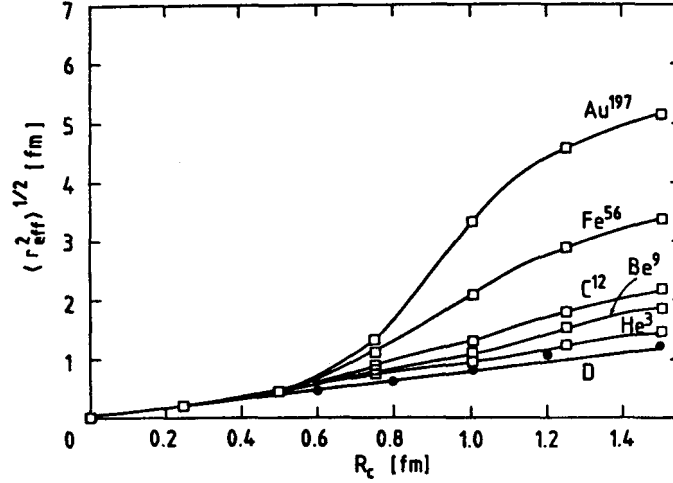


Fig. 7 The effective confinement radius $(r_{\text{eff}}^2)^{1/2}$ in nuclei as a function of the hard core size of the nucleon R_c .

bag. We have calculated in this more extensive model the nuclear Nolen-Schiffer anomaly. This anomaly is traced back to the deconfinement of up and down quarks in the nuclear environment which is different owing to their different masses. A crystal of nucleons (soliton matter) does not, of course, correspond to our understanding of nuclei as a liquid or a gas of nucleons. Therefore we have investigated ²⁸ what kind of structures develop when we allow a random distribution of the colour dielectric field with two values χ_{in} and χ_{out} characterizing the values of χ inside and outside the nucleons. Such a problem is analogous to percolation studies of electrical conductivity in amorphous materials. We have used such a model before to study the interaction radius in ³He. ²⁹ It has now been applied by many others to nuclei and deep inelastic scattering. Summarizing the results of this work, we show in Fig. 7 how the effective confinement zone (r_{eff}^2) increases with the hard-core nucleon radius R_c , inside which $\chi = \chi_{\text{in}}$. For the deuteron we see the physics of a dilute bound state, whereas for the heavier nuclei around $R_c = 1$ fm a transition takes place from the nuclei as colour insulators to colour conductors.

To conclude, I would like to summarize the three main theoretical points of this paper:

- i) Integration over high-frequency gluon configurations generates an effective colour dielectric model Lagrangian, characterized by a few coupling constants, which can be calculated numerically.
- ii) This Lagrangian generates bag-like solutions for the nucleon.
- iii) This Lagrangian allows a first attempt at determining quark wave functions in nuclei.

I think that progress along these lines in non-perturbative QCD will make experimental efforts done at institutions such as CEBAF accessible to a theoretical interpretation. Here I am mostly thinking of experiments done at the upper edge of the available energy in the scaling region with large x . Existing inclusive lepton-nucleus experiments in the intermediate x -region have given us a glimpse of quark degrees of freedom in nuclei.

Nuclear binding and a softening of the valence quark distribution are connected. New experiments at small and large x will illuminate the role of sea quarks, soft gluons, and the motion of quarks in nuclei.

I should like to thank the organizers of this Workshop for their kind invitation. Our work on colour dielectric lattice theory has greatly benefited from initial discussions with Michael Ilgenfritz. It has continuously been helped by Guido Martinelli, who also supplied us with an SU(3) version of the block transformation program.

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REVIEW OF THE NUCLEAR PHYSICS PROGRAM AT SLAC

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Abstract

A brief review is given of the recent activities and plans for future developments in the program of Nuclear Physics at SLAC. The SLAC facility offers unique opportunities to do electromagnetic nuclear physics experiments above one GeV that provide important data at large momentum transfer and shape ideas about the physics to be explored eventually at CEBAF.

The NPAS Program and the SLAC Experimental Facilities

The program of Nuclear Physics at SLAC began with a series of proposals by the American University Group starting in 1979 to add a new injector to use 1/6 of the two mile long linear accelerator to produce high intensity electron beams at energy from 1 to 5 GeV. In 1983 approval was obtained for construction of the new injector (called the Nuclear Physics Injector - NPI) and for the creation of a new program of nuclear physics measurements using the new beam.¹ Operation of the NPI began in late 1984 and the first experiment (NE3) took data in January 1985.

SLAC is presently the only place where a high intensity electron beam in the GeV energy range available for nuclear physics experiments until CEBAF is operating. The beam is generated with a pulsed structure with pulses 1.6 μ s long and a repetition rate of up to 180 pulses per second. At the maximum of 50 to 60 mA peak pulse current, the average current is about 15 μ A. This is 10 to 20 times more intense than can be obtained from the full SLAC linac when operated at energies below 5 GeV.

The NPAS program is presently conducted in the facilities of SLAC End Station A, which is a large experimental hall containing a complex of beam lines, spectrometers, shielding, cryogenic targets, particle detectors, electronics and computers. Most of this equipment was built for the electron scattering experiments in the SLAC High Energy Physics Program conducted over the past decade and can be readily turned to use in medium energy nuclear physics measurements. The principle spectrometers use by the NPAS program are the 1.6 GeV/c and the 8 GeV/c spectrometers. These are small aperture (1.8 msr and 0.8 msr respectively) devices of modest resolution (0.01 % in $\Delta p/p$). They are

* This research was supported by NSF Grant PHY85-10549 (American University), and DOE Contract No. DE-AC03-76F00515 (SLAC).

designed with sufficient bending power and shielding for detecting at most one high energy particle per pulse in the environment of high room background of soft photons and neutrons produced by the high powered beam (up to 60 kW) on thick targets (up to 2 g/cm²). The detectors available include lead glass and gas Cherenkov detectors for identifying scattered electrons, multiwire proportional chambers for particle tracking, and scintillator hodoscopes for triggering and time-of-flight measurements. Targets of liquid hydrogen and deuterium, gaseous helium, and various solid materials able to withstand the heat and radiation generated by the high current beam can be readily produced by the SLAC target Group. In addition to these standard facilities, End Station A offers a large experimental area that can be used for special experimental arrangements. One recent NPAS Experiment (NE4) used a large collection of quadrupole and dipole magnets and shielding blocks to construct a special 180° spectrometer system on the open floor space.

The End Station A facilities are now available essentially full time for nuclear physics experiments. Beam is delivered according to the constraints of the overall SLAC schedule and the available NPAS budgets. For the last two years there have been about 6 weeks of beam time per year. The Nuclear Program Advisory Committee meets approximately twice per year to make program decisions. To date there have been 13 proposals to the NPAS program and four experiments have been completed. Two experiments are currently in the development stages.

Recent and Planned Experiments

Completed NPAS Experiments

NE1—"Electroproduction of the Delta Isobar in Nuclei," by a collaboration from Argonne and Northwestern, H. Jackson, (Argonne) spokesman.

This experiment measured delta electroproduction in carbon and iron at momentum transfer around 0.1 (GeV/c)² to look for modifications due to the nuclear medium. The data included a rough longitudinal-transverse separation using the 1.6 GeV/c and the 8 GeV/c spectrometers. The collaboration, in conjunction with the experimenters from NE5, reinstrumented the 1.6 GeV/c spectrometer with a new gas Cherenkov counter, scintillator trigger counters, and a set of drift chambers for tracking. Data taking was completed in February 1986, and results are not yet available.

NE3 - "Inclusive Scattering from Nuclei," by a collaboration from the University of Virginia, Caltech, and Basel, D. Day (U.Va) spokesman.

The purpose of this measurement was to obtain the response function of several nuclei in the region of the inelastic threshold and quasi-elastic peak out the largest momentum transfer accessible with the SLAC beam. This was the

first NPAS experiment to obtain data with the new NPI beam. The 8 GeV/c spectrometer, with detectors recently upgraded by the high energy program, and a target array from a previous experiment were used. Data taking began in January 1985 and was completed in a second installment in January 1986.

Data were obtained scattering from ^4He , C, Al, Fe, and Au targets spanning six orders of magnitude in cross section. The data are analysed assuming quasi-free scattering from moving nucleons in terms of the nuclear scaling variable y , which is approximately the component of momentum of the struck nucleon along the direction of momentum transfer. The cross sections analyzed in terms of the y variable yield structure functions $F(y)$, which are interpreted roughly as the probability to find a nucleon at momentum y in the nucleus. Previous experiments on deuterium and helium at high Q^2 have shown that the data yield a universal function $F(y)$ for each nucleus approximately independent of momentum transfer. The new data will extend our knowledge of $F(y)$ to higher Q^2 and to more nuclei, and will give important information about momentum distributions in nuclei. An example of the preliminary results² for ^{12}C are shown in Figure 1.

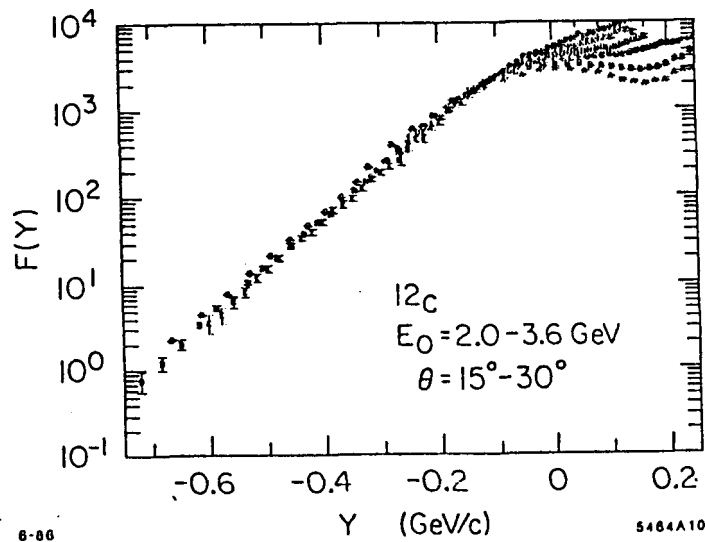


Figure 1. Preliminary results (Ref. 2) from NPAS Experiment NE3 on ^{12}C inelastic scattering analyzed to give $F(y)$ versus the scaling variable y . The data were taken at angles 15° to 30° and beam energies 2 to 3.6 GeV, corresponding to momentum transfer 0.2 to $2.2 (\text{GeV}/c)^2$.

NE4 –“Electron scattering from deuterium at 180° at Large Momentum Transfer,” by a collaboration from American University, Georgetown University, Lawrence Livermore National Laboratory, University of Massachusetts, SLAC, and

Tel Aviv, P. Bosted (AU) spokesman.

The primary goal of this experiment was to extend the measurement of the deuteron magnetic form factor $B(Q^2)$ out to the highest possible Q^2 . Another goal was to measure threshold and quasi-elastic scattering at backward angles to determine the inelastic structure function $W_1(\nu, Q^2)$. Because the magnetic form factor $B(Q^2)$ was expected to be much smaller than the forward angle function $A(Q^2)$, it was necessary to measure at 180° to determine $B(Q^2)$. In addition, due to resolution smearing in thick targets, the recoil deuterons must be detected in coincidence to pick out the elastic signal. This required the construction of a special double arm 180° spectrometer system, shown in Figure 2.

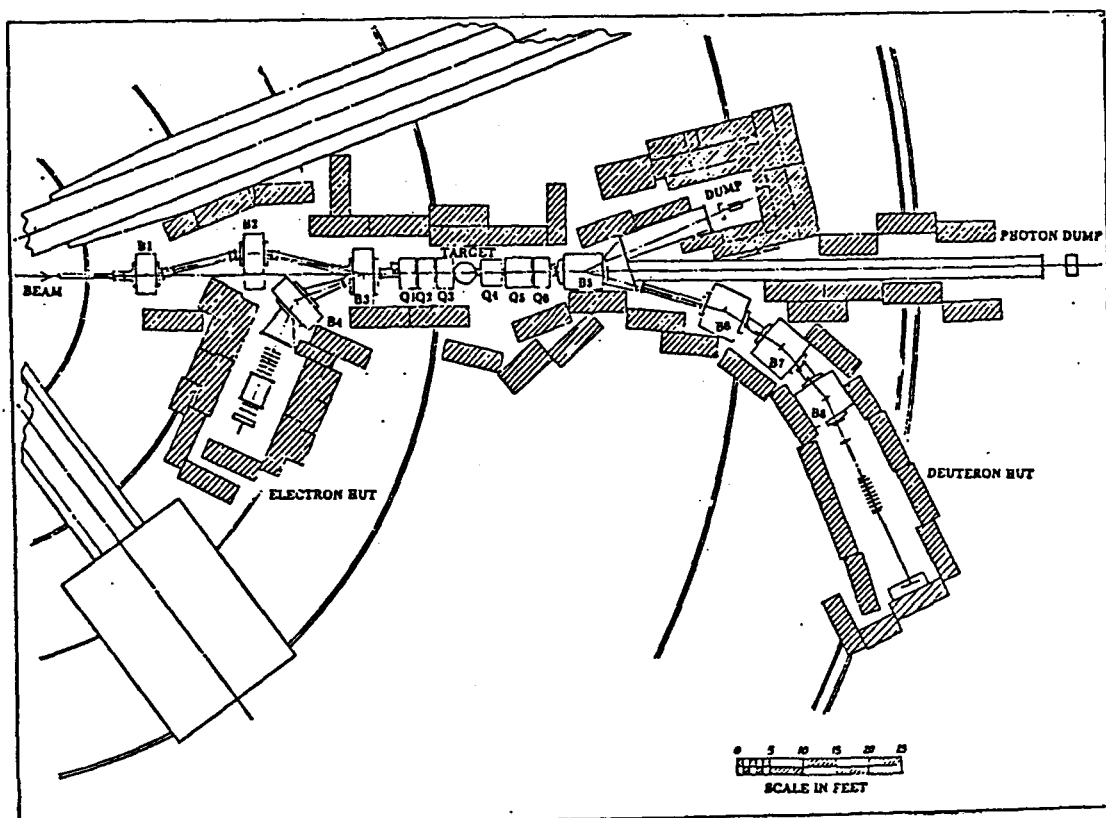


Figure 2. Experimental layout for Experiment NE4 on electron-deuteron scattering at 180° . A double arm spectrometer system was constructed on the floor of End Station A downstream of the traditional target location. The incident beam entered through chicane magnets B1 and B2 and was transported to the target through splitting magnet B3 and the quads Q1-Q3. The electrons scattered backwards were gathered up by quads Q1, Q2, Q3 and transported back through the bending magnets B3 and B4 into the electron detectors. The recoil deuterons were transported in the forward direction through a quad triplet

and into a spectrometer composed of 4 bend magnets B5-B8. The transmitted electron beam was split off by magnet B5 into a well shielded dump.

The first installment of data was obtained in June-July 1985, and the final run was completed in June 1986. Preliminary results³ from the 1985 data are shown in Figure 3. The new data extend $B(Q^2)$ out to $2.5 (GeV/c)^2$ and suggest the presence of a diffractive feature in that region. More data points were taken in 1986 in the region of the dip around $2 (GeV/c)^2$ and out to a maximum momentum transfer of $2.7 (GeV/c)^2$ that confirms the diffraction feature and sharpens the definition of its shape. The discovery of a diffractive shape for $B(Q^2)$ at such high momentum transfer will sharply distinguish between competing models of deuteron structure.

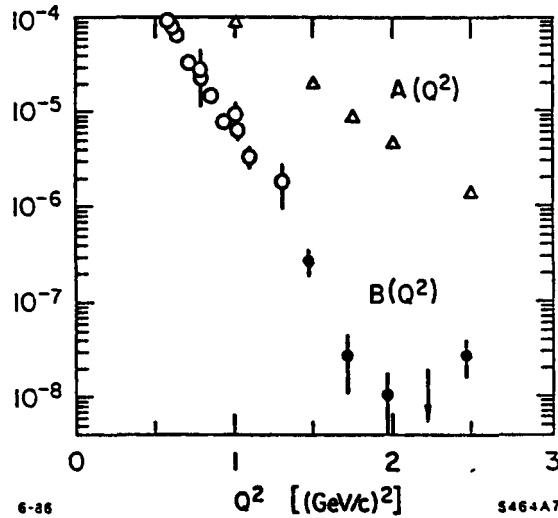


Figure 3. Preliminary results (Ref. 3) for $B(Q^2)$ from experiment NE4 and previous data for $A(Q^2)$.

NE5 – “Delta electroproduction in Nuclei,” by a collaboration from the University of Virginia, Caltech, Florida State, and CEBAF, S. Thornton (U.Va.) spokesman.

This experiment concentrated on the nuclear dependence and momentum transfer behavior of delta production in nuclei. It ran as a companion experiment to the similar measurement by NE1. Measurements were made on ^1H , ^4He , C, Fe, and W mostly at 37.5° and energies from 0.9 to 1.5 GeV using the 1.6 GeV/c spectrometer. The NE5 collaboration worked together with the NE1 group to upgrade the detector package of the spectrometer. Data taking was completed in February 1986 and results are not yet available.

NPAS Experiments in Progress

NE8 – “Two-Body Photodisintegration of the Deuteron,” by a collaboration from Argonne, Northwestern University, American University, Stanford, and the University of Virginia, R. Holt (Argonne) spokesman.

Disintegration of the deuteron into a proton-neutron final state will be measured out to photon energies up to 1.8 GeV using the upper end of a bremsstrahlung beam. Recoil protons will be detected over a range of angles from 35° to 135° in the 1.6 GeV/c spectrometer. The detector configuration is being modified to include an aerogel Cherenkov counter and better time-of-flight counters for proton detection. Preparation of the target and spectrometer equipment are underway for a data run sometime within the next year.

NE9 – “The Transverse and Longitudinal Response Functions for Several Nuclei at Momentum Transfer near $Q^2 = 1 \text{ (GeV/c)}^2$,” by a collaboration from the University of Virginia, Florida State University, Basel, Caltech, Saclay, CEBAF and Stanford, Z. Meziani (Stanford) spokesman.

The apparent suppression of the longitudinal response of nuclei in the quasi-elastic region seen in previous data⁴ up to Q^2 around 0.5 (GeV/c)^2 has attracted a lot of attention. It is thought that this effect could be a signature for loss of strength in the one-body electromagnetic process due to correlations in nuclei. It has also been noted that such an effect might be due to modification (swelling) of nucleons in nuclei. The aim of this experiment is to extend the separation of longitudinal and transverse cross sections out to larger momentum transfer.

The experiment will use the 8 GeV/c spectrometer to do a precise Rosenbluth separation at Q^2 around 1 (GeV/c)^2 in ^3He , ^4He , C and Fe. Construction of targets and calibration of the spectrometer acceptance is in progress in preparation for a data run within the next year.

Active Proposals

NE11 – “Proposal to Separate the Charge and Magnetic Form Factors of the Neutron at Momentum Transfer $Q^2 = 2$ to 5 (GeV/c)^2 ,” by a collaboration from American University, Lawrence Livermore National Laboratory, University of Massachusetts, SLAC, and Tel-Aviv, R. Arnold (AU) and P. Bosted (AU) spokesmen.

This is a proposal to test the hypothesis put forward recently by Gari and Krümpelmann⁵ that the neutron electric form factor G_{En} could be larger than the magnetic form factor G_{Mn} at large Q^2 . The experiment would use the 8 GeV/c spectrometer to measure quasi-elastic scattering from deuterium and elastic electron-proton scattering over a wide range of angles. The contribution from quasi-elastic scattering from the neutron would be extracted from the area of

the quasi-elastic peak by comparison to the (smeared) proton data. A separation of the neutron form factors at large momentum transfer would be important for testing quark models of nucleon structure and is needed as input to nuclear structure calculations.

Ne12 – “Proposal to Measure Elastic Electron Tritium Scattering at Large Momentum Transfer,” by a collaboration from American University, Lawrence Livermore National Laboratory, SLAC, Tel-Aviv, MIT, and Georgetown University, F. Dietrich (LLNL) and S. Rock (AU) spokesmen.

This proposal outlines a plan for testing leading to construction of a novel tritium target made of the compound LiT that would be used to measure the tritium elastic form factors F_{ch} and F_{mag} . The measurements would be made in specially constructed double arm spectrometers, one based on the 180° spectrometer used in NE4 to measure F_{mag} , and another one at forward angles for F_{ch} . The tritium form factors at large momentum transfer are fundamental quantities needed to continue progress in understanding the three nucleon system.

Recent High Energy Physics Experiments in End Station A.

The high energy physics program which operated in End Station A for more than a decade has wound down in the past year as the activities in high energy physics at SLAC have moved to other areas. The NPAS experiments over the past two years have operated alternately with the high energy experiments, and there has been a fruitful interplay between these activities. The experiments were carried out in some cases by the same physicists, and they are exploring similar questions in physics – probing the internal structure of nucleons and nuclei at high energy. The following is a description of the recent high energy experiments that have overlapped with the NPAS program and contributed to the development of the experimental equipment and the stimulation of physics ideas.

E136 – “Measurement of Elastic Electron Scattering from the Proton at High Momentum Transfer,” by a collaboration from American University, SLAC, Georgetown University, and Saclay, R. Arnold (AU) spokesman.

This was a measurement⁶ of elastic ep scattering with good precision over the Q^2 range from 2 to $31.3 (GeV/c)^2$ from which the proton magnetic form factor G_{Mp} was extracted. The experiment was motivated by the prediction of perturbative QCD that G_{Mp} at large Q^2 should fall with a power of $1/Q^4$. Deviations from this form would be expected from the additional Q^2 dependence of the strong coupling constant $\alpha(Q^2)$, and from the influence of the poorly known valence quark wave functions. A central open question in understanding hadron structure is: At what Q^2 do the contributions from soft nonperturbative processes die away and the perturbative QCD regime begin? The previous data

for the proton at large momentum transfer was not precise enough to test for deviations from $1/Q^4$ behavior. The new experiment, carried out using the 8 GeV/c spectrometer, obtained enough data to dramatically improve the precision at high Q^2 . The results in Figure 4 show G_{Mp} falling slightly faster than $1/Q^4$ for Q^2 above 10 $(\text{GeV}/c)^2$. This is consistent with the onset of the perturbative regime around 10 $(\text{GeV}/c)^2$, and leaves room for some Q^2 dependence from $\alpha(Q^2)$ and from the quark wave functions.

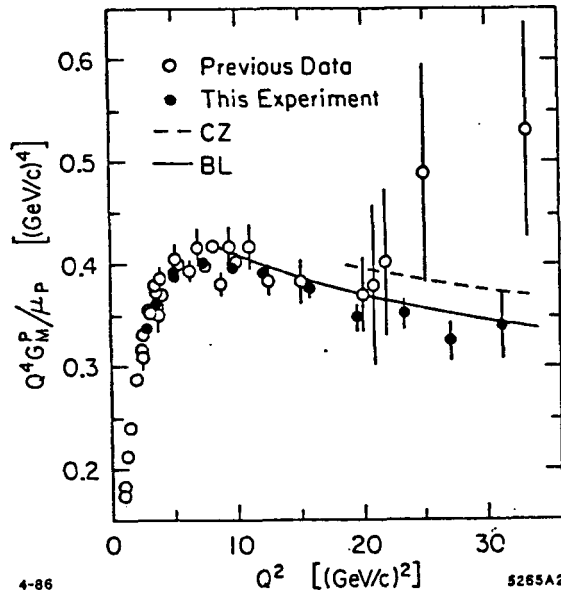


Figure 4. Results for $Q^4 G_{Mp}$ from experiment E136 (Ref. 6) along with perturbative QCD curves from Brodsky and Lepage (Ref. 7) and Chernyak and Zhitnitsky (Ref. 8).

There is now vigorous activity by various people trying to understand the interrelated questions of the applicability of perturbative QCD, the form of the proton wavefunctions, and the variation of the virtual photon-nucleon interaction versus momentum transfer.^{5,8,9} Understanding these issues for nucleon structure is vital to understanding nuclear structure at high momentum transfer and the answers will influence our ideas about the role of quark degrees of freedom at CEBAF energies.

E139 – “Measurement of the A-Dependence of Deep Inelastic Scattering,” by a collaboration from American University, SLAC, Saclay, Georgetown, University, University of Rochester, R. Arnold (AU) spokesman.

This measurement¹⁰ of the A-dependence of deep inelastic scattering in nuclei from helium to gold compared to that in deuterium was prompted by the EMC

discovery¹¹ of a modification of the cross section for high energy muon scattering on iron. The experiment showed convincingly that the EMC effect exists and its size increases with A . The data was taken on short notice using the 8 GeV/c spectrometer as instrumented for the E136. This data is widely regarded as one of the key pieces of evidence for modification of the quark distributions for nucleons embedded in nuclei. Observation of this effect has prompted a prolific literature describing a wide range of models. The discrepancy between the original EMC data and the SLAC data at x below 0.3 has apparently been cleared up. Preliminary results of new data¹² from the (slightly modified) EMC apparatus now agrees with the E139 data at low x , thus eliminating the need for models to explain a large value for the ratio of iron to deuterium at low x . There are still many open questions, particularly the x , Q^2 , and A dependence at x below 0.3, that hopefully will be addressed in future experiments at CERN and Fermilab. A major outstanding problem is to obtain a reliable separation of longitudinal and transverse cross sections for deep inelastic scattering on the nucleons and on nuclear targets. This is the subject of the following experiment.

E140 – “A Measurement of the x , Q^2 and Nuclear dependence of $R = \sigma_L/\sigma_T$,” by a collaboration from American University, University of Rochester, Caltech, Lawrence Livermore National Laboratory, University of Massachusetts, Tel-Aviv, Fermilab, SLAC, and Stanford, S. Rock (AU) and A. Bodek (Rochester) spokesmen.

For more than a decade there has existed a need for high quality separations of σ_L and σ_T in the nucleons (proton and deuteron) for purposes of testing QCD and ideas of nucleon structure. Following the discovery of the EMC effect, and the E139 results showing the A -dependence, it became important to separate the structure functions in heavy nuclei as well. The longitudinal cross section for deep inelastic scattering is expected to be small due to helicity conservation in the interaction of vector photons with the (massless) spin-1/2 quarks. Other interactions, such as scattering coherently from diquarks (quasi-pions), and QCD effects from quark transverse momentum, could enhance the longitudinal cross section. Some models of the EMC effect propose an enhancement of the pion content in nuclei, which might show up as a larger value for R in a nucleus.

The new experiment E140 measured R at a few x and Q^2 points for hydrogen, deuterium, iron, and gold. The experiment was conducted in the 8 GeV/c spectrometer following an upgrade of the lead glass shower counter and the gas Cherenkov counter used for pion-electron discrimination. Data taking was completed in January 1986, and so far only preliminary results for iron and deuterium are available.¹³ When the full results are known, they should help distinguish between the wide variety of models proposed for the EMC effect.

E141 – “An Electron Beam Dump Experiment for Light Short-lived Particles,” by a collaboration from University of Rochester, American University, Caltech, Fermilab, Georgetown, University of Massachusetts, and SLAC, E. Riordan (Rochester) spokesman.

This experiment¹⁴ was proposed, prepared, and completed within the space of a few months. It was designed to search for the signature of a new type of elementary particle (axion) suggested to explain the curious peaks observed in electron-positron production from heavy ion collisions at the GSI laboratory in Darmstadt. The 8 GeV/c spectrometer was arranged to detect high energy positrons that would emerge from a thin beam dump if the axions were produced by an incident electron beam and then subsequently decayed to electron-positron pairs after the dump. No signal of positrons above that expected from punch through of traditional electromagnetic shower particles was observed, and stringent limits on the mass and lifetime of possible new particles were obtained. The new limits eliminate the standard axion as an explanation for the phenomena observed at the GSI.

Future Developments

The SLAC facilities will continue for some time to provide opportunities for electromagnetic nuclear physics measurements at energies above one GeV. Several projects in various stages of planning could extend this capability to new areas.

A project to generate polarized electrons from the Nuclear Physics Injector has been discussed for several years, and is now receiving more active attention. A new electron gun using a laser driven photo cathode similar to that being developed now at SLAC for the Single Pass Collider would be used. Polarized electrons would be used in a series of measurements of spin transfer and parity effects in nuclear targets.

The possibilities for a new type of experiment using the high current stored beam in the PEP ring to interact with low density gas targets are now being vigorously pursued. A short test was conducted during the last run cycle in which nuclear electromagnetic interactions were measured using a distributed gas jet target in the TPC detector at one of the interaction regions on the ring. (See the talk at this workshop by K. Van Bibber.) The combination of high circulating current and low density targets offers possibilities for unique experiments, for example, using polarized targets, and detecting and identifying many particle is the final state. The physics and technical opportunities for internal target experiments at electron rings will be explored in detail at a workshop at SLAC in January 1987.

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PRESENT STATUS AND EXPERIMENTAL PROGRAM FOR THE ELECTRON STRETCHER RING ELSA

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A concise survey of the Bonn activities in the intermediate energy regime is presented. It is focussed on three topics: the Stretcher Accelerator, some Perspectives in Physics, and the Experimental Facilities.

INTRODUCTION

Since 1958, when the 500 MeV alternate gradient electron synchrotron came into operation, our institute has a continuous tradition in photon induced experiments. These experiments have started in 1967. In the late 1970's it became obvious, that a steady continuation of the experimental line required an improved accelerator and new experimental facilities. There were mainly two demands. The electron scattering experiments as well as real photon reactions with multiparticle final states needed a crucial increase of the duty cycle for counting rate limitations due to accidentals. A second demand was to reach higher energies for investigations on the formation and decay of heavy baryon resonances with high spins and for the excitation of vector mesons. In 1979 we submitted a proposal to the authorities in order to build a new laboratory and the Electron Stretcher Accelerator ELSA. Due to the bad financial situation this proposal had to be reduced to the machine itself and the tunnel. Now, in 1986, the construction is close to completion, and we hope to accelerate the first electrons within the next months.

1. THE ELECTRON STRETCHER ACCELERATOR

THE CONSTRUCTION

The ELSA ring is constructed for the storage of electrons up to 2.4 GeV, which have been accelerated in the synchrotron and, in addition, for accelerating electrons up to 3.5 GeV. It has a separate function magnet system, which provides radiation damping of the horizontal betatron oscillations and allows a wide range tunability of the betatron working point. The latter point is of importance for the resonant extraction. The ring has long straight, dispersion free sections, which are used for the RF cavities, the injection, one of the extractions and for wigglers, which improve the beam quality for synchrotron radiation experiments. ELSA has 16 FODO cells, each 10.28 m long, from the 32 half cells only 24 contain a bending magnet, which is inserted asymmetrically between the F- and D-quadrupoles. This yields straight sections in the arc cells. The whole configuration, together with the experimental facilities is displayed in fig. 1. ELSA has a very regular lattice with almost equal horizontal focussing power per cell, from which a very small variation of the radial β -function maxima results. In addition, since the bending magnets have a rectangular shape, they do not destroy this behaviour. Further elements are 8 sextupoles for the chromaticity compensation, 4 quadrupoles and 4 sextupoles for the extraction, and additional septum magnets.

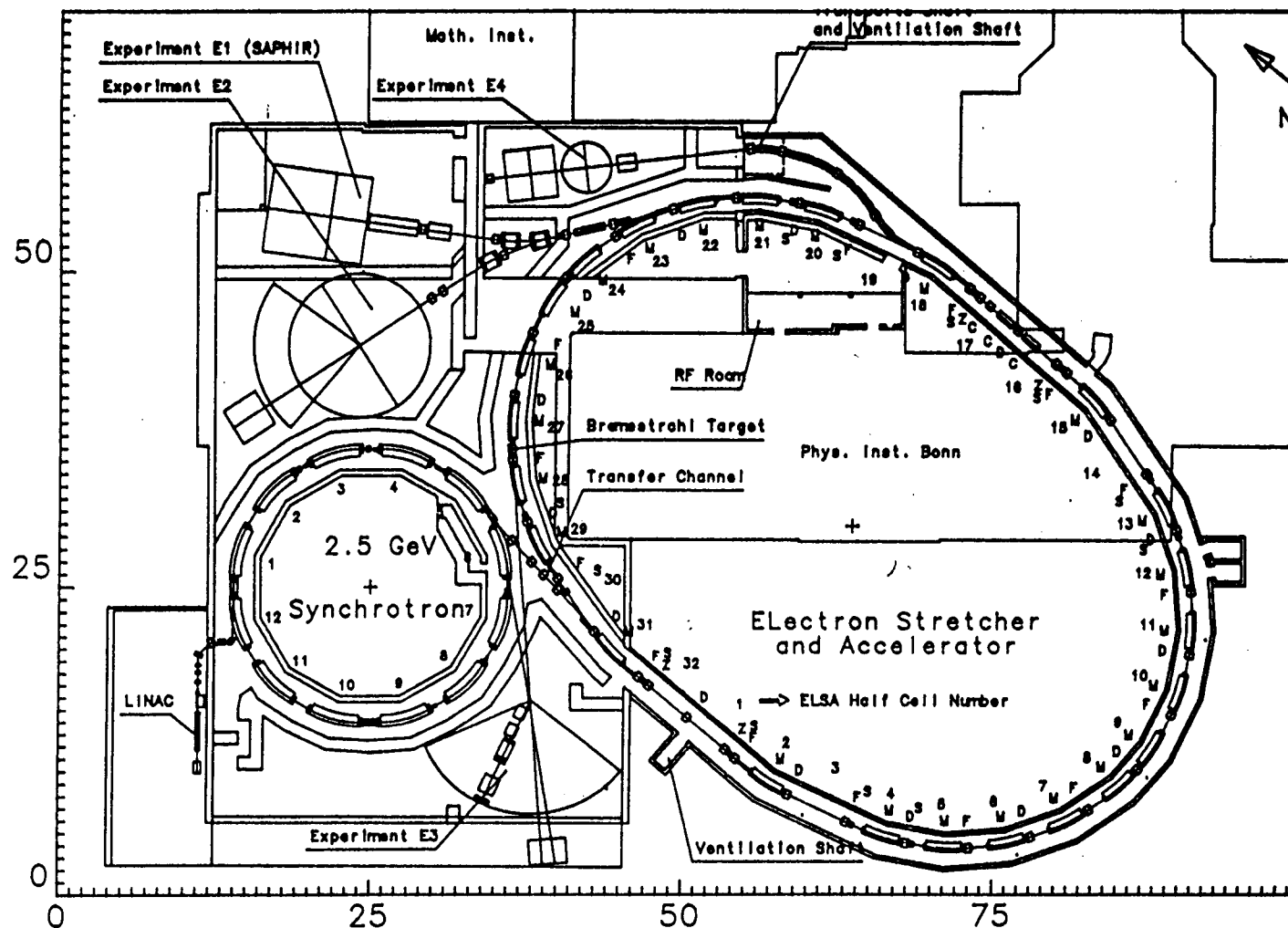


Fig. 1: Floor map of the Bonn facilities. E1 - SAPHIR, E2 - electron scattering experiment, E4 - PHOENICS

TABLE I: LATTICE AND OPTICS DATA

bending radius	10.84	m
circumference	164.4	m
number of FODO-cells	16	
length of FODO-cells	10.28	m
number of dipoles	24	
number of quads	32	
length of dipole core	2.84	m
length of quad core	0.45	m
<u>at 3.5 GeV and $Q_x = Q_z = 5.4$:</u>		
magnetic field strength	1.08	T
strength of F-quad	0.78	m ⁻²
strength of D-quad	0.75	m ⁻²
momentum compaction factor	4.6	%
damping times: horizontal	3.1	msec
vertical	3.1	msec
energy	1.6	msec
maximum horizontal β -fct.	22.2	m
minimum horizontal β -fct.	1.54	m
maximum vertical β -fct.	24.6	m
minimum vertical β -fct.	1.26	m
maximum dispersion fct.	3.42	m
natural energy width	0.09	%
natural horiz. emittance	0.76	mm x rad
natural chromaticity: horizontal	-9.1	
vertical	-8.4	

OPERATION MODES

ELSA can be operated in two different modes, the 'Stretcher Mode' for energies below 2.4 GeV, and the 'Accelerator Mode' with energies up to 3.5 GeV. In the first case the RF power is only needed to balance the synchro-

tron radiation losses. In this mode electron pulses from the synchrotron are injected at a maximum rate of 50 Hz (for low intensity requirements it can be reduced to 2 Hz). Immediately after injection electrons are extracted at constant rate for 20 ms or longer. It takes 1 ms to switch from injection to extraction, so the duty cycle is 95%. In this mode the internal beam intensity is 12 to 20 mA, the extracted current 200 to 400 nA.

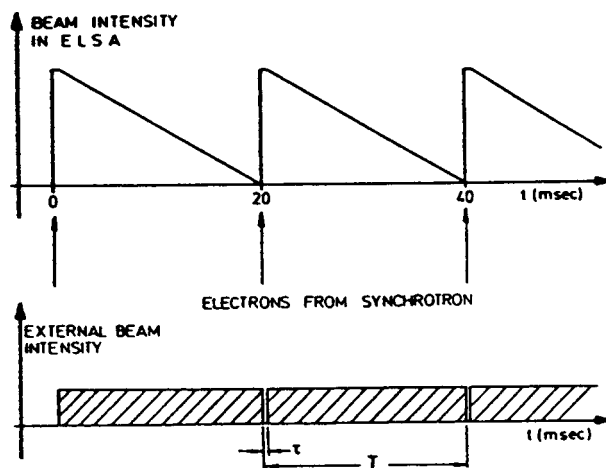


Fig. 2: Internal and external current of ELSA in the stretcher mode.

In the accelerator mode electrons are transferred to ELSA at 1.8 GeV, thus the field changes about a factor two at maximum. This demands 150 ms ramping time in both directions. In order to obtain a high duty cycle the extraction is done during a long flat top. At maximum energy 500 ms are foreseen, corresponding to a 60% duty cycle. Here the external current is about 10 to 20 nA. Higher intensities up to 100 nA are possible at a decreased duty cycle (see fig. 3).

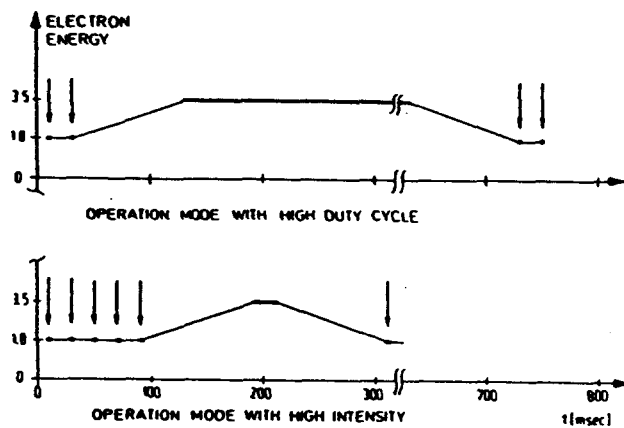


Fig. 3: Electron energy in the accelerator mode.

THE INJECTION

Since the circumference of the ELSA ring is 2.35 times larger than that of the synchrotron, a multiturn extraction from the synchrotron is necessary. We hope to get a homogeneous filling with a three turn extraction. The extraction from the synchrotron with kicker- and septum magnets is quite conventional, so a transfer efficiency of 50% seems feasible.

THE EXTRACTION

The extraction is one of the most serious problems at ELSA. A slow resonant method will be applied. To achieve a stable and controlled operation the working point is moved towards a third integer betatron resonance. The shrinkage of the stable phase space area is obtained via four additional quadrupoles at constant sextupole strength, then the electrons are kicked out of the hori-

zontal closed orbit plane using four extra sextupoles and two successive septum magnets. Calculations for this scheme predict an efficiency of 75% /1/.

THE RF SYSTEM

Two different RF systems will be used depending on the operation mode of ELSA. In the 'Stretcher Mode' only a moderate RF power between 500 W and 40 kW is required, here a 40 kW transmitter and a DORIS type cavity with low shunt impedance will be utilized. At the lower energies this system works more economical and is better matched to the beam stability requirements. In the 'Accelerator Mode' a 250 kW transmitter together with two PETRA type cavities will be used.

POLARIZED ELECTRONS

At Bonn the first polarized electron beam has been accelerated up to above 2 GeV in 1978 already /3/, but the intensity was so low, that no experiments have been performed. In the meantime, however, a new GaAs source on the basis of the photoeffect with polarized laser light has been developed and tested. With this source we hope to get an extracted current of 7 nA /4/.

BEAM LINES

ELSA has two external beams for electrons up to 3.5 GeV and 2 GeV, respectively. The first one will provide electrons to the e^- -scattering facility and to SAPHIR, the second one supplies the PHOENICS experiment. Both lines are foreseen to handle polarized electrons too.

STATUS

All magnets are in place and have been powered, the vacuum chamber is being installed. The RF system for the 'Stretcher Mode' is ready, the second one will come in during spring 1987. The ELSA injection as well as the first beam line are under installation, they will be ready in Sept. 86. We hope to have the first internal beam by the end of 1986.

For references concerning ELSA see /5/.

2. PERSPECTIVES IN PHYSICS AT ELSA

In the intermediate energy region structure seems to be mingled nearly inextricably with dynamics. Therefore, up to now, all QCD motivated models are in rather serious difficulties in this region. All the more a very precise experimental input is needed to check and improve the theories. Some progress in this direction seems possible with two kinds of experiments, precise measurements of spectroscopic data over a wide kinematical and energetic range and investigations on specific questions of dynamics. We hope to contribute our part to these problems with the new generation of experiments at ELSA. In this report emphasis is put mainly on such experiments that can be performed within a first period of three to five years.

ELECTRON SCATTERING EXPERIMENTS

One of the points of main effort at Bonn during the last years has been the investigation of the deuteron structure. Here electromagnetic probes can deliver complementary information to results from experiments with

strong interacting particles, like protons or pions; in addition the pointlike lepton probe is advantageous for the interpretation of the results. Of special interest in this energy region is the transition from conventional nuclear models to contributions from the quark degrees of freedom inside the nucleus. The main goal of our experiments is to disentangle contributions from meson exchange currents, isobar excitation, and from possible multiquark clusters to the deuteron wave function. The deuteron is especially suited for such investigations, because of its principally simple structure, which, however, is far from being understood in detail.

The first e^- -scattering experiments at ELSA will, using the improved experimental conditions, continue the investigations on the deuteron electrodisintegration and on the elastic scattering measurements on tensor-polarized deuterons.

DEUTERON DISINTEGRATION $e + d \rightarrow e' + p + n$

Kinematic variables:

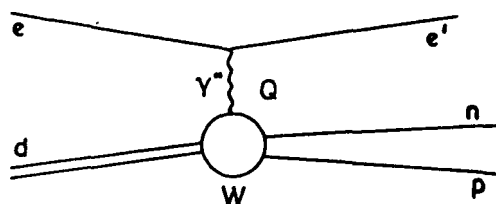


Fig.4: Deuteron electrodisintegration.

$$Q^2 = (e - e')^2, \quad \nu = E - E',$$

$$\epsilon = [1 + 2 \cdot (1 + \nu^2/Q^2) \cdot \tan^2(\theta_e/2)]^{-1}$$

measured observables:

$$\frac{d\sigma}{d\Omega dE' d\Omega_p} = r_t \left[\frac{d\sigma_u}{d\Omega_p} + \frac{d\sigma_L}{d\Omega_p} + \frac{d\sigma_T}{d\Omega_p} \cos 2\phi + \epsilon (\epsilon + 1) \frac{d\sigma_{TL}}{d\Omega_p} \cos \phi \right]$$

For a complete information the separation of the four terms is necessary for fixed parameters Q^2 , W and θ_p^* . First measurements at Bonn have been performed at $Q^2 = .23 \text{ GeV}^2$, $\theta_p^* = 80^\circ$, and in a W range from 1.965 to 2.335 GeV /6/, they will be extended to other proton

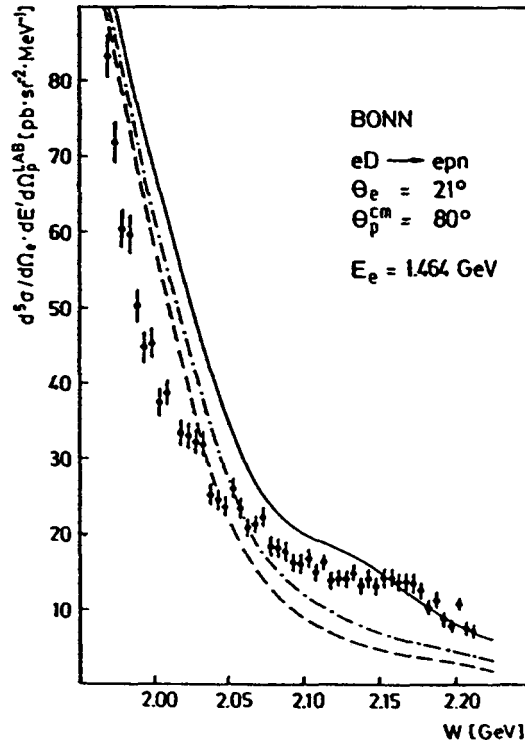


Fig. 5: Electrodissintegration of the deuteron at large excitation energies. Predictions from Arenhövel et al. /2/. Dashed line: without MEC or IC, dash-dot: with MEC, solid line: MEC and IC.

CM-angles in order to separate MEC and IC contributions. Additionally the Q^2 dependence will be investigated to get information about the off-shell behaviour of the nucleon formfactors and the $\gamma^*N\Delta$ transition formfactor. First results from the above mentioned experiment indicate a stronger Δ excitation in the deuteron than predicted (see fig. 5).

ELASTIC ELECTRON SCATTERING ON POLARIZED DEUTERONS

$e + d \rightarrow e' + d$

In electrodisintegration experiments /6/, /7/ the long range meson exchange currents coupling to the isovector part of the photon play an important role. In contrast to that, only the isoscalar coupling is of significance for the elastic e-d scattering, such providing supplementary information about the deuteron. One of the most interesting specific questions is the determination of the three elastic formfactors F_C , F_Q and F_M . The magnetic formfactor F_M has been measured at Bonn /8/ for Q^2 values between 0.5 and 1.3 GeV^2 as well as in many other laboratories. The results favour models, which include MEC contributions. The electric monopole formfactor $F_C(Q^2)$ is, besides the determination of the MEC contributions, of special interest in the search for effects from multiquark states. Even a small 6 q-admixture to the wave function might influence dramatically the F_C behaviour around $Q^2 = 0.6 \text{ GeV}^2$ /9/. First experiments have been performed at MIT-Bates /10/ for $Q^2 = 0.12$ and 0.16 GeV^2 by measuring the recoil tensor polarization T_{20} of the proton, and at Bonn for $Q^2 = 0.5 \text{ GeV}^2$. These measurements will be continued at ELSA under improved conditions.

Measured Observables:

Differential cross-sections for polarized (σ_{MEAS}) and unpolarized target (σ_{NP}).

$$\frac{d\sigma_{\text{NP}}}{d\Omega} = \sigma_{\text{NS}} \left[A(Q^2) + B(Q^2) \tan^2 \theta_e / 2 \right]$$

$$\text{with } A(Q^2) = F_C^2 + \frac{8}{3} \eta^2 F_Q^2 + \frac{2}{3} \eta F_M^2,$$

$$B(Q^2) = \frac{4}{3} \eta(1+\eta) F_M^2, \quad \eta = \frac{Q^2}{4m_d^2}$$

$$\frac{d\sigma_{\text{pol}}}{d\Omega} = \sigma_{\text{NS}} \left[A_T(Q^2) + B_T(Q^2) \tan \theta_e / 2 \right]$$

$$A_T(Q^2) = F_C^2 + \frac{2}{3} \eta F_C F_Q + \frac{10}{9} \eta^2 F_Q^2 + \frac{1}{2} \eta F_M^2$$

$$B_T(Q^2) = \frac{3}{2} \eta(1 + \eta) F_M^2$$

The cross-section of the elastic e-scattering on a tensor polarized ($T \approx 20\%$) target (σ_{MEAS}) is measured with the target vector polarization being perpendicular to the direction of the virtual photon in the scattering plane.

The measured cross-section σ_{MEAS} is connected to the cross-section for a 100% polarized target and the non polarized cross-section by the relation:

$$\sigma_{\text{MEAS}} = (1 - T) \cdot \sigma_{\text{NP}} + T \cdot \sigma_T,$$

where T is the tensor polarization strength,
or

$$\frac{\sigma_{\text{MEAS}}}{\sigma_{\text{NP}}} = T \cdot \left(\frac{\sigma_{\text{T}}}{\sigma_{\text{NP}}} - 1 \right) + 1$$

The formula allows a direct comparison between measurements and predictions. In fig. 6 calculations with different contributions are to be seen.

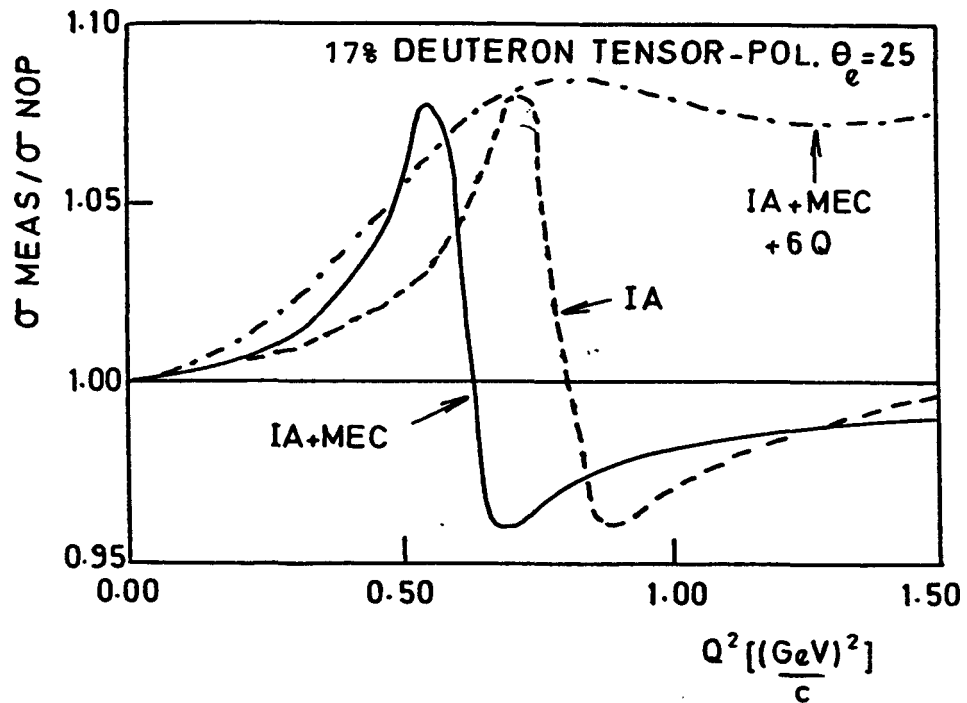


Fig.6: Predictions for GD_C with MEC and 6-quark (9) contributions.

At ELSA the following improvements are foreseen:

1. The cryostat power of the refrigerator will be increased from 20 mW to 200 mW. This allows an increase of the used beam current from 1 nA to 10 nA. The field strength of the target magnet will be slightly increased to 4 T.

2. For the electron detection under large scattering angles, a non magnetic shower counter detector system, covering a large solid angle will be built.

ELECTRON SCATTERING FROM NUCLEONS

During the last years, the investigation on excited nucleons has been an other topic of experimental activities at Bonn. One goal of these measurements is the determination of transition formfactors for the different resonant states. In a series of experiments /11/ distinct final states have been examined, and the dominant multipole amplitudes have been determined. These investigations will be continued up to a four momentum transfer $Q^2 \approx 3 \text{ GeV}^2$. Within the frame of single quark transition models /12/ certain relations between the resonance amplitudes of multiplet members are predicted, which can be tested by the experiment. In addition, such measurements can discriminate dynamical quark models /13/. Contrary to the situation in photoproduction experiments, there exist no measurements of polarization observables in e^- -scattering experiments. Hence these experiments are absolutely necessary for a complete multipole analysis of the data, in particular for the determination of weakly excited resonances. They are ideally suited for a low intensity machine like ELSA, and hence, will be of special interest for us during the next years.

In my short survey, I will restrict myself to two polarization experiments. For a more detailed and extended discussion see ref. /16/.

INCLUSIVE MEASUREMENTS $\gamma^*N \rightarrow X$

DETERMINATION OF THE HELICITY ASYMMETRIES $A_1(W, Q^2)$ and $A_2(W, Q^2)$

A_1 can be determined from the asymmetry of the inclusive cross-section $d\sigma/dE'd\Omega$ for parallel and antiparallel aligned spins of the virtual photon and the nucleon. The results are directly comparable to dynamical quark models, which make predictions on the Q^2 dependence of A_1 for individual resonances. From that and from a comparison with exclusive measurements one can also hope to get information on the helicity structure of the non resonant transition terms in the resonance region. Two points of special interest shall be mentioned. All dynamical quark models predict a strong Q^2 dependence of A_1 in the mass region of the $D_{13}(1510)$ and $F_{15}(1688)$ resonances. Whereas in the transition region for $Q^2 \leq 3 \text{ GeV}^2$ the different models give very distinct predictions, so that precise measurements can discriminate between the models.

A_2 can be measured by the asymmetry of the inclusive cross-sections for polarized virtual photons and polarized nucleons, where the nucleon spin is aligned perpendicular to the direction of the photon. A_2 is sensitive to the interference between transversal and longitudinal contributions to the γ^*N -coupling.

First inclusive polarization measurements were performed at SLAC /15/. The determined A_1 values, however, suffer from rather large statistical uncertainties. Hence, new and better data are desirable.

DETERMINATION OF THE ELECTRIC NEUTRON FORMFACTOR G_E^n

Using the Rosenbluth method, the electromagnetic proton formfactors were determined at many laboratories. The neutron magnetic formfactor was measured with the same method from quasielastic electron-deuteron scattering data, whereas the best values for G_E^n were derived from elastic electron deuteron scattering results /18/. This method suffers from the strong dependence on assumptions about the deuteron wave function. Therefore new independent measurements are desirable, since the electric neutron formfactor is essential for the test of microscopic nucleon models /30/, and will help to eliminate uncertainties in the interpretation of electron-deuteron scattering data /17/. At ELSA, the measurements will be performed by scattering polarized electrons on polarized neutrons (ND_3). The neutron spin is aligned in the scattering plane perpendicular to the photon direction. For this arrangement the elastic cross-section is given by:

$$\frac{d\sigma}{d\Omega} = \left(\frac{d\sigma}{d\Omega}\right)_{\text{NP}} (1 + P_E \cdot P_n \cdot A^n(Q^2))$$
$$A^n(Q^2) = \frac{2 G_E^n G_M^n \sqrt{n(1+n)} \tan \theta_e / 2}{\sigma_{\text{NP}} \cdot \sigma_{\text{Mott}}}$$

The asymmetry measurement for opposite spins of the initial electrons yields A^n , from which G_E^n can be determined. In order to keep the systematic uncertainties under control, the proton asymmetry A^p is measured simultaneously and, additionally, using the same method, from free polarized protons (NH_3). The spin-alignment of the polarized electrons will be flipped for each extracted spill, thus reducing the systematic errors considerably. With this method experimental asymmetries smaller than 10^{-3} can be detected.

REACTIONS WITH TAGGED PHOTONS

At ELSA, there will be two facilities for experiments with tagged photons. The PHOENICS detector will focus on reactions in the energy region below 1 GeV, whereas measurements at higher energies and those with many particle final states will be performed at the SAPHIR detector.

PHOTOPRODUCTION IN THE RESONANCE REGION

The electromagnetic properties of the nucleon resonances have been studied in N-scattering experiments as well as in the pion photoproduction. Whereas the most quantum numbers of resonant states are extracted from partial wave analyses of N-scattering data, the information about the NN coupling comes from pion photoproduction measurements. These coupling constants are of special importance, since they can be calculated in the framework of quark- and skyrmion models /19/. Therefore, unambiguous experimental results can discriminate between different models.

The present situation is complicated for various reasons:

1. The nucleon resonances are broad and overlap each other.
2. There are important non-resonant contributions.
3. The experimental information available is incomplete and does not allow an unambiguous analysis.
4. Even at relatively low energies, there are candidates with non negligible deviations between experiment and potential model calculations, e.g. P_{33} (1232), P_{11} (1440), D_{13} (1520).

5. At higher energies a tendency of decreasing elasticity with increasing mass is obvious. This means an increase of the importance of photoproduction data compared to results from pion scattering. Quark models predict a wealth of resonances, which are expected to couple weakly to πN and which have not yet been found.
6. Our knowledge about resonance-decays into strange particle final states is still marginal.

This situation leads very naturally to different experiments.

SINGLE PION PRODUCTION BELOW 1 GEV

A complete determination of the amplitudes requires the measurement of at least 7 independent observables (differential cross-section, single and double polarization quantities) for at least 3 of 4 possible reaction channels ($\gamma p \rightarrow \pi^0 p, \pi^+ n; \gamma n \rightarrow \pi^- p, \pi^0 n$). While there is a huge amount of cross-section data, and even the situation for single polarization observables from different laboratories (Bonn, Kharkov, Tokyo etc.) is not so bad, only a few double polarization experiments have been performed so far. One of the goals of the experimental program at ELSa is, to contribute new precise data in a wide kinematical range with emphasis on polarization observables. The experiments will begin with target asymmetry measurements. As soon as linearly polarized photons from an oriented crystal lattice and circularly polarized photons from polarized electrons will be available, also double polarization experiments will be performed at the PHOENICS detector system.

FORMATION AND DECAY OF BARYON RESONANCES

At the SAPHIR detector we intend to investigate the formation of baryon resonances and their decays in "inelastic" channels. This experiment will help to clarify the quest about the many, up to now not seen, candidates from potential models. Besides this, the investigation of cascade decays from resonances with high spin and mass will be examined, in order to test bag-string model ideas. An example for such a cascade decay could be:

$H_{3/2}^{+}(2420) \rightarrow F_{37}(1950) + \pi \rightarrow P_{33}(1236) + \pi + \pi \rightarrow N + \pi + \pi + \pi$.
In addition we will examine the channels $\pi\Delta$, ρN , ωN in particular by analyzing the final states $\pi^{+}\pi^{-}p$, $\pi^{+}\pi^{+}\pi^{-}n$ and $\pi^{+}\pi^{-}\pi^{0}p$. During the first period the measured quantities will be the differential cross-section and branching ratios; later investigations will extend to polarization observables from vector polarized nucleons as well as from linearly and circularly polarized photons.

STRANGE PARTICLE PRODUCTION

The production of strange particles is of special interest, since the s-quark is the first one not being a nucleon valence quark. Therefore it provides the possibility of an exemplary examination of a flavour threshold. The low multiplicities at the s-quark threshold allow a complete determination of the kinematics, and hence a separation of distinct reaction mechanisms. A comparison with data from pion induced reactions yields information about the photon dissociation into a $s\bar{s}$ -pair. Up to now no systematic investigations on the strange particle production in this energy region exist. The few data sets and analyses have been concentrated on the $K\Lambda$ and $K\Sigma$ channels. Measurements of these channels

have been performed at Bonn around 1970 /14/, we will continue these experiments under improved conditions with the SAPHIR detector, additionally the excitation of higher hyperons, like the $\Sigma(1385)$, will be investigated. Also a general examination of final states like $NK\pi\pi$, $NK\pi\pi\gamma$ or $NK\bar{K}$ is projected. For that purpose an efficient photon detection and a reliable kaon identification is necessary. In addition to this, specific questions of dynamics will be investigated, e.g. the $\Lambda^*(1405)$ compositeness as discussed by A.W. Thomas during this workshop. A further point might be the generation of bound hyperon states starting from the deuteron by the choice of an adequate kinematics.

CONTRIBUTIONS TO VECTOR MESON SPECTROSCOPY

Here the upper accessible mass is around 1.6 GeV, however, even with this restriction, a number of interesting questions is still to be answered. In general we have a good knowledge about the ground states ρ , ω , ϕ , but there are still uncertainties concerning the wave function at the origin $|\psi(0)|^2$, which can be reduced by precise Γ_{ee} and $\Gamma_{\mu\mu}$ measurements. Another useful information is a precise determination of the diffractive part of the ρ and ω production close to the threshold. Concerning excited states, we will concentrate on the $\rho'(1250)$ and $\rho'(1600)$ region, where also the determination of the wave function at the origin and the t -dependence of the cross-section are of some interest for the test of light quark potential models. A challenging experiment will be the determination of the VDM and GVD (generalized vector dominance) parts in the formation of excited states.

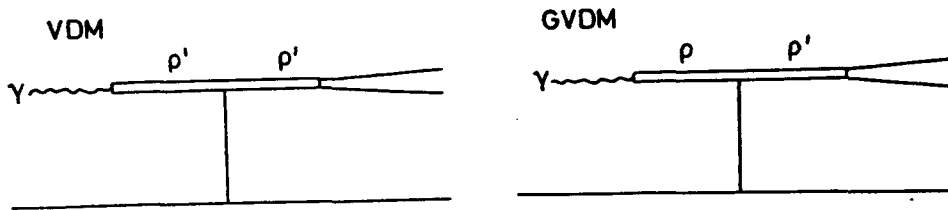


Fig.7: VDM and GVDM mechanism in the ρ' production.

The different contributions can be disentangled via shadowing measurements with different target lengths. However, this is certainly not an experiment for the very first run periods with the SAPHIR detector.

PHOTON-DEUTERON PHYSICS

Besides its in principle simple structure and its low binding energy, the deuteron is by no means simply a proton-neutron target. There are indeed many indications for a more complex structure, e.g. the relatively high frequency of quasi-deuteron reactions in nuclei or the problems concerning an unambiguous determination of the wave function. Investigations on the deuteron are an excellent starting point for many open questions in nuclear physics, like the connection between hidden colour and repulsive core, the $6-q$ contribution to the ground state or the Δ propagation in nuclear matter. In addition some theoretical tools for the analyses of light nuclei interactions may be developed and proved within this context, e.g. the treatment of the off-shell behaviour of vertices, like the $\gamma N \Delta$ vertex or the relativistic implementation of the fermi motion.

DEUTERON PHOTODISINTEGRATION $\gamma d \rightarrow pn$

The photodisintegration of the deuteron is one of the most fundamental processes in photonuclear physics. At low energies the NN-interaction plays a dominant role, whereas at higher energies nucleon resonances are excited. Therefore, ΔN - and NN^* - interactions are not to be neglected at intermediate energies. For a theoretical treatment a good knowledge of the deuteron wave function is required (including admixtures of excited nucleons and possibly 6-quark admixtures). Additionally, the photoexcitation of nucleon resonances, non resonant π -production and the interaction of Δ - and N^* - resonances with nucleons have to be described. Here the off-shell treatment of the γNN^* amplitudes from pion production off nucleons seems to be a crucial point. Due to the complicated spin structure of the photodisintegration process, 12 complex helicity amplitudes contribute and have to be determined. Therefore, a model independent analysis requires 23 measured observables in the complete kinematic range. Regarding the nucleon photoproduction results during the last two decades, this program seems not to be feasible, so that some model dependent assumptions are unavoidable. The experimental situation at photon energies $k > 100$ MeV is characterized by substantial discrepancies among different data sets for the differential cross section. On polarization observables /20/ only scarce information is available. At ELSA, a deuterium (ND_3) target with high (50%) vector polarization will be provided. With this equipment cross sections and polarization observables can be measured with the PHOENICS detector. The development of a frozen spin target, which is ideal in combination with low intensity beams, will also make possible measurements with tensor-polarized deuterons.

γ - DEUTERON INTERACTIONS WITH MULTIPARTICLE FINAL-STATES

Photon induced reactions with more than two particles in the final state will be examined at the SAPHIR-detector. There are, besides the general reasons mentioned before, two main motivations for these experiments, namely to investigate the production and to prove the existence of dibaryon candidates around 2 GeV, which have been observed at Bonn /21/. Fig. 8 shows the most pronounced enhancement at 2.014 GeV with a width $\Gamma \leq 5$ MeV and a statistical significance of 4.7 standard deviations. This signal has been seen in the invariant M_{pp} mass distribution from the reaction $\gamma d \rightarrow pp \pi^-$. Other final states will be examined too, like $d\pi^+\pi^-$ and $pp\pi^-\pi^0$. For these we will determine the n-fold differential cross-section and invariant mass distributions of the different sub-

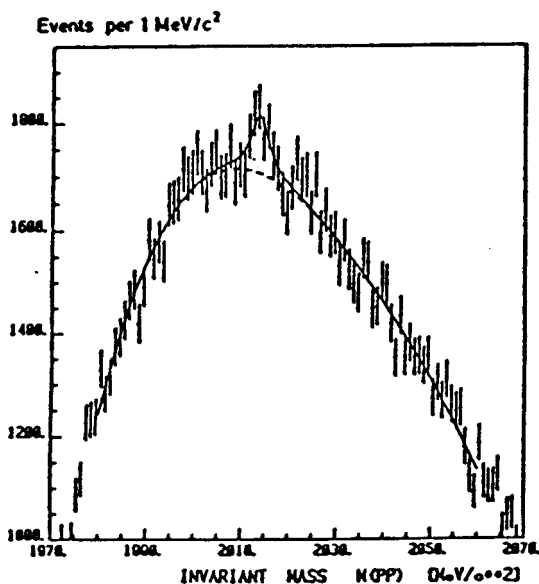


Fig. 8: The pp mass distribution from $\gamma d \rightarrow pp \pi^-$.

systems as well as cross-sections for special channels like $\Delta\Delta$ production. Since the question of the existence of 6 q-states is of vital interest, not only for an understanding of colour force saturation within the framework of QCD, but also for the nucleon-nucleon interaction in nuclei, this experiment will have a high priority at SAPHIR. If our results will be confirmed, angular distribution measurements and a helicity analysis, as well as an isospin determination must be the next step for an interpretation.

PHOTOREACTIONS ON LIGHT NUCLEI

Besides the investigations on the deuteron, the understanding of the properties of light nuclei, like ^3H , ^3He , ^4He is another important step in nuclear physics, since a reliable theory must be able to describe basic properties like binding energies, magnetic moments and electromagnetic formfactors as well as photon induced reactions. At intermediate energies the excitation and propagation of nucleon resonances and the influence of mesonic exchange currents have to be taken into account. As the 3-body problem can be solved exactly starting from a given NNpotential, the investigation of the $A = 3$ nuclei is an useful testing ground for theories. An exciting issue is the question, if 3-body force effects can be observed. The basic processes that can be studied with real photons are elastic and inelastic π -production and photodisintegration reactions.

COHERENT PIONPRODUCTION ON THE $A = 3$ SYSTEM

The theoretical description of pion production from nuclei starts with the elementary pion production on nucleons in combination with specific nuclear models. Off-

shell corrections to the production amplitudes, Pauli blocking and Fermi motion have to be taken into account as well as final state interactions. The requirement of elastic reactions, that the struck nucleon remains embedded in the residual nucleus leads to a dramatic decrease of the cross-section with increasing photon energy. The $A = 3$ system is an ideal candidate for these investigations, since ${}^3\text{He}$ and ${}^3\text{H}$ have a well known structure. A detailed comparison of both members of the isospin doublet can give valuable information on mesonic exchange effects and opens a possibility to exploit experimentally the isospin symmetry of strongly interacting particles.

4 elastic reaction channels are possible and measurable:

$$\begin{aligned} & {}^3\text{He}(\gamma, \pi^+){}^3\text{H}; \quad {}^3\text{H}(\gamma, \pi^-){}^3\text{He} \\ & {}^3\text{He}(\gamma, \pi^0){}^3\text{He}; \quad {}^3\text{H}(\gamma, \pi^0){}^3\text{H} \end{aligned}$$

From the experimental point of view it is a great advantage that the light recoil nuclei have energies high enough to be detected outside the target. An existing ${}^3\text{H}/{}^3\text{He}$ target /22/ previously used at the normal bremsstrahlung beam will be prepared for the use in the tagged photon beam at ELSA.

PHOTODISINTEGRATION OF LIGHT NUCLEI

Besides the pion production reactions, photodisintegration processes also have to be investigated and compared to theoretical calculations. At present, only the 2-body disintegration $\gamma {}^3\text{He} \rightarrow \text{pd}$ has been studied over a wide photon energy range, but the different data sets show sizable discrepancies. No data are available for the breakup reaction $\gamma {}^3\text{He} \rightarrow \text{ppn}$ at energies above pion threshold. It is planned to measure photodisintegration

processes on light nuclei using the PHOENICS detector. For low photon energies a high pressure drift chamber /23/ is under development, which is designed to detect very low energy particles in a large solid angle.

PHOTON REACTIONS AT MEDIUM AND HEAVY NUCLEI

Due to the difficulties to calculate structure and properties of heavy nuclei from first principles, an interpretation of photonuclear reactions on heavy nuclei is much more complicated. There are, however, basic problems which can be studied even without a detailed theory.

Some of the problems we will investigate are:

1. The determination of the total hadronic cross-section as a function of the photon energy for a broad range of nuclear masses. Of special interest is the question, if higher nucleon resonances are visible.
2. The measurement of correlations in energy and angle between NN or π N pairs following the photon absorption is a first step for a deeper understanding of the photon absorption mechanism in nuclei.

3. THE EXPERIMENTAL FACILITIES AT ELSA

According to the main physics goals, there are three main experimental facilities at ELSA:

1. the Electron Scattering Experiment
2. the PHOENICS detector system for physics on nuclei and nucleons at energies below 1 GeV
3. the SAPHIR detector for investigations on multiparticle final states off the nucleon and the deuteron

Besides these, there are available a pair-spectrometer set-up, mainly used for test purposes, and a forward magnetic spectrometer system for measurements with bremsstrahlung photons.

THE ELECTRON SCATTERING EXPERIMENT

This arrangement has already been used at the 2.5 GeV synchrotron, but has been improved and will be subject to further improvements for the use at ELSA. Depending on the distinct demands of the measurement, variable combinations of magnetic and non-magnetic spectrometers will be utilized.

The following spectrometers are available:

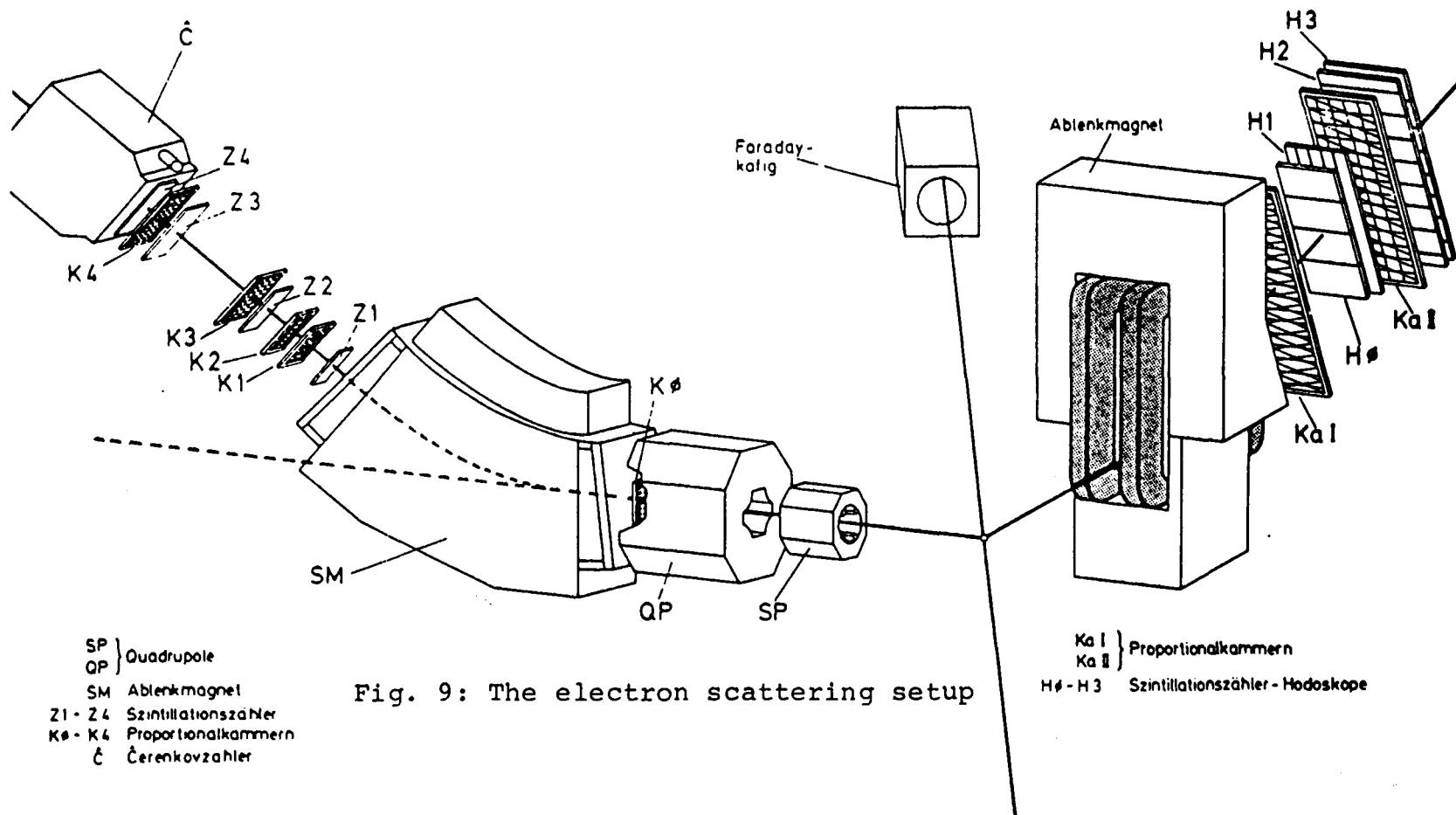
a) Electron spectrometers

1. Focusing magnet spectrometer

Momentum Range		0.2 - 1.9 GeV/c
Acceptance	$\Delta p/p$	20%
Resolution	$\delta p/p$	0.6%
Solid angle acceptance	$\Delta\Omega$	1.8 msr
Angular resolution	$\delta\theta = \delta\phi$	1.5 mrad

2. Non-focusing magnet spectrometer

Momentum range		0.1 - 0.8 GeV/c
Acceptance	p_{MAX}/p_{MIN}	2
Resolution	$\delta p/p$	1.2%
Solid angle acceptance	$\Delta\Omega$	40 msr
Angular resolution	$\delta\theta = \delta\phi$	8 mrad



b) Hadron spectrometer

Non focusing magnet spectrometer

Momentum range		0.1 - 1.4 GeV/c
Acceptance	$p_{\text{MAX}}/p_{\text{MIN}}$	4
Resolution	$\delta p/p$	1.2%
Solid angle acceptance	$\Delta\Omega$	50 msr
Angular resolution	$\delta\theta = \delta\phi$	8 mrad

For the future, a non magnetic, large solid angle detection system consisting of a lead glas or BGO array is planned for the detection of electrons under large production angles, as well as a large solid angle neutron detector for measuring the electric neutron formfactor G_E^n . A typical e-scattering arrangement, as it has been used at the 2.5 GeV synchrotron, is shown in fig. 9.

Two other important ingredients to the electron scattering experiment are the polarized electron source and the polarized target /24/. The latter one uses superconducting coils in Helmholtz configuration, which makes possible horizontal and vertical spin alignment. The accessible deuteron polarization values at 4 T are 50% for the vector polarization and 20% for the tensor polarization.

THE PHOENICS DETECTOR SYSTEM

The setup planned for the investigation of photonuclear reactions at ELSA, consists of a tagged photon beam and a scintillation counter array for the detection of photon induced reactions. Polarized H and D targets and a $^3\text{H}/^3\text{He}$ target are available, polarized beams with linear and circular photon polarization will be available in the near future.

THE TAGGING SYSTEM

The tagged photon beam facility /25/ has been optimized for operation at photon energies below 1 GeV. The magnetic spectrometer for the momentum analysis of the recoil electrons from the bremsstrahlung process consists of a synchrotron-type C-magnet with curved pole faces and a strong field gradient ($dB/dr = 5 \text{ T/m}$). This field provides a point to point focusing in the deflection plane. Electrons between 50 and 800 MeV can be detected with an average dispersion of 1.3 mm/MeV. The deflection of the primary beam at 1 GeV is 35° . The detection system is built up of 128 small scintillation counters and 16 thick backing counters, which serve for timing, background suppression and trigger purposes. The energy resolution is limited by the size of the counters to 6 MeV at higher and 12 MeV at low photon energies. The complete system will be ready in summer 1986. A scintillation counter module has been tested at the Mainz microtron; after off-line corrections the timing resolution obtained was better than $\sigma_t = 150 \text{ ps}$. Fig. 10 gives a survey of the tagging system.

It is planned to use a diamond crystal radiator to produce linearly polarized photons. Since the polarization is highest at approximately $0.3 \cdot E_0$, the tagging system was designed to cover an energy range from $(0.2-0.95) \cdot E_0$ simultaneously. When polarized electrons will be available at ELSA, they will be used to produce circularly polarized photons too. Therefore, the transverse spin orientation inside ELSA has to be turned into a longitudinal one. A separate beam-line containing 4 dipoles, 3 superconducting solenoids and several quadrupoles has been designed to achieve a spin rotation in a wide energy range.

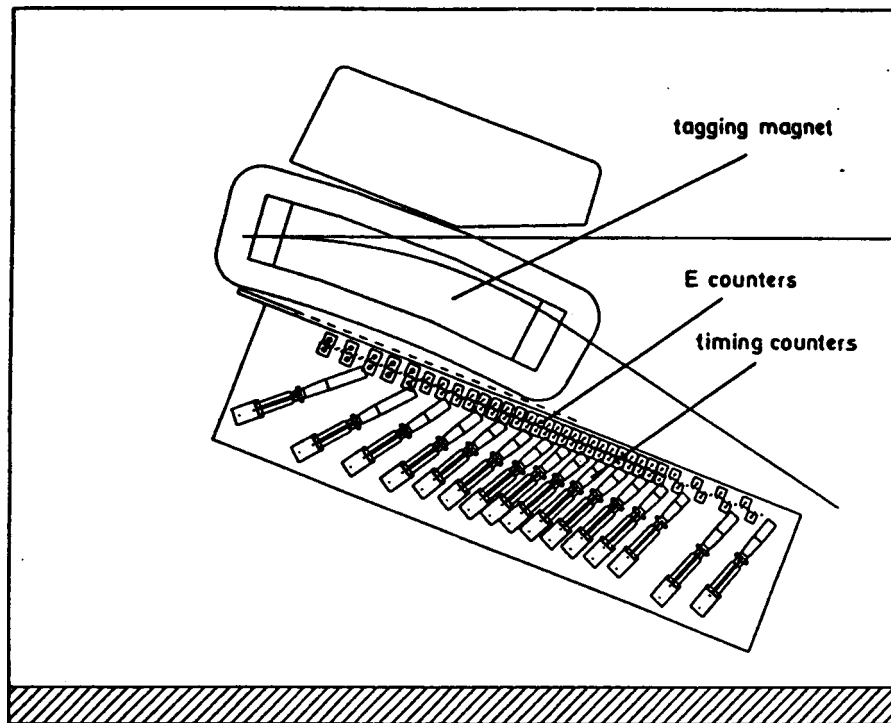


Fig.10: The tagging system for the PHOENICS detector.

THE SCINTILLATION COUNTER SYSTEM

The detector arrangement (fig. 11) consists of an inner ring of scintillation counters close to the target and an outer one at larger distance. All counters have two multiplier tubes each for the vertical position resolution.

The counters are:

- 30 start- or veto counters ($50 \times 0.5 \times 10$) cm^3
- 31 stop counters for charged particles ($170 \times 20 \times 20$) cm^3
- 26 stop counters for neutrons or charged particles ($100 \times 20 \times 20$) cm^3

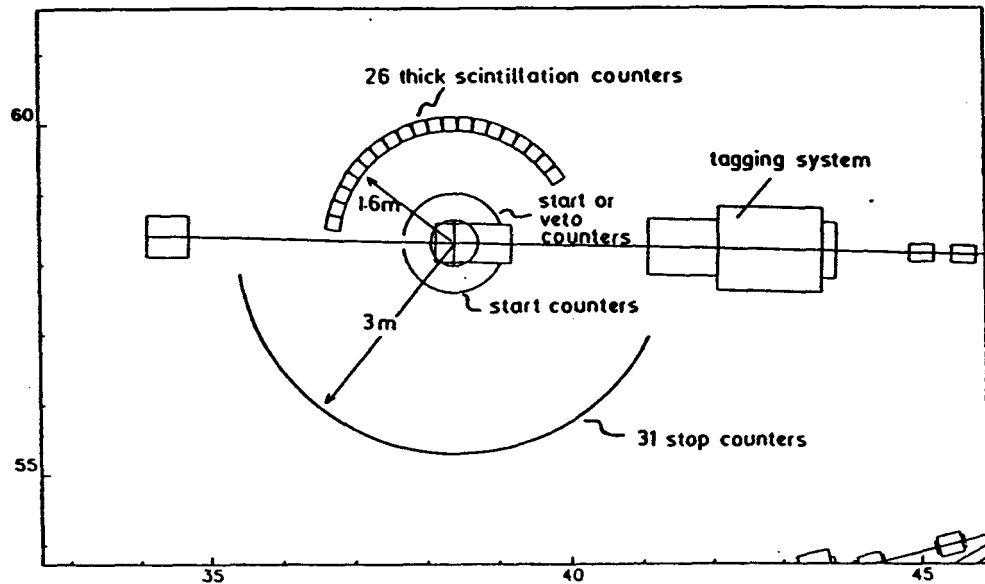


Fig.11: The PHOENICS time of flight hodoscope system.

Particle separation can be achieved in a wide kinematic range by a combination of the informations on the energy deposited in the counters and the time of flight. Timing and pulse-height are supervised via a computer controlled LED system.

Status: at present the start counters are being built, the stop counters are ready and being tested.

THE HIGH-PRESSURE DRIFT CHAMBER DETECTOR

An alternate approach to a detector design for the detection of low energetic charged particles has been realized. The principle is a time projection chamber around a long gaseous target, both operated at the same high pressure [23]. The long high-pressure target yields a high amount of material in γ -beam direction and a low one perpendicular to the beam, allowing the detection of

low energy particles. The drift chamber, separated from the target by a thin foil, consists of a central cathode cylinder, one outer layer of 32 signal- and potential wires and cathode strips. The chamber is surrounded by a 8 element scintillation counter hodoscope. The whole setup is housed in a stainless steel vessel. Fig. 12 shows the arrangement. The track reconstruction is achieved by measurements of the drift-time to the signal wires, charge division on the wires and by the induced signals on the cathode strips. Since there is only one layer of wires, a measurement of the signal development as a function of the time in flash-ADC's is required. First test measurements at a pressure of 10 atm have been performed with cosmic rays and for a limited number of channels. A test with the complete electronic equipment is projected for later this year.

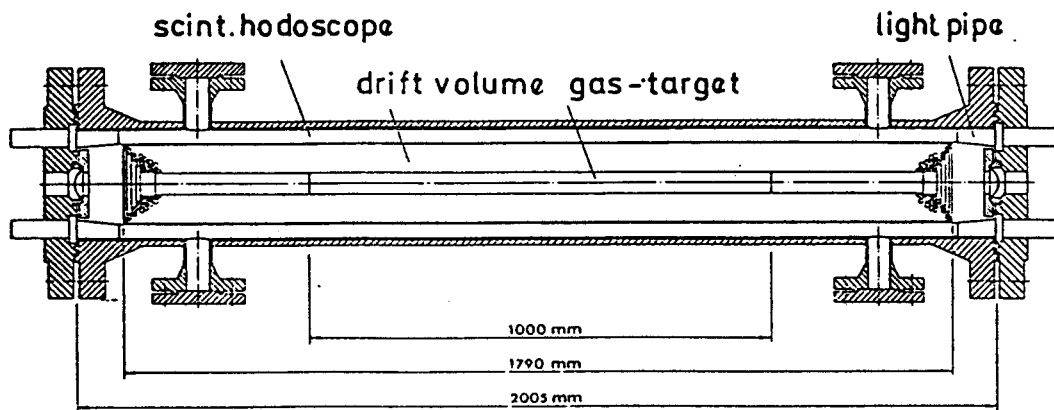


Fig.12: The high pressure drift chamber detector.

TARGETS

A polarized NH_3 and ND_3 target /27/ is available. It is operated at 2.5 T and at a temperature of 0.2 K, which is maintained by a $^3\text{He}/^4\text{He}$ dilution refrigerator. The obtained polarization for protons is larger than 0.8, the one for deuterons is about 0.35. A method to provide a tensor polarization $\geq 25\%$ is under development. The low intensity of a tagged photon beam makes the use of a frozen spin target possible, which can be maintained in a much weaker holding field, thus causing less problems with respect to the pair-production background. This target type will also be used in the SAPHIR detector. Standard liquid H_2 , D_2 and ^3He targets are available from previous experiments.

THE $^3\text{H}/^3\text{He}$ TARGET AND DETECTOR SYSTEM

A $^3\text{H}/^3\text{He}$ gas-target /22/ has been used at the bremsstrahlung beam of the former 500 MeV synchrotron /26/ and will be modified for the use with the tagged photon beam.

The modifications mainly involve:

1. concentration of the tritium (now 2000 Cu corresponding 800 cm^3 NTP) in a cell with smaller diameter,
2. the construction of a detection system inside the vacuum vessel with large energy- and angle acceptance, which is optimized for the detection of recoil nuclei.

Charged pions will be detected in the standard scintillation counter setup.

THE SAPHIR DETECTOR

In the intermediate energy range, hitherto mainly experiments with two particle final states have been performed, albeit even in reactions on the nucleon many particle final states dominate. This was caused mainly by the existing set-ups with small solid-angle spectrometers, which have been used in order to avoid accidentals due to the bad duty factors of the accelerators. Now, a new generation of machines with 100% duty factor makes possible experiments with large acceptance detectors together with tagged photon beams.

DEMANDS TO THE DETECTOR SYSTEM

The proposed experiments require the following detector properties:

- a large solid angle for the detection of multiparticle final states with good acceptance for charged particles and for photons,

- a good track definition for the event reconstruction, demanding a dense packing of the detector planes inside the magnetic field,

- a minimized amount of material along the path of the particles, since the majority of produced particles has momenta below 500 MeV,

- a positive particle identification is needed,

- a tagging system, that covers a wide energy range and that delivers a good energy determination.

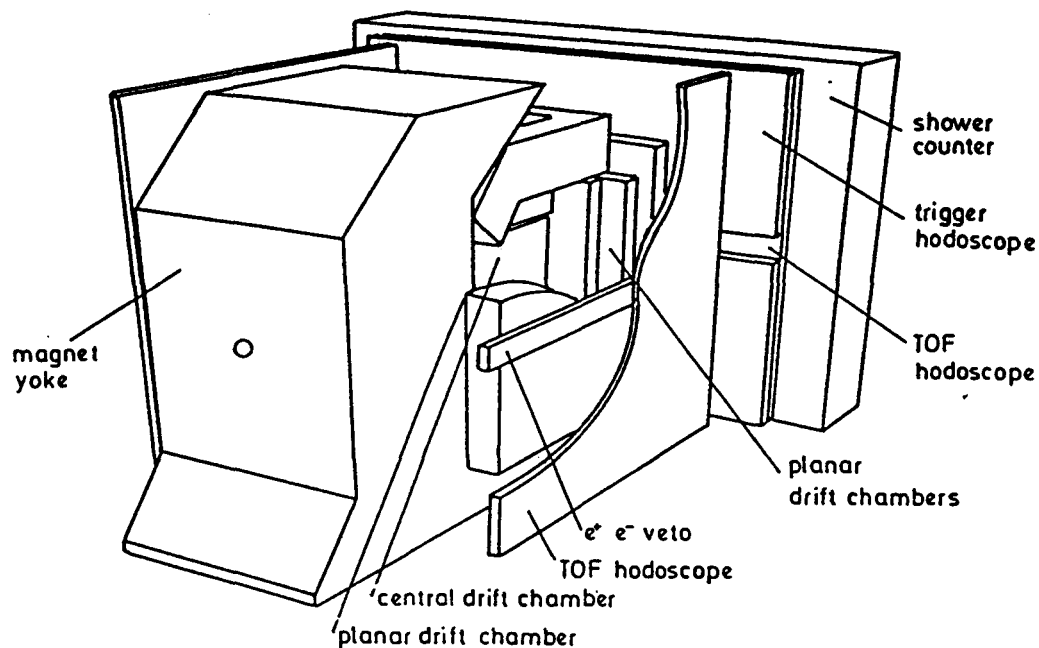


Fig.13a: The SAPHIR detector, perspective view.

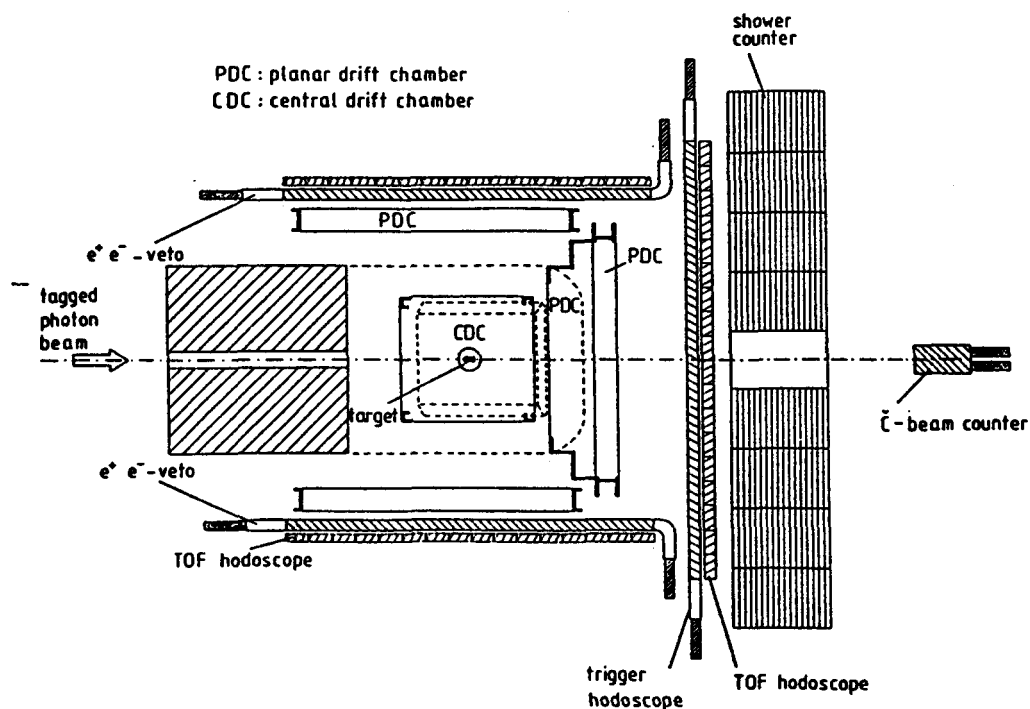


Fig.13b: Top view of the SAPHIR detector.

THE SET-UP

The general lay-out is displayed in fig. 13. The system provides a solid angle $\Delta\Omega = 0.6 \times 4\pi$ sterad and a momentum resolution σ_p/p between 1 - 2% in the forward cone, which becomes worse by a factor of 1.5 for side-ward emitted particles.

THE MAGNET

The C-type magnet has pole faces of $1 \times 1 \text{ m}^2$ and a gap height of 1 m. It is operated at a magnetic field of 0.6 T at maximum.

THE TARGET

The target is centered in the middle between the pole faces. It consists of a horizontal cylinder of 8 cm length, which can be filled with liquid H_2 or D_2 . The cylinder is connected to the cryostat through a hole in the upper pole face.

THE TRACKING SYSTEM

The tracks of charged particles are determined with the help of the central drift chamber, one planar chamber on either side of the magnet and two planar chambers in the forward direction. The track coordinate informations from this set-up will be used for the event reconstruction and to evaluate the event vertex, secondary vertices of V-particles and the momenta of the particles at their vertices.

All SAPHIR driftchambers have a hexagonal cell structure to approximate a rotational field symmetry, since there is no predominant direction of the trajectories. Accor-

ding to the expected higher particle flux close to the target and in the forward cone, the drift cells in the central- and in the forward chambers have a slightly smaller diameter than the ones in the side chambers (see fig. 14).

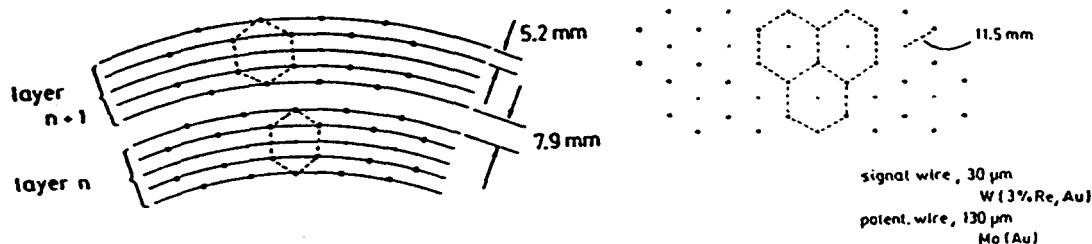


Fig. 14a: Cylindrical layers. Fig. 14b: Double layers.

In the central chamber the drift cells form concentric layers (fig. 15) with a spacing of 28.6 mm between the signal wires of neighbouring layers. The cylindrical inner part of the chamber has eight layers with vertical wires and six "stereo" layers with an inclination angle of $\pm 5^\circ$ with respect to the vertical direction. They are used to determine the vertical coordinates of track-points between the pole faces. Additionally, the central chamber has layers of horizontal cells at its edges. In order to keep the effects of Coulomb scattering small, a He/Ne/CH₄ gas mixture will be used.

The planar chambers on either side of the magnet (fig. 16) have the drift cells arranged in double layers, in order to resolve "left-right" ambiguities (fig. 14). They have two such double layers for measuring horizontal and vertical track coordinates with the same resolution. One "stereo" layer was introduced to avoid multi-track ambiguities. The two planar front-chambers will have vertical, horizontal and "stereo" layers at

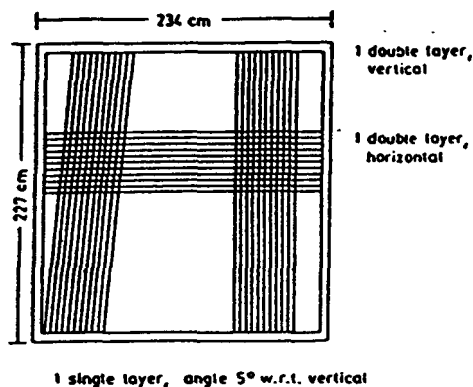
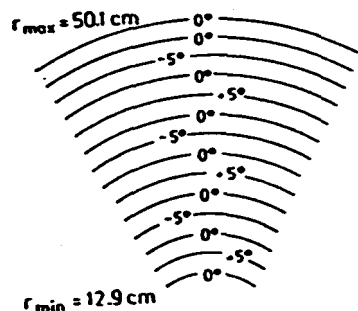


Fig. 15: Layer arrangement in the central drift chamber

Fig.16: The planar side chambers.

angles of $\pm 5^\circ$. Normally both will be mounted together and positioned very close to the front edge of the pole faces. However, for a better momentum resolution, the second one can be positioned in front of the trigger hodoscope.

A summary of the data is given in table II:

Central Drift Chamber

dimensions	1.14 x 1.19 x 0.98 m ³
layers	8 vertical, 6 stereo, 1 horizontal at each edge
number of cells	1540
distance signal-, potential wire	10.4 mm
gas mixture	He/Ne/CH ₄ (17/49/34)
resolution	$\sigma_x = 180 \mu\text{m}$

Planar Side Chambers

dimensions	2.24 x 2.30 x 0.20 m ³
layers	1 horizontal double layer, 1 vertical double layer, 1 stereo layer
number of cells	560/chamber
distance signal-, potential wire	11.5 mm
gas mixture	Ar/CH ₄ (90/10)
resolution	$\sigma_x = \sigma_y = 180 \mu\text{m}$

Planar Front Chambers

dimensions	1.65 x 1.62 x 0.38 m ³
	and 2.35 x 2.55 x 0.20 m ³
layers	5 horizontal, 5 vertical, 8 stereo
number of cells	1762
distance signal-, potential wire	10.4 mm
resolution	$\sigma_x = \sigma_y = 180 \mu\text{m}$

For a more detailed description of the drift chamber system and further performed tests see ref. /28/.

Status: The central drift chamber and the two side chambers are assembled and are being tested. The front chambers have been designed, the frames are in production.

THE TIME OF FLIGHT- AND TRIGGER HODOSCOPES

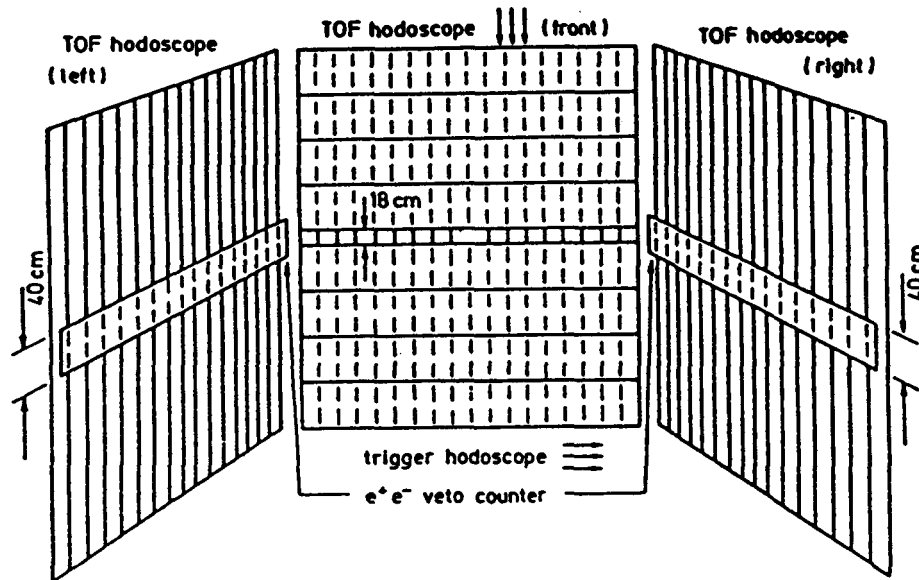


Fig.17: The TOF- and trigger hodoscopes.

The TOF hodoscopes are provided for particle separation by means of a time of flight measurement. The start signals come from thick scintillation counters in the tagging hodoscope.

One of the problems of a magnetic detector system, where the target is situated inside the field volume, is the electron and positron background due to pair-production. The expected e^+, e^- -rate of 23 kHz (at a luminosity $L = 3.4 \cdot 10^{29} \text{ cm}^{-2} \text{ s}^{-1}$) is rather high compared to the hadronic rate of 60 Hz. Therefore the fast trigger logic has to reject e^+, e^- -pairs with an efficiency of at least 99.9%. A Monte Carlo simulation has shown, that 94% of the electrons (and positrons) hit the forward TOF hodoscope (fig. 17) within approximately $\pm 7 \text{ cm}$ around the horizontal symmetry plane, the remaining 6% hit the sideward hodoscopes with a somewhat broader distribution. Therefore in front of the forward TOF hodoscope an

additional trigger hodoscope will be set up, leaving a gap of 18 cm at beam height. The fast trigger logic will only accept coincidences between TOF- and trigger counters. For the side hodoscopes a signal in one of the TOF counters is used only if there is no coincidence with the 40 cm high veto counter. The pulse height and timing stability of the TOF counters is supervised with a laser calibration system via glass fibre cables.

Status: the TOF hodoscope counters are assembled and have been tested partially. The obtained time resolution is $\sigma \approx 150$ ps. The trigger and veto counters are being built.

THE SHOWER COUNTER

With respect to the large area, which has to be covered, a comparatively inexpensive system for the photon detection has been designed. It is a modular system in a metal-gas type sandwich construction with proportional wire read out. It has been developed and tested at Bonn /29/. The design is to be seen from fig. 18.

The technical data and the measured resolutions are:

size of the front area of one module	35 x 51 cm ²
depth of a module	12 x λ_0
converters	24, 0.5 x λ_0 each
signal wires gold plated tungsten	50 μ m diameter
energy resolution	$\sigma_E/E = 13.3\% / \sqrt{E \text{ (GeV)}}$
spatial resolution	15 mm at 1 GeV, if 3x3 signals are added and read out by one electronic channel
detection efficiency	> 90% at 50 MeV

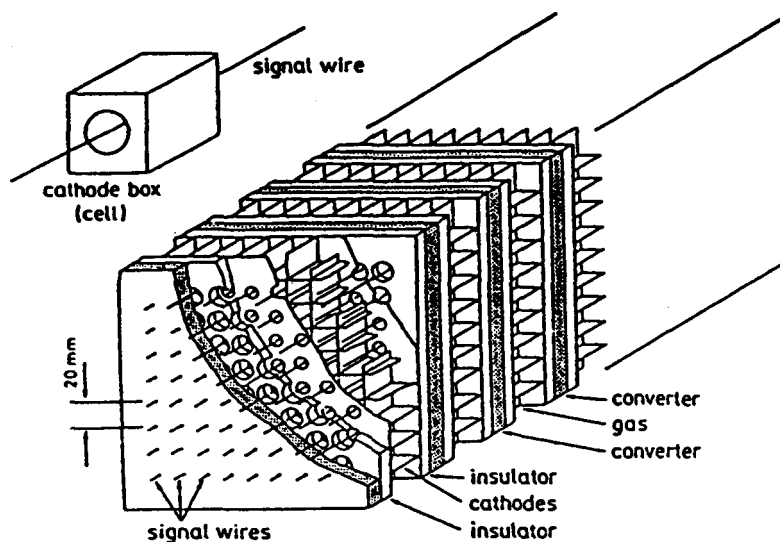


Fig.18: A shower counter module.

Status: The forward shower detector consists of 98 modules, of which 60 have been assembled, the remaining 38 modules will be ready late this year.

THE AEROGEL ČERENKOV-COUNTER

For the pion kaon separation in the momentum range 0.8 - 2.4 GeV/c an aerogel Čerenkov-counter with $n = 1.03$ is in preparation. This counter is needed for the forward direction only, since particles emitted towards the sides are slow enough for a separation via a time of flight measurement.

THE TAGGING SYSTEM

At SAPHIR, during a first period, the scintillation counter system designed for the PHOENICS tagging system will be mounted at a tagging magnet with homogeneous

field. The focusing in the vertical deflexion plain is gained with the fringe field only. With this arrangement the energy range covered is limited to $(0.5-0.95) \times E_0$, providing an energy resolution of 5 - 10 MeV. At a later stage, a new system, using a multiwire proportional chamber and scintillation counters for the electron determination, as well as a focusing magnet will be built. This system has been designed, it will yield an energy range $(0.2 - 0.95) \times E_0$ up to the endpoint energy of 3.5 GeV. This system has been postponed to a later period for financial reasons.

DATA AQUISITION AND PATTERN RECOGNITION

Last, but not least I would like to mention this essential part of any large solid angle detection system. The data transfer from the different experimental components and the supervision of the detector components is done via an intelligent system of front-end computers on 68000 basis, which have been developed in our institute. There is a threefold parallel readout with data buffers before the information is put together in an "event-composer". From here the event data are transfered to the online computer. In addition, this system supervises the performance of all components and delivers information about the event statistics.

The off-line track-finding is done according to a 3-dimensional "principle component method". If a set of measured points has been associated to a track, these points are fitted with a quintic spline in order to extract the physical parameters. Here Coulomb scattering and energy losses are taken into account. Besides these pattern recognition and geometry programs, there exists a simulation chain for the preparation of specific expe-

riments, the interpretation of displayed patterns and for a general understanding of the detector system.

ACKNOWLEDGEMENT

My thanks go to my colleagues at Bonn, especially to Dr. J. Arends, Dr. D. Husmann and Prof. G. Knop, for helpful discussions and for providing parts of the material presented here. I would also like to thank the organizers of the CEBAF workshop, in particular Prof. F. Gross, Dr. V. Burkert, Prof. B.A. Mecking, Dr. R.R. Whitney and Miss R. Rickman for their efforts in organizing a very instructive and pleasant meeting.

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MULTIPARTICLE FINAL STATES FROM PION-INDUCED REACTIONS

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ABSTRACT

In this talk I will review multiparticle final states originating from pion-induced nuclear reactions. I will cover both quasi-elastic knockout and true absorption using recent experimental and theoretical results to demonstrate our level of understanding at the present time.

I. INTRODUCTION

In spite of rumored complaints by several of the theorists in the audience, I am here to discuss the "dirty" business of pion physics. In this talk I will limit my discussion to multi-particle final states produced in pion-nucleus reactions. I will further restrict my discussion to pions with kinetic energies corresponding to the region of the $\Delta(1232)$ resonance (~ 100 - 250 MeV). In this energy range two reactions dominate the reaction cross section -- inelastic scattering, which is primarily quasielastic scattering, and true absorption (no pion in the final state). This is demonstrated in Fig. 1 which shows the A-dependence of the reaction, inelastic, and absorption cross sections for 165 MeV π^{+} .^{1,2} At this energy for medium mass nuclei the inelastic and absorption cross sections each represent $\sim 50\%$ of the total reaction cross section. Note that the A-dependence suggests that inelastic scattering is more peripheral than absorption which is roughly proportional to the area.

Also shown in Fig. 1 are the energy dependences of the reaction cross sections for ^{12}C . The absorption cross section clearly reflects the Δ -resonance and dominates at very low energies. Above resonance the inelastic scattering becomes more and more dominant.

In the following I will attempt to review the present status of pion quasielastic scattering and in-flight absorption. Where possible I will discuss relationships to electromagnetic processes.

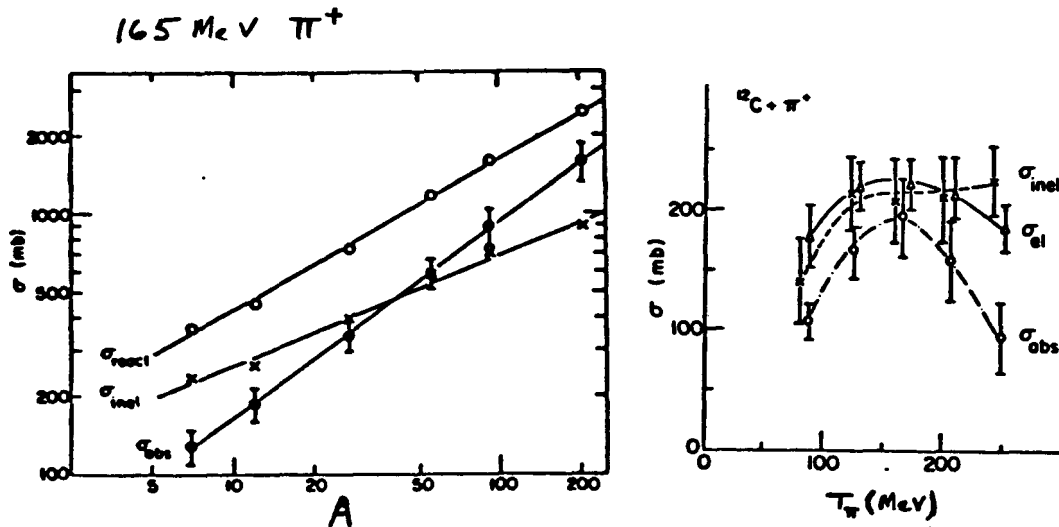


FIG. 1. Total cross sections for π^+ -nucleus interactions. Left: A-dependence of the total inelastic, absorption, and reaction cross sections for 165 MeV π^+ . Right: Energy dependence for $\pi^+ + {}^{12}\text{C}$. (From Refs. 1 and 2.)

II. QUASIELASTIC $\pi^\pm$ SCATTERING

In this section I present results for quasielastic scattering in order of increasing sophistication of the experiments, starting with inclusive measurements and then turning to coincidence studies.

A. Inclusive $\pi^\pm$ Quasielastic Scattering

At intermediate energies the inelastic cross section is essentially composed of quasielastic π -N scattering. Shown in Fig. 2 are large angle inclusive (π^+, π^+) spectra at 160 MeV.³ A clear quasi-free peak representing the scattering of pions from individual nucleons bound in the nucleus is observed, even for nuclei as heavy as Pb. The quasifree peak lies below the free π -N location (arrow) due to the binding energy and has a width determined by the Fermi motion of the struck nucleons.

I think that this is a remarkable result. We have a very strongly interacting particle showing a predominance of single π -N collision processes. Similar studies of intermediate energy (p, p') only show such a pronounced quasifree peak for the lightest targets at small angles. There are several reasons why we might expect the

strongly interacting pions in this energy range to show such a clean quasifree peak. Firstly, at these energies a major component of the reaction cross section is pion absorption. These pions are removed from the inelastic scattering channel and therefore do not produce "background" in the (π, π') spectra. Secondly, in this energy range after an initial π -N scattering, on average the pion energy is below resonance thereby reducing the π -N cross section and thus the probability of a second scattering. It is interesting to note that at higher pion energies ($T_\pi \approx 400$ MeV) the quasifree signature is gone for backward angles.⁴ Thirdly, due to the strong absorption the reaction is localized to the nuclear surface where the nucleon density is lower and the pions have less material to traverse.

To show the angular dependence of quasifree scattering, in Fig. 3, I present 163 MeV inclusive $^{16}\text{O}(\pi, \pi')$ spectra at various angles.⁵ These spectra show the centroid of the quasifree peak decreasing in energy with increasing angle in accordance with π -N scattering. The curves are DWIA calculations by Thies.⁶ The solid curves include the effects of formation and propagation of the Δ in the nucleus using the Δ -hole model and a phenomenological Δ -nucleon potential. These calculations provide a rather good description of the experimental data. Not surprisingly, calculations neglecting Δ formation and propagation (dashed curves) do a poor job in describing the data.

Examination of the small angle data in Fig. 3 shows that the

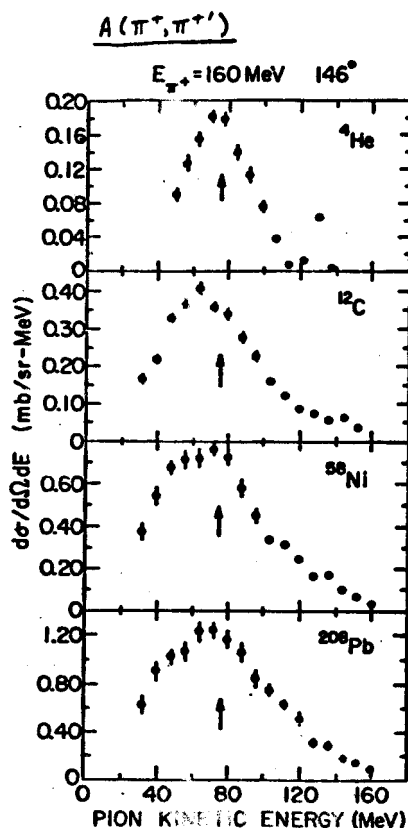


FIG. 2. Inclusive 160 MeV (π^+, π^+) cross sections at 146° (Ref. 3).

theory greatly underpredicts the data at low outgoing pion energies ($T_\pi < 75$ MeV). This is not surprising and we presume that the data in this region contain a preponderance of events from multiple scattered pions. For pions, unlike other projectiles, one has a extremely nice method of quantifying the yield from pion double scattering. In particular, one can measure the continuum yield produced in pion double charge exchange ($\pi^\pm, \pi^\mp$) which must involve two π -N interactions.

Figure 4 shows $^{16}\text{O}(\pi, \pi')$ inclusive data for two angles compared to $^{16}\text{O}(e, e')$ at the same momentum transfer normalized to the pion data. Also shown are $^{16}\text{O}(\pi^+, \pi^-)$ data normalized at low energies.⁷ This is a very informative figure and demonstrates two essential features. (1) The agreement with (e, e') indicates that the bulk of the (π, π') quasifree peak arises from single π -N scattering. (2) The low energy tail of the quasifree peak arises from double scattering, the shape of which can be described by pion double charge exchange data.

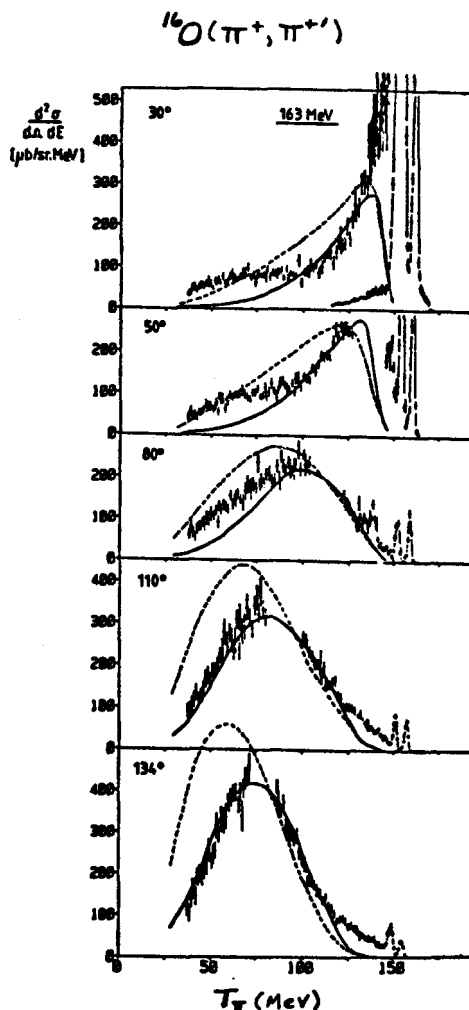


FIG. 3. Inclusive 163 MeV $^{16}\text{O}(\pi^+, \pi^+')$ cross sections (Ref. 5). The curves are the theoretical calculations with (—) and without (---) the effects of Δ -propagation (Ref. 6).

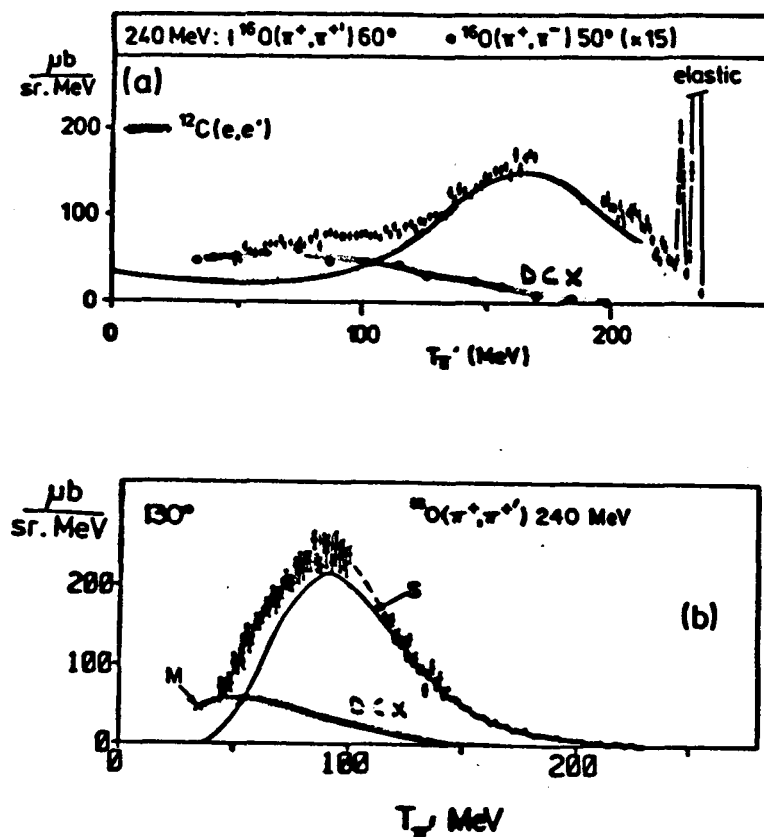


FIG. 4. Inclusive 240 MeV $^{16}O(\pi^+, \pi^+)$ and $^{16}O(\pi^+, \pi^-)$ DCX cross sections. The solid curve corresponds to $^{12}C(e, e')$ at the same momentum transfer. The DCX data have been renormalized by about a factor of 15 (Ref. 7).

B. Exclusive Quasielastic Scattering

The next level of sophistication and detail in quasifree scattering is direct measurement of the π -N coincidences. A very nice coincidence experiment for $^{16}O(\pi^\pm, \pi^\pm p)^{15}N$ has been carried out at SIN.^{8,9} This experiment achieved sufficient energy resolution to separate the $1/2^-$ and $3/2^-$ states (see Fig. 7) and sufficient coverage of phase space for the coincident protons to allow accurate integration for the total yields. Results for two pion angles are presented in Fig. 5. The top panels (a) show the inclusive (π^+, π^+) yield at 60° and 130° . The solid line represents the $(\pi^+, \pi^+ p)$ yield integrated over all proton angles, but subject to a ~ 35 MeV proton energy threshold. Also shown are breakdowns of the integrated

yields into p-shell and core contributions. The coincidence pion spectra strongly resemble the inclusive pion spectra, but contain $\lesssim 50\%$ of the total inclusive yield. The ratios of the exclusive to inclusive yields as a function of T_π are shown in panel (b). Note that addition of $(\pi^-, \pi^- p)$ to $(\pi^+, \pi^+ p)$ effectively includes the $(\pi^+, \pi^+ n)$ contributions to the inclusive yield. At the maximum of the quasifree peak the coincident yield (subject to the proton energy threshold) accounts for $\sim 55\%$ (70%) of the inclusive yield at 60° (130°). At lower outgoing pion energies this ratio decreases as expected, since a significant component of the inclusive yield results from double (or more) scattering.

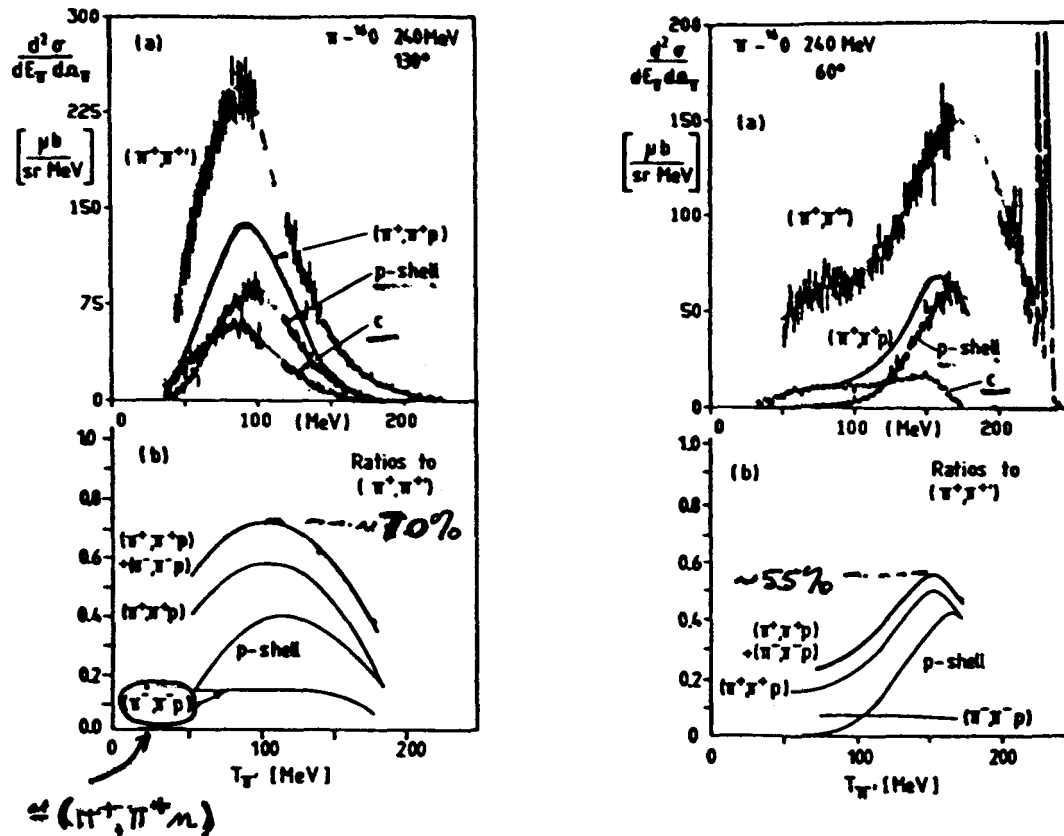


FIG. 5. Inclusive $^{16}\text{O}(\pi^+, \pi^+)$ and exclusive $^{16}\text{O}(\pi^+, \pi^+ p)$ cross sections. The $(\pi^+, \pi^+ p)$ data are separated into p-shell and "s-shell" knockout. The lower panels show the ratio of the exclusive to inclusive yields (Ref. 7).

The fact that at the peak we miss $\sim 30\%$ to 50% of the yield presumably reflects the fact that roughly $1/3$ to $1/2$ of the quasi-

elastically scattered protons interact with additional nucleons in traversing the nucleus, losing sufficient energy so that they fall below the 35 MeV threshold. Note that at the peak for 60° the pion kinetic energy is ~ 160 MeV leaving $\lesssim 65$ MeV for the emergent proton. At 130° the proton has $\lesssim 135$ MeV, which because of its higher penetrability would account for the higher probability of being detected in the proton detector.

As a possible check on this hypothesis I compare the 60° data to factorized DWIA calculations based on the "cartoon diagram"¹⁰ shown in Fig. 6. In the factorized DWIA the cross section can be written as¹¹

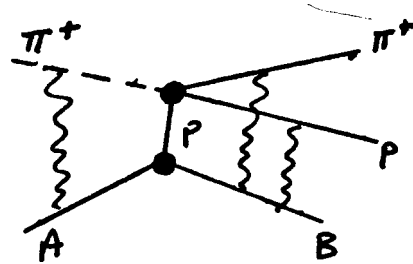


FIG. 6. "Cartoon" diagram for the $A(\pi, \pi p)B$ reaction in the DWIA.

$$\frac{d^3\sigma}{d\Omega_{\pi} dE_{\pi} d\Omega_p} = KF \cdot \left. \frac{d\sigma}{d\Omega} \right|_{\pi^+p} \cdot \sum_{\Lambda} |T_L^{\Lambda}|^2 \quad (1)$$

where KF is a known kinematic factor, $\left. \frac{d\sigma}{d\Omega} \right|_{\pi^+p}$ is the two-body π^+p elastic scattering cross section, and

$$T_L^{\Lambda} = (2L+1)^{-1/2} \int \chi_{\pi^+}^{(-)*}(\vec{r}) \chi_p^{(-)*}(\vec{r}) \langle \psi_B | \psi_A \rangle \chi_{\pi^+}^{(+)}(\gamma \vec{r}) d^3r. \quad (2)$$

The χ 's represent distorted waves obtained from optical model potentials fit to elastic scattering and $\langle \psi_B | \psi_A \rangle$ is the overlap of the initial and final nuclear states which is proportional to the bound proton single particle wave function.

Since I have already shown the importance of Δ -propagation in Fig. 3, one might ask why I am doing this factorized calculation. Empirically we have found that factorized DWIA calculations using phenomenological potentials provide a rather good description of exclusive (π^+, π^+p) data.¹² Presumably the use of phenomenological pion optical-model potentials includes many of the major effects from Δ -propagation.

By integrating Eq. (1) over the full proton solid angle I can

calculate $\frac{d^2\sigma}{d\Omega dE_\pi}$ for each state in the residual nucleus. DWIA calculations summed over the $p_{1/2}$, $p_{3/2}$, and $s_{1/2}$ shells in ^{16}O , with and without the absorptive part of the proton optical-model potential, should represent the exclusive $(\pi, \pi p)$ and inclusive (π, π') results, respectively. The calculations compared to the experimental data⁸ are displayed in Fig. 7. The DWIA calculations have not been renormalized; i.e., they are normalized by the nucleon occupancy. In the region of the peak the DWIA calculations provide a rather good description of the data. In addition, they reproduce the factor of approximately two reduction for the coincidence measurements through the absorption of emergent protons.

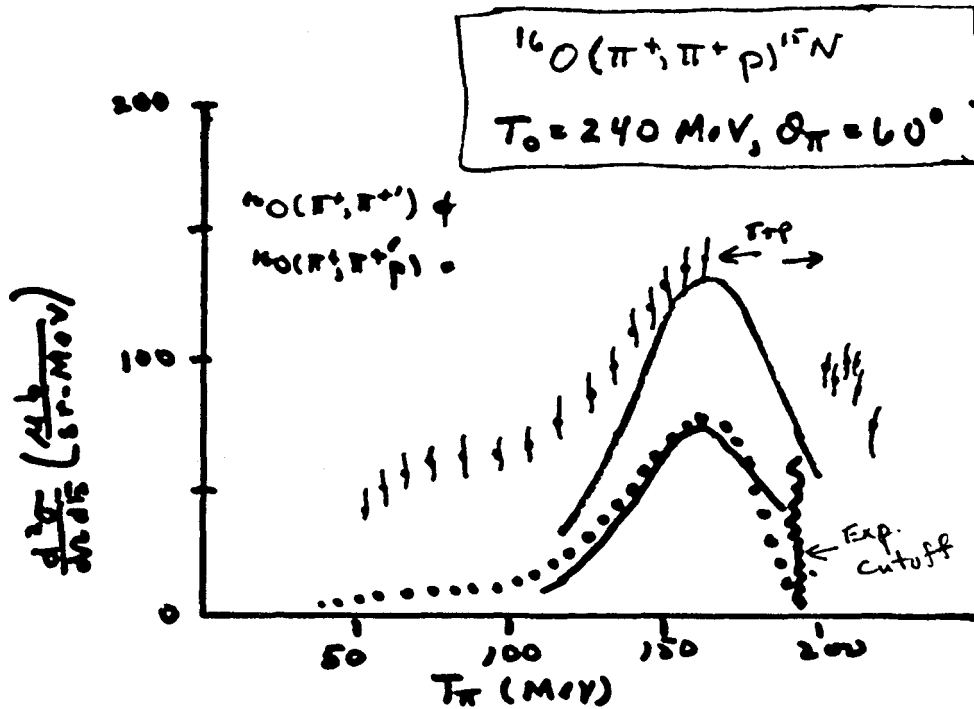


FIG. 7. Inclusive $^{16}\text{O}(\pi^+, \pi^+')$ and exclusive $^{16}\text{O}(\pi^+, \pi^+ p)$ cross sections. The curves are DWIA calculations as discussed in the text.

As a final topic on quasifree π -N scattering I consider the same SIN experiment^{8,9} in which comparisons were made between $^{16}\text{O}(\pi^+, \pi^+ p)^{15}\text{N}$ and $^{16}\text{O}(\pi^-, \pi^- p)^{15}\text{N}$ for individual final states. Presented in Fig. 8 are the excitation energy spectra demonstrating the separation of the $1/2^-$ and $3/2^-$ states, and the ratio

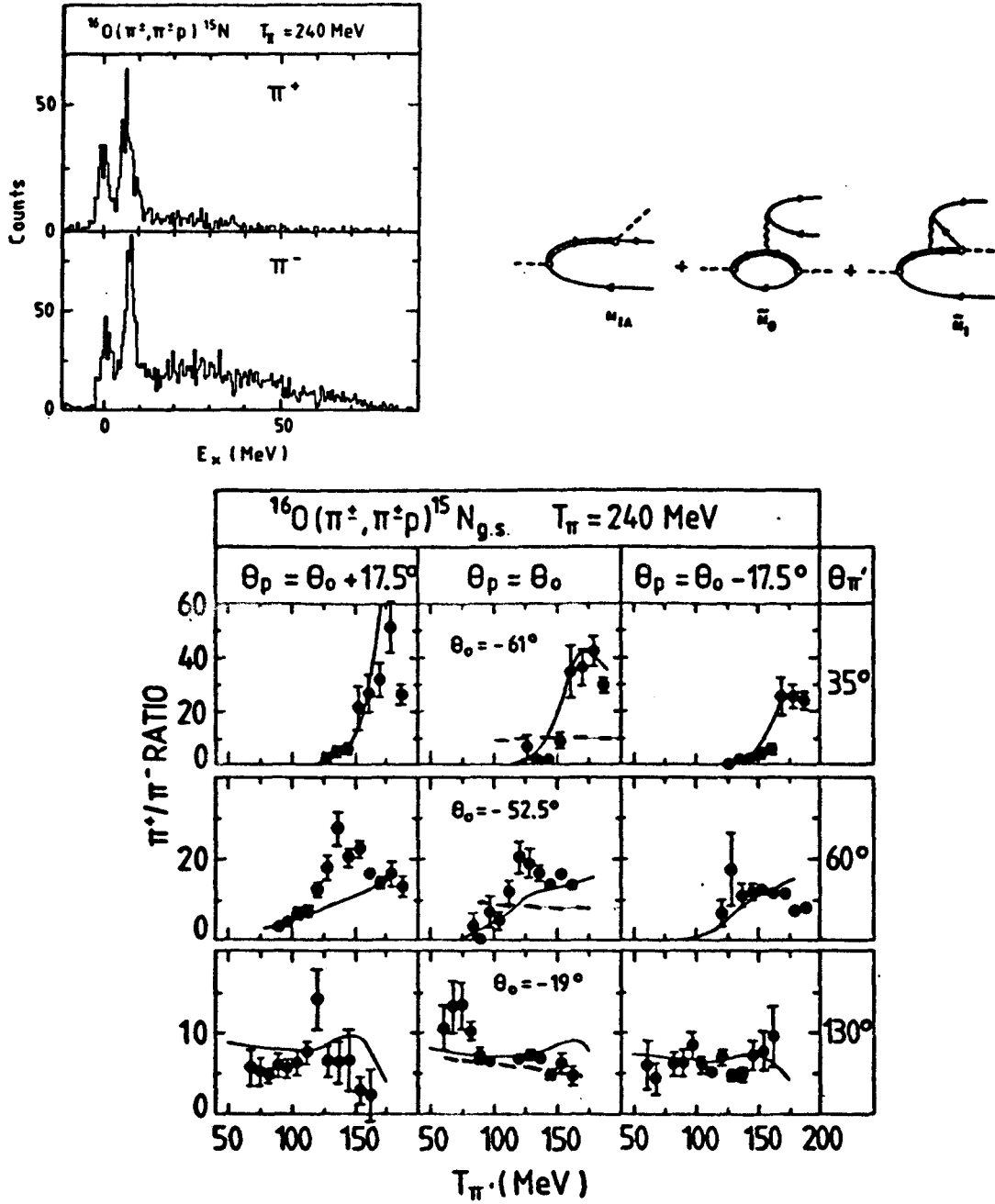


FIG. 8. $^{16}\text{O}(\pi^\pm, \pi^\pm p)$ at 240 MeV. Top Left: Typical excitation energy spectra. Bottom: Ratios of the $(\pi^+, \pi^+ p)$ to $(\pi^-, \pi^- p)$ cross sections for the ^{14}N ground state for various pion and proton angles (Ref. 9). The solid curve is a theoretical calculation based on the Δ -hole model using the diagrams indicated at the top right (Ref. 13). The dashed curve represents the ratio based on the impulse approximation.

$(\pi^+, \pi^+p)/(\pi^-, \pi^-p)$ for the ^{15}N ground state. The ratios are shown for three pion angles and three proton angles with θ_0 corresponding to the quasifree π -N angle. The absence of the expected minimum at θ_0 for p-shell knockout merely reflects the large solid angle (~ 250 msr) effects. Based on the impulse approximation we would expect the ratio to reflect the free π^+p/π^-p cross section ratio of about nine. Shown for the $\theta_p = \theta_0$ data are the expected ratios based on a factorized DWIA calculation (dashed curves). The experimental ratios differ greatly from these values ranging from $\sim 1/4$ to ~ 4 times the DWIA predictions.

The origin of this discrepancy in ratio has been satisfactorily explained by Hirata, Lenz, and Thies¹³ by including Δ -N interactions in which the nucleon is ejected from the nucleus. The diagrams they considered are shown in Fig. 8 and predictions for the ratios are shown as solid lines. These calculations quantitatively reproduce the major features of the experimental ratios. The inclusion of the Δ -N knockout terms primarily affects the (π^-, π^-p) calculation, since the impulse approximation (IA) term is small. The effect on (π^+, π^+p) is $\lesssim 10\%$ compared to the IA term alone, again providing some small justification for the use of the factorized DWIA in the analysis of (π^+, π^+p) . Clearly, such calculations would be totally inappropriate for (π^-, π^-p) .

Note that such Δ -N knockout terms will certainly play an important role in intermediate energy electromagnetic reactions. For example (γ, n) reactions at energies around 300 MeV can be expected to have important contributions from such terms.

Overall, I would say that most of the qualitative features of quasifree π -nucleus reactions are understood. However, a good description of the experimental data requires a quite sophisticated and difficult theoretical calculation. Thus, I personally do not see pion quasielastic scattering as providing new nuclear structure information, but rather it provides a testing ground for reaction theories, particularly those involving Δ -formation. Much of what is learned from such reaction dynamics studies can then be carried over to reactions with other probes such as electromagnetic probes.

III. PION ABSORPTION

As stated in the introduction, the other major pion-nucleus reaction channel is that of absorption. In the absorption process the pion gives up its total energy to the nucleus, or at least to some of the nuclear constituents. This process cannot happen on the free nucleon so that the elementary process is that of $\pi^+d \rightarrow p+p$. In addition, although kinematically the pion can absorb on a single nucleon bound in a nucleus, the probability is extremely small, requiring the presence of a bound nucleon with a very high momentum ($k \sim 2k_F$). Thus, little of the absorption cross section comes from a single nucleon process, and must arise from absorption on two or more nucleons.

Shown in Fig. 9 is the diagram for two-nucleon absorption generally expected to dominate at pion energies near the Δ resonance. Kinematically, absorption of an energetic pion by two nucleons nearly at rest leads to back-to-back protons with quite high momenta ($\gtrsim 500$ MeV/c). If the

interaction were short-ranged, absorption studies would clearly probe very short distances (< 0.5 fm), and studies of $(\pi, 2N)$ reactions have often been considered as a possible method for probing short-range correlations in nuclei. Unfortunately,

estimates based on the diagram in Fig. 9, including recoil effects, indicate an effective range of 1.5 to 2.0 fm.¹⁴ Therefore, we expect $(\pi, 2N)$ reactions to be quite sensitive to longer range correlations, and hopes of extracting information on short-range correlations rest on a thorough understanding of the reaction dynamics.

As in the case of quasielastic scattering, I will discuss absorption by considering representative experiments with increasing

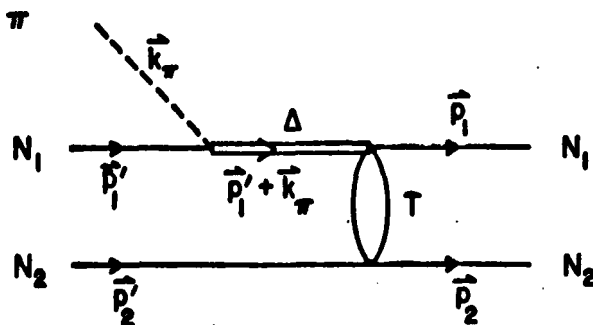


FIG. 9. Diagram expected to dominate two-nucleon absorption near resonance.

levels of detail and sophistication. Each of these experiments provides information on the basic absorption mechanism, including the important question of how many nucleons are involved directly in the absorption of pions by nuclei.

A. Global Studies of Pion Absorption

An extensive study of the $A(\pi^\pm, p)$ reaction has been carried out by McKeown et al.^{15,16}. They measured inclusive proton cross sections for a series of nuclei (C to Ta) with both π^+ and π^- at several energies between 100 and 220 MeV. In Fig. 10 are shown typical spectra for Ni at two angles. Note that at small angles protons with energies above the beam energy of 160 MeV must arise from pion absorption, while for lower energies quasifree knockout can contribute.

The authors used these data to plot contours of the invariant cross section on a rapidity plot, a typical example of which is also shown in Fig. 10. These contours tend to lie on a circle with its center displaced from zero rapidity, although close inspection of the rapidity plot indicates that the contour is flattened in the region of low rapidity. If the data do lie on a circle, one can attribute the source of the protons to the decay of hot matter in thermal equilibrium moving with a velocity corresponding to the central rapidity. With this assumption McKeown et al. fit circles to each data set to find the source rapidity. They next assumed that the pion transferred all of its energy and momentum to a group of N nucleons in the nucleus creating the hot moving source. It is then a simple matter to extract the number N of nucleons involved in the absorption process. Based on these assumptions they find 3 to 5 nucleons are involved in the absorption process depending on mass and energy. They further argue that because of the relatively long mean free path of protons in nuclei (~ 5 fm), the additional nucleons, beyond the expected pair, cannot arise from final state interactions but must originate directly from the absorption process. They also point out that the idea of many-body absorption ($N > 2$) is consistent with the unexpectedly large yield of protons

from π^- induced reactions — much more than would be expected from a two-nucleon absorption process with final state interactions.

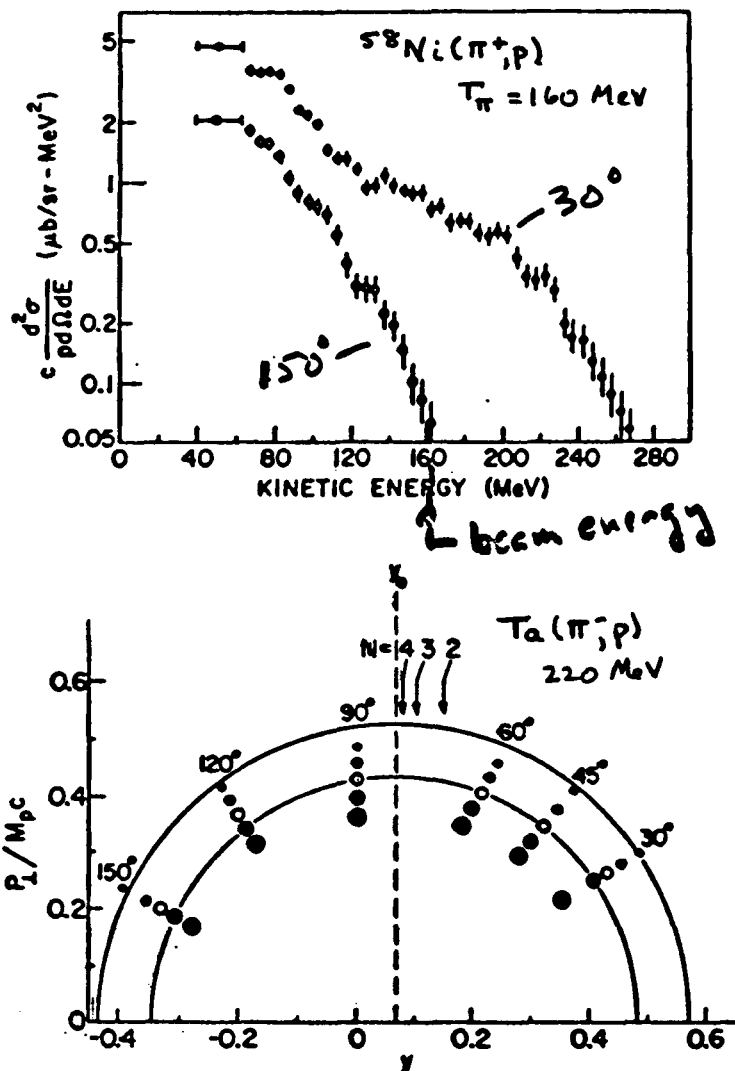


FIG. 10. Results for inclusive $A(\pi, p)$ measurements (Ref. 15). Top: $^{58}\text{Ni}(\pi^+, p)$ cross sections at 160 MeV for two angles. Bottom: Rapidity plot for $\text{Ta}(\pi^-, p)$ at 220 MeV.

Although one can disagree with the conclusions from several standpoints (e.g., the assumption of thermal equilibrium since the contours deviate from circles, the assumption that initial state pion quasifree scattering or final state interactions are not very important,¹⁷ etc.), this work has raised a number of interesting questions and has stimulated great interest in the pion absorption

mechanism. Furthermore, the use of rapidity plots is an excellent method of presenting extensive experimental results in a simple graph. Finally, in contrast with my opinions nine months ago, I now believe that multi-nucleon absorption ($N > 2$) is, in fact, a significant component of absorption. I will show further evidence of this in the remainder of my talk.

$^{12}\text{C}(\pi^+, 2p)$
@ $T_\pi = 165 \text{ MeV}$

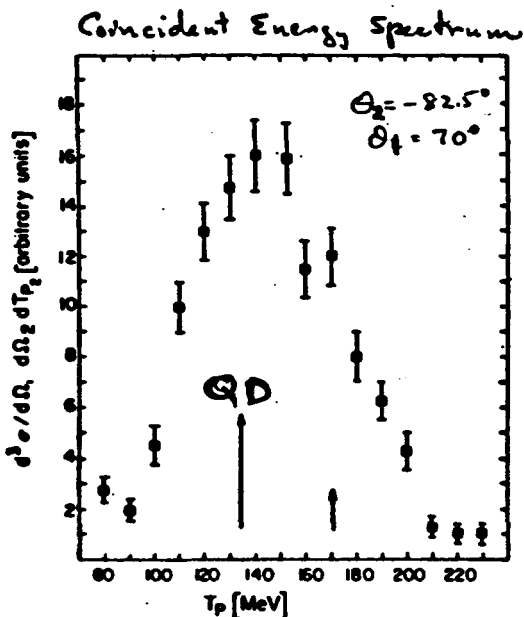
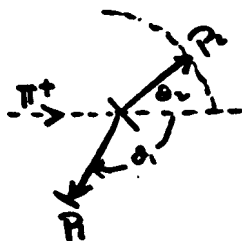
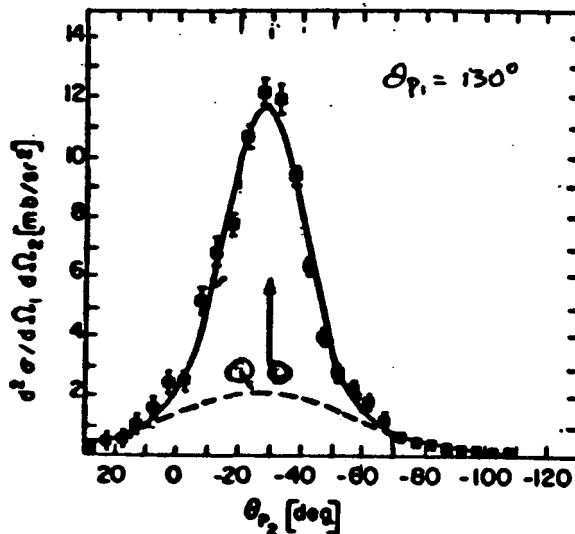


FIG. 11. Results from $^{12}\text{C}(\pi^+, 2p)$ measurements at 165 MeV (Ref. 18). Top: Coincidence energy spectrum for one angle pair. Bottom: Angular correlation at 130° from integrating the energy spectra at each angle pair. The curve is a fit with the sum of two Gaussians.



The obvious extension of this work is the study of N-N correlations. An extensive study of the $(\pi^+, 2p)$ reaction has been made by the Tel Aviv group.^{18,19} The type of experimental measurements are schematically depicted in Fig. 11. One detector was fixed at θ_1 and a second detector (actually arrays of detectors) was moved over a large range of θ_2 . A typical energy spectrum for $^{12}\text{C}(\pi^+, 2p)$ at one pair of angles is shown in Fig. 11. These data show a clear peak located at the quasideuteron energy indicated by the arrow labeled QD; i.e., the energy corresponding to the two-body $\pi^+ + d \rightarrow 2p$ reaction.

The energy spectrum for each angle pair was integrated over the acceptance of the detectors, subject to some energy cut, and for each θ_1 the double differential cross section $\frac{d^2\sigma}{d\Omega_1 d\Omega_2}$ was plotted as a function of the angle of the second proton θ_2 . One such result is also shown in Fig. 11. Again the data are strongly peaked at the quasideuteron angle (indicated by the arrow), the angle corresponding to $\pi^+ + d \rightarrow 2p$. The authors fit the angular correlation data with the sum of two Gaussians, one narrow and one broad. They assume that the narrow Gaussian represents two-nucleon absorption and the broad Gaussian represents some more complicated absorption process.

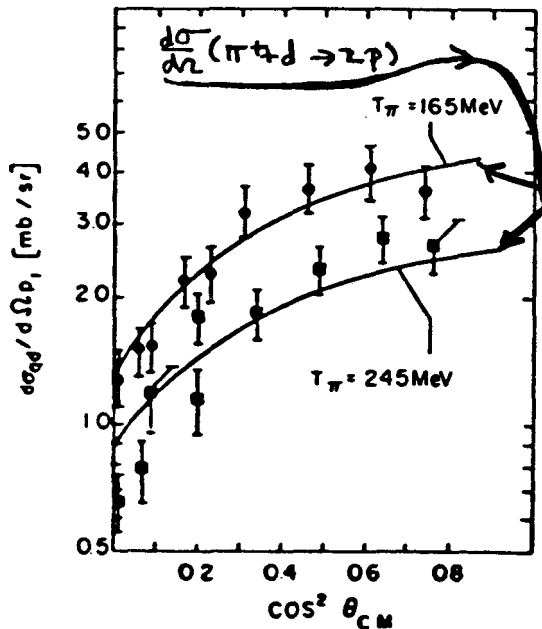


FIG. 12. Differential cross sections for $^{12}\text{C}(\pi^+, 2p)$ obtained by integrating the narrow Gaussian component of the angular correlations such as that shown in Fig. 11. The solid curves represent $\pi^+ + d \rightarrow 2p$ normalized to the data (Ref. 18).

Using the narrow Gaussian to integrate over $d\Omega_2$ Altman et al.¹⁸ arrive at a single differential cross section $\frac{d\sigma}{d\Omega_1}$ for two-nucleon absorption. This result as a function of angle for two incident energies is shown in Fig. 12. The curves correspond to the free $\pi^+d \rightarrow 2p$ cross sections normalized to the data.

Based on the results in Figs. 11 and 12 one might conclude that indeed two-nucleon absorption, similar to the elementary $\pi^+d \rightarrow 2p$ process, is the dominant mechanism involved in absorption of pions by nuclei. The width of the distributions in both energy and angle then represents the momentum distribution of the quasideuteron in the nucleus. This would be in contrast to the conclusions reached in the inclusive studies. However, when one integrates the angular distributions of Fig. 12 one obtains only about 10% of the true absorption cross section. Estimates of the effects of final state interactions based on either simple mean free path arguments or DWIA calculations would only double this result to about 20%. Data for other nuclei are presented in Fig. 13 and the fractions of the absorption cross section, including final state interaction estimates, shown in Table I. Note that the width of the angular correlation increases by less than 50% from ^{12}C to ^{209}Bi suggesting that the quasideuteron has similar momentum distributions for all nuclei. I will return to this point later in my talk. The work presented in Refs. 18 and 19 thus suggests that two-nucleon absorption accounts for only about 20% of the total absorption cross section, independent of the target nucleus when one includes estimates of the effects of final state interactions. This is more or less consistent with the inclusive measurements, and suggests that two-

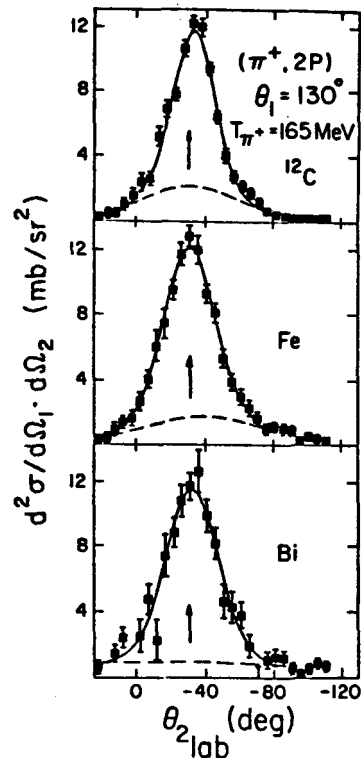


FIG. 13. Angular correlation data for exclusive $A(\pi^+, 2p)$ measurements at 165 MeV (Ref. 19).

nucleon absorption is a relatively small component of the absorption process.

TABLE I. Yields for $A(\pi^+, 2p)$ at 165 MeV obtained from the "narrow" Gaussian fit to angular correlation data.

Target	$\sigma(\pi^+, 2p)$ [Narrow] $\sigma_{\text{absorption}}$	Ratio Corrected for FSI
^{12}C	0.09 ± 0.02	$0.19 \rightarrow 0.28$
Fe	0.05 ± 0.01	$0.14 \rightarrow 0.21$
Bi	0.02 ± 0.004	$0.15 \rightarrow 0.23$

I have some reservations concerning the $A(\pi^+, 2p)$ results and conclusions presented by these authors. Firstly, the assumption that the broad Gaussian contains no two-nucleon absorption contributions is suspect. Ritchie et al.²⁰ using a quasideuteron model have shown that absorption on an L=2 quasideuteron leads to a broad distribution comparable to the broad Gaussian shown in Fig. 11. Recent high resolution data for ^{16}O confirm the shape predicted by the quasideuteron model for L=2 deuterons.²¹ However, even including the contributions from the broad Gaussian, it is difficult to see that one could get more than a factor of two increase in cross section still leaving more than 50% of the cross section missing. Secondly, if these data represent a large fraction of the $(\pi^+, 2p)$ yield, then in what channel is the absorption cross section? We know that the (π^+, pn) yield²² and the $(\pi^+, p \text{ cluster})$ yields²³ are small compared to $(\pi^+, 2p)$. This suggests that much of the yield lies below the experimental energy thresholds. Finally, in my opinion, one of the problems with these data is that the energy thresholds lie significantly higher than indicated in the papers, leading effectively to a rather high missing mass cut. This has led to a great deal of confusion amongst those scientists working in the field. I will provide experimental results which support this conclusion when I present more detailed experimental results for $^{58}\text{Ni}(\pi^+, 2p)$.

B. More Detailed Studies of Pion Absorption

Now I would like to turn to $(\pi^+, 2p)$ results which contain additional detail in terms of energy information. Burger et al.²⁴ have studied the $^{58}\text{Ni}(\pi^+, 2p)$ reaction at 160 MeV using an array of NaI detectors. This system provided a missing mass resolution of about 6 MeV. Shown in Fig. 14 is a contour plot of the yield for $^{58}\text{Ni}(\pi^+, 2p)$ at $\theta_1/\theta_2 = +75^\circ/-75^\circ$ (the quasideuteron angle pair) as a function of the energies of the outgoing protons. The diagonal line corresponds to an axis of excitation energy in the residual nucleus. One observes a large yield in the region of low excitation energy corresponding to bound or nearly bound states in ^{56}Co . This yield must arise predominantly from two-nucleon absorption based on energetics. In addition, we see a significant yield for low outgoing proton energies ($T < 100$ MeV) corresponding to many-particle final states.

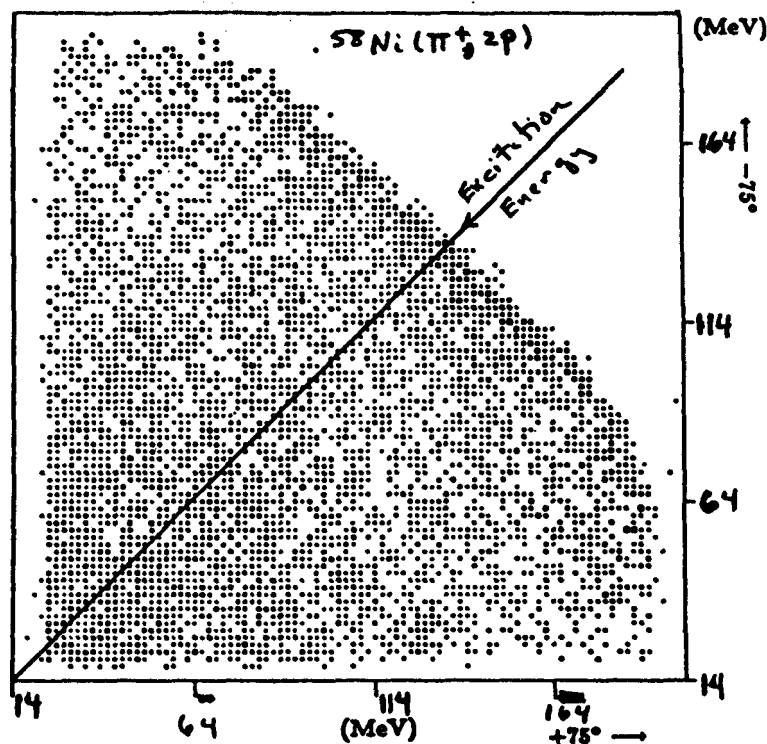


FIG. 14. Energy contour plot for $^{58}\text{Ni}(\pi^+, 2p)$ at 160 MeV. The two protons were detected at $\pm 75^\circ$. The diagonal line corresponds to an excitation energy axis in ^{56}Co (Ref. 24).

Projecting such data onto the excitation energy axis we arrive at the results presented in Fig. 15 for three angle pairs with θ_1 fixed at 75° and θ_2 at the quasideuteron angle (-75°) and 25° and 50° away from this angle. We immediately observe that the region of low excitation energy (< 50 MeV), which we believe arises from two-nucleon absorption, is strongly angle dependent, peaking at the quasideuteron angle. The region of high excitation energy (~ 160 MeV) on the other hand is nearly angle independent and clearly indicates more complicated reaction mechanisms. Unfortunately, these results do not distinguish between various more complicated mechanisms such as multi-nucleon absorption ($N > 2$) and initial or final state interactions.

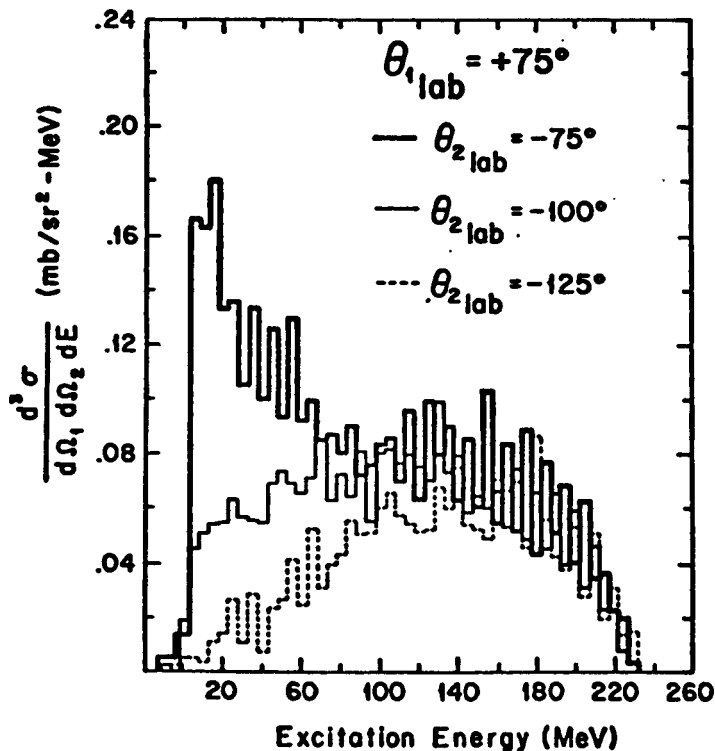


FIG. 15. Excitation energy spectra for $^{58}\text{Ni}(\pi^+, 2p)$ at 160 MeV for $\theta_1 = +75^\circ$ and three θ_2 . The $\pm 75^\circ$ setting corresponds to the quasideuteron angle (Ref. 24).

Binning these data into three regions of excitation energy we obtain the angular distributions in Fig. 16. This figure simply details the results indicated by Fig. 15, and shows the strong angular dependence for low excitation energy ($T_1 + T_2 > 230$ MeV) versus the rather flat angular dependence for high excitation energies ($T_1 + T_2 < 160$ MeV). The intermediate excitation energy region,

chosen to insure true pion absorption ($T_1 + T_2 > \text{beam energy}$), gives rise to a significantly broader angular correlation. It is clear from these results that the width of the angular correlation will depend reasonably sensitively upon the excitation energy cut.

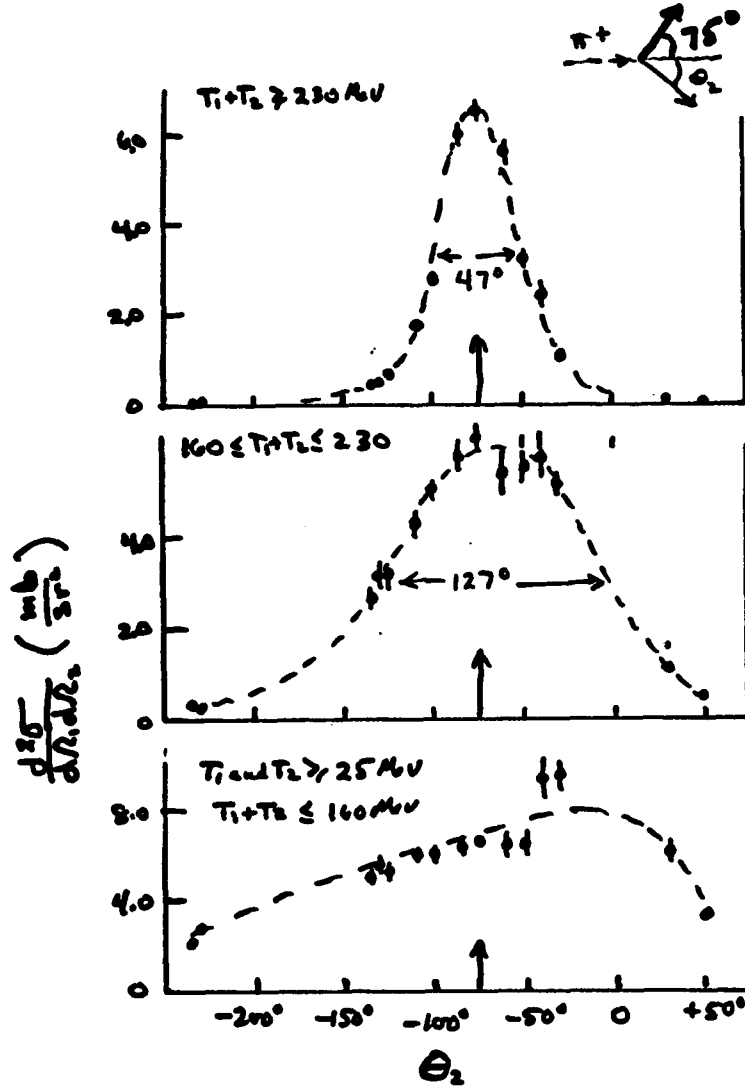


FIG. 16. Angular correlations for $^{58}\text{Ni}(\pi^+, 2p)$ at 160 MeV with $\theta_1 = +75^\circ$. The three distributions correspond to various non-overlapping excitation energy cuts (Ref. 24).

For comparison purposes with other published results (and disregarding what I stated previously about the validity of such an analysis) we also analyzed the angular correlation data with a sum of two Gaussians. This analysis is shown in Fig. 17 where for this

display all data above a specified (T_1+T_2) have been summed. The two Gaussians provide a good fit to the data and a convenient method of integrating the results. Burger et al.²⁴ find that the area of the narrow Gaussian is nearly constant, independent of the excitation region. The broad Gaussian, of course, increases greatly as the excitation energy is increased. Not surprisingly, considering the previous figures, the narrow Gaussian arises almost solely from the low excitation energy region, which we have already suggested should be two-nucleon absorption.

Before continuing the discussion of these data I would like to compare the present ^{58}Ni angular correlations with those obtained by Altman et al.¹⁹ for ^{56}Fe (see Fig. 13). From comparisons of the width, the magnitude of the cross section, and the relative contributions of the narrow and broad Gaussians, I am convinced that the data of Altman et al. represents only about the first 50 MeV or so of excitation energy in the residual nucleus, in contrast to what is implied in the paper.

Returning to ^{58}Ni , integrating the fitted Gaussians Burger finds that the narrow Gaussian follows the $\pi^+ + d \rightarrow 2p$ cross section. Integrating this angular distribution he concludes that the narrow Gaussian represents only $9 \pm 2\%$ of the total absorption cross section. These results are contained in Table II. This fraction is in reasonably

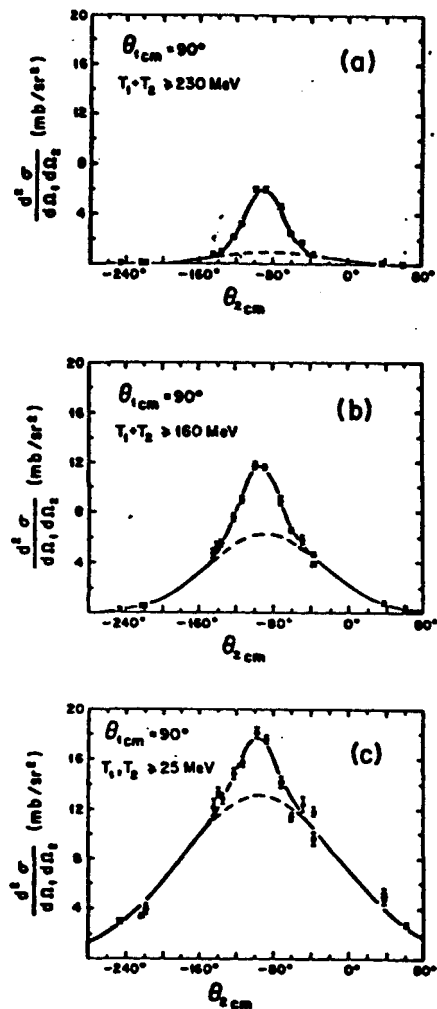


FIG. 17. Angular correlations for $^{58}\text{Ni}(\pi^+, 2p)$ at 100 MeV with lower cuts on excitation energy. The solid curve results from fitting the sum of two Gaussians to the data. The dashed curve shows the contribution from the broad Gaussian (Ref. 24).

good agreement with the Fe results of Altman et al., again suggesting a relatively small contribution from the two-nucleon absorption -- at least in the sense that the yield is represented by the low excitation energy or narrow Gaussian region.

TABLE II. Yields for $^{58}\text{Ni}(\pi^+, 2p)$ at 160 MeV obtained from "narrow" plus "broad" Gaussian fits to angular correlation data.

Energy Cut	$\frac{\sigma_{\text{narrow}}}{\sigma_{\text{absorption}}}$	$\frac{\sigma_{\text{broad}}}{\sigma_{\text{absorption}}}$	$\frac{\sigma_{\text{broad}} + \sigma_{\text{narrow}}}{\sigma_{\text{absorption}}}$
$T_1 + T_2 > 230 \text{ MeV}$	7%	3%	10%
$T_1 + T_2 > 160 \text{ MeV}$	10%	29%	39%
$T_1, T_2 > 25 \text{ MeV}$	9%	97%	106%

Integrating the broad Gaussian, however, one obtains approximately a flat angular distribution and the resultant yield compared to the total absorption cross section is given in Table II. Note that the measured yield is equivalent to the total absorption cross section. Given the offsetting effects of the fact that the present yield contains some events of higher multiplicity versus the fact that the experiment excludes part of the phase space (T_1 or $T_2 < 25 \text{ MeV}$),²⁵ it is reasonable to assume that $(\pi^+, 2p)$ can account for the bulk of the total absorption cross section. This is consistent with the observed small yields for other reaction channels.

We are still left with the fact that the data which one can reasonably safely assign to two-nucleon absorption represents $\lesssim 10\%$ of the total absorption cross section. We can then ask what is the role of initial and final state interactions in two-nucleon absorption. Burger attempted to estimate the magnitude of these effects. Figure 18 shows schematically the estimates he carried out. For the two-body absorption he assumed the presence of quasideuteron in the nucleus with a Gaussian momentum distribution of width 300 MeV/c and that the absorption was given by the free $\pi^+d \rightarrow 2p$ cross section (Fig. 18b). For initial state interactions he allowed a specified fraction of the pions to scatter once according to free π^+p scat-

tering, retaining only those pions which lost less than 50 MeV of kinetic energy (Fig. 18a). These two processes were then combined in a Monte Carlo calculation to compute the angular correlation. The final state interactions depicted in Fig. 18c were estimated by both mean free path arguments and DWIA calculations. Both indicate about a factor of three reduction from final state interactions.

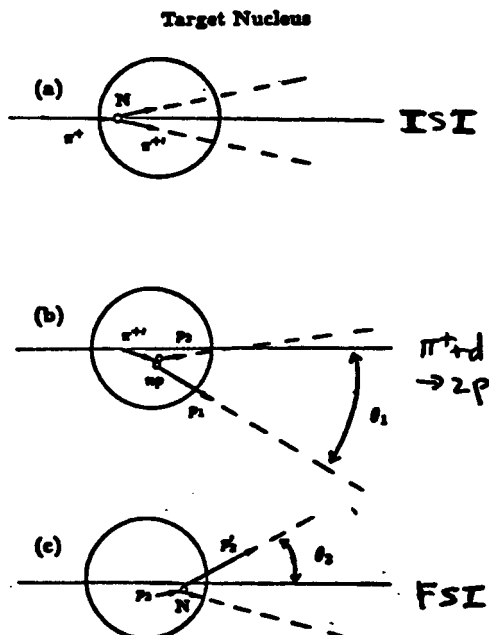


FIG. 18. Schematic representation of (a) an initial state interaction, (b) the absorption on a quasideuteron, (c) and final state interactions (Ref. 24).

The calculated angular correlations for 0% and 30% of the pions scattered before absorption are shown in Fig. 19 compared to the experimental data. We note that initial state interactions lead to a pronounced asymmetry in the angular correlation which is not present in the experimental data for the low excitation energy cut. The physical origin of this asymmetry has been explained by Schiffer²⁶ and arises in part from the fact that absorption after a π -N scattering takes place at a lower energy thereby increasing the opening angle. More sophisticated INC calculations show the same effect.²⁷ From this comparison we conclude that initial state interaction effects are small for the low excitation energy region. It is interesting to note that the angular correlations for higher

excitation energy regions (see Fig. 16) show the development of an asymmetry and may therefore contain events arising from initial state interactions. However, the asymmetry is relatively small. Overall, using the estimates of initial and final state interactions it is difficult to argue that more than approximately 50% of the $(\pi^+, 2p)$ yield arises from two-nucleon absorption.

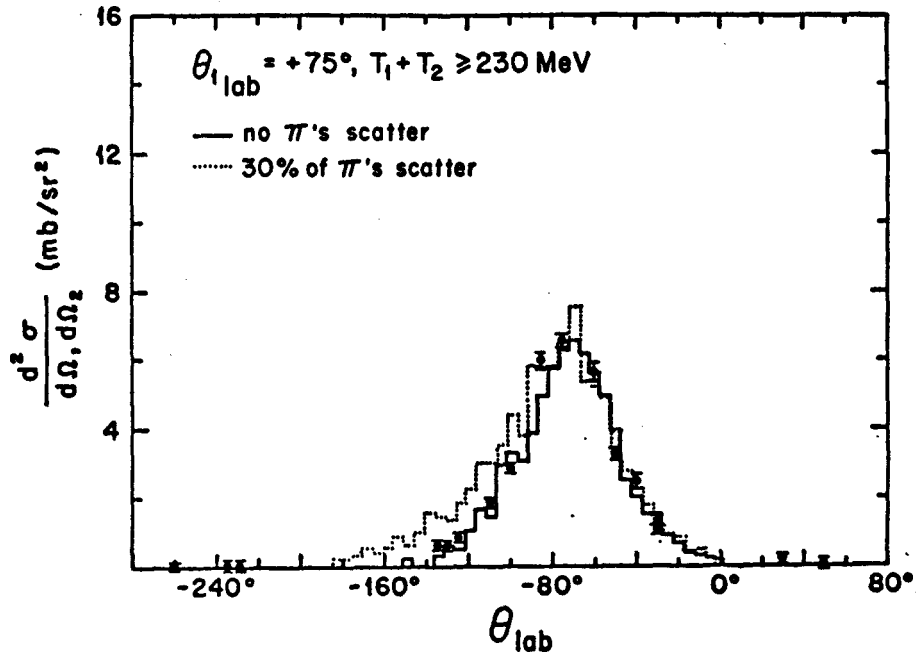


FIG. 19. Angular correlation data for $^{58}\text{Ni}(\pi^+, 2p)$ from the highest energy cut. The histograms are Monte Carlo calculations representing various probabilities for pion scattering before absorption (Ref. 24).

The above study again implies the probable importance of multi-nucleon ($N > 2$) absorption. Two other experiments which suggest similar conclusions are the works of Tacik et al.²⁸ and Backenstoss et al.²⁹ Tacik et al. measured the $^{12}\text{C}(\pi^+, 3p)$ cross sections over a very limited range of phase space. Firstly, their data show no evidence for strong ISI effects. Secondly, integrating their $(\pi^+, 3p)$ data over all phase space (an integration which requires extensive extrapolation), they estimate that the $^{12}\text{C}(\pi^+, 3p)$ reaction contains approximately 39% of the total absorption cross section. Although there is great uncertainty in this value, the results do

suggest a large component of $(\pi^+, 3p)$.

Backenstoss et al.²⁹ have measured a detailed angular correlation for ${}^3\text{He}(\pi^+, 2p)n$. This result is presented in Fig. 20. The angular correlation shows the strong peak expected for two-nucleon absorption. However, in addition there is an underlying component which is consistent with three-body phase space and may well represent three-nucleon absorption. The integrated yield for this component represents 30% of the total absorption cross section. Similar results for ${}^3\text{He}$ have been obtained at TRIUMF.³⁰

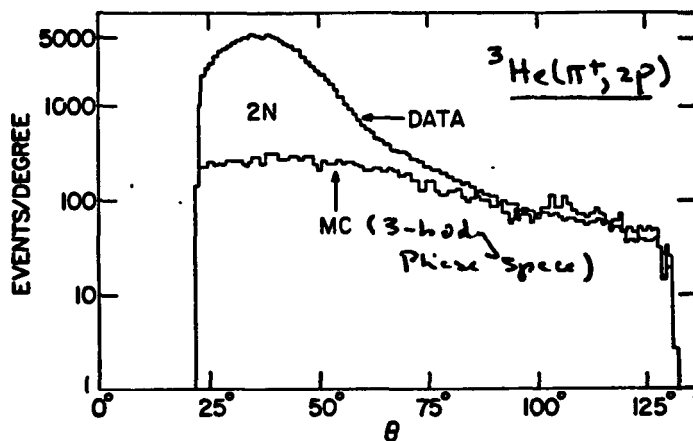


FIG. 20. Angular correlation for the ${}^3\text{He}(\pi^+, 2p)$ at 120 MeV. The curve labeled MC corresponds to a Monte Carlo simulation for three-body phase space (Ref. 29).

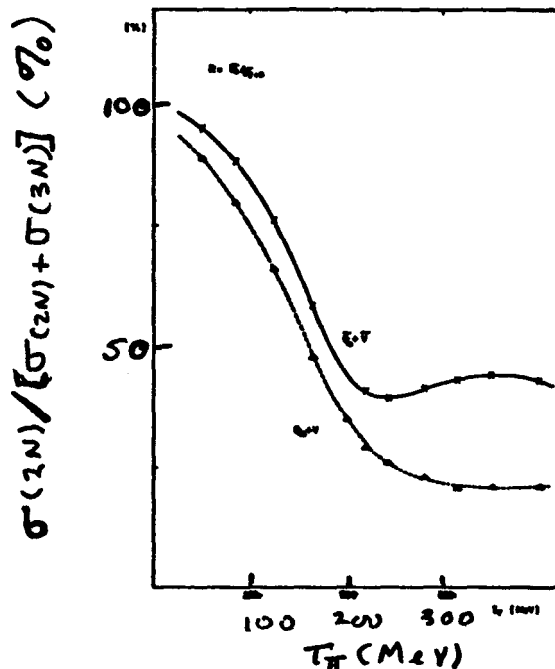
In addition to the experimental support for multi-nucleon absorption, a recent theoretical paper by Oset et al.³¹ reports calculations of the relative importance of two- and three-nucleon absorption. The diagrams considered by these authors are shown in Fig. 21. They find the sequential Δ production dominates the three-body absorption. The calculated ratio of two-nucleon to two- plus three-nucleon absorption for ${}^{12}\text{C}$ is shown in Fig. 21 as a function of the pion bombarding energy. These calculations indicate a large component of three-nucleon absorption, representing approximately 50% of the absorption cross section at resonance energies. This result would be reasonably consistent with present experimental results, although they may underestimate the effects at low energies.

There are two interesting sidelights to these calculations. Firstly, estimates of four-nucleon absorption by Oset et al., pro-

cesses which contain the double Δ -production mechanism proposed by Brown et al.³² to explain the absorption process, indicate that less than 15% of the absorption cross section is contained in such processes. Much of this suppression arises from Pauli blocking. Secondly, such considerations will be of crucial importance for the understanding of $(\gamma, 2N)$ or $(e, e2N)$ reactions. One expects that three-nucleon processes will be of equal importance in such reactions.



FIG. 21. Calculated ratio of three-nucleon to two- plus three-nucleon absorption cross sections as a function of pion energy for the diagrams shown in the figure (Ref. 31).



The present experimental data and theoretical calculations discussed seem to indicate roughly equal contributions from two-nucleon and three(or more)-nucleon absorption. This presents two avenues for continued research on pion absorption. Based on kine-

matics, one can experimentally isolate regions in which two-nucleon absorption dominates. Continued research into this aspect will require experiments with better energy resolution in which individual final states are resolved. In this way one can use particular nuclear states as filters in order to examine the reaction mechanism, and then quantify the role of two-nucleon correlations in the absorption process.

The second avenue of research is the explicit study of three or more nucleon absorption components. These experiments will require explicit measurements of higher multiplicity events using large acceptance devices. Multiplicity measurements for pion absorption are already underway at NIKHEF and Los Alamos using detector systems with $\Omega > 2\pi$ sr. I would hope that these measurements will have sufficient energy and angular resolution to be able to isolate three-body, four-body, etc. final states. I believe that the multiplicity measurements primarily allow the study of reaction dynamics, and will not only clarify the pion absorption mechanism, but will also aid in the understanding of electromagnetic interactions in which hadrons are detected in the final state.

C. Transitions to Specific Nuclear States

As a final topic, I will discuss some $(\pi^+, 2p)$ reactions leading to specific final states (or groups of states) in the residual nucleus. By isolating regions in which the residual $(A-2)$ nucleus is bound, we expect two-nucleon absorption to be dominant.

It is instructive to compare the data to factorized DWIA calculations based on an extreme quasideuteron model. In Fig. 22 we show another "cartoon" diagram which depicts the DWIA calculation. In this model the triple differential cross section for the $A(\pi^+, 2p)B$ reaction to a specific final state in B can be written³³ as

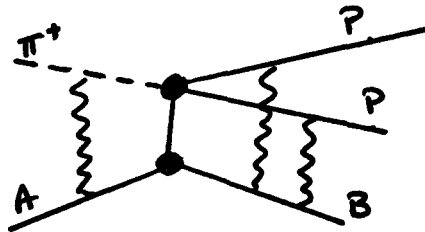


FIG. 22. "Cartoon" DWIA diagram for pion absorption on a quasideuteron.

$$\frac{d^3\sigma}{d\Omega_1 d\Omega_2 dE_1} = KF \cdot \left. \frac{d\sigma}{d\Omega} \right|_{\pi^+ + d \rightarrow 2p} \cdot \sum_L |T_L^\Lambda|^2 \quad (3)$$

where KF is a known kinematic factor and $\left. \frac{d\sigma}{d\Omega} \right|_{\pi^+ + d \rightarrow 2p}$ is the two-body $\pi^+ + d \rightarrow 2p$ cross section. We take this two-body cross section to be equal to the free $\pi^+ + d \rightarrow 2p$ cross section corresponding to either the $\pi^+ + d$ initial state (IEP) or the $2p$ final state (FEP). Due to the Δ -resonance the two prescriptions can lead to quite different values for $\frac{d\sigma}{d\Omega}$. The quantity

$$T_L^\Lambda = (2L+1)^{-1/2} \int \chi_{p_1}^{(-)*}(\vec{r}) \chi_{p_2}^{(-)*}(\vec{r}) \langle \psi_B | \psi_A \rangle \chi_\pi^{(+)}(\vec{r}) d^3r \quad (4)$$

where the χ 's represent distorted waves calculated with phenomenological optical model potentials. The overlap integral for the initial and final nuclear states is taken to be a bound deuteron wave function calculated in a Woods-Saxon potential; i.e.,

$$\langle \psi_B | \psi_A \rangle = S^{1/2} \Phi_{NL}(\vec{r}). \quad (5)$$

The quantum numbers N,L are based on a oscillator shell model in which all oscillator quanta are placed in the n-p center-of-mass motion wave function, the relative motion wave function containing zero oscillator quanta; e.g., for $(f_{7/2})^2$ we would have $4S(L=0)$, $3D(L=2)$, $2G(L=4)$, and $1I(L=6)$ bound deuteron states. The spectroscopic amplitude contains the shell model wave function information (amplitudes, two-particle CFP's, etc.) as well as the ls - jj recoupling coefficients and Moshinsky brackets.

One can again ask why we carry out such simplified calculations -- other than the fact that these are the calculations which we experimentalists know how to do. Firstly, $A(\pi^+, 2p)B$ calculations carried out by Ohta, Thies, and Lee,³⁴ based on the Betz-Lee³⁵ π -NN phenomenological Hamiltonian, indicate that the quasideuteron component dominates the reaction. Shown in Fig. 23 are calculated cross sections for $\pi + 2N[{}^{2S+1}L_J] \rightarrow 2p$ as a function of pion energy.

In their model, the cross section for the $S=1, l=0$ relative n-p wave function is significantly larger than any other n-p component allowed for two p-shell nucleons. For ^{12}C at 170 MeV, the absorption on $^3\text{S}_1$ pairs produce approximately 80% of the absorption cross section calculated with all possible pair configurations.

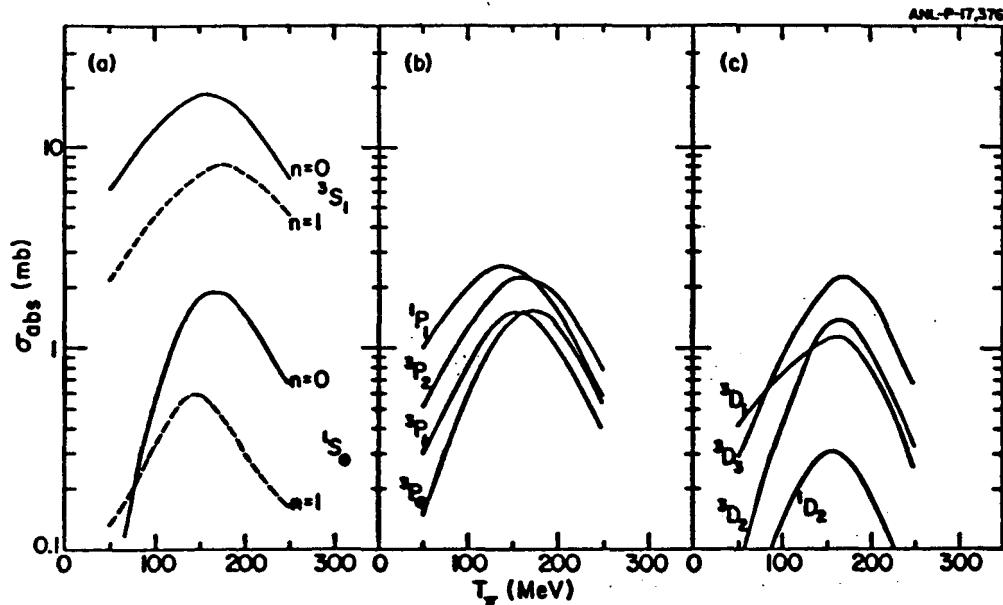


FIG. 23. Absorption cross sections as a function of pion energy for various n-p configurations in the initial state (Ref. 34).

Secondly, in spite of problems related to the simplicity of the model, this calculation at least provides an estimate of the role of initial and final state interactions. To see these effects, in Fig. 24 we present the ratio of distorted wave (DWIA) to plane wave (PWIA) quasideuteron calculations as a function of the pion energy. The calculations are for absorption on an $L=0$ deuteron with equal energies and equal angles for the outgoing protons, chosen such that the momentum of the residual nucleus is zero. On resonance the distortion effects lead to roughly a factor of 5-7 reduction in cross section. In the lower panel calculations for ^{40}Ca are presented in which distortions are included for only the incoming π or only the outgoing protons. One clearly sees the role of the Δ -resonance in the initial state. Near resonance one obtains roughly equal attenuation from the initial state and final state interac-

tions. It is worth pointing out that the case with no pion distortion provides an estimate of the magnitude of attenuation one might expect in a $(\gamma, 2p)$ reaction from final state interactions.

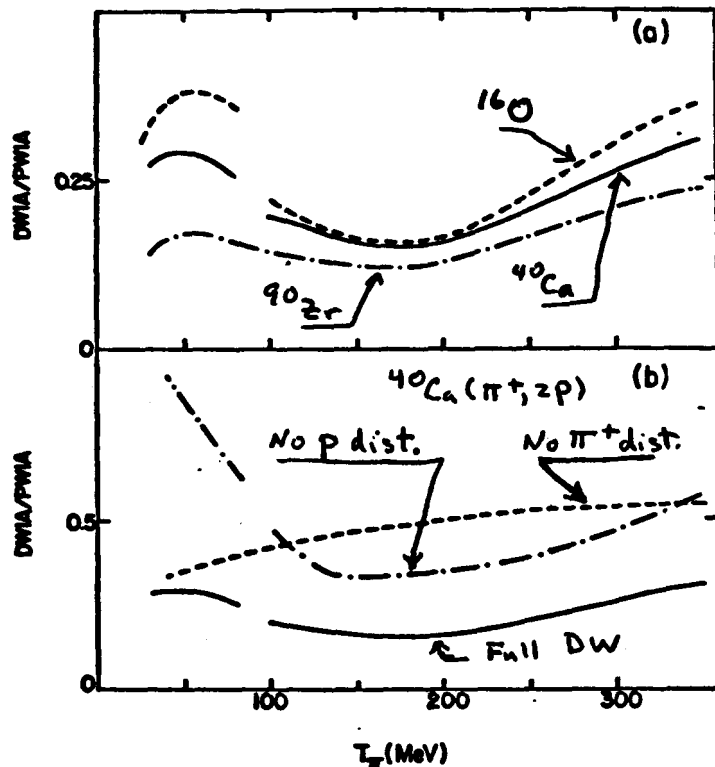


FIG. 24. Ratio of factorized DWIA to PWIA cross sections as a function of pion energy. The calculations are for $L=0$ with $\theta_1=-\theta_2$, $T_1=T_2$ and $\text{Precoil} = 0$. The lower panel exhibits the results for ^{40}Ca when either the pion or proton distortions are excluded (Ref. 33).

Having at least rationalized why we might expect these DWIA calculations to have some relevance, let us now compare them to experimental data. In Fig. 25 are the cross sections³⁶ at 60 MeV for $^6\text{Li}(\pi^+, 2p)$ leading to the ground state of ^4He (upper panel) and the 25 MeV excited state region of ^4He (lower panel). Not surprisingly, considering the $\alpha+d$ cluster structure of ^6Li and weak binding, the DWIA quasideuteron calculation provides an excellent fit to the ground state data both in terms of shape and magnitude (spectroscopic factor of 1.17). For the excited state the two curves and normalizations arise from the two different prescriptions, FEP and IEP, for $\frac{d\sigma}{d\Omega}$ in Eq. (3). In this case the effects are large since the binding energy is about 25 MeV out of the 60 MeV incident energy, an deuteron energy located on the side of the Δ -resonance. We see that the calculations reproduce the shape quite well indicat-

ing that such a model contains the proper kinematic behavior. Unfortunately, the extracted spectroscopic factors are uncomfortably large. I can present an argument in the same DWIA framework for why these should be reduced, but rather than that let us look at other comparisons.

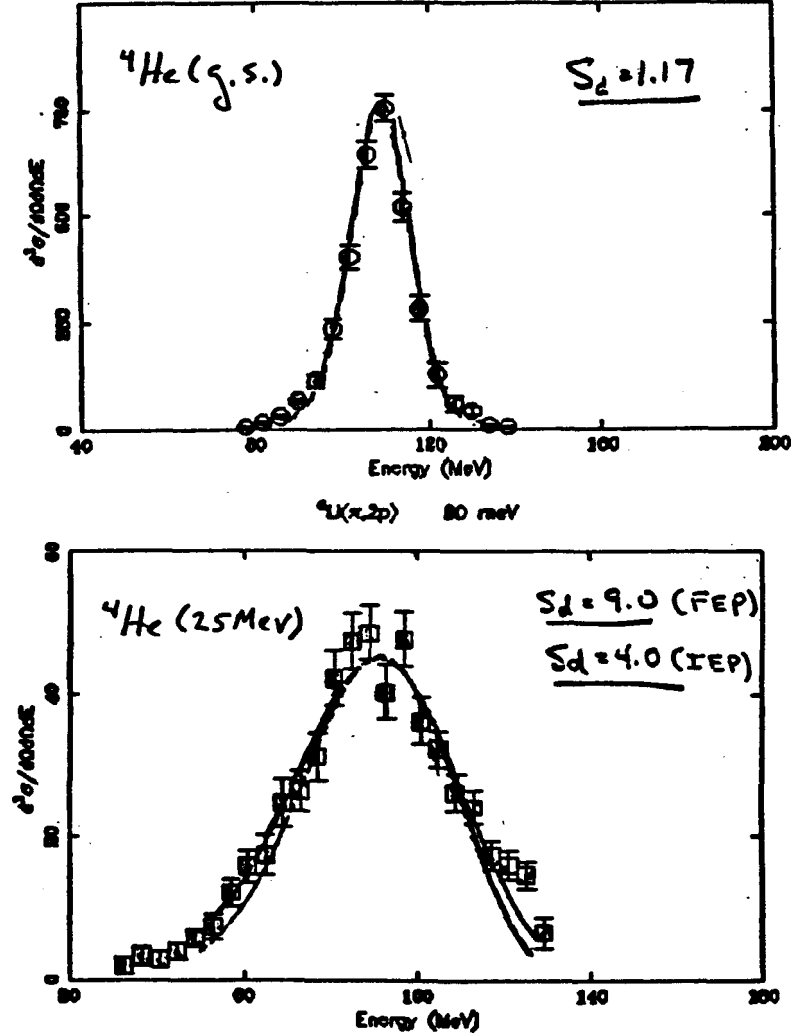


FIG. 25. The ${}^6\text{Li}(\pi^+, 2p){}^4\text{He}$ energy sharing cross sections at 60 MeV (Ref. 36). The curves are factorized quasideuteron DWIA calculations with the normalizations specified.

In Fig. 26 recent angular distributions²¹ for ${}^{16}\text{O}(\pi, 2p){}^{14}\text{N}$ leading to the 3.95 MeV, 1^+ (upper panel) and the 7.03 MeV, 2^+ (lower panel) states are shown. The calculations are DWIA calculations using Cohen-Kurath³⁷ wave functions and must be renormalized by a factor of 6 to 7. Again, these calculations provide a reasonable

fit to the shape, including the asymmetry in the $L=2$ transition which arises from the behavior of the π^+d cross section. (The asymmetry arises from the fact that the π^+d interaction takes place at an effective higher energy for backward angles, and since we are below resonance the cross section is larger.) However, the magnitude is much too small. This would seem to indicate that the model provides a poor description of the data, and perhaps one is seeing the effects of additional two-nucleon correlations. Before one reaches a definite conclusion, I would point out that

this factor of 6 to 7 is not so unlike the "unhappiness" factor that one has used in the past to renormalize two-nucleon transfer calculations. Thus, an understanding of the results will not only require additional data and tests of the $(\pi^+, 2p)$ reaction, but also should be coupled to analyses of other two-nucleon data such as $(p, {}^3\text{He})$.

Finally, let me show the worst case. In Fig. 27 are shown Burger's data for the ${}^{58}\text{Ni}(\pi^+, 2p){}^{56}\text{Co}$ reaction for the first 50 MeV of excitation in ${}^{56}\text{Co}$. The upper panel shows the energy sharing

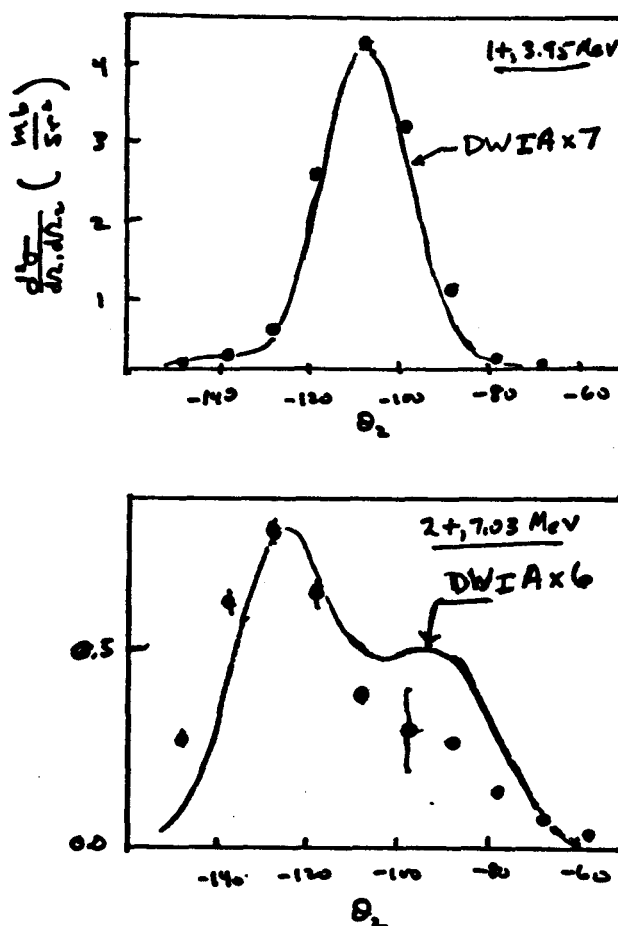
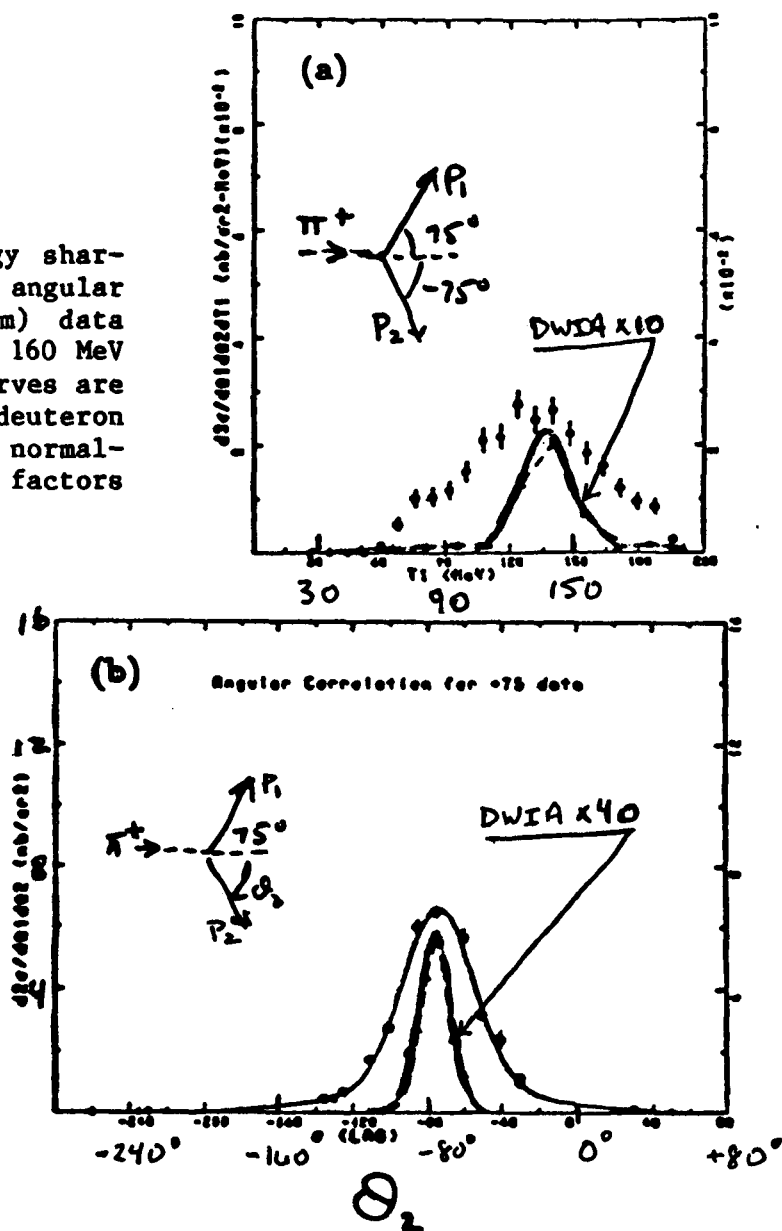


FIG. 26. Angular correlation data for ${}^{16}\text{O}(\pi^+, 2p){}^{14}\text{N}$ at 116 MeV and $\theta_1 = 50^\circ$ leading to specific states (Ref. 21). The curves are factorized quasideuteron DWIA calculations using Cohen-Kurath wave functions renormalized by the factors indicated.

data for $\theta_1/\theta_2 = +75^\circ/-75^\circ$ and the lower panel shows the angular correlation for $\theta_1 = 75^\circ$ where the data have been integrated over T_1 .

FIG. 27. Energy sharing (top) and angular correlation (bottom) data for $^{58}\text{Ni}(\pi^+, 2p)$ at 160 MeV (Ref. 24). The curves are factorized quasideuteron DWIA calculations normalized by the factors indicated.



Quasideuteron DWIA calculations were carried out using the full sum rule limit for the $f_{7/2}$ shell. The $L=0, 2, 4$, and 6 contributions were added incoherently. The results of these calculations, renormalized by the indicated amounts, are shown in Fig. 27. We see that the predicted shape is a factor of two to three too narrow, and

the magnitude is a disaster. The inclusion of additional shell model orbitals would not be sufficient to correct these major disagreements. Thus for medium weight nuclei, the simple quasideuteron model, assuming absorption on $S=1$ n-p pairs in a relative $1s$ state, is totally inadequate.

How does one proceed at this point? Clearly we need a significant theoretical effort, even for this simpler case of two-nucleon absorption, if we hope to extract useful physics information. Work is in progress on this problem.^{34,38} Two simplifications in the present quasideuteron calculations which must be assessed are the following. Firstly, is the limitation of the absorption to $1s$ relative motion for the n-p pair. Even if $l=0$ relative motion dominates, other principle quantum numbers may provide significant contributions. For ^{16}O the p-shell nucleons permit either $1s$ or $2s$ relative motion, and for ^{58}Ni f-shell nucleons permit four possible $l=0$ configurations ($1s$ through $4s$). If configurations other than $1s$ are important one would expect both an increase in cross section from the coherence and an increase in width from the lower oscillator quanta then present in the center-of-mass motion. Both changes increase as we increase the single particle oscillator quanta. From that standpoint I would point out that the extreme quasideuteron model of placing all of the oscillator quanta in the n-p c.m. motion leads to a decrease in the angular correlation width as one proceeds to higher major shells, in contrast to the experimental results of Altman et al.¹⁹ Since the discrepancies in both shape and magnitude between the data and quasideuteron calculations increase with increasing mass, the inclusion of additional terms at least provides corrections which go in the right direction.

Secondly, one must assess the possible coherence between the two different L-transfers possible to an unnatural parity state. Although one can show that these contributions are incoherent if the $\pi^+d \rightarrow 2p$ vertex is diagonal in the channel spin, we know that this is an incorrect assumption. If the reaction is coherent in L we again might expect an increase in the calculated width of the distributions.

A prerequisite to understanding both of the above considerations is the inclusion of a proper treatment of the π -NN vertex, perhaps along the lines of Ref. 34. In addition, our understanding would be enhanced by the availability of additional data in which individual states, or at least groups of states, are isolated so as to provide a more detailed test of the reaction theory. These data should cover a range of nuclei from different major shells.

VI. SUMMARY

I have attempted to review the status of two major pion-induced reactions leading to multiparticle final states using recent experimental results. In my opinion quasifree knockout reactions are reasonably well understood, at least qualitatively. However, the theoretical calculations are very difficult and at a minimum must properly include the Δ . The extraction of new information from such studies will require both improved experiments with good energy resolution and detailed, difficult theoretical calculations.

Pion absorption on the other hand remains poorly understood. We are still uncertain as to the role of multi-nucleon absorption, initial state interactions, and final state interactions. Experimental advancements can be made by work with improved energy resolution, phase space coverage, and clever choice of targets. However, the whole area is in desperate need of theoretical support and calculations, if we are to extract the important underlying physics available in these reactions.

ACKNOWLEDGMENTS

This work was supported in part by the National Science Foundation.

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ELECTRON SCATTERING FROM NUCLEONS AND NUCLEI IN THE SKYRME MODEL

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Abstract

The Skyrme model of the baryons as topological solitons of meson fields may be viewed as an approximate boson representation of QCD that is applicable at low energies. The model and its generalizations permit complete construction of the electromagnetic current operators of nucleons and nuclei - including the exchange current operators - as well as the nuclear force itself. The isovector current is a Noether current of the model Lagrangean, whereas the transverse isoscalar current is topological and model independent. Because of this model independence the Skyrme approach is particularly well suited for investigation of nuclear electromagnetic isoscalar observables.

1.Introduction

The Skyrme model for the baryons, in which the nucleons and nuclei are formed as topological solitons of pion and scalar meson fields [1] describes many of the important features of the nucleons and nuclear force in a qualitatively correct way [2,3,4,5]. As the Skyrme model may be viewed as an effective boson representation of QCD in the large color limit at low energies [6] it provides a theoretically appealing starting point for the calculation of nuclear properties at low and possibly intermediate energies.

The Skyrme model and its generalizations are based on a nonlinear Lagrangean of meson fields [1,7,8] and an anomalous current which carries the baryon number as a topological charge. The latter is not a Noether current of the Lagrangean and hence independent of the detailed dynamical model. It is related to the Wess-Zumino interaction [6,9], which does not contribute to the energy in the case of SU(2). As it is not quite clear which interaction terms the Lagrangean should contain in order to represent the low energy large color limit of QCD the anomalous baryon current operator, which depends on the fields but not on their interactions will be of exceptional value in settling the utility of the whole approach.

In this talk I shall therefore focus the discussion on the consequences of the anomalous baryon current on nuclear electron scattering observables. The

anomalous baryon current determines the transverse isoscalar current operators of nuclei, including the exchange current components, in a model independent way. This is particularly interesting as hitherto a theoretical basis for such currents has been lacking. The transverse isoscalar exchange currents have been constructed from a variety of more or less adhoc meson exchange mechanisms. The topological soliton approach makes firm predictions for these current operators which can be confronted with experiment. This will be illustrated by a model independent calculation of the charge and magnetic form factors of the deuteron.

2. The isoscalar current of the nucleon

In the Skyrme approach with a single SU(2) field U made up of a scalar σ and isovector pseudoscalar pion field π :

$$U = \sigma + i\vec{\tau} \cdot \vec{\pi}, \quad (1)$$

the anomalous baryon current operator has the form [1]

$$B^\mu = \frac{\epsilon^{\mu\nu\alpha\beta}}{24\pi^2} \text{Tr} \{ U^\dagger \partial_\nu U U^\dagger \partial_\alpha U U^\dagger \partial_\beta U \}. \quad (2)$$

For single nucleons the topological soliton field has the hedgehog form

$$U(r) = \exp i\vec{\tau} \cdot \hat{r} \theta(r). \quad (3)$$

Here $\theta(r)$ is a function (the chiral angle), which satisfies an Euler-Lagrange equation of motion the form of which depends on the Lagrangean. The form of the current (2) on the other hand does not depend on the Lagrangean but is determined by θ alone irrespective of the details of the Lagrangean. To avoid the need to specify the Lagrangean one may determine θ directly from the electromagnetic structure of the nucleon and then use that chiral angle to predict other isoscalar electromagnetic observables. This approach will be illustrated here for the case of the electromagnetic form factors of the deuteron.

In order to construct an e.m. current operator from (2) that has the proper spin and isospin structure the spin-isospin degeneracy must be broken by rotating the hedgehog field (2) according to [2] as

$$U(r) = A(t)U(r)A^\dagger(t). \quad (4)$$

Here $A(t)$ is a time dependent SU(2) matrix, the elements of which may be described in terms of 3 Euler angles, which are the coordinates of the baryon

wavefunctions. To obtain a current operator that has the conventional form in terms of Pauli spin and isospin matrices for the nucleons instead of the cumbersome Euler angles we utilize the projection theorem [10]

$$\langle N' | A \underline{\tau} A^\dagger | N \rangle = -\frac{1}{3} \langle N' | \underline{\sigma}^N (\underline{\tau} \cdot \underline{\tau}^N) | N \rangle. \quad (5)$$

Here N and N' represent single-nucleon states and $\underline{\sigma}^N$ and $\underline{\tau}^N$ are the spin- and isospin operators for a single nucleon.

The isoscalar charge and (spin) current operators for the nucleon then take the form

$$\rho(\vec{r}) = -\frac{1}{4\pi^2 r^2} \sin^2 \theta(r) \theta'(r), \quad (6a)$$

$$\vec{j}(\vec{r}) = \frac{1}{8\pi^2 \lambda r^2} \sin^2 \theta(r) \theta'(r) \underline{\sigma} \times \vec{r}. \quad (6b)$$

Here λ is the moment of inertia of a single soliton. It can be calculated from the Lagrangean if that is known, but here we treat it as a parameter to be determined so as to obtain the correct value for the isoscalar magnetic moment of the nucleon.

The isoscalar electric and magnetic form factors of the nucleon are obtained from the Fourier transforms of the operators (6) as

$$G_E^S(q) = -\frac{2}{\pi} \int dr j_0(qr) \sin^2 \theta(r) \theta'(r), \quad (7a)$$

$$G_M^S(q) = \frac{4m_N}{3\pi\lambda} \int dr [j_0(qr) + j_2(qr)] \sin^2 \theta(r) \theta'(r). \quad (7b)$$

The chiral angle θ may be determined by inverting the Fourier transform (7a) if $G_E^S(q)$ is known. With the conventional dipole form factor

$$G_E^S(q) = \frac{1}{(q^2 + \Lambda^2)^{-2}}, \quad (8)$$

with $\Lambda^2 = 0.71 \text{ GeV}^2$, one obtains the following equation for θ

$$\theta(r) - \frac{1}{2} \sin 2\theta(r) = \pi e^{-\Lambda r} (1 + \Lambda r + \frac{1}{2} \Lambda^2 r^2). \quad (9)$$

The chiral angle is then found by numerical solution of this equation. In Fig.1 we compare the values for θ determined in this way to the result given by the Skyrme model Lagrangean, when augmented by a pion mass term [11]. The results are qualitatively similar, but should not be expected to be too close as the Skyrme model result does not include the vector meson degrees of freedom, which are implicit once the dipole form factor is used as the starting point. For large distances $\theta(r)$ behaves as $r^{2/3} e^{-\Lambda r/3}$ when using the dipole form factor rather than as $e^{-m_\pi r}/r$ as suggested by the Skyrme model.

In Figs.2 and 3 we show the calculated isoscalar form factors as obtained in the Skyrme model, and compare them to the 5 parameter dipole fit of Ref.[12]. The agreement is reasonably good. Similar form factor calculations have recently been presented in Ref.[13].

It is worth noting that the form factors (7) are related as

$$G_M^S = - \frac{2m_\pi}{\lambda} \frac{dG_E^S}{dq^2}, \quad (10)$$

a relation that is amenable to experimental test.

3. The magnetic form factor of the deuteron

In order to construct the isoscalar current operator for the two-nucleon system one may use the product ansatz

$$U(\vec{r}_1, \vec{r}_2, \vec{r}) = U(\vec{r} - \vec{r}_1)U(\vec{r} - \vec{r}_2), \quad (11)$$

for the soliton field with baryon number 2 [1]. Here $\vec{r}_1$ and $\vec{r}_2$ are the centers of the two solitons and $\vec{r}$ is the point of interaction with the e.m. field. With this ansatz the baryon current splits into a sum of two single current operators and an exchange current operator as

$$B^\mu(\vec{r}_1, \vec{r}_2, \vec{r}) = B_1^\mu(\vec{r} - \vec{r}_1) + B_2^\mu(\vec{r} - \vec{r}_2) + B_{ex}^\mu(\vec{r}_1, \vec{r}_2, \vec{r}). \quad (12)$$

Here B_i^μ is the baryon current for a single nucleon considered in section 2 above and B_{ex} the irreducible two-body exchange current operator.

The spatial part of the exchange current operator B_{ex} in (12) may be brought into the form [10]

$$B_{ex}^\mu(\vec{r}_1, \vec{r}_2; \vec{r}) = \frac{1}{g_A^2} \nabla \times \text{Tr} \{ \dot{U}_1 U_1 U_2 \nabla U_2^\dagger - \dot{U}_2 U_2^\dagger U_1^\dagger \nabla U_1 \}, \quad (13)$$

where we have used the abbreviation

$$U_k = U(\vec{r} - \vec{r}_k), (k = 1, 2). \quad (14)$$

Here $U(\vec{r} - \vec{r}_k)$ is the hedgehog field for nucleon k . A nontransverse term that does not contribute for nucleonic states has been left out from (14). The spin and isospin dependence of the exchange current operator may be obtained by rotating both solitons separately according to eq.(4). This procedure finally yields the exchange current operator

$$\vec{j}_{ex} = -\frac{i}{36\pi^2\lambda} \vec{\tau}_1 \cdot \vec{\tau}_2 e^{i\vec{q} \cdot \vec{r}/2} \vec{q} \times \int d^3R e^{i\vec{q} \cdot \vec{R}} \sin^2 \theta(R) \left[\frac{\sin 2\theta(R)}{2R} (\vec{\sigma}^1 + \vec{\sigma}^2) + \left(\theta(R) - \frac{\sin 2\theta(R)}{2R} \right) (\vec{\sigma}^1 + \vec{\sigma}^2) \cdot \hat{R} \hat{R} \right]. \quad (15)$$

Here $\vec{r}$ is the internucleon separation.

To evaluate the matrix elements of the combined current operator (single soliton currents and the exchange current) one has to use a realistic model for the deuteron wavefunction. Here I have used the wavefunctions that correspond to the parametrized Paris potential [14]. The results are very similar when the Reid soft core interaction wavefunctions [15] are used instead. In Fig.4 the calculated magnetic form factor of the deuteron is shown when obtained using the dipole form factor for G_E^d to determine the chiral angle [16]. At large values of momentum transfer the two-body contribution dominates the form factor. The agreement with the empirical form factor values from refs. [17,18] is remarkable. For comparison it should be mentioned that when using the Skyrme model [11] to determine the chiral angle one overpredicts the magnetic form factor (Fig.5). This may be an indication for the need to include explicit vector meson degrees of freedom in the Skyrme model along the lines suggested in the refs.[7,8].

4. The charge form factor of the deuteron

The time component of the exchange current term in the separation (12) is [10]

$$B_{ex}^0(\vec{r}_1, \vec{r}_2; \vec{r}) = \frac{i}{8\pi^2} \vec{q} \cdot \int d^3R e^{i\vec{q} \cdot \vec{R}} \text{Tr} \{ A_1 U_1^\dagger \nabla U_1 A_1^\dagger \times A_2 U_2 \nabla U_2^\dagger A_2^\dagger \}. \quad (17)$$

Here the A 's are the time dependent rotation operators for the two solitons

introduced in eq(4). In terms of conventional Pauli spin and isospin variables for the nucleons the exchange charge operator is

$$\begin{aligned} \rho_{ex} = & -\frac{i}{36\pi^2} \vec{q} \cdot \vec{x}' \cdot \vec{x}'' e^{i\vec{q} \cdot \vec{r}/2} \int d^3R e^{i\vec{q} \cdot \vec{R}} \frac{\sin^2 \theta(R)}{R} \left\{ \frac{\sin^2 \theta(\xi)}{2\xi} [\vec{\sigma}^2 (\vec{\sigma}' \cdot \hat{R}) \right. \\ & + \vec{\sigma}^4 (\vec{\sigma}^2 \cdot \hat{R}) - 2\hat{R} \cdot \vec{\sigma}' \cdot \vec{\sigma}^2] + [\theta'(\xi) - \frac{\sin^2 \theta(\xi)}{\xi}] \\ & \left. [\vec{\sigma}' \cdot \hat{R} \vec{\sigma}^2 (\hat{R} \cdot \hat{R}) + \vec{\sigma}' \cdot \hat{R} \vec{\sigma}^2 \cdot \hat{R} - 2\vec{\sigma}' \cdot \hat{R} \cdot \vec{\sigma}^2 \cdot \hat{R}] \right\}. \end{aligned} \quad (18)$$

Here $\vec{\xi} = \vec{R} - \vec{r}$.

The matrix elements of this exchange charge operator will contribute to both the charge and quadrupole form factors of the deuteron. The predicted charge form factor is shown in Fig.6 as obtained when the chiral angle has been determined from the dipole parametrization for the electric form factor of the nucleon. In this case the deuteron wavefunction that corresponds to the Paris potential has been used [14]. The importance of the result lies in the fact that the exchange current contribution draws the node in the form factor towards smaller values of momentum transfer. Although the correct point for the node has not yet been determined by experiment, as the measured form factor A is dominated by the quadrupole contribution in the vicinity of the node, it is known from the charge form factors of the three- and four-nucleon systems that such an exchange charge correction is needed.

5. Discussion

The study of the utility of the Skyrme model and its generalizations as a theoretical basis for nuclear physics at low and intermediate energies is still in its infancy. The relevant form of the Lagrangean of the meson fields is not yet known although much progress has recently been made in this direction [7,8]. Nevertheless there are within the approach certain predictions that do not rely on the details of the Lagrangian which can already be confronted with experiment. All operators that are determined by the topological baryon current lead to such model independent predictions. Hence the Skyrme approach is particularly well suited for the study of the isoscalar e.m. observables in nuclei, and it does in fact provide the first coherent theoretical basis for the transverse isoscalar current operators.

The model independent predictions of the (isoscalar) form factors of the deuteron suggest because of the good agreement with the empirical values that the Skyrme approach is a useful one. The theory of the exchange current contributions to these has hitherto had to rely on various meson exchange current mechanisms that involve a number of assumptions and parameters, and it has thus had little predictive power [19].

The isovector e.m. current and charge operators are Noether currents of the Lagrangean used in the Skyrme approach and thus model dependent. A reliable construction of those will therefore have to await the appearance of a Lagrangean that is realistic enough to lead to a quantitatively acceptable model for the nucleon-nucleon interaction.

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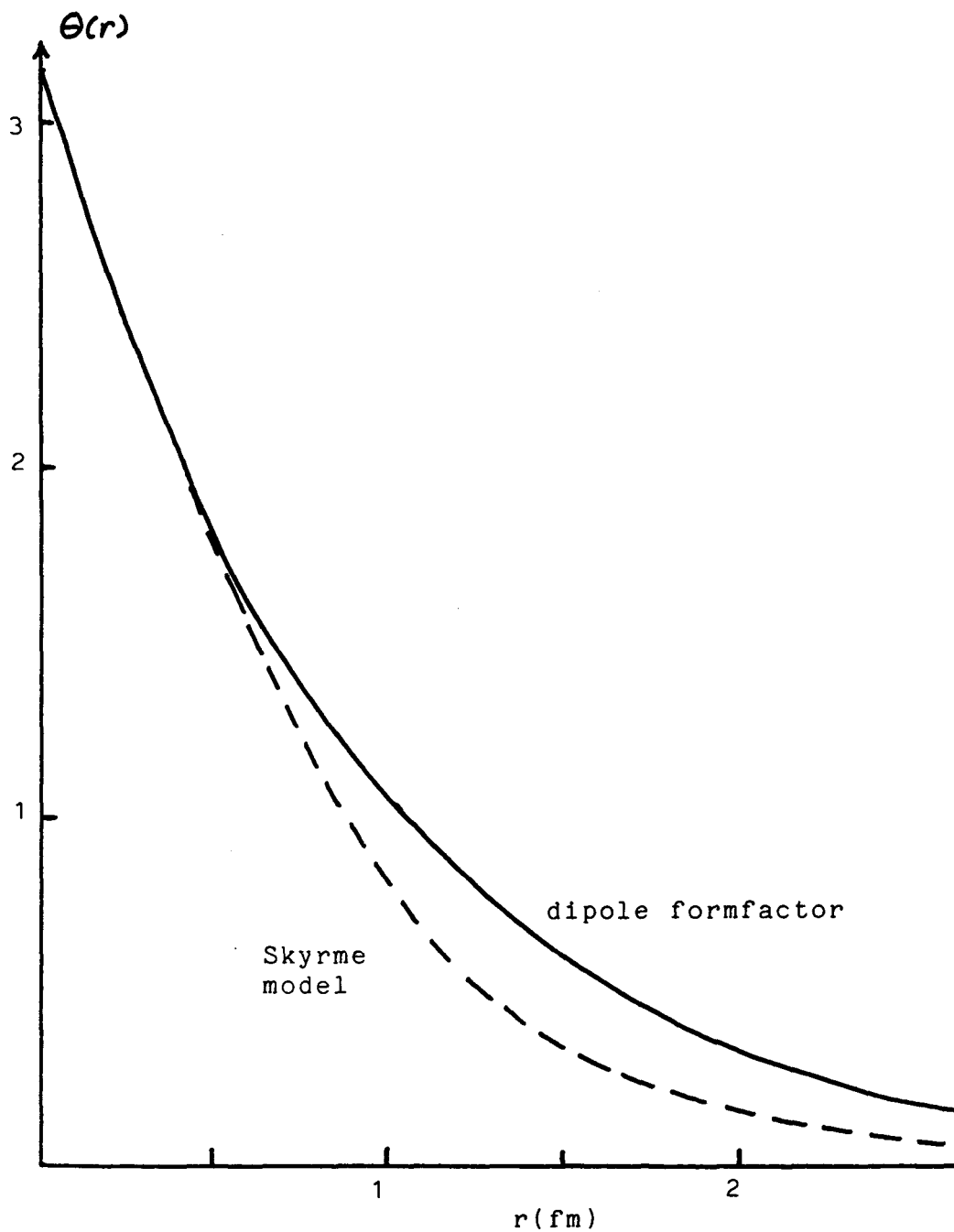


Fig.1

The chiral angle as determined from the dipole form form the isoscalar nucleon electric form factor and from the Skyrme model.

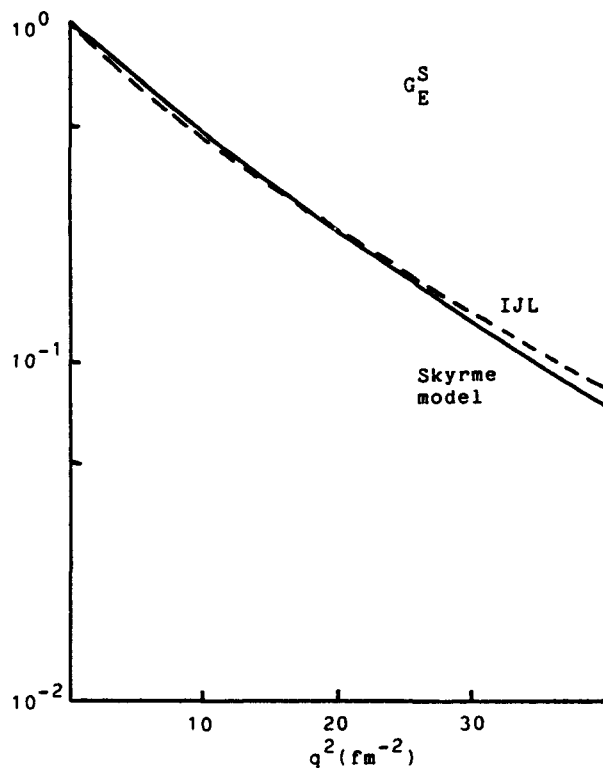


Fig.2

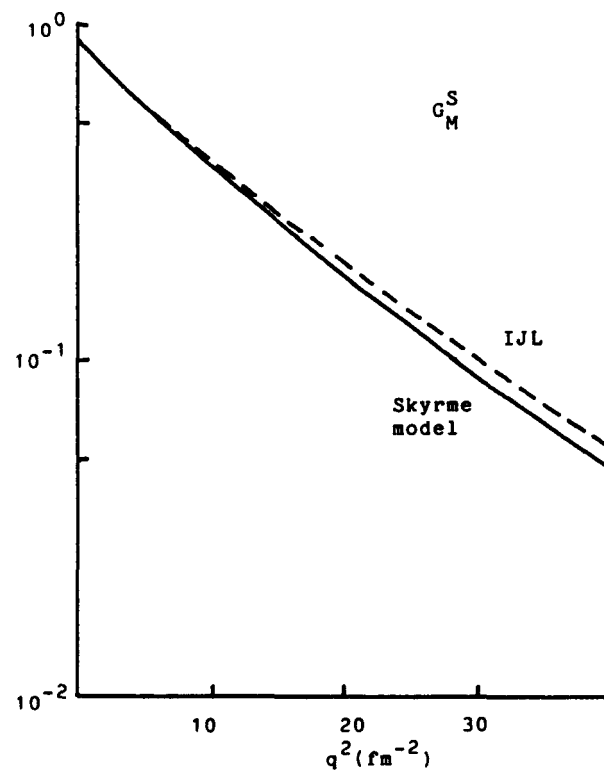


Fig.3

The isoscalar electric (Fig.2) and magnetic (Fig.3) form factors of the nucleon in the Skyrme model and as given by the 5-parameter dipole fit of Ref.12 (IJL).

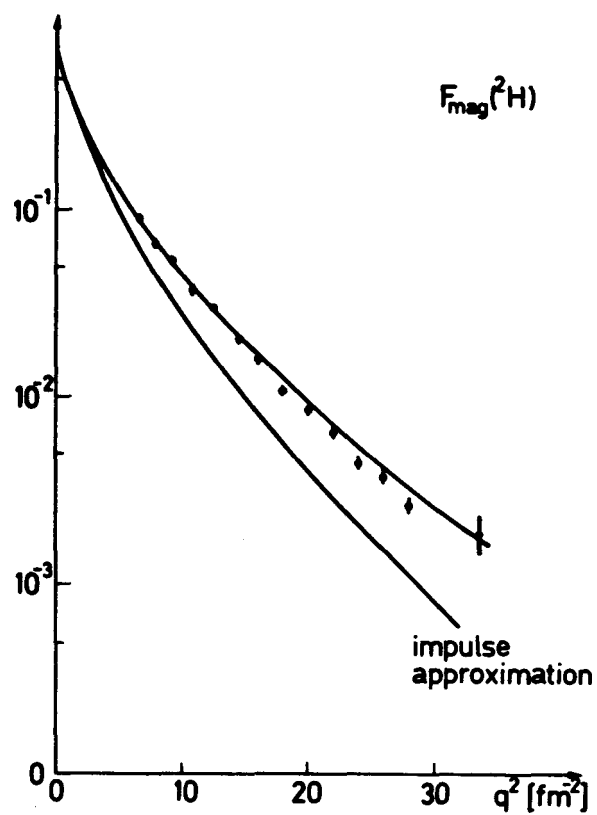


Fig. 4

The magnetic form factor of the deuteron as obtained with the chiral angle determined from the dipole form for the isoscalar electric form factor of the nucleon. The impulse approximation result is shown separately.

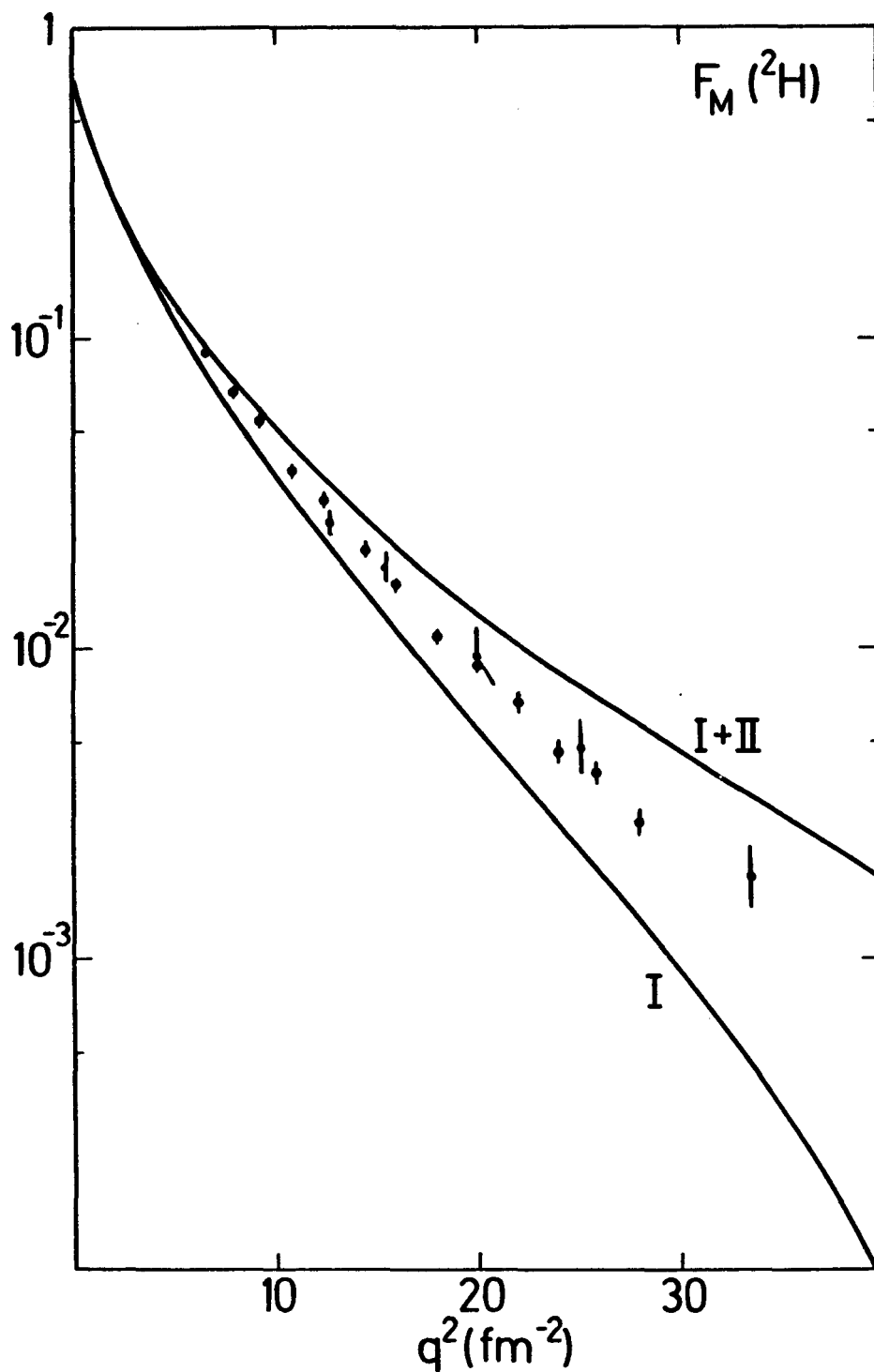


Fig.5

The magnetic form factor of the deuteron as obtained with the Skyrme model. The impulse approximation (I) is shown separately. The data points are from Refs.17 and 18.

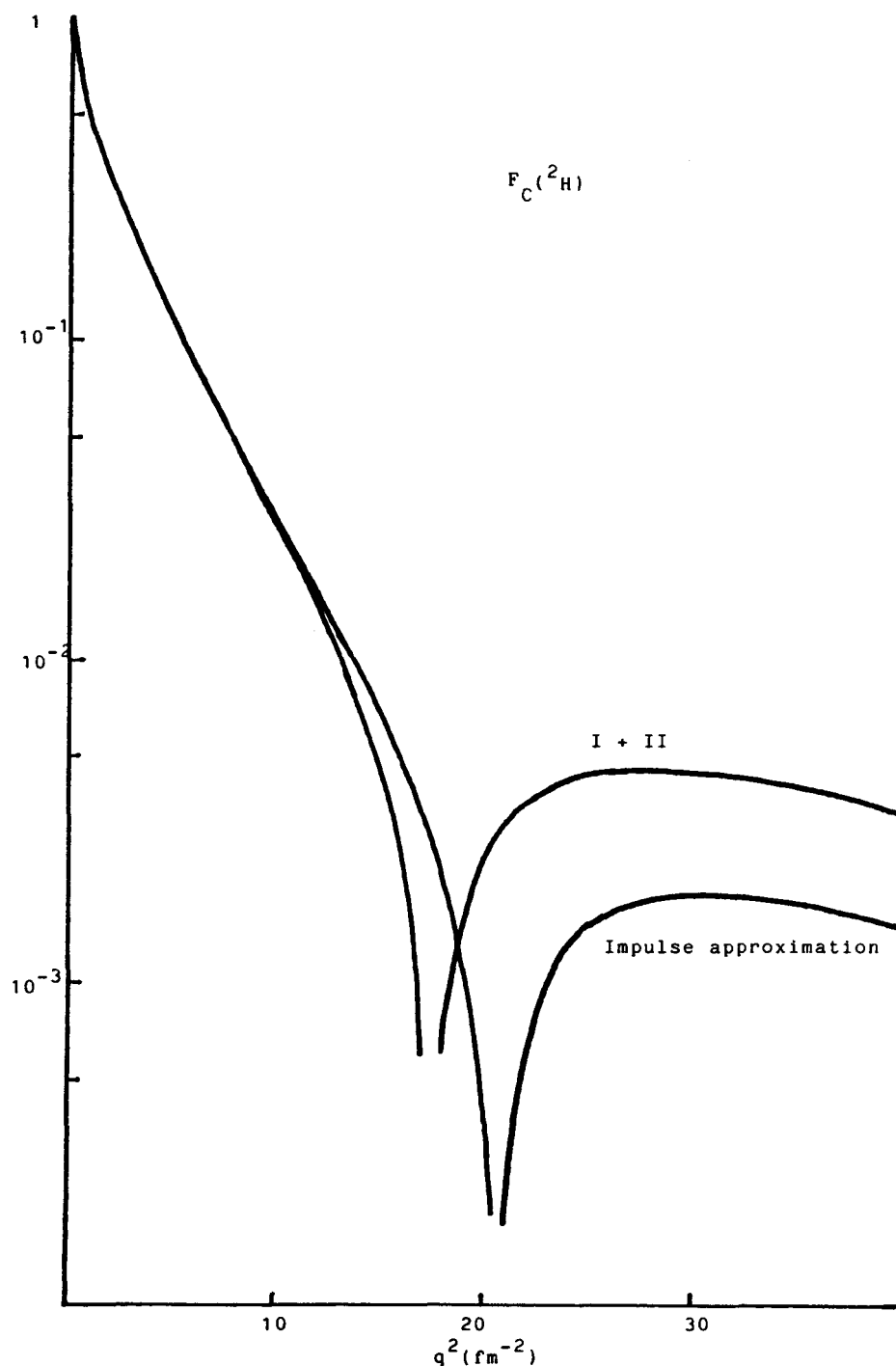


Fig.6

The charge form factor of the deuteron as obtained with the chiral angle determined from the dipole form for the isoscalar nucleon electric form factor.

THE PROTON ELECTROMAGNETIC FORM FACTORS IN THE TIME - LIKE REGION

J. Duclos, Saclay

Abstract

The proton form factor in the time - like region is measured at LEAR using the electromagnetic annihilation of antiprotons: $\bar{p}p \rightarrow e^+e^-$. The experiment is described and preliminary results are given.

Introduction

The electromagnetic form factors (EMFF) of the nucleons in the space-like region are rather well known and further experiments are planned, particularly at CEBAF, to complete the measurements. But a good knowledge of the nucleon electromagnetic structure requires precise EMFF measurements in all the available regions, and few data exist in the time - like region.

In the time - like region the mass square of the intermediate photon is positive and the form factors are function of the total energy square (s). Concerning the proton, two reactions can be used, both with a threshold at $s = 4Mp^2$:

$$\bar{p} + p \rightarrow e^+ + e^- \quad (1)$$

$$e^+ + e^- \rightarrow p + \bar{p} \quad (2)$$

The reaction (2) is limited by the luminosity of the e^+e^- colliders, particularly near threshold where the phase space is small and the detection of the low energy proton is difficult. The reaction (1) was limited until recently by the intensity of the antiproton beams. Not more than 200 events were detected from both reactions before the present LEAR experiment [1 - 5].

Various interpretations of the EMFF exist, depending on the region which is considered. The popular interpretation as a space distribution of

electric and magnetic multipoles is valid only at very low momentum transfers in the space-like region. In the asymptotic region, the quark counting rule [6] predicts a t^{-2} (s^{-2}) dependence. QCD calculations are in progress to get quantitative results in the intermediate space-like region [7]. Few QCD predictions exist in the time-like region [8].

The data are generally compared with the predictions of the Vector Dominance Model (VDM) assuming that the photon is coupled to the hadronic matter through intermediate vector mesons. In that scheme the EMFF in the time-like region, particularly near threshold, is sensitive to the nearest heavy vector mesons. More generally any hadronic object, having the photon quantum numbers ($J^{PC} = 1^{--}$) and coupled to the $p\bar{p}$ system, should affect the EMFF behaviour. The cross section of reaction (2) is given by:

$$d\sigma/d\Omega = (\alpha^2/8k\sqrt{s}) \cdot \{|G_m|(1 + \cos^2\theta) + |G_e|(4Mp^2/s)\sin^2\theta\}$$

where θ is CM lepton-antiproton angle. In the time-like region (above $4M\pi^2$) the FF are complex and in the absence of polarized beam or target only $|G_e|$ and $|G_m|$ are accessible. A $|G_e|/|G_m|$ separation can be obtained from an angular distribution measurement. At threshold we have strictly $|G_e| = |G_m|$ and the VDM predicts only small deviations from this equality up to 6 GeV^2 [9,16]. Nevertheless, if a 1^{--} state exists near the $p\bar{p}$ threshold or in the region of the measurement, important local effects could appear on the $|G_e|/|G_m|$ ratio [11]. Until now, the small number of events did not allow to perform any angular distribution and the data have been analysed assuming $|G_e| = |G_m|$.

The LEAR PS170 experiment

The new beams of antiprotons available at LEAR:

- Momentum: 100 – 2000 MeV/c
- Intensity: $10^6 \bar{p}/\text{sec}$.

allow for the first time a systematic study of the reaction (1). Measurements can be done from threshold up to a total energy: $s = 5.8 \text{ GeV}^2$, with enough statistics to perform a $|G_e|/|G_m|$ separation. The PS170 experiment [12 – 15] has been installed at LEAR in that purpose. The experimental set-up must satisfy the following requirements:

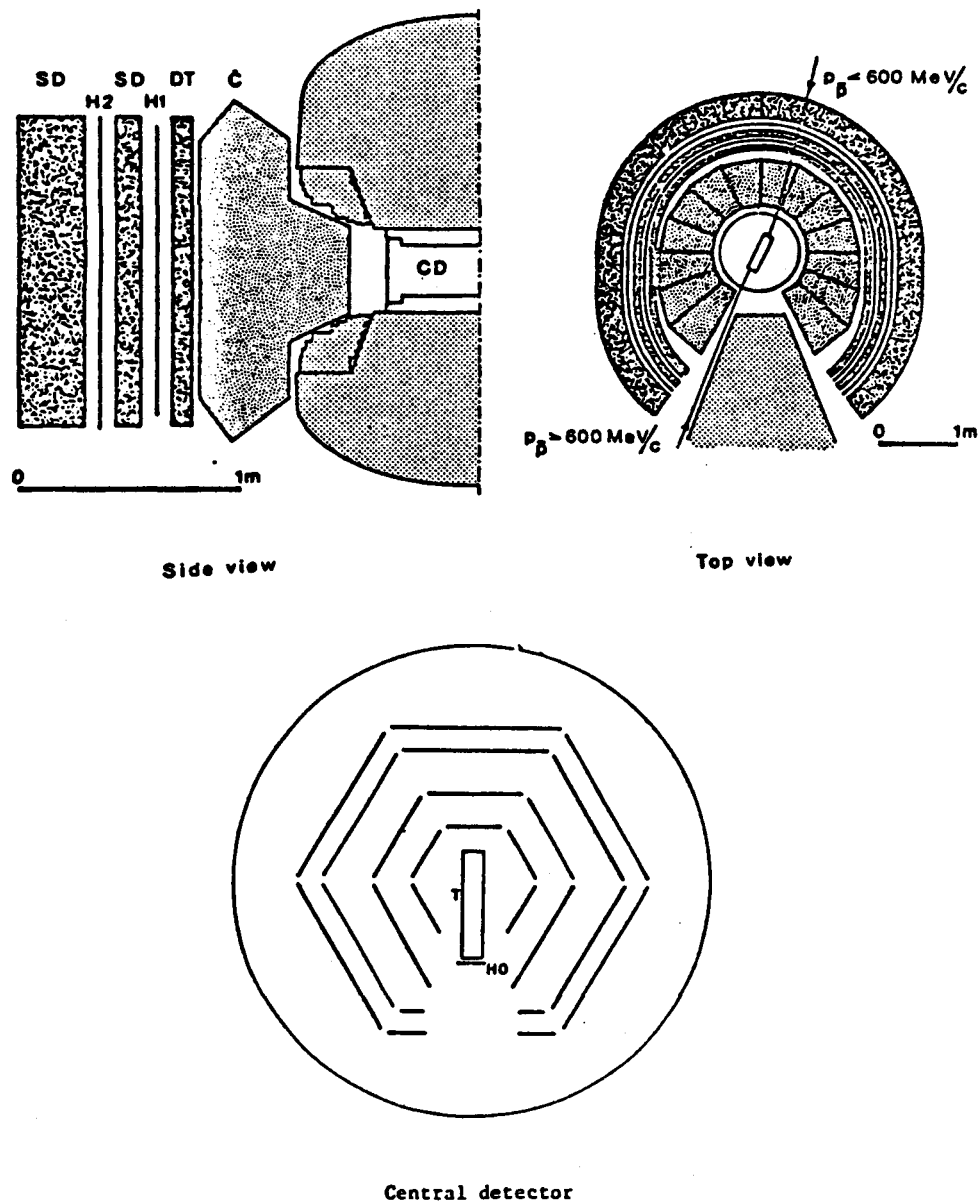


Fig. 1. Experimental set-up

- a powerful selection of e^+e^- pairs against the huge hadronic background (the reaction (1) has a branching ratio between 10^{-7} and 10^{-8})
- a large solid angle acceptance to compensate the low cross section.
- a reasonable momentum resolution (3%) to select the 2-body events

Such performances allow to measure also the following reaction:

$$p\bar{p} \rightarrow \pi_0 + V_0 (\rightarrow e^+e^-) \quad (3)$$

in which the mass of the vector meson V_0 is given by the reconstruction of the e^+e^- pair. A large amount of such events is expected in the data from annihilation of antiprotons at rest and then a contribution to the vector meson spectroscopy between 1 and 1.7 GeV can be obtained as a by-product.

The figure 1 is a sketch of the apparatus. A 30 cm long liquid hydrogen target is located in the center of a (1.3 tesla) circular magnetic field ($\phi = 1\text{m}$, $h = .4\text{m}$). The target is surrounded by several shells of detectors:

- a central detector, in the magnetic field, made of 4 plans of wire chambers with cathode read-out (3000 analog channels).
- a 28-cell gas Cerenkov counter (C).
- 100 drift tubes (DT).
- two hodoscopes (H1, H2) of 90 scintillators.
- an electromagnetic calorimeter (SD) of lead and streamer tubes (11000 digital channels).

The $\bar{p}$ beam comes in through a hole in the apparatus and crosses the wire chambers and a small hodoscope: H_0 (scintillators and silicon strips) before entering the target. A "hadron" trigger is given by the incoming beam associated with 2 charged particles in the H1-H2 hodoscopes at a minimum opening angle. An "electron" trigger requires in addition a signal in the corresponding cells of the gas Cerenkov counter. All the "electron" events and a fraction (1/1000) of the "hadron" events are registered in order to make a relative measurement of the reaction (1) versus the reaction:

$$p\bar{p} \rightarrow h^+h^- (h^+h^- = \pi^+\pi^- \text{ or } k^+k^-) \quad (4)$$

The absolute cross section and angular distribution of (4) are known from other experiments [17]. The momentum resolution for the 2-body events

(1) or (4) is about 2%. Using (off-line) both the gas Cerenkov counter and the calorimeter responses, the rejection factor against the h^+h^- events in the e^+e^- analysis is at a level of 10^{10} .

Data have been taken for various $\bar{p}$ momenta from 300 to 600 MeV/c (table 1). At 300 MeV/c the antiprotons stop in the target. The annihilations in flight, during the $\bar{p}$ slowing down from 300 to 100 MeV/c can be selected from the large amount of annihilations at rest. The 300 MeV/c data are used to get a EMFF value at threshold and near threshold, to study the vector mesons and to calibrate the apparatus.

Table 1: Present status of the data taking

Momentum (MeV/c)	Number of $\bar{p}$ on target (10^{11})	Events h^+h^-	Events e^+e^-
300	0.7	20000	2600 (250 in flight)
540	0.75	2530	52
600	4.9	21000	400

A total of $6 \cdot 10^{12}$ antiprotons were requested in the proposal. Assuming the expected LEAR intensity ($10^6 \bar{p}/s$), it corresponds to about 100 days of data taking. Unfortunately the running conditions and the heavy program of LEAR in the period 1984–86 did not allow to get more than 10% of original request. The table 1 gives the number of antiprotons received on the target during that period at different incident momenta and the corresponding number of hadronic ($\pi^+\pi^-$, k^+k^-) and electronic (e^+e^-) 2-body events. These data will give, after complete analysis, 4 points of the form factor (threshold, near threshold, 540, 600 MeV/c). A significant angular distribution will be obtained at 600 MeV/c.

Preliminary results

On figure 2 are shown, for the annihilation at rest, the angular correlation between the 2 charged particles and their momentum distribution.

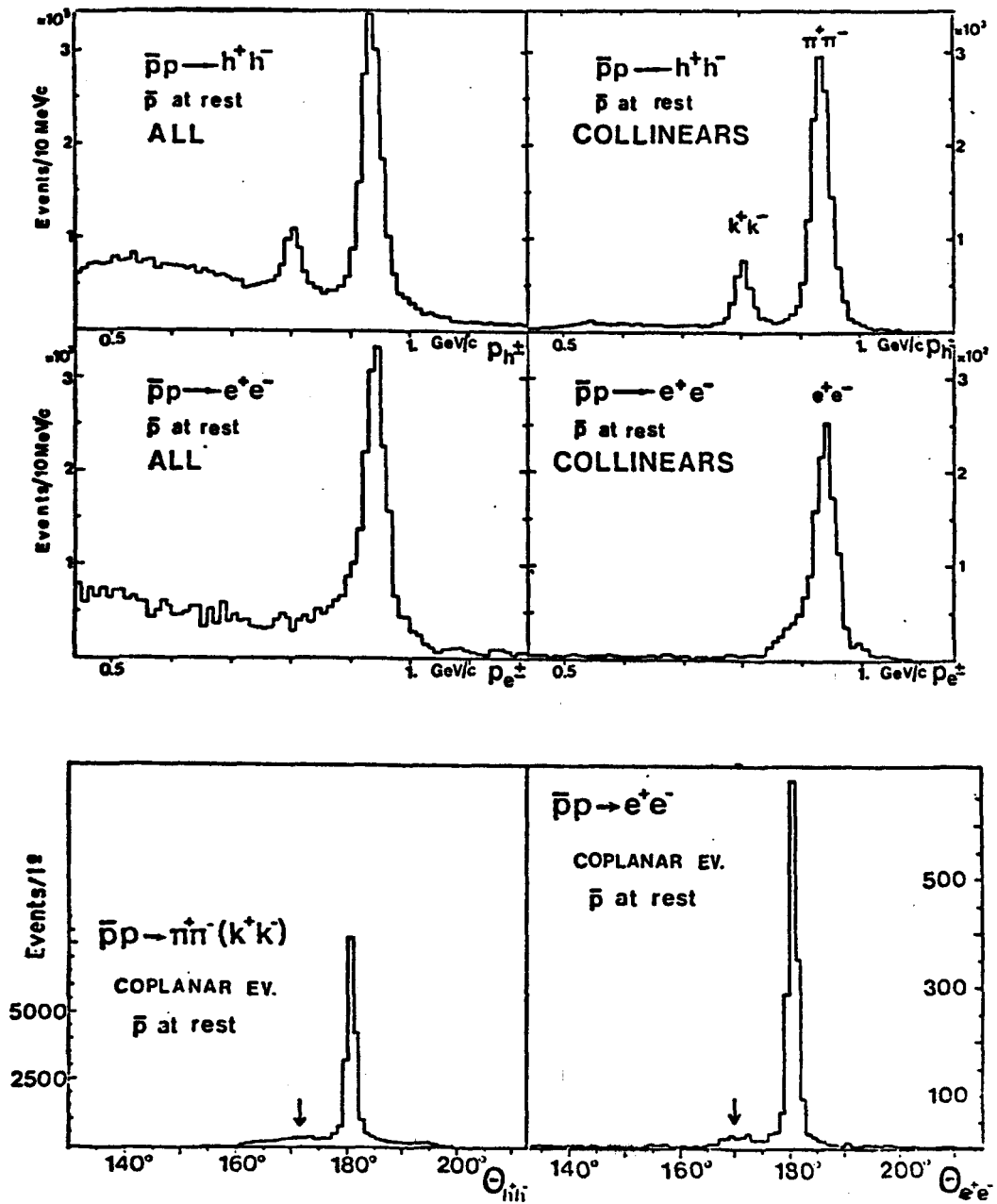


Fig. 2. Momentum distribution of the 2 outgoing particles (h^+h^- , e^+e^-) from the $\bar{p}p$ annihilation at rest (top). Distribution of their opening angle (bottom). The arrows show the in-flight events.

The 2-body reactions (1) and (4) appear at 180° . In the momentum distribution of the hadronic data, the $\pi^+\pi^-$ and the k^+k^- events are clearly resolved. In the corresponding e^+e^- distribution, the absence of peak at the kaon momentum demonstrates the good rejection of the (h^+h^-) hadron events. The following branching ratios have been found [10]:

$$\Gamma(\bar{p}p \rightarrow k^+k^-)/\Gamma(\bar{p}p \rightarrow \pi^+\pi^-) = 0.32 \pm 0.05$$

$$\Gamma(\bar{p}p \rightarrow e^+e^-)/\Gamma(\bar{p}p \rightarrow \pi^+\pi^-) = (1.5 \pm 0.3)10^{-4}$$

Similar performances are obtained from the analysis of the in-flight data. The angular acceptance of the apparatus is obtained from a Monte-Carlo calculation. The result is checked with the experimental distribution at rest which is isotropic. The relative cross section of the e^+e^- and h^+h^- reactions is free from systematic errors, excepted for the efficiency of the Cerenkov counter and the calorimeter that play differently in the 2 channels. In the present state of the analysis, no definite result can be given. On figure 3, preliminary values of the form factors are plotted, assuming $|G_e| = |G_m|$. One can make two remarks.

the form factor at threshold agrees with a VDM fit [16] of the parameters of 5 vector mesons: $\rho(773)$, $\omega(782)$, $\phi(1019)$, $\rho'(1600)$, $\phi'(1660)$ to the space-like and time-like existing results. A fit to the space-like data only would give a larger value of the time-like form factor near threshold. This region is more sensitive to the heavy mesons ρ' (isovector) and ϕ' (isoscalar). A correct fit to both the space-like and time-like data can be obtained by choosing the ρ' and ϕ' coupling constants to give opposite effects so reproducing the small value of the proton FF. But, then, the two effects will add-up in the neutron FF and the surprising result is a very large cross section of $n\bar{n} \rightarrow e^+e^-$ compared to $p\bar{p} \rightarrow e^+e^-$.

The value of the form factor at 600 MeV/c seems smaller than the VDM prediction. In other words the slope of the form factor does not seem compatible with the effect of the 5 vector mesons considered in the fit and suggests that some other contributions exist.

In conclusion, it is clear that a precise systematic study of the time-like form factor is possible at LEAR and is very useful for the knowledge of the proton electromagnetic structure. The PS170 experiment will continue after the installation of an additional ring: ACOL, that should increase the antiproton production at CERN by a factor 5.

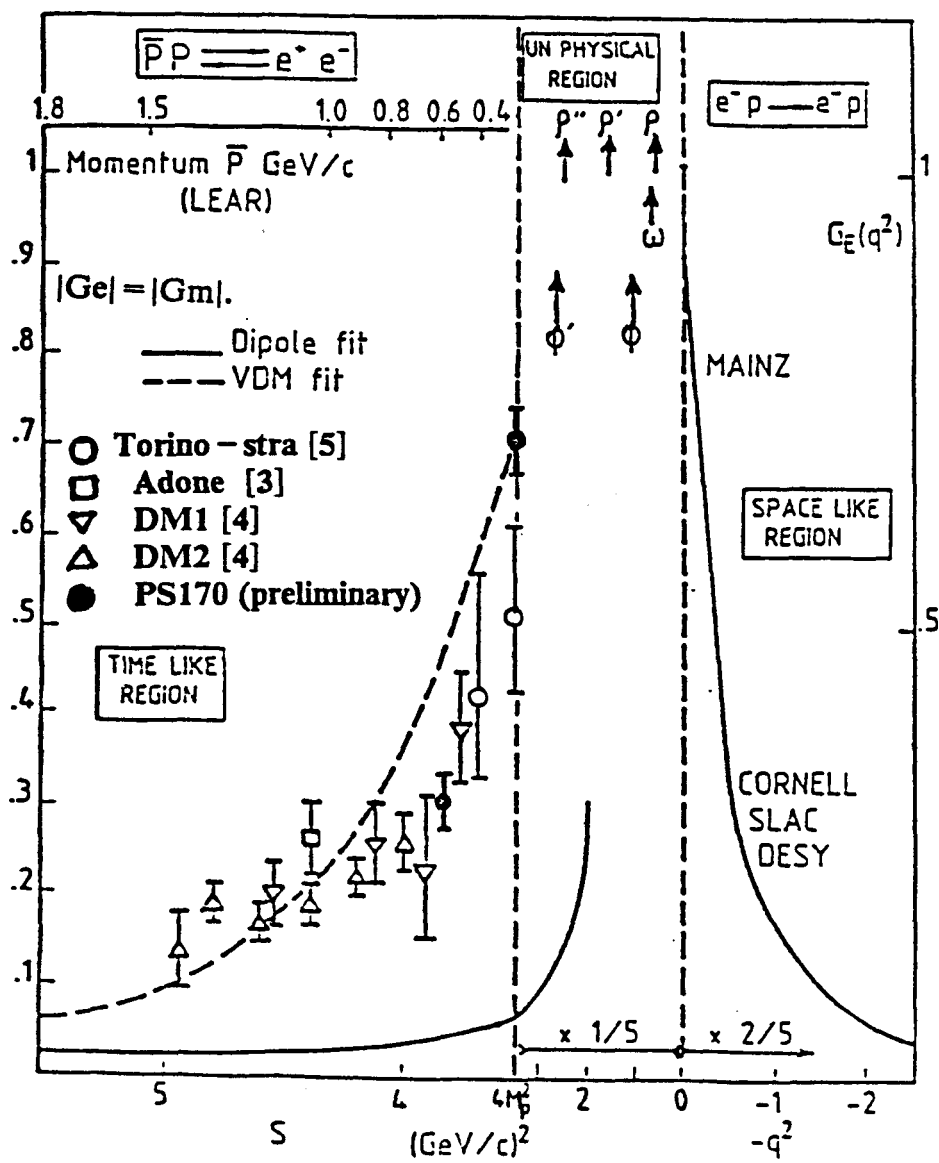


Fig. 3. Proton form factor in the time-like region assuming: $|G_E| = |G_M|$.

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LEAR experiment PS183

Comments on the CEBAF experimental program

J. Duclos, Saclay

Abstract

The characteristics of CEBAF will open a new era in the electromagnetic program of nuclear physics. The qualities of the beam expected from the superconductive linac should also allow the use of new techniques of detection.

CEBAF: a new machine

The CEBAF project represents a considerable improvement compared to the present electron machines of the nuclear physics community.

- The energy is increased by almost a factor 10 and the high intensity is maintained.
- The duty factor is multiplied by 100 and the possibility of splitting the beam between three simultaneous experiments provides in fact an additional factor 3.
- The emittance is considerably improved and, also very important, the beam is expected to be extremely clean. The absence of tail (or halo) in the emittance, expected from the dc acceleration, is indeed a condition of the superconductivity technique.

These extremely favorable conditions will open of course a new program of research but should also allow the use of new techniques beyond the simple extrapolation of the present experiments.

The physics program has been extensively discussed by the various working groups. As the main goal of the workshop is to define the CEBAF experimental equipment, I will make some personal comments

in that domain. Most of the experiments of course use the main characteristics of the machine: high energy, high intensity, and good duty factor. But particular performances are needed by some experiments.

- The experiments which include some nucleus spectroscopy require high resolution spectrometers. Typical examples are the coincidence measurement of single particle densities and the study of hypernuclei.
- Some nuclear experiments need only moderate energy resolution. This category includes deep inelastic reactions, pion electro – production, observation of 2 – body correlations by $(e, e' - 2N)$ experiments... For such experiments the energy resolution is mainly used to exclude the production of additional pions.
- The experiments of particle physics (nucleon form factors, nucleon resonances) or others like parity violation, don't need either a high resolution but for most of them polarized electrons would be extremely useful or even necessary

A high resolution experimental room

An impressive study has been made of all the possible spectrometers to be installed at CEBAF. They are necessary to achieve the good resolution required by the nuclear spectroscopy. But they have also many advantages that made them, until now, the traditional equipment of the electron machines:

- Possibility of high peak luminosity in a noisy environment
- Relatively simple detectors easily accessible to everybody
- Rapid analysis of the data.

There is no doubt that one experimental room at CEBAF should be totally equipped with two very performant spectrometers. One at least should have the best possible resolution. Such an equipment will guarantee high quality results as soon as the first beam of electrons is delivered.

Open geometry

But spectrometers have also inconvenients:

- Limited solid angles and limited angles
- No detection of neutral particles (neutron, π_0)
- Difficulty for out of plan measurements.

- Small momentum and angular acceptance
- Large encumbrance, absence of flexibility.

The totally new conditions of CEBAF may tolerate the use of detectors in open geometry eventually in direct view of the target. The pile-up of particles is a severe limitation to such techniques on the present electron machines. In that respect the situation at CEBAF could be improved by a factor 1000 (100% duty factor – beam sharing – clean beam). If the beam has no halo, the background is due only to interactions in the target. The maximum luminosity, acceptable for an open geometry, will be obtained by using a high current and a thin target. It could be measured on existing machines or eventually calculated by the EGS code.

The Large Acceptance Detector

Following this line, the LAD is a 4π solid angle, multipurpose detector. It is mainly foreseen for low intensity electron beams associated with polarized targets, and photon beams. Though expensive, the design is very attractive. A collaboration of several universities has been launched during this workshop in view of supporting the LAD experimental program. Such a detector indeed requires an enormous amount of work. The various steps: – building, setting up electronics and triggers, calibration, data taking, data analysis – are strongly connected and have to be undertaken by the same large group of physicists.

A flexible experimental room

For many experiments of moderate resolution a large flexibility would allow to optimize the apparatus for the specific performances required by some experiments. In that view, an experimental room could be equipped with only one spectrometer to detect the scattered electrons, generally in the forward sector, the hadrons detection being left totally open.

Elastic scattering: $e+n \rightarrow e+n$ (from deuterium) Solid state detectors may be used to measure directly light recoil nuclei (α – particles). The experiments involving neutrons or π_0 , the polarimeters, the out-of-plan measurements cannot be done easily with classical spectrometers but should be possible at CEBAF, using detectors in open geometry.

One may wish of course to equip the second experimental room with three spectrometers for the $(e, e' - 2p)$ experiments but such an heavy equipment could dangerously affect the flexibility for other experiments and techniques. A good intermediate solution may be the versatile magnet VAS. Nevertheless, a unique device will never fit all the experiments.

An important program of CEBAF will certainly be the pion electroproduction in the direction of the virtual photon. The longitudinal component of the cross section in that case is dominated by the one pion exchange diagram and probes the nucleus as a pion source. The pions have to be detected at very small angles from the beam direction ($< 10^\circ$) in a strong electron background. This could be done for example by using a non focussing magnet (borrowed from some other place?), with specific shielding and detectors (wire chambers, Cerenkov counters). A similar equipment has been used by equivalent experiments in particle physics [3]

The CEBAF experiments could benefit in the future of the enormous work done presently by the high energy groups to improve the performances and lower the price of all kind of detectors (BGO, scintillator glasses, silicon calorimeters...) [1]. The segmentation of the detectors becoming easier with the modern read-out systems can help solving the pile-up problem [2].

Polarized beams

Polarized beams are not considered in the first step of the CEBAF construction. However many experiments require polarized electrons. (Deuterium form factors, pion electroproduction, $E2/M1$ measurement...). A polarized source, based on the existing GaAs sources should be studied by a CEBAF-users laboratory. Several years are indeed necessary to build a reliable device. Provision should be made on the machine for using polarized electrons. It may be difficult to include a polarized beam at all energies in the present project, but even a facility for discrete energies is worth considering.

A fourth experimental room

The possibility of 3 simultaneous users is a great advantage of CEBAF. however the gain will be reduced by the limited number of experimental rooms, making the scheduling extremely difficult. On the other hand, because of the unprecedented characteristics of CEBAF, non-tra-

ditional (exotic) experiments may be proposed. An exemple is given in the next section. Then the possibility of having a fourth end station with no detection equipment should be considered.

Parity violation

The parity violation in electron – hadron scattering has been observed in atoms and, using polarized electrons, at SLAC. A proposal has been approved at Bates. A model independent description of the effect depends on 4 parameters obtained by combinations of the various weak currents (vector, axial, isoscalar, isovector). Several different experiments are necessary to solve the problem. The Standard Model is now well established particularly by the neutrino experiments. However, for such a fundamental question it may be interesting to check all the accessible channels, with eventually a precision useful for the theory. A possibility is to measure the parity violation using a polarized beam, in the following reactions:

- Elastic scattering: $e + p \rightarrow e + p$
- Elastic scattering: $e + n \rightarrow e + n$ (from deuterium)
- Delta excitation: $e + N \rightarrow e + N^*$

A high momentum transfert gives a large violation effect, but a medium energy is necessary to resolve the N^* exitation and to get a reasonable cross section. A good compromise is the scattering of 1.2 GeV electrons at angles around 90° . A possible apparatus is shown on fig 1. The scattered electrons are identified in 30 cells including a gas Cerenkov counter and a lead glass calorimeter. The rate assymetry between forward and backward polarizations is about 10^{-4} . A 20 μ A beam in a 10 cm liquid hydrogen target produces 10^5 events/sec. A strong shielding and eventually an axial magnetic field would reduce the single rates. CEBAF is the ideal machine for such an experiment.

A storage ring in a parasitic mode

An inconvenient of the superconducting linac is the desapearence of the ring. A considerable amount of work has been done by the CEBAF study groups about the internal targets, in the stretcher ring project. In the new project, let us consider a ring as an additional device working in a parasitic mode. The ring is filled up to 100 mA in less than 1 hour. The beam is cooled down by synchrotron radiation. The energy is maintained by an RF system. Relatively thin jet targets (10^{14} cm^{-2}) provide a useful

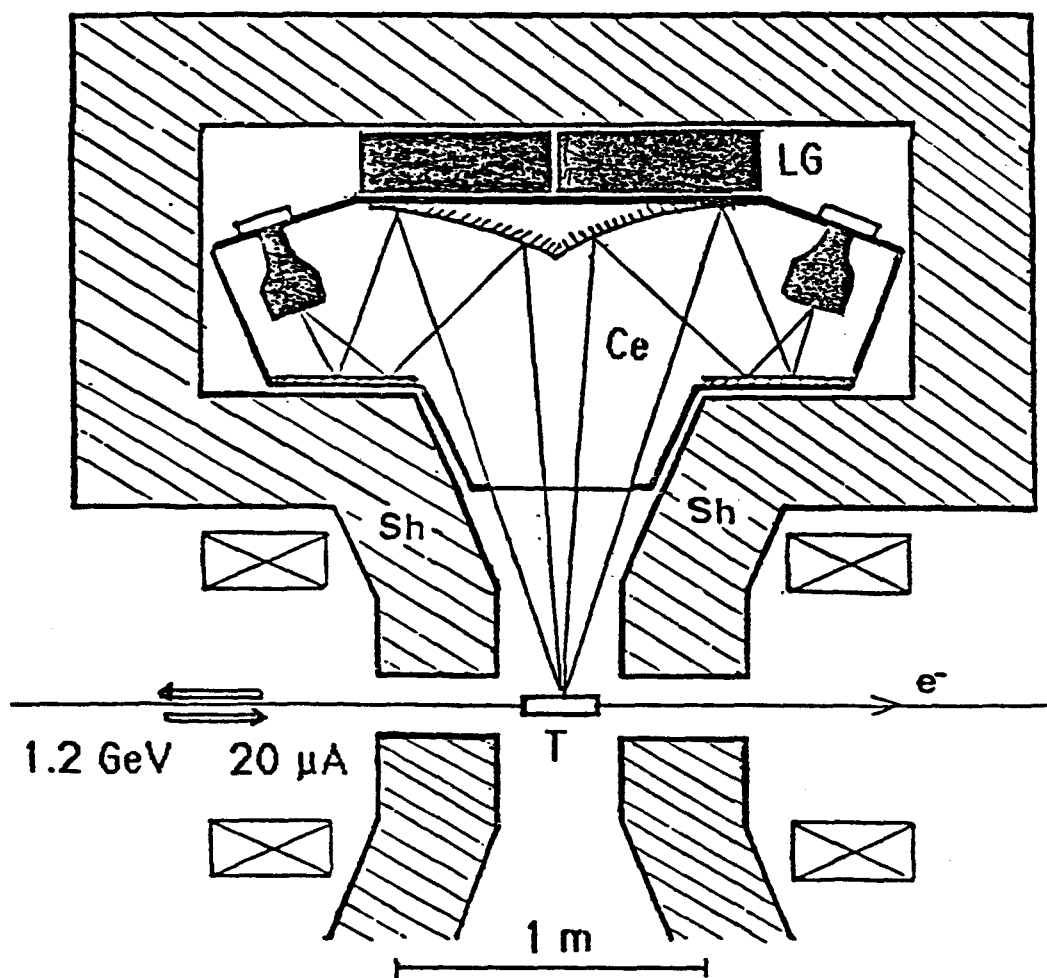


Fig. 1 Apparatus suggested for a parity violation experiment Ce: gas Cerenkov counter, LG: lead glass counters, T: hydrogen or deuterium target, Sh: shielding.

luminosity (10^{32} cm^{-2}) during about 10 hours. The experiments may include polarized target, and large solid angle detectors. The 2 – body reactions for example could be easily identified by a very compact detector using silicon micro – strips [4].

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INTRODUCING THE WORKING GROUP ON EXCITED BARYONS
AT THE CEBAF 1986 WORKSHOP: "SON ET LUMIERE"*

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ABSTRACT

I outline here important issues that needed to be debated at sessions of the working group on excited baryons in the 1986 CEBAF Workshop. Our primary charge was to examine the experimental concepts outlined in the RPAC-85¹, published by CEBAF, in the context of research to be done on the excited baryons.

1. INTRODUCTION

I am going to give here a preview of the sessions arranged for the working group on excited baryons at the 1986 CEBAF Workshop. The chief aim of our group is to review the concepts of proposed experiments at CEBAF in this area, contained in the RPAC-85¹, and to see if there is room for improvement, or even radical changes, before the freeze of the designs of principal experimental facilities. This is the charge of my working group. Already Franz Gross and Jean Mougey have set the scene for the discussion this morning. Plenary talks by Thomas, Pirner, Riska and the broad overview by Langacker will give you many snapshots of our current knowledge of hadron structure.

In order to avoid duplication, let me remind you of some references that already provide us with some of the pertinent information concerning this area of CEBAF physics. Let me emphasize the CEBAF materials: past reviews by Isgur² and Cassel³, and contributions in RPAC-85 by Burkert⁴, Stoler⁴, Carlson⁵, Mukhopadhyay et al.⁶, Robson⁷, Weber and Weyrauch⁸. Opening chapters of Part One of RPAC-85 form

* Invited Talk at the CEBAF 1986 Summer Workshop, Newport News, VA.

[†] Supported by the U.S. Dept. of Energy Contract #DE-AC02-83ER40114-A003.

convenient points of departure for us. For more extensive material on baryon structure and spectroscopy, you can consult the proceedings of the conference BARYON 1980⁹, and the latest particle data book, just published¹⁰. I would urge caution in taking the electromagnetic helicity amplitudes for baryons in the last reference at face value, a topic we shall revisit during our working sessions.

2. MAIN ISSUES

You have received a long list of questions to be discussed here in our working group. [I guess Volker Burkert is the author of them.] I have a set of questions and topics as well, not necessarily identical to those compiled by Volker. Here is my partial list of my issues in the context of Baryon Resonances (B^* for short, B being the nucleon):

- (1) Why electromagnetic excitation of B^* 's?
- (2) Which ones in particular?
- (3) For what new physics?
- (4) Existing data: γBB^* amplitude ambiguities.
- (5) Extraction of the resonant amplitude: theoretical problems.
- (6) B^* in nuclei.
- (7) Future experiments: before and at CEBAF.

Unless stated otherwise, γ refers to both real and virtual photons. You can add your own items to either of the lists mentioned above! Let me introduce some of these from my short list. I do not have time now for all of them. I shall return to them in my summary talk.

3. WHY ELECTROMAGNETIC EXCITATION OF B^* 's?

Let me recall the following lines from Feynman¹¹:

"One very powerful way of experimentally investigating hadrons is to look at them with a known particle, in particular, the photon - no other is known as well."

Here, of course, both real and virtual photons are implied. Feynman then compares the following two reactions:

$$\pi + p \rightarrow \pi + p, \quad (a)$$

and

$$\gamma + p \rightarrow p + \pi. \quad (b)$$

In the reaction (a), he points out, two hadrons with complicated structures are hitting each other; thus, it is like hitting one "unknown" with another. The q^2 of the pion is fixed; it is m_π^2 . The reaction (b) has the advantage of hitting an "unknown", here proton, with a "known" object, γ . The q^2 can be varied by using virtual photons. "A pion far off its mass shell may be a meaningless - or at least highly complicated - idea", Feynman concludes. Notice that, in reaction (b), the photon can hadronize, say, into vector mesons, at the appropriate kinematic region, but even that is clear and selective.

Thus, we are after γBB^* , the three-point function that should be understood from fundamentals of hadron structure, and some elusive hadrons that are predicted, but not yet seen. The electromagnetic transition amplitudes between B and B^* 's are vital for our understanding of the baryon multiplet structure.

4. WHICH ONES?

To do efficient baryon bookkeeping, we may start with the SU(6) quark model, introduced first by Gürsey and Radicati in the sixties. Writing the elementary six-dimensional column vector

$$\vec{\xi} = \begin{pmatrix} p \uparrow \\ n \uparrow \\ \lambda \uparrow \\ p \downarrow \\ n \downarrow \\ \lambda \downarrow \end{pmatrix} \quad (1)$$

as the basic $[6]$ multiplet of SU(6), with $\uparrow$ and $\downarrow$ meaning spin-up and spin-down configurations for the appropriate quark flavor, we can recall the identity

$$[6] \otimes [6] \otimes [6] = \underline{56}_S \oplus \underline{70}_M \oplus \underline{70}_M \oplus \underline{20}_A, \quad (2)$$

with the subscripts indicating the symmetry of the representation. Those of you, who want to refresh your memories of the meaning of the above SU(6) direct product identity, may consult standard works on quark model [for a pedagogical review, see Gasiorowicz and Rosner¹²]. It may be useful to remember the SU(3) content of the SU(6) multiplet. We give this in Table 1, with some examples of our familiar baryons. The relevant orbital angular momentum is L, N is the degree of the tensor in reduced coordinates $(\vec{\lambda}, \vec{\rho})$ that can be appropriately defined in terms of the three coordinates of the three quarks making the baryon.

Table 1: Some members of the SU(6) Baryon multiplets

N	Parity	SU(6), L	SU(3) Content	Example of States
0	+	<u>56</u> , 0	<u>8</u>	N (939) $\frac{1}{2}^+$
			<u>10</u>	Δ (1232) $\frac{3}{2}^+$
1	-	<u>70</u> , 1	<u>1</u>	Λ (1405) $\frac{1}{2}^-$
			<u>8</u>	N (1535) $\frac{1}{2}^-$
			<u>10</u>	Δ (1700) $\frac{3}{2}^-$
2	+	<u>56</u> , 2	<u>8</u>	N (1720) $\frac{3}{2}^+$
			<u>8</u>	N (1680) $\frac{5}{2}^+$
			<u>10</u>	Δ (1910) $\frac{1}{2}^+$
	+	<u>56</u> , 0	<u>8</u>	N (1440) $\frac{1}{2}^+$
			<u>10</u>	

The study of electromagnetic excitation of all these baryons to measure the γBB^* amplitudes remains the goal of this physics.

5. WHY BOTHER ABOUT B^* 's AT CEBAF?

S. Koonin¹³, while reviewing the CEBAF Workshop in 1985, came to the following conclusion concerning the possibility of studying excited Baryons at CEBAF:

"It is clear that much work in this neglected field could be done at CEBAF, but I am not yet convinced that it should be. What new might be learned?"

The answer to this question is going to be the main theme of our working group and its successors. If the answer to this question is in strong affirmative, it will make sense then to go ahead.

The short answer to Koonin's question is obvious: non-perturbative quark-gluon structure of hadrons is still unknown. It is only for this that the entire CEBAF enterprise, in particular, the experiments on excited baryons, will be worthwhile. These experiments can help future understanding of it.

Consider the large number of models for hadrons: old quark models, new quark shell model, bags of various varieties, topological and non-topological solitons, lattice gauge theory, flux tubes, strings, and so on. In the "so on" categories lie all of the exotic "super" models, which somehow will not only solve all the diseases of other non-super models, but also give the correct properties of the hadrons as we know them. All these models also purport, at low-energies, to give the correct limit of QCD.

We are clearly observing in this instance, blind men surveying an elephant. Only sense that we can bring to this enterprise is to provide accurate experimental constraints that the theory must reproduce. Imagine studying atomic physics without knowing the atomic spectroscopy! How can we study hadron physics without precise hadron spectroscopy?

6. INTERPLAY BETWEEN THEORY AND EXPERIMENT

The experimental baryon spectroscopy alone is not enough. We must learn how to tackle the theoretical problem of extracting the resonant γBB^* amplitudes from the data. Feynman¹¹ has put his finger on the crux of the problem:

"How much is resonances and how much is background? Can the background below a resonance be simply tails of other resonances?"

We have learned a lot, since Feynman's analysis of the Walker approach of extracting the γBB^* amplitude, for the $N \leftrightarrow \Delta$ (1232) transition. However, even there, the model dependence¹⁴ of extracting the resonant amplitude is considerable.

Can we surmount this vicious problem for the baryon spectrum beyond Δ (1232)? Can we do any better than what has been done thus far?

A capital question for our theorists in this workshop!

7. EXCITED BARYONS IN NUCLEI

So far the electron and photon factories have addressed the question of photo- and electroproduction of Δ (1232) in nuclei. Recent Mainz-RPI-MIT-Pittsburgh collaboration experiments at Bates¹⁵ have yielded spectacular results for charged pion photoproduction, as have the Bates-MIT-Boston experiments¹⁶ on the neutral pions. The role of a nuclear medium modifications have been found to be very large¹⁷, both in coherent and incoherent π^0 photoproduction, requiring one to go beyond the standard distorted wave impulse approximation (DWIA) theory. Evidences are mounting on the nuclear surprises for the charged pion photoproduction¹⁸. Our understanding of the role of production of neutral pions, which can then lead, by charge exchange, to the same nuclear final state as the charged pion photoproduction reaction, is very low.

The question arises as to how we would escalate these works at the CEBAF energies¹⁵. How do we go beyond the Δ (1232) excitation in nuclei? The answer has to start with the study of the electro- and photoexcitation of the excited baryons at CEBAF kinematics, from free nucleons. Our theoretical understanding of this will be the key to the modifications of the elementary γBB^* amplitude in the rich environment of the complex nuclei.

7. AU REVOIR

Above are some of the issues for our working group.

Let us look for their answers in planning future experiments, both at CEBAF and in the interim, in our sessions ahead¹⁹.

I thank Franz Gross for his kind invitation to the CEBAF Workshop, and Volker Burkert for critical discussions that influenced our working group on excited baryons.

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Report on

WORKING GROUP ON SINGLE NUCLEON EMISSION

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It has already been recognized that $(e,e'N)$ experiments will be a major part of the CEBAF physics program^{1,2}. With the present generation of electron facilities, they have proved to be the most direct and powerful tool to obtain informations about the nucleons bound in a nucleus. Two of the most relevant physics issues are:

1. The determination of single nucleon densities. In particular, the spectral function $S(E_m, p_m)$, which contains the momentum and separation energy distribution of the nucleons, can be extracted from these measurements. This is the "classical" aspect of $(e,e'N)$ experiments, and the one which has been mostly exploited during the last decade.

2. The study of the electromagnetic properties of the bound nucleon. In principle, these experiments can tell us how the elementary interaction of an electron with a bound nucleon differs from that with a free nucleon. It must be emphasized that a good knowledge of this elementary interaction is preliminary to any reliable extraction of momentum distributions, spectroscopic factors, etc..

Since doubts have appeared from recent measurements on the validity of the standard approach to quasi-free scattering, the participants to this working group choosed to focus their attention on this problem, in relation with the more general and presently popular question: do nucleons keep their intrinsic properties in the nuclear medium ?

This problem is addressed somewhere else in these proceedings^{3,4,5}. In Ref. 3, the present experimental status and

recent theoretical developments are reviewed. A deeper discussion of some theoretical aspects can be found in Ref. 4,5 and several contributed papers to CEBAF 1986 Summer Study Group. In this report I only recall briefly the questions which remain open at the present time and draw the future research axes which have appeared specific of a multi-GeV CW electron facility like CEBAF.

QUASI-ELASTIC SCATTERING. PRESENT STATUS

Impulse approximation. Off-shell effects

In the standard approach to quasi-free scattering it is assumed that the nuclear electromagnetic current J_μ is the superposition of individual one-body currents:

$$J_\mu = \sum_i j_\mu^N(i)$$

In this framework, the target nucleus is described as a collection of independant nucleons. Three ingredients are necessary to define the nucleon current:

1. The nucleon spinor.
2. The nucleon form factors.
3. The current operator.

In Impulse Approximation the spinor and form factors of the free nucleon are used. Only the current operator has to be extrapolated to take into account off-shell effects: some prescription is needed for this extrapolation. It is commonly believed that the kind of prescriptions given by De Forest let little ambiguity in the choice of the current operator. Recent (e,e'p) experiments are usually analyzed following such prescriptions.

Break-down of Impulse Approximation: experimental evidences

Separation of the transverse R_T and longitudinal R_L response functions in inclusive measurements at MIT-Bates and Saclay Linac have revealed an unexpected behaviour of the longitudinal response⁷. Classical calculations (relativistic Fermi Gas⁸ or relativistic shell-model⁹) overestimate R_L by 20 to 40% for medium nuclei. This results in a non saturated Coulomb Sum Rule even at the highest momenta (550 MeV/c) reached in those experiments. However it should be remarked that in recent Bates measurements on ^{238}U the quenching of R_L is found smaller.

Recently, the separation of transverse and longitudinal response functions has been extended to $(e,e'p)$ experiments^{10,11}. Let me recall that, in one-photon exchange approximation, the $(e,e'p)$ six-folded cross section writes:

$$d^6\sigma = \Gamma_V [\sigma_T + \epsilon\sigma_T + \epsilon\cos 2\alpha \sigma_{TT} + (\epsilon(\epsilon+1))^{1/2}\cos\alpha \sigma_{TL}]$$

Detecting the knocked-out proton in the direction of the three-momentum transfer, the interference terms σ_{TT}, σ_{TL} vanish.

This technique has been adopted in the NIKHEF¹² and Saclay¹³ experiments. The ratio σ_L/σ_T extracted from these experiments on ^{12}C and ^{40}Ca deviates significantly from the DWIA prediction. The order of magnitude of the effect is very similar to that observed in inclusive measurements and the tendency to increase with A is confirmed.

The consistence of this now rather extensive set of data tends to indicate that the classical Impulse Approximation picture breaks down when one studies into details the behaviour of the nuclear response functions. Various explanations to this break-down have been proposed.

Break-down of Impulse Approximation: proposed explanations

Different theoretical viewpoints are reviewed in Ref 3. and references are given. I briefly summarize the different kinds of effects which have been invoked:

1. Nuclear effects in classical framework.

They include:

- Collective effects, which are treated through RPA calculations.
- Mesonic Exchange Currents and Isobar contributions, which should affect mainly the transverse response functions.
- Ground state correlations. At momentum transfers greater than $2k_F$ (k_F = Fermi momentum of the nucleus) no Pauli blocking occurs and several theoretical estimates conclude that correlations are unlikely to contribute more than a few %.
- Final State Interaction effects. Handling these effects in the Impulse Approximation framework is a serious problem. Important corrections due to the distortion of the outgoing nucleon wave are to be considered. This problem is discussed by S. Boffi in these proceedings⁶.
- Distortion of the incoming photon wave. The "effective q^2 approximation" generally used to take into account effects beyond the one-photon exchange approximation is likely to become too crude for large A nuclei. Whether it is sufficient for medium nuclei like Ca and Fe remains to be proved on a rigorous basis. Several works are in progress on this subject.

2. Relativistic effects.

Treatment of the nucleon as a Dirac spinor in a Hartree mean field approach is in progress. In these models, the nucleon motion obeys the Dirac equation. The effect of the intense scalar σ and vector ω fields interacting with nucleons would be to shift down the mass and energy scales in the Dirac equation. The wave function of the nucleon as well as the current operator are affected: the combined effect could produce a quenched longitudinal response function.

3. Pionic effects on the proton electric form factor.

In this approach, the increase of the proton charge radius would be, similarly to the increase of the electromagnetic polarizability in the nuclear medium, due to distortion of the pion cloud.

4. Increase of the quark confinement volume in bound protons.

This hypothesis has received much attention because of its possible connection with the EMC effect. Change of a single mass or size scale, controlled by some QCD order parameter appeared an attractive unified explanation of low energy and high energy nuclear phenomena.

OPEN QUESTIONS

Present day data do not allow to draw definitive conclusions on the validity of all these working hypotheses. However, there are increasing evidences against the quenching of the proton magnetic form factor, as predicted by the "swollen" nucleon hypothesis:

1. The transverse response functions in inclusive measurements are fairly well reproduced without modifying G_M .

2. The y -scaling behaviour of the inclusive cross-section at high four-momentum measured at SLAC on ^3He , Al, Fe is destroyed by a " q^2 rescaling" of more than a few per-cent.

3. The q^2 behaviour of the elementary ep cross-section, extracted from the Saclay $^{40}\text{Ca}(e,e'p)^{39}\text{K}$ measurement in the range .3 to .6 GeV^2 appears quite incompatible with more than a few per-cent change in the proton magnetic radius.

Nevertheless, a number of questions remain open and controversial:

- The longitudinal response functions in inclusive measurements and the ratio σ_L/σ_T in $(e,e'p)$ measurements are still not understood. The limited (q^2, ϵ) domain and the accuracy of present experiments do not allow to rule out a quenching of the proton charge form factor.

- The effect of short-range correlations. In higher momentum measurements they should be minimized. This problem has been discussed by the working group on two-body correlations. See the report by J. Lightbody.

- The effect of Mesonic Exchange Currents and virtual Isobar excitations. They affect mainly the transverse components and it is very important that they be well under control, especially in view of possible $(e,e'n)$ measurements.

- Handling of Final State Interaction. This problem, crucial for the analysis of one or more nucleon knock-out experiments must not be addressed only from a theoretical point of view. Specific measurements will be necessary to check the calculations.

It is obvious that more complete and more accurate data on response functions are needed in order to isolate specific effects and distinguish between various theoretical predictions. An inclusive experiment with σ_T vs σ_L separation is planned at SLAC in the momentum range near 1 GeV² (Experiment NE9). Further $(e,e'p)$ experiments on ⁴He and ⁶Li are in progress at Bates, NIKHEF and Saclay. But, because of too low energy and duty cycle, the needed extension of these experiments towards higher q^2 is out of reach of present experimental facilities.

FUTURE RESEARCH AXES

From the discussions and contributions to the working group on "Single nucleon emission", several recommendations have emerged, which I summarize as follows:

1. Be as exclusive as possible.

This condition is necessary to single-out specific channels and disentangle various mechanisms: correlations, MEC, FSI, ... A good separation of the electric response function is of primary interest. Separation of the interference terms could bring valuable information on FSI. Some counting rates and accuracy estimates have been given by Morgenstern¹⁴ in the case of Deuterium. They showed that it is not easy to achieve a good accuracy in extracting σ_L and

σ_{TT} , even with a 2% accuracy on the cross-section measurement. Uncertainties in the backward angle vs forward angle comparison are amplified in the longitudinal vs transverse separation. Therefore measurements at large ϵ (e.g. small angles) are essential to obtain valuable information on the Coulomb response. They require high energy incident electrons.

2. Study systematically the effects of nuclear medium parameters.

The average nuclear density and the binding energy of the nucleons are believed to play a role in several theoretical approaches. Therefore, not only the dependance with A should be systematically studied, but also the dependance with the separation energy of the knocked-out nucleon. For that kind of investigations a good energy resolution (a few 10^{-4}) would be helpful.

3. Study the behaviour of the separated response functions in as large a momentum range as possible.

Clearly, a comparison of the "bound nucleon form factors" with the free nucleon ones necessitates to study the behaviour of σ_T and σ_L in a wide momentum range. This test is certainly the most powerful to make clear the question of a possible deformation of the nucleon in nuclei. However above $\sim 3 \text{ GeV}^2$ the longitudinal response becomes a few per-cent of the total response and thus, as stated above, very difficult to extract with accuracy, even with a 4 GeV incident beam.

CONCLUSIONS

It seems unquestionable to the participants to this working group that a part of the future (e,e'N) program at CEBAF should be devoted to deepen our understanding of the coupling of virtual photons to the elementary nucleonic current in nuclei.

The kinematical flexibility, the gain in counting rates and in true/accidental ratio brought by a high intensity, multi-GeV, CW accelerator like CEBAF will considerably open the field of this kind

of studies.

This field of research is indeed highly connected with other topics which have been discussed during the 1986 workshop. I already mentioned the working group on two-body correlations. But I would like to emphasize that the study of the nucleon resonances in the medium appears probably the most promising continuation of these investigations. Covering a large (q^2, ω) domain in exclusive measurements will certainly open an entirely new chapter of nuclear physics and allow to build a bridge with the high energy phenomena where quarks and gluons are expected to become the relevant degrees of freedom.

ACKNOWLEDGEMENTS

I would like to thank all the participants to the working group on "Single Nucleon Densities" for their active and stimulating participation to the discussions, and for having provided numerous valuable and helpful contributions. I am especially indebted to F. Gross and J. Mougey for the warm hospitality and constant encouragement I received during my stay at CEBAF.

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(e,e'2N) WORKING GROUP - SUMMARY

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Background

The (e,e'2N) reaction is one of the important areas that the CEBAF facility will be able to access by virtue of its high energy and CW beam character. Experiments of this type are to serve as part of the core program for CEBAF and, because of the need for three spectrometers, will have a strong impact on facility planning. Now is the time to focus on the details of this experiment and develop theoretical models of the process. It is very useful to try to project the impact, breadth, and feasibility of such research in order to prioritize it with respect to other research planned for CEBAF.

Up to the present time only the single nucleon knockout reaction has been studied extensively, and these studies have been confined mostly to quasifree kinematics due to large accidental coincidence rates. These experiments have revealed information on the one-body momentum distribution, in contrast to inclusive electron scattering studies which reveal integral properties of this distribution. By extrapolation, we expect that two-nucleon knockout experiments will reveal details of the two-nucleon momentum distribution and, in turn, information about the nature of two-body correlations in nuclei. A number of very important questions need to be addressed, including: what is the relation between perceived high momentum components of the one-body momentum distribution and short range N-N interactions; what is the role of the delta and explicit meson

exchange in virtual photon absorption at large momentum transfers; more fundamentally, how are momentum and energy transfer shared by the nucleons and constituent quarks; are there six quark, hidden color configurations manifest in two-nucleon knockout reactions?

The conventional "diagrammatic" understanding of the two-nucleon knockout process is shown in Fig. 1. The first of these diagrams (Fig. 1a) represents the effect of initial state correlations. It has long been recognized that, in the absence of final state interactions and meson exchange currents (MEC), the initial state two-body density distribution can be obtained from two-nucleon knockout cross sections. The relation is given very crudely at the bottom of Fig. 1. The importance of the pair correlation function is that it reflects how two nucleons interact in the nuclear medium, and it is intimately connected with the Coulomb sum rule and y -scaling, both of which are currently subjects of widespread interest. (The non-relativistic Coulomb sum rule is directly proportional to the Fourier transform of the two-body density, and the onset of y -scaling with increasing momentum transfer corresponds to the disappearance of high Fourier components in the two-body density.) It is, however, a terrible over-simplification to think that initial state correlations dominate the other processes. For example, delta production dominates the transverse virtual photon absorption spectrum in the 300-500 MeV excitation range, and it is known from pion absorption measurements in this energy region that there is a strong coupling to the two-nucleon channel. Therefore, a substantial fraction of what is observable in that region should be connected to pion absorption data that already exists. We should take advantage of this information in planning what new information can be derived from a program of two-nucleon knockout studies.

The choice of kinematics is particularly important. For example, one can exploit the insensitivity of the longitudinal response

function to the delta and other MEC effects. The explicit processes depicted in Fig. 1c-f are absent (at least non-relativistically) from the longitudinal part of the $(e, e'2N)$ cross section. Electrons can be detected at forward angles, where the longitudinal response dominates, in order to maximize our sensitivity to initial state correlations. In contrast to the pion absorption data, which is more a surface phenomenon because of the strong pion-nucleon coupling at resonance, photo-absorption is intrinsically a bulk process. We must, of course, recognize that because the delta has a finite lifetime, processes involving delta propagation in some sense average over very short range phenomena. Nevertheless, a comparison of transverse electron scattering to the two-nucleon final state, which is enhanced at larger scattering angles, with pion absorption to the same two-nucleon channel, may provide more complete understanding of medium effects on the process.

We must additionally deal with rescattering processes and final state interactions (Fig. 1b). Again, the choice of kinematics is important. We can enhance or suppress some of these effects by proper choice of outgoing nucleon momenta: using the fact that there is a broad minima in the N-N cross sections at about 200 MeV kinetic energy, one can minimize the effect of final state interactions with the residual system; and, choosing the relative momentum of the outgoing nucleons to be greater than 100 MeV/c, one can limit rescattering effects to of order several percent. Flexibility in positioning spectrometers close in angle to each other may also be important to achieving a full measurement program in two-nucleon emission, since we could then enhance or diminish final state interaction effects at will to calibrate our understanding.

Kinematics must also be taken into careful consideration when optimizing count rates. Forward angles for the electron spectrometer enhance the count rate dramatically. Clearance of the

spectrometers and beam line must be adequate in order to take advantage of this gain. In addition, we should consider the circumstances under which a single spectrometer can be used for detecting both protons in the case of $(e, e'2p)$ studies. We should follow the measurements of this latter type which are planned at Saclay^[FR85].

Theoretical predictions for the two nucleon knockout process are few in number. Laget^[LA86] has made a calculation for the $^3\text{He}(e, e'2p)$ reaction, using a Faddeev calculation for the ground state, taking MEC and rescattering effects into account, however not using a proper continuum calculation for the final state. There are other less well founded calculations or estimates of the $(e, e'2p)$ cross sections by Cioffi degli Atti and Mougey^[CI85], as well as by myself^[LI85], in the CEBAF 1985 RPAC Report. These calculations start with the single nucleon knockout spectral function and $(e, e'p)$ formalism. Phase space arguments are used to estimate the probability for detecting a recoiling second nucleon, assuming at large recoil momenta that the spectral function is dominated by short range two-nucleon interaction effects. Although some effects of final state interactions have been included in these latter estimates, they have less dynamics in them than the Laget calculation, including essentially only initial state correlations. Clearly, a great deal more work in this area is needed. Nevertheless, these works serve as our only guides for count rate estimates. The rates themselves are very low, and the quality of $(e, e'2p)$ or $(e, e'np)$ data will rest on the size of the accidental coincidence rates; i.e., $(e, e') \times (e, p) \times (e, p)$.

Average accidental rates for an n -fold coincidence reaction can be obtained from the expression $R_{\text{acc}} = (2T/DF)^{n-1} R_1 R_2 \dots R_n$, where R_i is the average singles rate in channel i . Fig. 2 shows the true and accidental rates for ^3He and $^{16}\text{O}(e, e'2p)$ at $E_e = 1$ GeV, an electron

scattering angle of 20° , at varying proton energy ($T_1=T_2$), and for proton scattering angles indicated in square brackets. These conditions correspond to zero recoil momentum of the daughter system. The results were obtained for a luminosity of $2.4 \times 10^{38} \text{ e}^-/\text{sec}/\text{cm}^2$, solid angles of 50 msr for the protons and 4 msr for the electrons, an energy interval of 20 MeV for the protons and 40 MeV for the electrons, a coincidence resolving time $2T=5 \text{ ns.}$, and a duty factor $DF=.01$. The true rates shown were based on my estimates for the cross section, which essentially agree with those of Cioffi degli Atti and Mougey, which, however, exceed the estimates of Laget. Without going into great detail on the interpretation of Fig. 2, the prime message here is that independent of duty factor the rates we are going to have to live with will be of order tens to hundreds of true events per hour! We should be able to obtain true rate to accidental rate ratios greater than one at 100 percent duty factor. We should also be able to increase the beam energy by of order 2X and decrease the electron scattering angle by 2X to obtain roughly an order of magnitude rate increase. Even with this increase, however, the separation of the cross section into the various structure functions will not be easy. We must therefore be sure that the spectrometers are designed to give us maximum possible solid angle and momentum acceptance, and with flexibility to get to far forward electron scattering angles.

Discussion

There was a wide spread in sentiment within the group regarding direct access to the pair correlation function, ranging from enthusiastic statements that this is the most important physics that we can address, to a deep skepticism that we will learn anything new. Within the optimistic camp, the great interest is in determining the nature of the N-N interaction inside the nuclear medium and examining the question of whether or not the interaction, or indeed the

nucleon itself, is altered in form from that determined by on-shell scattering measurements. Naively, one might expect the pion cloud surrounding the nucleon should be polarized at the 5-10% level by the presence of other nucleons, simply taking the free nucleon size and comparing it to the volume occupied per nucleon based on nuclear density. It is also interesting to ask how momentum is shared at the [(virtual photon)-(nucleon)] vertex when the nucleon is imbedded in nuclear matter. Are there, for example, meson exchange diagrams coupling more than two nucleons together at this vertex?

From the viewpoint of those holding more conservative, perhaps more realistic, attitudes towards $(e, e'2N)$ reactions, they see our lack of understanding of conventional ideas as a barrier to the unambiguous interpretation of two nucleon knockout data. Included in this area are the electron-nucleon off-shell behavior, relativistic effects, the nature of the pion-nucleon form factor vis-a-vis the cut-off mass in exchange current processes, and the lack of a theoretical framework with which to explicitly take pionic degrees of freedom into account.

One positive point was reiterated several times. The longitudinal structure function is the most promising to access short range correlations. Contributions to the longitudinal cross section from the two-body MEC matrix elements depicted schematically in Figs. 1(c-f) are absent non-relativistically and expected to be small in a relativistic treatment. Furthermore, the low energy loss region, below roughly 200 MeV, is the best kinematic region in which to look because the delta resonance, which dominates the transverse cross section, is still small. As a result, the extraction of the longitudinal response function, either by Rosenbluth separation or by measurement of the LT interference structure function, is least complicated. The low energy loss, high momentum transfer region is also interesting from the standpoint of y -scaling because that is

the region where interesting deviations from the best calculations have shown up. If one can isolate the longitudinal structure function, theoretical interpretation of the data must, however, include both initial state and final state correlations, as depicted in Figs. 1a and 1b.

One goal of studying correlations would be to understand the nature of the inclusive longitudinal (e, e') cross section. Recently, there has been a great deal of work on the Coulomb sum rule. The integrated longitudinal cross section has been found to be less than that expected from closure arguments. To explain this phenomenon, a number of people suggested that nucleon structure functions are modified in nuclei, corresponding essentially to larger nucleons. Fig. 3 roughly depicts the results obtained by Barreau et al.^[BA85]. Their data do not rule out a high energy loss tail on the one-body momentum distribution. If the sum rule were to be experimentally fulfilled, it would have to come from data taken at these higher energy losses. Within current understanding it is hard to generate much strength at large energy loss except by introducing very strong short-range two-body correlations. We have, however, not ruled out some less radical options to explaining the data: failure of the radiative corrections or failure of the one photon exchange mechanism. I believe this is what one would call a long shot, however it would not be hard to test at least the one photon exchange postulate by comparing e^+ and e^- scattering in the quasielastic region. It is certainly within the capabilities of the Saclay linac to probe this question and they are already contemplating such an experiment for elastic scattering. As a word of caution it does not take much of a discrepancy to alter the results of a Rosenbluth separation.

Fig. 4 shows transverse (e, e') data of Barreau et al. for ^{12}C in the quasi-elastic region. The calculations are from Shakin et

al. [SH86]. The dashed curve corresponds to single nucleon knockout with unmodified nucleon structure functions. The solid curve corresponds again to single nucleon knockout but with nucleon structure functions which essentially are needed to fit the longitudinal data. Clearly there is a discrepancy with the transverse inclusive data. However, the calculations can be reconciled with experiment if there is a substantial component of two-nucleon knockout in the inclusive cross section, indicated by the broad curve. This is another reason why it is important to have more and more exclusive data, even if it is aimed solely at understanding inclusive results. Fig. 5 schematically shows some Bates (e,e'p) data for ^{12}C from Finn et al. [FI86]. These data were taken in exactly the region shown in Fig. 4 where two nucleon processes may be important. The Finn et al. data reveal the normal 1p- and 1s-shell knockout peaks. They also show a broad flat component which Finn et al. speculate might be due to two nucleon emission. With actual two nucleon emission data we will be able to give definitive answers to the various theoretical questions posed.

The transverse structure functions are much more sensitive to the MEC effects depicted in Figs. 1(c-f). However, since these effects are largely isovector in nature, with a current operator antisymmetric in isospin, they couple much more strongly to the isospin-antisymmetric pn channel than the isospin-symmetric pp channel. By calculation, MEC effects in the np channel are roughly a factor of ten greater than MEC effects in the pp channel. We therefore have some handles to decompose various of the mechanisms at play in the transverse virtual photon sector. MEC effects are interesting in their own right as they clearly relate to the forces between nucleons. However, as discussed by Riska in his talk at this conference, we do not yet fully understand the nature of MEC effects and the N-N potential in a fully self-consistent formalism. In a study of MEC effects via the (e,e'2N) reaction it is essential

to have the capability of detecting neutrons as well as protons. Neutron detectors must be developed for this purpose.

Earlier, we discussed the subject of delta dominance in photo-reactions. As indicated in Fig. 1c, if we make a cut along the delta propagator line, there is no real distinction between this diagram and one in which a pion is absorbed leading to the delta. A great deal of pion absorption data leading to two nucleon-emission has already been taken. Roos has discussed this subject in his talk at this conference. Fig. 6 shows some interesting results from this pion absorption sector. In the 3-3 resonance region fifty percent or more of the absorption goes into the two-nucleon channel. It is not, therefore, fully clear that this type of data can not be used to the same end as $(e, e'2N)$ data. We must be careful to understand which of the isospin channels we are addressing by such studies; however, in pion absorption we have a complementary probe with which to study nucleon pairs. According to Van Orden^[VA86] the delta can be viewed as having a substantial lifetime, and therefore will average over short range phenomena, perhaps reducing the sensitivity to details of the two-body density. This averaging occurs during the process of delta propagation through the nucleus. We have learned a great deal about pion nucleus interactions through the concept of delta-hole dynamics. Hopefully, the $(e, e'2N)$ reaction, through its essentially bulk production amplitude compared to the strong surface absorption of pions through the delta resonance, will provide new information to test and refine these concepts.

Little has been said about hidden color, six-quark clusters in connection with two nucleon knockout. No one has come forth with predictions as to what the important signatures of such a pre-existent cluster would be. Depending on what the quantum numbers are for this system, there may be specific angular distributions which are favored in the final two-nucleon cm system. Hwang^[HW86] sug-

gested that there may be some interesting questions to address in this area. One troublesome problem is that the energies of the protons required to examine such questions are obviously not known. The criteria that we are currently using to design the spectrometers to study two-nucleon emission are not particularly well defined. For one of them, one chooses the momentum which minimizes final state interactions, roughly 1 GeV/c. We must be very sure not to exclude the capability for higher energy ($e, e'2N$) reaction studies from the experimental facilities. The only way I know of doing this is to preserve flexibility in the spectrometer designs, which means creating modular magnetic elements that can be reconfigured or added to, for extending the energy. This design philosophy was advocated by Duclos in one of his talks at this conference.

Experimental Considerations

In the study of two-nucleon knockout we must consider the possibilities not only of pp and pn emission, but also the possibility of deuteron emission. This certainly is easier from the standpoint of spectrometer requirements than triple coincidence studies. Deuteron and more complex particle emission was discussed separately by Berman and by Comfort. Fig. 7 shows data from an ${}^6\text{Li}(e, e'd)$ experiment at NIKHEF by Berman et al. [BES6]. Similar data for carbon exists that demonstrates the need for good resolution in order to resolve the final states of ${}^{16}\text{B}$. Knowledge of the final state of the residual system can tell details about the initial state of the two nucleons that may be missed in a low resolution pair knockout leading to the continuum. The only reservations regarding such studies concern the count rate. Targets must be very thin for deuterons and more complex systems to escape without serious attenuation. Experience with such studies at Saclay and at NIKHEF will be valuable to further planning. As with all the two nucleon knockout studies, there will be serious ambiguities regarding the pre-

existence of such clusters in complex nuclei, or their formation through final state pick-up reactions. Nevertheless, pp, pn, and deuteron knockout by electrons have different sensitivities and therefore should be considered as complementary aspects of a larger program.

There is a clear need for developing a neutron counter for the neutron part of pn knockout studies. It may be possible to operate a large scintillator array in a time-of-flight set-up to achieve large solid angle. There should be plenty of room for such a system in the planned halls. In addition, one can contemplate out-of-plane measurements with such a system because of the relative ease with which it could be moved vertically compared to a large acceptance magnetic spectrometer. This point was emphasized by both Lindgren and Berman during group discussions. The construction of a large solid angle neutron detector with modest energy resolution, whether it be of the time-of-flight variety or recoil proton spectrometer, would appear to be an essential piece of equipment for the lab to have. Such a device would also be useful in $(e, e'n)$ reaction studies.

In another detector related issue, one must consider the trade-offs of using magnetic spectrometers versus large acceptance detectors based on range telescopes. This is very important when counting rates are low and accidentals may represent a serious limitation. The true-to-accidental ratio for two-nucleon knockout is proportional to $(DF/2T/\text{luminosity})^2$. When the DF and coincidence resolving time are at their limits, the only remaining degree of freedom to control this ratio is the luminosity. If necessary to operate with reduced luminosity, one can then contemplate the use of large acceptance, non-magnetic detectors. The expense of reduced beam intensity or target thickness can then be made up for in increased detector acceptance. This is certainly an area that, like the neu-

tron detector, should receive serious consideration once the bulge in magnetic spectrometer design is passed, or the budget is exceeded.

Out-of-plane measurements received little attention. The spectrometer requirements are still not resolved, largely for lack of theoretical input as to the importance of the interference structure functions that these measurements address. The use of an LAD detector, such as the one designed by Mecking, to operate along the momentum transfer direction was discussed at the CEBAF Summer Workshop of 1985. A full ϕ -range would be accessible with such a device for particles moving substantially along the q -direction. Nothing has changed in this regard. However, Chang^[CH85] has discussed the importance of such a device in connection with possible detection of jet-like events as seen in high energy experiments.

Finally, it is important to consider the structure of accidental coincidence rates in more detail. As I discussed earlier, the accidental rate is proportional to the product of n -singles rates. Each of these rates is roughly proportional to the atomic weight. For example, electron scattering cross sections for evaluating singles rates in the quasi-elastic region are proportional to the sum of charges or the magnetic moments of the constituent nucleons. In the production of proton or neutron singles rates by the quasideuteron mechanism, the cross section is proportional to NZ/A . Similarly, for the production of nucleons by direct knockout or by production and decay of the delta, the cross sections are proportional to A . In this way, the true-to-accidental rate is proportional to A^{-2} , and the study of two-nucleon knockout for increasingly heavy nuclei becomes problematic. What really needs to be done to check this is (i) careful modeling of backgrounds and processes which contribute to uncorrelated singles rates in the various detectors and spectrometers planned for CEBAF, (ii) some realistic calculations of true

event rates, and then (iii) a simulation of the results of a triple coincidence measurement. It could be unfortunate should extensive experimental programs be planned without serious simulations aimed at detecting fundamental problems.

Recommendations

1. Foremost amongst our recommendations are the need for calculations. Primarily, we need these calculations for count rate estimates, and to see if there is a reasonable basis for an extensive experimental program. We also need these calculations to evaluate spectrometer options insofar as acceptancies and even angular ranges are concerned. The type of calculations that are needed can not be ordained, however it appears that little effort is going the direction of new quark degrees of freedom. The CEBAF project is rooted in the quark-nuclear physics interface and work in this direction should be encouraged. For the $(e, e'2N)$ core program, if it exists, some attention should be given to exploring the signature of exotic six-quark bags.
2. Evaluate the requirements on neutron detectors. They properly should be of major interest in both the single-nucleon and two-nucleon knockout programs. Such detectors have the potentiality of large solid angle and also of operating out-of-plane without the major laboratory resource commitment required for a magnetic spectrometer.
3. Attempt to preserve the flexibility in the experimental halls by using modular and demountable spectrometer systems, where possible. This follows the pattern of high energy physics experiments. Also, consider the more extensive use of large

acceptance scintillator range telescopes for high multiplicity events.

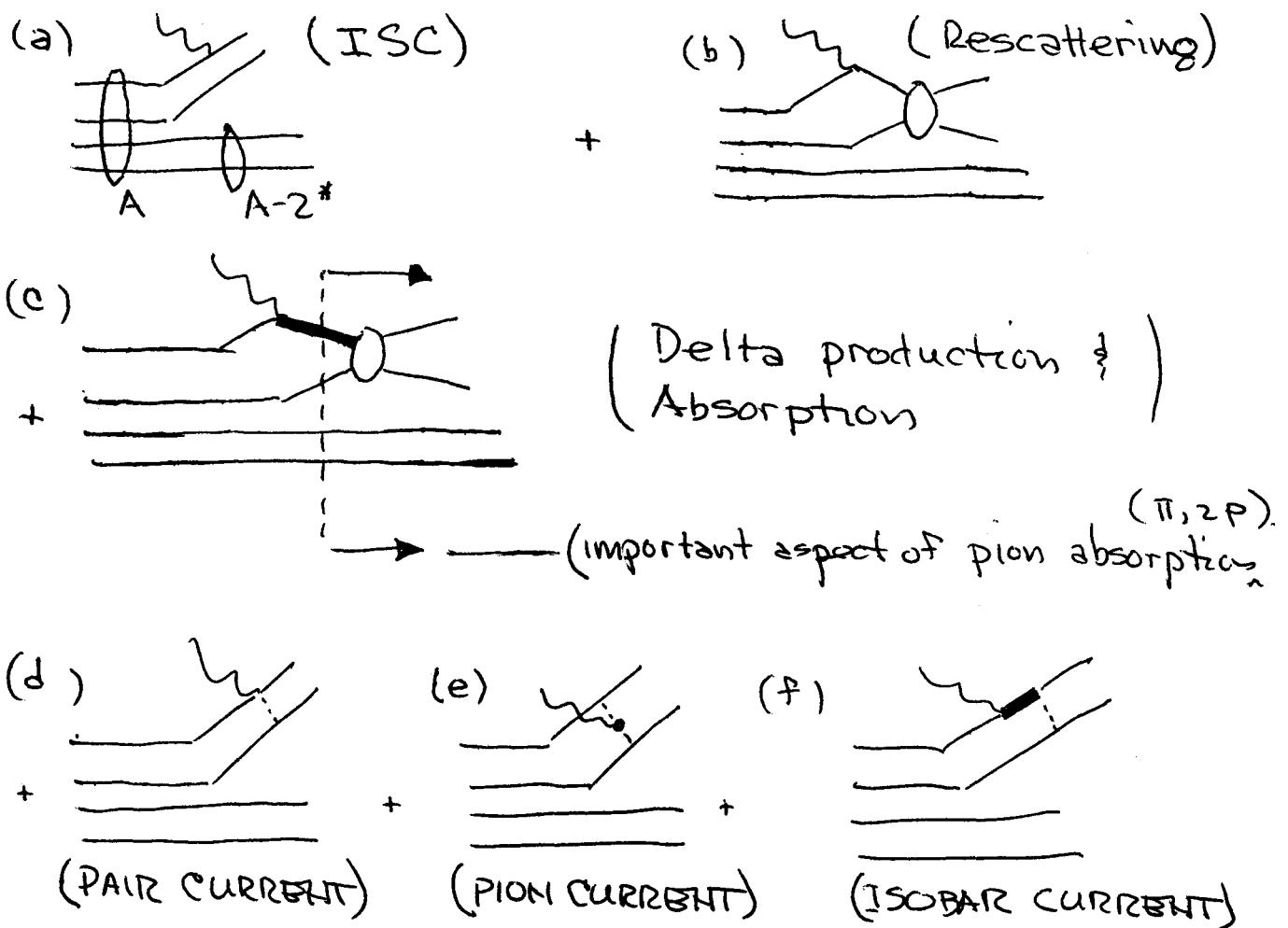
4. Change the focus of $(e, e'2N)$ reaction studies to include more of the systematics of $2N$ and higher multiplicity knockout, its relation to inclusive electron scattering, and its general behavior with respect to variations in momentum and energy transfer. This should be done for nuclei extending from few nucleon systems to complex nuclei. In addition, there should be an attempt to correlate these studies with existing pion absorption data vis-a-vis delta propagation and bulk versus surface delta production and decay to multinucleon final states.
5. Begin realistic simulations of two-nucleon knockout measurements. Modeling of both uncorrelated singles rates and true rates for light and heavy nuclear targets are essential to such simulations.

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Relation to 2-body densities & other processes



- Explicit MEC (d,e,f) small for pp knockout
" " pn knockout
- (c-f) Absent from longitudinal part of $T_{e,e'2N}$
- Initial state correlations (a)

$$d^{(9)}T \sim T_{eN} P_1 E_1 P_2 E_2 n(P_{em}) n(P_{rel}) \Gamma^{-1}$$

$$n(P_{rel}) = \int e^{i \vec{P}_{rel} \cdot (\vec{r}_1 - \vec{r}_2)} \rho(\vec{r}_1, \vec{r}_2) d^3 \vec{r}_1 d^3 \vec{r}_2$$

" related to Coulomb sum rule, γ -scaling

Figure 1.

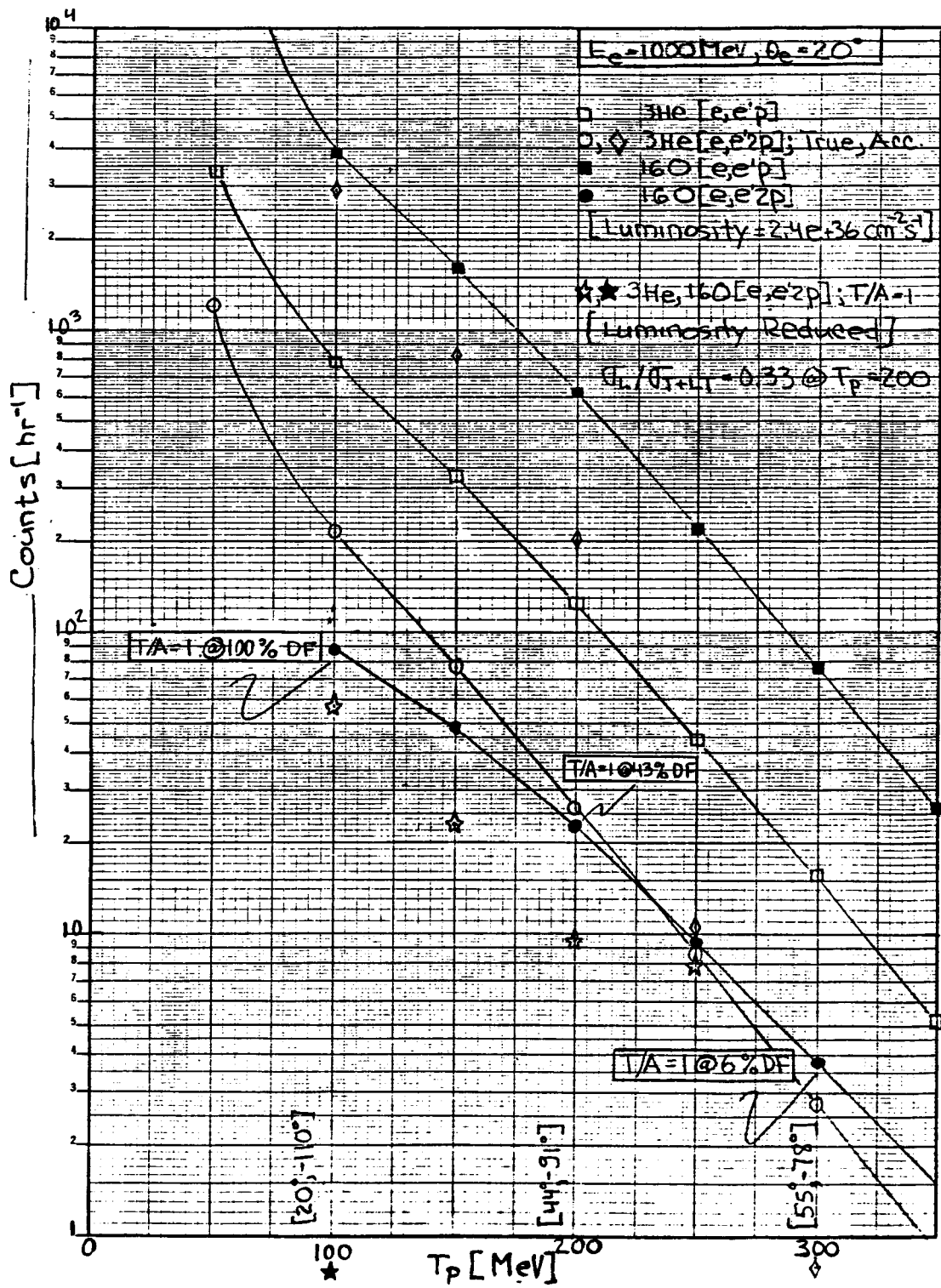


Figure 2.

Sadley (ee') Barreau et al



Source of tail

- (1) Correlations
- (2) Failure of radiative corrections
- (3) " " one-photon exch.

Figure 3.

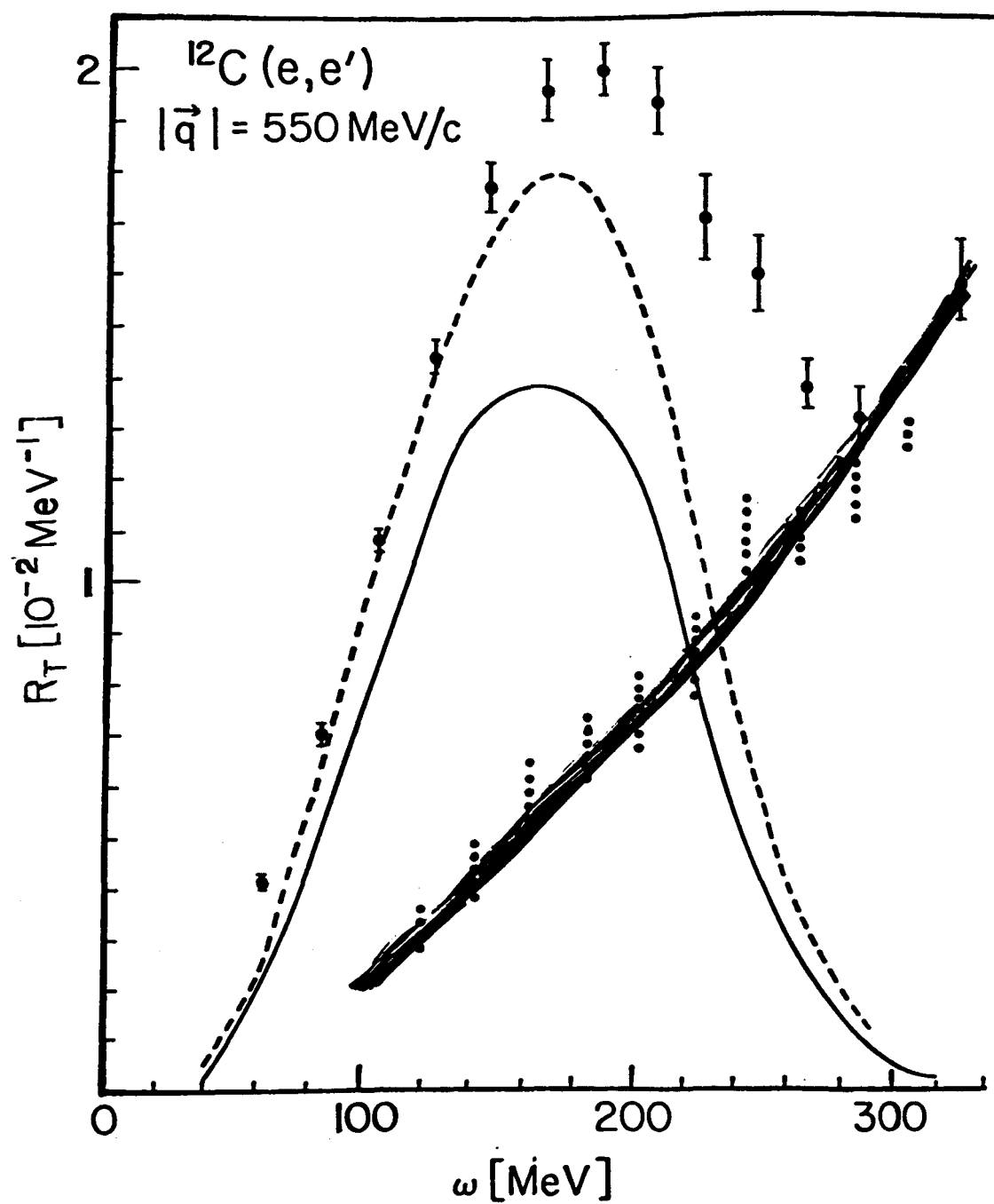


Figure 4.

Bates (e,e'p) Finn et al

^{12}C

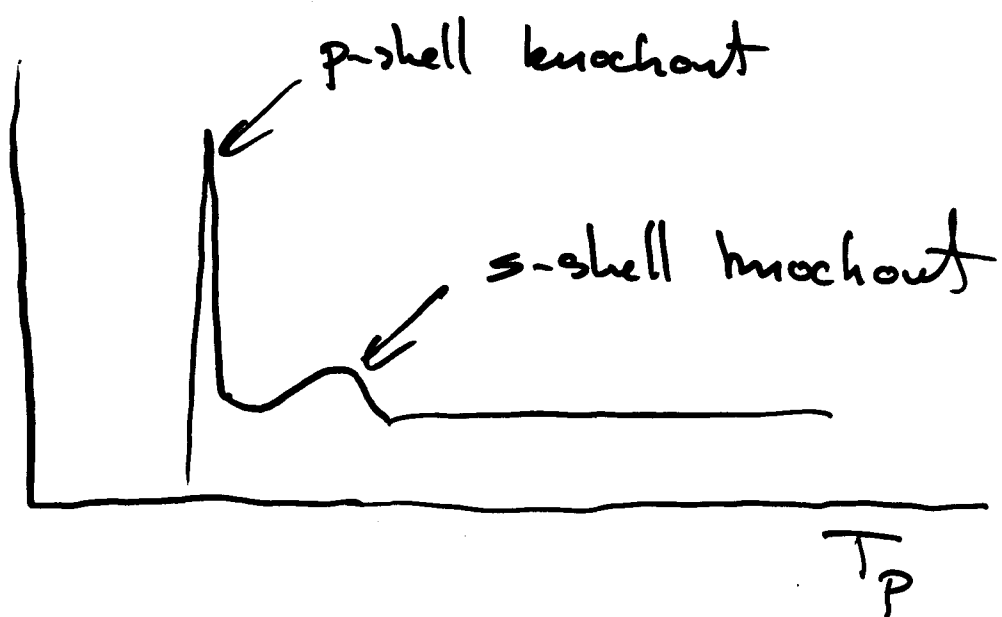
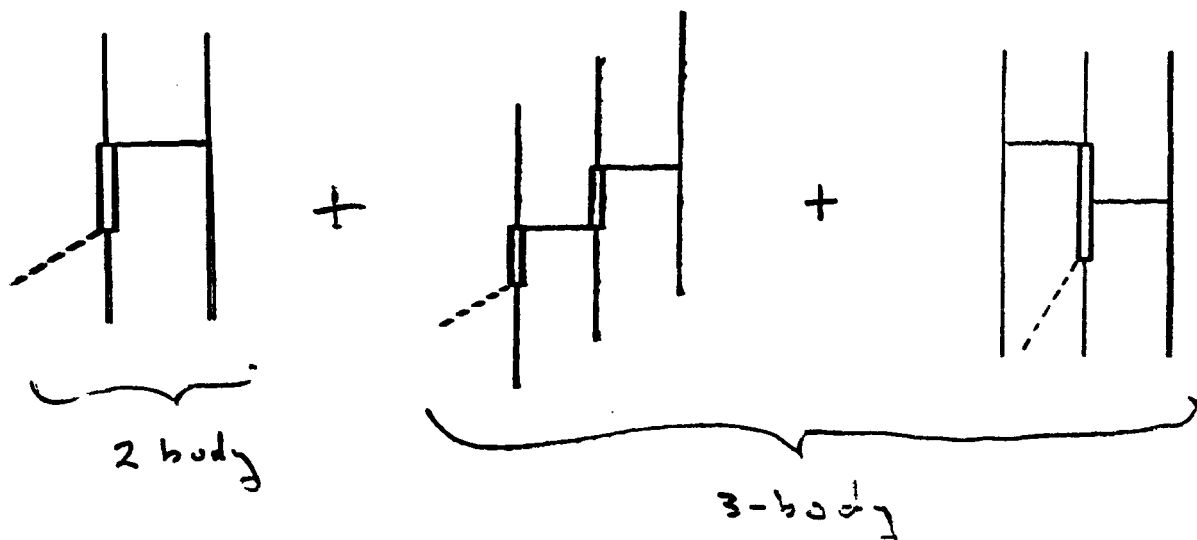
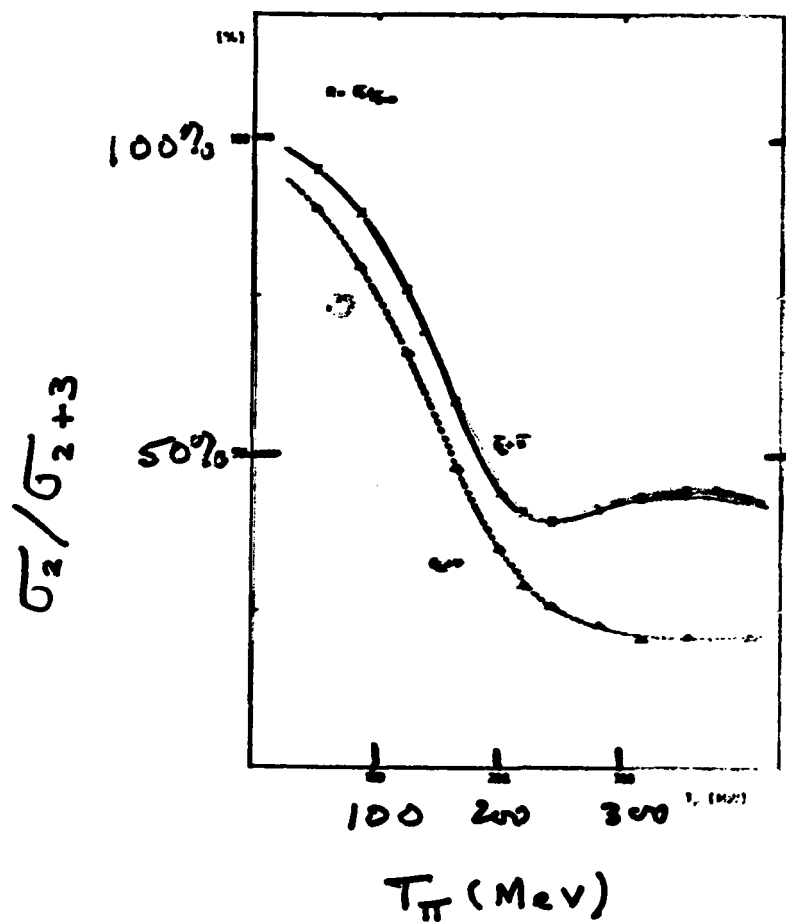


Figure 5.



70 Absorption cross Section



$$\sigma_{4B} < 15\% \sigma_{Abs}$$

Figure 6.

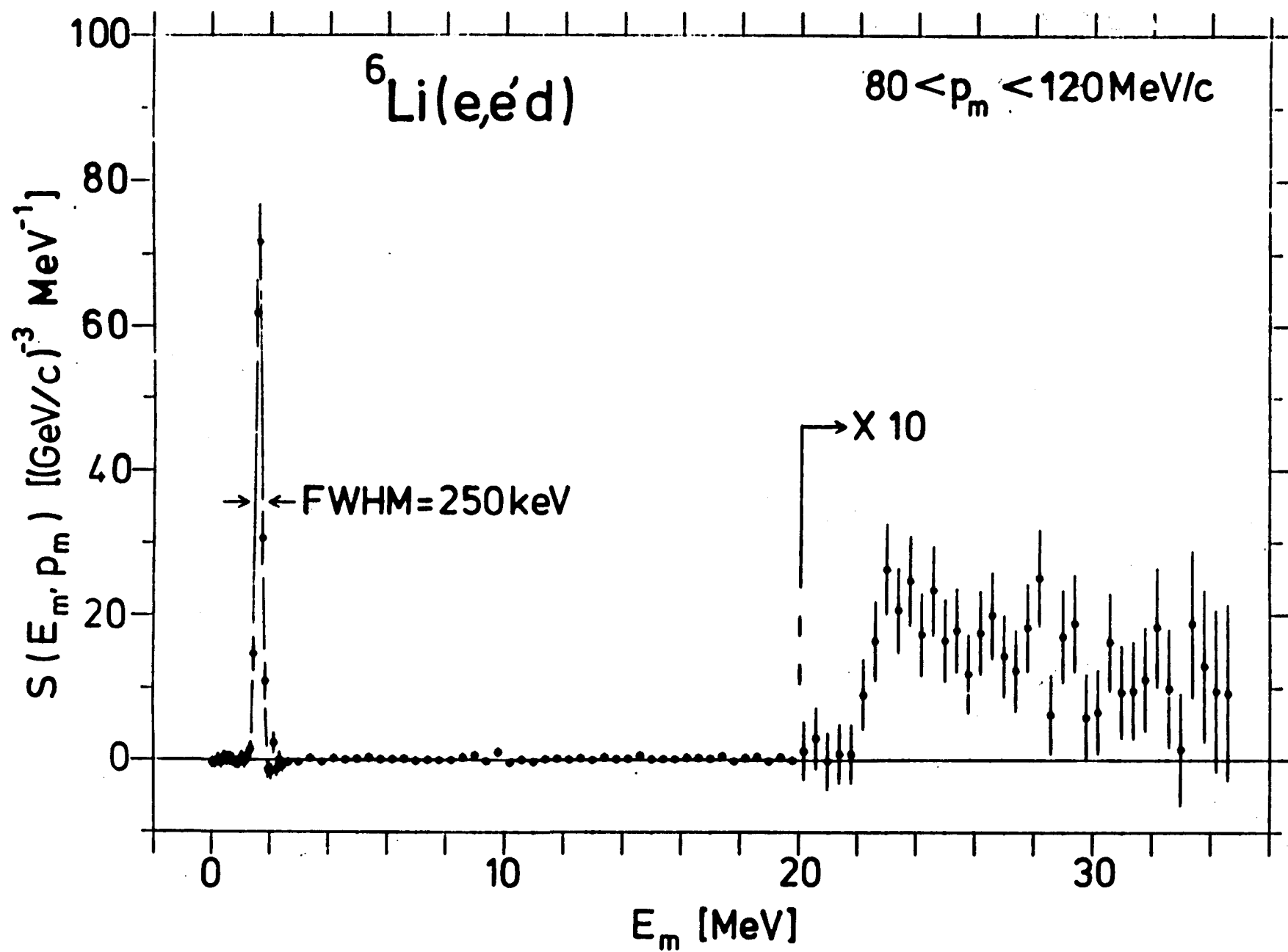


Figure 7.

EXCITED BARYONS*

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ABSTRACT

I give below a summary report on sessions of the working group on excited baryons at the CEBAF 1986 Workshop. It contains a tour of the theoretical horizon, some examples, experimental specifics and recommendations for interesting experiments in this area.

1. INTRODUCTION

Theoretical frontiers in particle physics can be visualized in two distinct horizons at the high-energy and low-energy ends. In the former, the main goal is to go beyond the standard $SU(3) \otimes SU(2) \otimes U(1)$ model, which is comprised of the demonstrably correct Glashow-Weinberg-Salam theory and hopefully correct quantum chromodynamics (QCD), ultimately to a grand unification of all interactions. Here one wants to find new experimental signatures which would hint at the need for new theoretical enterprise or confirm extant theoretical expectations. Examples of such discoveries would be finding top quarks, Higgs particles, evidences for techniparticles, discoveries of supersymmetric remnants and so on. At the other end lie possible experimental discoveries of nonzero neutrino mass, lepton flavor-violating decays, proton decay, etc. A truly remarkable opportunity at the low-energy end is the development and testing of the nonperturbative version of the QCD. From this, we would have to understand the structure of hadrons in their "free" state, or as aggregates in

* Invited talk at the CEBAF 1986 Workshop, Newport News, VA.

the environment of the atomic nucleus, where modifications of properties of free hadrons remain distinctly possible. Out of this should emerge an ultimate understanding of the atomic and molecular matter, from which stars, galaxies and the large-scale structure of the universe evolve. The Continuous Electron Beam Accelerator Facility, or CEBAF, now being built at Newport News, VA, would provide us with the opportunity to study critically the electromagnetic structure of hadrons in the excitation energy region of a few GeV, which is rich in resonances. Thus, for particle physics, it will provide an ample supply of electroweak data to test the nonperturbative domain of the QCD. In the context of nuclear physics, the hope is to build on the quark-gluon picture of free hadrons and uncover new phenomena and understanding of the multi-hadron environment. Finally, by duality, all of these should be intimately connected with the high-energy aspect of the QCD.

The task of our working group at the CEBAF 1986 Workshop was to examine experimental concepts, discussed in the book entitled Research Program at CEBAF (1985) [RPAC-85], in the area of excited baryons; we were also asked to suggest new ideas or modifications that might have a strong bearing on the design of future experimental halls and equipment. On the theoretical side, our desire was to take stock of theoretical issues that might motivate future works on the electromagnetic structure of baryons at CEBAF. We had five theoretical contributions: from Pfeil, Thomas, Browniowski, Manley and this author. Experimental contributions were twice as many: from Burkert, Mecking, Crannell, McCarthy, Stoler, Gastaldi, Schwille, Miskimen, Perdrisat and Minehart. I discuss some of these as I go along.

Due to lack of time, I would not be able to do justice to these talks here. The reader should consult contributions of these authors in these proceedings or in RPAC-86. An important achievement of our working group was the formation of an ambitious collaboration of several universities and institutions (ten at the time of this Conference). This was called the LAD (Large Acceptance

Detector) collaboration. Interested readers may contact B. Mecking, R. Whitney (both of CEBAF), or H. Crannell (Catholic U.) for further information on this collaboration and its current status. For this development alone, our working group has justified its long working hours. I shall return to the role of LAD in the excited baryon spectroscopy, a little later.

In the next few pages, I shall try to give an overview of the main themes of our working group and its recommendations. I have to avoid here too much technical details. My hope is to help people, who were not at our sessions, understand its issues, and to impress on Ingo Sick, our conference conscience, on its importance. It should also serve, for members of our working group, as an interesting point of reference for further discussions.

Let me divide the remainder of my summary in the following fashion: theoretical horizon, some examples, experimental specifics, suggestions and recommendations. Some overlap with my introductory talk is unavoidable, and is, indeed, necessary!

2. THEORETICAL HORIZON

At present, theorists firmly believe that the correct theory for strong interaction is QCD. It is a non-Abelian gauge theory, which is thought to provide a fundamental basis to understand all strong interaction phenomena in nature. At high energy, the effective gauge coupling α_s of the theory becomes small, and deep inelastic scattering of leptons from hadrons reveals point-like quark structure of hadrons. The vanishingly small "strong" coupling at high energy incorporates the phenomenon of "asymptotic freedom". At the low-energy side of this theory, the coupling constant becomes large, providing for the "infrared slavery" of quarks. In effect, hadrons become prisons from which quarks cannot be liberated in a free, isolated state. This, at once, "explains" why colored states are not produced in nature. [One, of course, discounts here experiments of Fairbank and collaborators¹

at Stanford, who seemed to have detected fractional charges in niobium.]

QCD allows one, at high energy, to use perturbative methods to test the theory by confronting it with experiments. For example, one can look for processes in which photon (or gluon)-quark Compton scattering or its crossed reactions, occur as parton subprocesses. Another example is to test scaling in the Drell-Yan process, $q\bar{q} \rightarrow \gamma^* \rightarrow \ell^- \ell^+$, seen in $pp \rightarrow \ell^- \ell^+ X$. Such tests² are thought to be quite successful, even though they can never match the stringency of tests of quantum electrodynamics.

Unfortunately, the situation at low energy is less than satisfactory. Here we are in the nonperturbative regime of QCD. Lattice methods³ are currently believed to be the only reliable way to test QCD ideas in this difficult region. The lattice approach starts with the problem of dealing with the Feynman path integral for the thermodynamic partition function involving the Euclidean action $S(A_\mu, \psi, \bar{\psi})$, where A_μ is the glue field and $\psi, \bar{\psi}$ are quark fields. It then imposes a short-distance cut-off by introducing a space-time lattice with $N_t \times N_s^3$ sites for the lattice spacing, a . One then evaluates expectation values of operators, using a Monte Carlo strategy. Eventually, a physics result is hoped for in the continuum limit in which $a \rightarrow 0$. Due to limitations of the present-day numerical calculations, we are a long way from the real world of excited baryons. According to a recent observation of a specialist⁴, one needs "Cray Millenia" before truly reliable results can be obtained and confronted with experiment.

In the interim, a variety of models are being developed as candidates for the desirable QCD-motivated low-energy approach to hadron structure. We may start with the old nonrelativistic $SU(6)$ picture known since the sixties⁵. Here nearly all meson and baryon resonances observed are embraced in a few multiplets of representation of $SU(6)_q \otimes SU(6)_{\bar{q}} \otimes O(3)$. Thus, most baryons are in the S- and D-wave

56-plets and a P-wave 70-plet. Most mesons are in the S- and P-wave 36-plets. Color provides here the antisymmetry needed by the Fermi-Dirac statistics of quarks. In the limit of mass degeneracy in a given multiplet, the p, n, λ quarks are degenerate and inter-quark forces are spin-independent.

Beyond the simple group theory, a variety of hadron models exist that emphasize at least some aspects of QCD relevant at the nonperturbative level. Thus, bags of different types⁶⁻⁸ have quark and gluon confinement by fiat and they may or may not incorporate chiral symmetry. Quark shell-models also emphasize confinement, and strong spin-dependent color magnetic forces suggested by QCD. Solitons⁹, such as Skyrmions, arise in a natural way in theoretical approaches in which QCD is equivalent to a meson theory in some well-defined context, as proved¹⁰ by 't Hooft and Witten, following pre-QCD ideas of Skyrme¹¹. Skyrmions thus serve as interesting dynamical models for baryons. Finally, the theory of strings¹² presents yet another new frontier, with many theoretical attractions. At present it begs the question: "Where are the quarks, gluons, etc., in these models?" It is a bit like a set of violins waiting for a Mozart or a Beethoven!

Out of this bewildering variety of hadronic pictures at low energy, perhaps an interesting one for the excited baryon spectroscopy is the quark shell model. In its most "modern" version, advocated by De Rújula, Georgi and Glashow¹³, and refined by Isgur, Karl and collaborators¹⁴, the baryon Hamiltonian is of the form

$$H = \sum_i m_i + H_o + H_{hyp}, \quad (1)$$

where

$$H_o = \sum_i \frac{p_i^2}{2m_i} + V_{conf}, \quad V_{conf} = \sum_{i < j} v_{conf}^{ij}, \quad (2)$$

$$v_{conf}^{ij} \approx \frac{1}{2} K r_{ij}^2, \quad (3)$$

$$H_{\text{hyp}}^{ij} = \frac{2\alpha_s}{3m_i m_j} \left\{ \frac{8\pi}{3} \vec{S}_i \cdot \vec{S}_j \delta^3(\vec{r}_{ij}) + \frac{1}{r_{ij}} \left[\frac{3\vec{S}_i \cdot \vec{r}_{ij} \vec{S}_j \cdot \vec{r}_{ij}}{2} - \vec{S}_i \cdot \vec{S}_j \right] \right\}, \quad (4)$$

with i being the quark index. The parameters of the theory are m_i ($i=u,d,s$), K and α_s . The interaction (4) is the well-known hyperfine interaction, familiar in atomic and nuclear physics; in the quark context, it arises in the nonrelativistic limit of the one-gluon exchange between a quark pair. A nuclear physicist looking at Eq. (4) is immediately reminded of the rich nuclear literature on the D-state contribution to the ground-state wave function of deuteron, due to the tensor force arising out of the one-pion exchange between nucleons. The color magnetic tensor force in (4), not surprisingly, also gives rise to the SU(6)-violating D-states in the nucleon and delta, in an entirely analogous fashion.

The power of the Isgur-Karl approach is in its ability to correlate an enormous amount of baryon spectroscopic data (the corresponding meson theory is equally powerful for mesons). It provides us with a natural framework to discuss, say, the electromagnetic transition properties between a pair of baryons. It also predicts baryon states with definite electromagnetic (and strong) properties that can be confronted with the powerful experimental tools to be available at CEBAF. In short, the Isgur-Karl theory is to baryon spectroscopy, what the nuclear many-particle shell model is to nuclear spectroscopy.

Before concluding this section, I would like to give a pictorial connection, due to G. Karl, between various models of hadrons as a way of orienting your ideas on excited baryon spectroscopy at CEBAF (Fig.1).

Hadron Spectroscopy

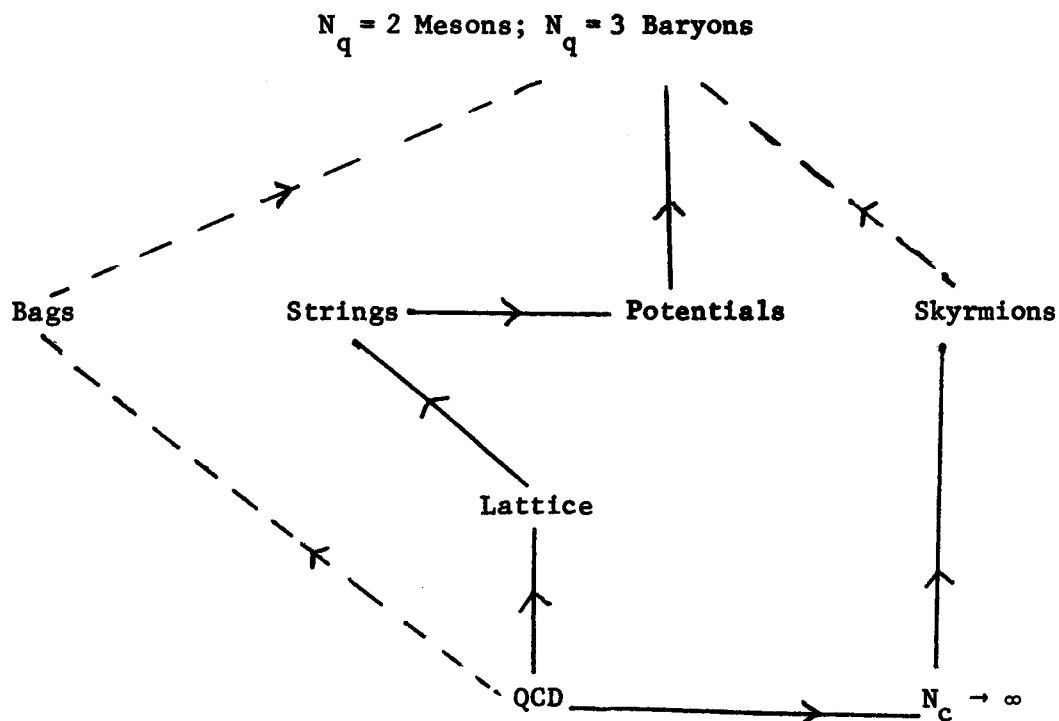


Fig.1: Schematic relation between some approaches to hadron structure and its connection to hadron spectroscopy [after G. Karl¹⁵]. Solid lines are direct connections, dashed lines are suggestive linkages.

Concerning the hadronic modelling in the direct context of CEBAF physics, we had two contributions at our working group. Broniowski¹⁶ reported on the work of the Maryland School. This group treats the structure of $N(939)$, $\Delta(1232)$ and the Roper resonance in the chiral soliton model. An interesting point of this work is the explicit inclusion of quarks, thus differing from the Skyrme approach in its pure form. Broniowski showed that spectroscopic properties of the N and Δ are successfully reproduced, within 20%, in their work, but those of the Roper resonance are not. The contribution of Manley¹⁷ examined implications of the quark shell model on the spectrum of non-strange baryon resonances. It also explored their inelastic decay modes. Manley¹⁷ emphasized some resonances predicted in this model that decouple from the πN channel. These would be prime candidates for study via $\gamma N \rightarrow 2\pi N$ and $\gamma N \rightarrow 3\pi N$ reactions, where γ is a real or virtual photon. More on this a little later.

3. CONCRETE CASES OF INTEREST FOR CEBAF EXPERIMENTS - SOME EXAMPLES

From the survey of recent theoretical approaches to the hadron (specifically, baryon) structure given above, it is clear that a complete theory of the hadron structure based on the interactions of quarks and gluons is unknown at this time, even as there are many models of this with interesting features. Thus, in the non-relativistic quark shell model, the relativistic effects are ignored, and the confinement mechanism is oversimplified. The correctness of the residual interaction is not proved, it is only tested through its phenomenological implications. In bags, the quark confinement is by fiat, and there is not enough dynamics. Quantum corrections to the soliton theory are uncertain. How to go from the large number of colors to only three in the "real-life" QCD is not clear. To quote Witten¹⁰, "Although QCD is equivalent to a meson theory, this is not a simple theory such as the nonlinear sigma model or Skyrme model. It is an extremely complex and unknown theory with an infinite number of elementary fields." With the present state of

computation, the lattice gauge theory is far away from predicting realistic answers for baryon spectroscopy. Strings carry ideas for future developments. In short, critical experimental tests are needed at CEBAF energies in order to sort out predictions of different models pertaining to the quark-gluon structure of hadrons. In this section, we shall discuss a few concrete examples of such tests possible at CEBAF. Some of these were emphasized in contributions of Pfeil, Broniowski, Thomas and this author to the excited baryon working group.

1) The Nucleon to Delta (1232) Transition Amplitudes

The photon transition matrix elements between the nucleon and delta (1232) have been a subject of intense theoretical and experimental research over many years. In simple quark models ($SU(6)$, $SU(6)_W$), the nucleon and delta are spherical. Hence the electric quadrupole (E2) transition amplitude is zero. This is the statement of the theorem due to Becchi and Morpurgo¹⁸. In Table 1, we display predictions of different hadronic models for the ratio $E2/M1$, $M1$ being the dominant magnetic dipole transition amplitude. Notice the wide spread of this quantity in different models. It is amusing to recall a result due to Carlson¹⁹ from the asymptotic QCD. He has shown that the hadronic helicity conservation constraint of the theory alone is enough to predict this ratio: it is unity! While this result has no obvious bearing at CEBAF energies, it demonstrates the importance of measuring the electric quadrupole transition amplitude accurately.

Even more interesting is the prospect of measuring the E2 and the scalar quadrupole $S2$ (or equivalently, the longitudinal quadrupole $L2$) amplitude for arbitrary mass squared of virtual photons in the nucleon to delta transition. The $SU(6)$, $SU(6)_W$ arguments again imply these to be zero, while more realistic hadronic models predict their nonzero values as functions of virtual photon mass squared.

Table 1. Predictions for the Ratio E2/M1 for the Nucleon to Delta photon Transition in Various Models

Model	Prediction
SU(6), SU(6) _W [Ref.18]	0
Quark shell model (Ref.20]	$\sim -3 \times 10^{-3}$
MIT Bag [Ref.21]	0
Chiral Bag [Ref.22]	$\sim -9 \times 10^{-3}$
Skymion [Ref.23] & Maryland Soliton [Ref.16]	$\sim -5 \times 10^{-2}$
Asymptotic QCD [Ref.19]	1

Using the presently available experimental information on the M1 and E2 multipole amplitudes for real photons in the nucleon to delta transition, extracted by Pfeil and Schwela²⁴, and by Berends and Donnachie²⁵, from cross-section and polarization data, we²⁶ have tried to extract the ratio E2/M1. Our primary concern in this analysis has been the model-dependence of the separation of the background effects from those of the resonance. Using an approach pioneered by Olsson, we have fitted two gauge couplings g_1 and g_2 pertaining to the $\gamma N \Delta$ vertex, by computing the background contributions at the tree level, and by insisting on unitarity via the Watson theorem. We have obtained $E2/M1 = (-1.5 \pm 0.2) \times 10^{-2}$. The error obtained here has been found to be too optimistic to reflect all theoretical uncertainties of the effective Lagrangian theory. A more thorough analysis²⁷ has subsequently given us this ratio between - 0.5% to - 1.6%.

A theoretical investigation of the longitudinal/scalar multipole amplitude in the nucleon-delta transition in the context of the quark shell model by Bourdeau and this author²⁸ has revealed an

interesting problem. While the L2 and S2 amplitudes are not independent and related by gauge invariance, the actual model calculation in the truncated Isgur-Karl basis does not obey this constraint, a problem known²⁹ in other areas of many-body physics. Our contribution²⁸ to the RPAC-86 emphasizes this point. We believe the S2 amplitude is less sensitive to the truncation of the quark model space and should be confronted with experiment at CEBAF.

In the CEBAF 1986 workshop, Pfeil³⁰ clarified some of the problems in extracting multipole (helicity) amplitudes from the existing data. Burkert³¹ and McCarthy³² dealt with the importance of polarized targets and beams at CEBAF in extracting interesting amplitudes like E2 and S2. Miskimen³³, Perdrisat³⁴ and Minehart³⁵ discussed recent proposals at Bates, SLAC-NPAS and other labs on measuring these amplitudes. It became clear from these discussions how important CEBAF will be in attacking problems of this kind.

2) The Electromagnetic Structure of the Roper (1440) Resonance

The Roper resonance belongs to the $N = 2$, 56, 0 multiplet in the SU(6) scheme. Its mass is between 1400 and 1490 MeV and full width 120 to 350 MeV, with sizeable $N\pi$ and $N\pi\pi$ decay channels. It has been recently suggested by Arndt that it might be actually split into two resonances. [In our conference, Pfeil³⁰ criticized this line of argument.] Its mysterious electromagnetic properties have been discussed at length in the RPAC-85. From the analysis of the older DESY and NINA data at $Q^2 = 1 \text{ GeV}^2$, there is a strong suggestion³¹ that the coupling of this resonance to the longitudinal photon is appreciable. We heard at the CEBAF-1986 Conference new proposals from Perdrisat³⁴ and Minehart³⁵ on the possibility of studying this resonance at Bates and SLAC-NPAS, but these pre-CEBAF prospects did not look very encouraging.

Burkert³¹ presented arguments at this conference on how the measurement of target asymmetries under special experimental conditions would be sensitive to the strength of the Roper excitation.

We refer the reader to the contribution³¹ of Burkert to the RPAC-85 for details. Clearly, the possibility of utilizing beam and target polarizations at CEBAF makes these experiments front-line possibilities.

From the theoretical point of view, the Roper resonance will be a test case for many models. Thus, soliton models have difficulty in handling properties of this resonance, as we heard from Broniowski¹⁶. Quark shell model calculations³⁰, used by Burkert, require serious scrutiny in the context of the role of the non-resonant background, and the truncation of the quark model space, in determining the longitudinal transition amplitude. We need to understand large differences³⁰ among predictions of the helicity amplitude $A_{1/2}$ as a function of Q^2 , by different authors, all utilizing quark shell models.

3) Where are Missing B^* 's?

Isgur has noted recently that quark shell models predict more baryon resonances than seen experimentally so far. For example, positive parity ($N=2$) N^* and Δ^* states predicted exceed those that are seen. One possibility of this is that the quark shell model is not counting correctly the relevant degrees of freedom; the other is that these resonances are indeed there, as yet to be discovered. Manley, in his contribution, emphasized the CEBAF interest in the excitation energy region around 2 GeV where such states are copiously predicted. Isgur and collaborators have stressed that many of these states are mixed between and within SU(6) multiplets (reminiscent of the complex configuration mixing in nuclear physics) and hence they decouple from the πN channel. These are good candidates to be studied by the $\gamma N \rightarrow B^* \rightarrow \Delta\pi$, ρN , ωN , etc., avoiding the πN channel which is dynamically unfavorable. Isgur has given a helpful table of photoproduction amplitudes of precisely such resonances in the 2 GeV excitation energy region. We reproduce here this table for ready reference of our readers.

Table 2. Quark shell model estimates of the photoproduction amplitudes of some baryon resonances around 2 GeV by Isgur and collaborators. Coupling amplitudes to the strongest channels are also shown. (From Isgur³⁹).

State	Photoproduction Amplitudes				Strong Coupling Amplitudes			
	$A_{3/2}^P$	$A_{1/2}^P$	$A_{3/2}^n$	$A_{1/2}^n$	$N\pi$	$\Delta\pi$	$N\rho$	$N\omega$
$N \frac{5}{2}^+(1955)$	+9	- 67	- 45	- 22	0	- 7	- 8	+12
$N \frac{5}{2}^+(2025)$	- 3	+ 2	+51	+15	1	- 7	+7	-11
$\Delta \frac{5}{2}^+(1975)$	+76	+61			1	+ 6	- 18	na
$N \frac{3}{2}^+(1870)$	+ 6	- 19	- 6	- 21	3	- 4	+1	- 10
$N \frac{3}{2}^+(1955)$	+4	- 16	- 39	+ 6	1	- 9	+ 6	- 9
$N \frac{3}{2}^+(1980)$	- 5	+19	+ 22	- 20	1	+ 9	- 7	+ 3
$N \frac{3}{2}^+(2060)$	0	- 1	- 15	+ 2	1	+ 5	- 3	+ 9
$\Delta \frac{3}{2}^+(1975)$	- 7	+18			0	- 8	+ 5	na
$N \frac{1}{2}^+(1890)$		- 20		- 1	4	+ 3	- 5	+ 6
$N \frac{1}{2}^+(2055)$		+ 7		+ 3	1	+ 2	- 1	- 5

Manley³⁷ and Mecking³⁸ emphasized that the predicted photoproduction amplitudes (Table 2) for a few of these missing resonances could be very large and extremely favorable for exploration at CEBAF. Multiparticle final states make these experiments very suitable for exploration by the large acceptance detector (LAD).

Clearly, discovery and identification of the "missing" resonances and measurements of their photo- and electroproduction amplitudes are of high priority at CEBAF. The effective excitation region between

2 to 3 GeV should be very valuable to explore. It should provide critical tests for the nonrelativistic quark shell model and bring out new insights on the effective quark interaction.

4) B^* in Nuclear Environment

During the last ten years, lepton and meson factories have provided us with insights on the production of $\Delta(1232)$ inside the nucleus and its propagation in the nuclear medium. From our experiences in this field thus far, we are apt to ask the following questions for excitation of any baryon resonance B^* in the nuclear medium, by real or virtual photons to be available at CEBAF:

- a) Is the γBB^* vertex modified in the nuclear medium?
- b) How does a particular B^* propagate in the nuclear medium?
What is the effect of the nuclear medium on its decay width?
- c) How do we handle theoretically and experimentally the production of many hadrons in the final state in the nuclear interior?
- d) Are there any interesting physics for nuclear production of the $\Delta(1232)$ at CEBAF energies?
- e) Is there an early onset of scaling inside the nucleus?

Questions like these were asked at the workshop, but answers were not plentiful. Undaunted by this, Stoler went ahead and produced computer simulations for CEBAF kinematics. Reproduced here is a typical rate estimate (Table 3) from his study concerning ${}^4\text{He}(e, e' p \pi^-)X$ reactions. Stoler concludes³⁸: "It appears that charged pion photoproduction experiments involving delta and B^* production in nuclei will be feasible at CEBAF." Amen, but we need more work in the meantime, both on theory and on experimental feasibility.

Table 3. Laboratory kinematic conditions and counting rates for the experiment ${}^4\text{He}(e, e', p\pi^-)X$ designed to separate the longitudinal and transverse components of the differential response. The pion decay angle in the delta frame is 160° . Pion and proton angles are relative to the incident electron direction. Γ is the virtual flux factor, ϵ , virtual photon polarization. The cm $p\pi$ energy is 1250 MeV. [After Stoler³⁶].

E_e (GeV)	E'_e (GeV)	θ_e (deg)	ϵ	Γ	P_p (MeV/c)	θ_p (deg)	P_π (MeV/c)	θ_π (deg)	$(dN/dt)p\pi$ (hr ⁻¹)
1.9	1.3	20	0.9	3	842	29°	144	159	500
0.87	0.35	60	0.44	.3	842	15°	144	145	50

My feeling is that the production of excited baryons, beyond the well-known $\Delta(1232)$, in complex nuclei at CEBAF kinematics is a virgin research area. We do not expect to get immediate guidance from theorists as to what to look for. Survey experiments may be needed to get first insights into this new area of research. Physics Advisory Committees may have to allow expeditions into uncharted waters before new physics becomes visible on this research horizon!

Anhang: In my summary, I have been avoiding any discussion of production of strange baryons and their resonances at CEBAF. This interesting subject had been partially covered in the working group on hypernuclear physics chaired by Ed Hungerford. Thomas⁴⁰ contributed a paper to our session discussing the possibility of $\Lambda^*(1405)$ as "inter-truder" state. He suggested that experiments at CEBAF could test whether this is a pure 3q state or a $\bar{K}N$ bound state. Schille⁴¹ reported on dibaryons.

Due to lack of time, we could not take up the topic of the search for pre-existing nucleon resonances in nuclei. An example of this would be the question of the $\Delta\Delta$ and other BB^* components of the deuteron wave function. We also could not discuss the possibility of exciting multi-quark exotic clusters in complex nuclei and electroweak physics connected to the test of the standard model at CEBAF kinematics. Clearly more work is needed to clarify these interesting issues. I apologize for any other unintentional omissions in the physics of excited baryons in the CEBAF context.

4. EXPERIMENTAL SET-UP AT CEBAF FOR EXCITED BARYON RESEARCH: SOME OPTIONS

Given the physics objectives discussed above, we can suggest some experimental options that are worthy of further considerations by interested physicists. [Here the present author is completely out of his elements! Volker Burkert and other experimental colleagues in attendance at the CEBAF 1986 Workshop are prime movers for this section.] These options are by no means exhaustive; they are presented with the hope that they would serve at least as talking points for future discussions.

(a) Physics objective: Measurement of electromagnetic transition form factors for $W \leq 2 \text{ GeV}/c^2$, $Q^2 \leq 3(\text{GeV}/c)^2$.

Process: $\gamma_V B \rightarrow B' \eta$, $B' \gamma$ (B: Nucleon, B': baryon of interest).

Target: cryogenic H_2 or D_2 with high cryo power. Polarized solid state targets (NH_3 , ND_3) with high cooling are also desirable.

Temperature range: $T \sim 0.25$ to 1°K . Magnetic field needed is 3.5 to 5.0 Tesla. Beam current I_e : Unpolarized $I_e \sim 30 \text{ nA}$ to $\sim 100 \mu\text{A}$. Polarized $I_e \sim 30 \text{ nA}$ to $100 \mu\text{A}$.

Detectors: 4 - 6 GeV e^- spectrometer of moderate resolution, large solid angle, large momentum acceptance. One hadron spectrometer with $p_m \leq 3.5 \text{ GeV}/c$. Particle identification, out of plane capability desirable. Resolution $\leq 5 \times 10^{-3}$. The spectrometer called VAS I is a suitable candidate for the hadron spectrometer.

(b) Physics objective: Study of $\gamma_V B \rightarrow B' \pi\pi$. Here an additional spectrometer is needed to detect the second hadron, of lower momentum compared to that in the case (a). The spectrometer called VAS II is a candidate for use in this case.

(c) Physics objective: $\gamma B \rightarrow B' \pi\pi$: Search for unseen states. One needs here the tagged photon facility. Large acceptance detector (IAD) is needed for the detection of hadrons in the final state.

(d) Physics objective: Study of B^* in nuclei.

Target: D_2 , light nuclei. Beam: Unpolarized $I_e \sim 10 \mu A$.

Detectors: For the process $A(e, e', p\pi)X$, one needs one e^- spectrometer of moderate to poor resolution, one hadron spectrometer for the large-momentum proton and another hadron spectrometer for the low-momentum pion, both of moderate resolution.

Contributions of Burkert³¹ and Stoler³⁶ to the RPAC-86 should be consulted for further details. Burkert and Mecking discuss in depth the designs of the VAS and LAD devices respectively.

5. CONCLUDING REMARKS

In the early nineties, experiments will begin at CEBAF. From careful considerations of the state of particle and nuclear physics at this time and with reasonable extrapolation of this to the not too distant future, it is safe to expect that the nonperturbative quark-gluon structure of hadrons either in free states, or as constituents of nuclear matter, will remain a subject of great intellectual interest at the time when CEBAF delivers its beams to experimental halls. Quoting Feynman again, "one powerful way of experimentally investigating the strongly interacting particles (hadrons) is to look at them, to probe them with a known particle; in particular, the photon (no other is known as well)." This is precisely what we have been discussing so far.

With the polarized targets and beams at CEBAF, a new era of exploring baryon structure will begin. It will produce new understandings of the role of color magnetism and other quark-quark interactions, suggested by QCD. It will allow us to explore the electro-weak structure of known baryon resonances to great precision and to find those resonances which are predicted to exist, but have eluded experimental discovery thus far.

We strongly recommend this new world to our future explorers.

I thank Franz Gross for his kind invitation to the CEBAF-1986 Workshop. Special thanks are due to all contributors to the truly exciting sessions of the excited baryons at this workshop. I am particularly grateful to V. Burkert, W. Pfeil and P. Stoler for many helpful discussions. I also thank J.S. O'Connell and A. Halprin of the Lewis Center for Physics. My research is supported by the U.S. Department of Energy.

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Report on the Workshop on Electromagnetic Production of Strange Nuclear Systems

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ABSTRACT

This paper reports the discussion of the Working group on Electromagnetic Production of Strange Nuclear Systems held at CEBAF, June 23-27. It contains a brief review of the status of photoproduction of strangeness and hypernuclei, and what one might expect to investigate with the CEBAF machine. Recommendations of the working group are presented in the conclusions.

INTRODUCTION

The electromagnetic production of strangeness -1 nuclear systems has been proposed and reviewed several times in the past.¹⁻⁵ In particular, the production of hypernuclei with energy resolution on the order of 200 keV offers the exciting possibility of exploring "new" physics. By this I mean physics beyond a conventional nuclear structure picture. How this might arise in experiments remains only reasonable speculation at present, however such speculation, can form the basis of new physical understanding.

A working group on the electromagnetic production of strange nuclear systems was held in conjunction with the CEBAF workshop, June 23-27. This review summarizes the results of the working group and interjects some additional comments. One immediately recognizes that the collective nature of a working group naturally elicits a distribution of excitation strength over the entire group, to the extent that some individuality cannot be ascertained. Thus it is not possible, everywhere, to acknowledge each particular individual, however material is referenced to the extent possible. A list of working group contributors is given in the appendix.

ELEMENTARY AMPLITUDES

The elementary amplitudes for $p(\gamma, K^+)\Lambda$ are not well measured. The cross-section between 0.9 to 4.0 GeV incident energy rises sharply from threshold (0.9 GeV) to a more or less constant value above 1.2 GeV.⁶ It falls slowly as a function of the kaon production angle for angles greater than 10° . Theoretical predictions of this cross-section rely on unknown coupling constants of the K and K^* to the Λ , Σ and their resonances.⁷ These constants can be derived under SU(3) symmetry using a quark model,⁸ but because of the large number of resonances that contribute, and the sensitivity of the calculation to the coupling constants, the calculated cross-sections are unreliable.

It would be useful to develop a "hard mass" limit as one approaches threshold for use in a similar way to soft pion theorems.⁹ To check the validity of such theorems, spin dependent data near threshold are necessary. Of importance is the spin dependent nature of the cross-section. Since the $p(\gamma, K)\Lambda$ amplitude contains a significant spin-flip strength, nuclear structures are excited that cannot be reached via (K, π) or (π, K) reactions. For example, in hypernuclear systems the (γ, K) reaction will excite non-natural parity states².

HYPERNUCLEAR CALCULATIONS

Uncertainties in the elementary amplitudes lead to uncertainties in the hypernuclear excitation strength. Therefore, without good elementary amplitudes it will be impossible to identify whether a deviation from the data is due to an error in the reaction mechanism or nuclear structure. However, because the reaction does lead to the excitation of non-natural parity states, spectroscopy different to the (K, π) and (π, K) reactions will occur. In the later reactions the weak spin flip probability at forward angles leads to the excitation of only natural parity states.

Because both the γ and K^+ weakly interact with other particles it is also possible, in principle, to excite hypernuclear levels

within the interior of the nucleus. This is in contrast to the (k,π) and (π,K) reactions which are peripheral in nature. However, it is noted that distortion of the K^+ in the final state may be significant for large A , particularly at higher momentum where pion production becomes important.

However, no matter how hypernuclear excitations are produced one has the usual problems associated with the theoretical treatment of a many body system. One could hope these problems would not be too large because of the weaker strength of the probes, but one always has ambiguities in the extension of the elementary amplitudes to "off-shell" values and in the incorporation of relativistic calculations or corrections.

The large momentum transfer of the reaction preferentially excites states reached through high angular momentum transfer. Momentum transfer at 0^0 as a function of incident momentum is shown in figure 1 and compared to the (K,π) and the (π,K) reactions. All (K,π) inflight hypernuclear transitions have occurred near the "magic" incident momentum, so that substitutional states are populated. That is, states are populated in which there is no change of single particle shell model quantum numbers between the nuclear and the λ . On the other hand the (π,K) reaction requires large momentum transfer and populates stretched angular momentum states. Both reactions reach only states of natural parity that are near the surface of the nucleus.

SUMMARY OF THE HYPERNUCLEAR PRODUCTION

There has been considerable effort in demonstrating that one can produce hypernuclei using the (e,e^1,K) reaction.^{1,2} In fact it has been shown that excellent resolution (< 250 KeV) can be obtained with reasonable counting rates for light hypernuclear systems.⁴ We now need calculations of cross-sections for heavier systems, par-

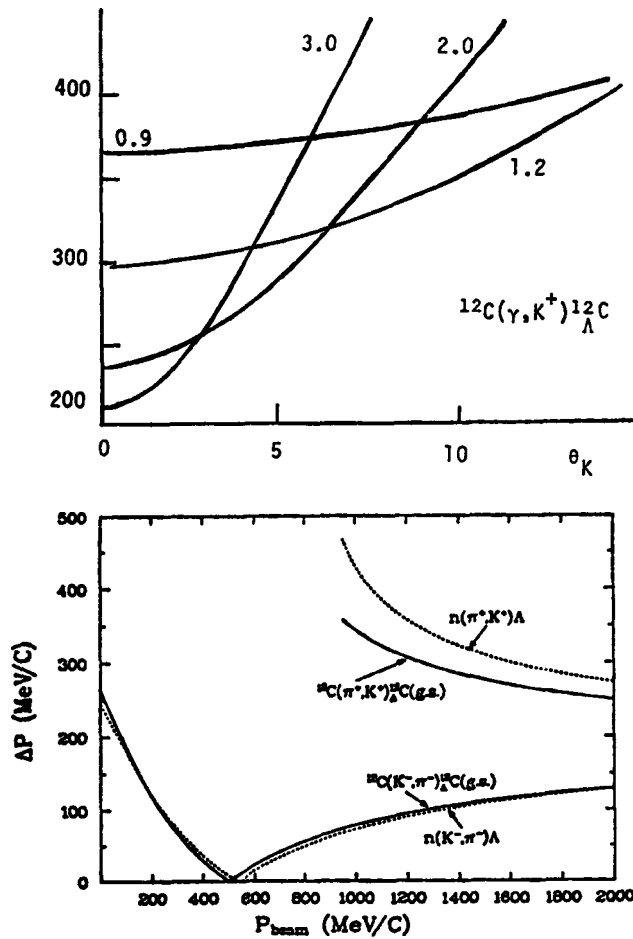


Fig. 1)
Momentum transfer for;
a) (γK^+) reaction as a
function of angle:
b) (π, K) reaction as a
function of the inci-
dent momentum

ticularly for states where the nuclear core is not excited and the lambda is deeply embedded within the nuclear medium.

Previous investigations of hypernuclear systems have reasonably mapped out the 1s and 1p shell hypernuclei using the (K, π) and (π, K) reaction.¹⁰⁻¹² The Λ effective interaction at least in the lower 1p shell is well determined, except for the tensor contribution which significantly contributes only for nuclei near the 1p shell closure.¹² There is however the open question regarding the charge symmetry breaking strength of the interaction.¹³ A large effect is observed in the $A = 4$ system that has been attributed to a number of factors including 3 body forces and Λ - Σ coupling. One needs to

investigate other light hypernuclear systems in order to shed more light on this subject. Because of the quality and intensity of the electron beam at CEBAF one can consider using separated isotope targets. Such targets would allow one to easily reach many hitherto inaccessible charge symmetric systems. In contrast, the mesonic reaction such as (π, K^0) ¹⁴ in comparison with the (π, K^+) reaction would be needed to address these issues.

A premiere hypernuclear program would involve the direct deposition of a Λ into the interior of a heavy nucleus, with the ability to identify the level into which it is placed. Such states could possibly lead to the observation of new physics. In these experiments one looks for the effects of partial deconfinement of the Λ into quarks. This could be more probable or easier to observe in strange systems because the strange quark is heavier and is distinguishable from the other quarks in the medium. For example, the binding energy of the hypernuclear ground states may be very different from what is to be expected in conventional models.¹⁵ In addition to the energy level structure the life time, and magnetic moment of the Λ may be altered. One might view the Λ as having a large effective size in the medium. In conventional terms a swelling of the Λ in the nuclear medium, may be one way of interpreting Λ - Σ type coupling. As has been mentioned, this can also lead to charge-symmetry breaking.

It is clear that the elementary amplitudes need to be carefully measured. This will require a measurement of both energy and spin dependence so that calculations can reliably predict the strength for the hypernuclear excitation. There are several other problems of somewhat technical interest. One needs measurement of the amplitudes to calibrate theoretical reaction calculations of hypernuclear systems, and the hypernuclear many body problems using the unique behavior of the reaction particles.

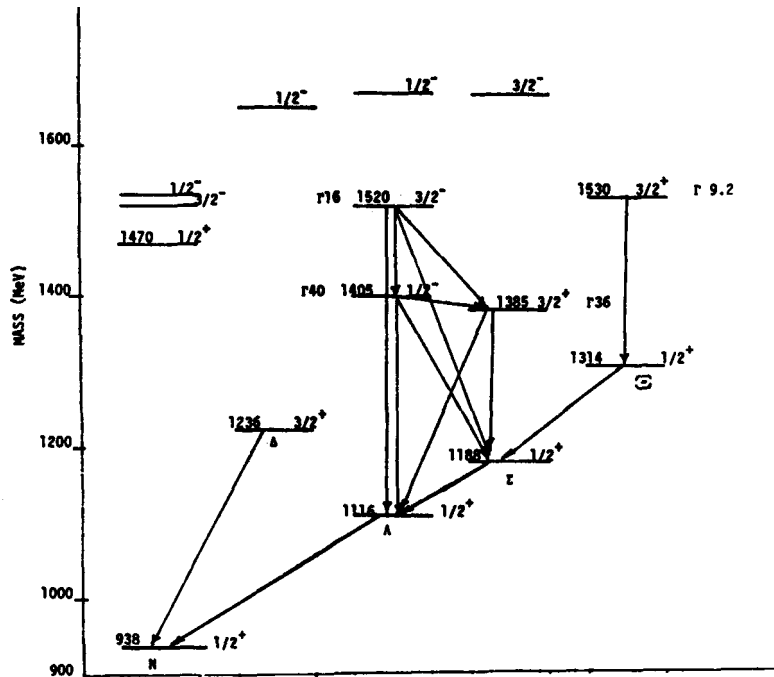


Fig 2) Possible Radiation Decays of Excited Hyperons

SUMMARY OF HYPERON RESONANCES AND DIBARYONS

There are a number of narrow hyperon resonances, and studies of their quark substructure particularly the admixtures of $q^4\bar{q}$ components are certainly important.¹⁶ Radiative decays of the resonances are sensitive to quark structure models. (Fig. 2) For example, the decay of $\Lambda(1405) \rightarrow \Sigma(\text{g.s.}) + \gamma$ is not allowed if the wave function has only a q^3 component. Only a few measurements have been made, and only two with a non-zero result.¹⁷ These experiments have low count rate and high background.

In the excitation of baryon number 2, $S = -1$ resonances, there are still open questions as whether any states are seen.¹⁸ The low lying states with $S = -1$ are below meson threshold although above two particle break-up thresholds. (Fig. 3) It is important to note that this is not the case for the $S = 0$ systems where the color-

magnetic force is repulsive instead of attractive. The lowest lying $S = -1$ dibaryon forms a spin doublet. However, only the triplet level which may lie near the ΣN threshold can be reached from a (K, π) reaction on the deuteron. Spin flip using the (γ, K) reaction should excite both states simultaneously. Thus, while the over-all mass scale cannot be calculated with accuracy, the spin splitting prediction which is based on the $N\Delta$ mass-splitting is much more reliable.¹⁹ This reaction thus offers a good test of quark bag models, where uncertainty remains in the quark baryon coupling.

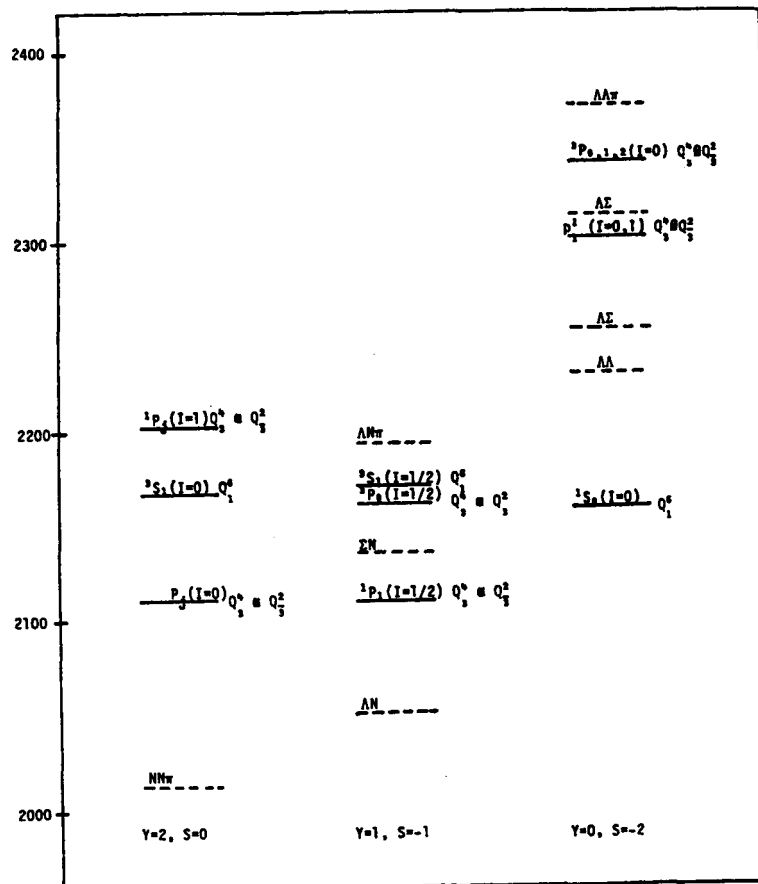


Fig 3) Predicted Dibaryon Spectrum

EXPERIMENTAL METHODS AND DETECTOR SYSTEMS

Previous work has developed a proposed layout for hypernuclear physics at CEBAF. The best geometry (i.e. highest resolution and count rate) uses virtual photons produced near 0° . This technique has been discussed by several people.^{2,3} The required spectrometer system has also been outlined.

The usable beam flux at zero degrees is limited by the number of electrons reaching the focal plane of the electron tagging spectrometer. This number is dominated by bremsstrahlung electrons in the target which increases as Z^2 . Thus, for a fixed radiation length target, the number of scattering centers decreases as a function of A . This is illustrated in Table 1. It is clear that the counting rates for the heavier hypernuclei fall dramatically. Since the "significant" physics occurs for heavier A , it is important to determine if the dependence on the beam flux (and thus the count rate) due to bremsstrahlung can be relaxed.

Table 1. Number of Nuclei/cm² in a Radiation Length

	$\chi(\text{gm/cm}^2)$	$\rho(\text{gm/cm}^3)$	# Nuclei $\chi(\times 10^{23})$
C	42.70	2.27	21.4
Si	21.82	2.33	4.7
Cu	12.86	8.96	0.81
SN	8.82	7.31	0.45
Pb	6.37	11.35	0.19

SUMMARY OF EXPERIMENTAL METHODS

The presently proposed layout should be optimized for heavier nuclei. In this case the count rate drops dramatically so one should investigate ways to use more of the available beam current. Of course one needs to develop particle ID for the focal plane detector. In addition one should further investigate the trigger rate for the experiment.

The backgrounds also need to be explored. Of interest is what hardronic flux is present near the target, and what flux will reach the rear of the spectrometers? It is clear that a good conceptual design is now available, but one needs detailed calculations in order to develop an actual proposal.

In a sense the experiments are background limited because the bremsstrahlung background limits the number of beam electrons on target. One needs a method of reducing this background.

RECOMMENDATIONS

Consideration of the above remarks leads to the following recommendations for a hypernuclear program at CEBAF:

1. Strange quark systems provide an exciting list of possibilities that may produce new physics.
2. Studies of the heavier hypernuclear systems with excellent resolution are a unique feature of CEBAF.
3. Two spectrometers with the following properties are necessary:

<u>TYPE</u>	<u>K⁺</u>	<u>e⁻</u>
Pmax	~ 1 GeV/c	0.3 GeV/c
$\Delta\Omega$	35 msr	1 msr
Angle range	10-25 deg	~ 0 deg
$\Delta P/P$	10 ⁻⁴	2 x 10 ⁻⁴
P max/P min	1.2	3.6
Length	10-15 m	-

4. The beam must be 100×10^{-9} amp at ~1.2 GeV/c with no halo. It should be capable of being dispersed on target. Long run times (500-1,000 hrs) must be expected.
5. Measurement of electromagnetic form factors in lighter systems over an amplitude range of ~ 1 to 10 are possible.
6. Measurements of (γK) elementary amplitude must be improved. Most of these studies are best done with equipment other than that proposed for hypernuclear spectroscopy.
7. Studies of the few body strange system should continue.
8. Space near the target will be needed to position detector systems to capture the break-up products.

Appendix
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October 1986

Report of the Working Group on x-Scaling

Third CEBAF Summer Workshop

June 23-27 1986

Newport News, Va. 23606

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1. General Remarks

The working group on x -scaling was a relatively late addition to the agenda of the third CEBAF Summer Study. While the charge to the other working groups (*e.g.* $(e, e'2N)$, hypernuclear physics, excited baryons) was well-circumscribed, our charge was somewhat more self-defining. We elected as a group to depart from a more restrictive interpretation ($Q^2 \geq 1(\text{GeV}/c)^2$ and $x > 0.1$, where hard, non-diffractive features of e - p scattering become manifest) to consider the more general topic of *the particle physics aspects of electronuclear physics at energies up to 6 GeV*. We emphasize that CEBAF, even without an upgrade beyond 6 GeV, is a particle physics machine. Even a decade ago forefront particle physics with electrons and photons was pursued at accelerators with similar energy (CEA, DESY, Cornell); significant lacunae still exist today even in the measurement of fundamental processes (extent of π , K form factors, spin structure functions $g_{1,2}$) not to mention the almost total lack of nuclear dependence ($A \geq 3$). It should be born in mind that 6 GeV is above $p\bar{p}$ threshold even for nucleon targets, and above J/ψ threshold for $A > 2$. It is our opinion that the physics program at CEBAF would be greatly enriched by giving a balanced emphasis to both nuclear physics from the baryon-meson standpoint, and under its particle physics aspects.

As a practical matter, the working group proceeded by enumerating ten experiments of high scientific interest, and of probable feasibility with CEBAF energies and luminosities.

2. Experiments

The experiments are ordered roughly in the following way. The first four represent those presently under active study. The middle three are those whose feasibility is clear, but which to our knowledge have not been examined in detail, regarding rates and specific experimental constraints. The last three are more speculative, but should be included within the program pending further experimental or theoretical developments.

2.1 SYSTEMATIC DETERMINATION OF THE QUARK DISTRIBUTION AND FRAGMENTATION FUNCTIONS $e + p, n \rightarrow e' + h + X$ at 6 GeV

An important question is how low in energy is the Quark-Parton Model (QPM) valid in exclusive deep-inelastic scattering. The normalized cross sections $(\frac{1}{\sigma_h} \frac{d\sigma_h}{dz})_{Q^2, x}$ (where $z \equiv E_h/\nu$, $h = \pi, K, \text{ etc.}$) are expressed as sums bilinear in the flavor distributions (weighted by $q_i^2 \frac{4}{9}u(x)$, $\frac{4}{9}\bar{u}(x)$, $\frac{1}{9}d(x)$ etc.), and their fragmentation functions *e.g.* $D_u^{\pi^+}(z)$, $D_{\bar{u}}^{\pi^+}(z)$, $D_d^{\pi^+}(z)$..., and within assumptions may be determined independently. This is an exercise of two-fold significance (and important both for nuclear physics both for PEP and CEBAF) as (i) the DESY/Cornell streamer chamber experiments of a decade ago were severely limited in particle identification, and (ii) upon making this investment to learn the $q, \bar{q}$ distribution functions from the nucleon, one may hope to measure the increase in the sea quark distribution at low- x in nuclei. This information is contained in the hadron ratios $(h^+/h^-)_x$, extrapolated to $z = 1$. A significant body of literature exists for charge retention in e, μ, ν scattering from nucleon targets which roughly validates the naive QPM picture that the leading hadron carries the struck quark¹⁻⁴. Figure 1 shows (a) the mean charge $\langle q \rangle$ as a function of x (for $x_F > 0$), and (b) in x_F for $x > 0.2$ ¹. Specifically in (a); the solid lines show the QPM prediction for two values of η (a universal constant shifting the observed charge distribution; $\eta = 0$ if $s\bar{s}$ pairs are produced with equal probability to $u\bar{u}$, $d\bar{d}$ in fragmentation, $\eta = -\frac{1}{6}$ if $s\bar{s}$ is totally suppressed). In (b) extrapolating to the fragmentation function endpoint (in the Feynman variable $x_F \equiv p_{||}^*/p_{max}^*$) shows that when a u quark is struck, a positive hadron results 80% of the time. (Note that at $\langle x \rangle = 0.28$, 87% of the total cross section is due to the interaction with u quarks and 11% with d quarks; for $x_F > 0.5$, 72% of the charged hadrons are positive.) Although at low- Q^2 the 'charge leakage' (overlapping of the target fragmentation into the jet fragmentation regions) is significant, the relative cleanliness of the signal in the DECO work gives hope that nuclear dependences may be discovered (see Ref. 5 for a hint of the increase in the $q\bar{q}$ sea with atomic weight.)

A full discussion of this topic is to be found in the excellent review of Janata in these Proceedings⁶. A study of particle fractions, multiplicities, fragmentation functions, *etc.* with the Lund Monte Carlo hadronization code, relevant to CEBAF energies will be found in the 1986 RPAC report⁷.

2.2 MEASUREMENT OF THE PIONIC SEA IN ATOMIC NUCLEI

BY $(e, e'\pi)$ IN QUASIELASTIC KINEMATICS

The importance of the pionic sea in understanding the EMC effect both in the high- and low- x regions was first underlined by Llewellyn Smith⁸, and Ericson and Thomas⁹. That the pion distribution function could be extracted for the nucleon by isolating the longitudinal component from the four structure functions contributing to exclusive $N(e, e'\pi)$ was demonstrated by Güttner¹⁰ *et al*, who urged a systematic study in A to exhibit the predicted increase in the pionic content of nuclei with A . We will not expand on this discussion as it has been done in more detail in a contribution to the 1986 RPAC¹¹, and elsewhere in these Proceedings¹².

2.3 MEASUREMENT OF THE SPIN STRUCTURE FUNCTIONS OF THE NUCLEON

One of the most important tests of our understanding of QCD in the nucleon is the measurement of the spin structure functions $G_{1,2}(Q^2, \nu)$ which are independent of the two spin-averaged structure functions $W_{1,2}(Q^2, \nu)$. This is also an area where CEBAF can make important contributions in the transition to scaling regime $Q^2 = 1 - 3 \text{ (GeV/c)}^2$. Measurements with polarized hydrogen have been made both in SLAC experiments E-80, E-130 as well as a more recent experiment by the European Muon Collaboration. The data so far are in agreement with the Ellis-Jaffe sum rule (on $2\int_0^1 g_1^p dx$, where $M^2\nu G_1(Q^2, \nu) \rightarrow g_1(x)$ in the scaling limit) and the Bjorken sum rule (which connects $\int_0^1 (g_1^p - g_1^n) dx$ to the ratio of the axial to vector weak coupling constants $|\frac{g_a}{g_v}|$), assuming the neutron asymmetry $A_1^n = (d\sigma^{\uparrow\uparrow} - d\sigma^{\uparrow\downarrow})/(d\sigma^{\uparrow\uparrow} + d\sigma^{\uparrow\downarrow})$ is small¹³. CEBAF is in an excellent position to both improve the statistics of the proton measurements, the relative experimental error on the test of the Ellis-Jaffe sum rule

being $\approx 30\%$, and furthermore measure for the first time the neutron asymmetry. Since the time of the SLAC-Yale experiments significant advancements have been made in the development of frozen spin targets NH_3 , ND_3 . Furthermore a working gaseous $^3\vec{H}e$ target may soon be practical¹⁴. This latter target is especially interesting in that unlike the case of $\vec{d}$, no proton subtraction needs to be done if the d-state admixtures in the ^3He wavefunction may be ignored.

A feasibility study for this program was carried out during the 1986 Summer Study. Details and time estimates for a 3% measurement of A_1^p and a 7% measurement of A_1^n at $Q^2 = 1$ and 3 $(GeV/c)^2$ are to be found in the report of Geesaman in the 1986 CEBAF RPAC report¹⁵.

2.4 ELECTROPRODUCTION OF VECTOR MESONS

$(e, e'\rho), (e, e'\phi)$ ON NUCLEON AND NUCLEAR TARGETS

Electroproduction allows the study of the transition from hard non-diffractive scattering (large x and Q^2) to the vector meson dominance (VMD) regime ($Q^2 \rightarrow 0$). In the VMD picture, the photon (real or virtual) may appear as a hadron with the same quantum numbers. The formation length of a particular vector meson $b_v = c\Delta\tau \approx \hbar c 2\nu [Q^2 + m_v^2]^{-1}$ where m_v is the mass of the vector meson and $\nu = E - E'$. This length may be made considerably shorter (high Q^2) or longer ($Q^2 \rightarrow 0$) than both the average internucleon separation ($\approx 2 fm$) or the mean free path of a meson in a finite nucleus (typically $1.5 - 2 fm$). For a maximum energy of 6 GeV, the ρ formation length will be $\geq 3 fm$.

It is thus of interest to perform these experiments on nuclear targets in addition to hydrogen as the interposition of nuclear matter in principle allows the determination of the forward vector meson scattering amplitude T_{vv} and cross section σ_v . As the decay angular distribution ($\rho \rightarrow \pi^+\pi^-$, $\phi \rightarrow K^+K^-$) contains the L-T information, those quantities may be determined separately for longitudinally and transversely polarized mesons. Several aspects of ρ production from hydrogen have never been understood adequately due to the relatively

poor quality of the data (associated with the duty factor, spectrometer solid angle, etc.); a thorough discussion of the limitations of the $\gamma, \gamma^* + p \rightarrow \rho + p$ data and call for future work is given both by the reviews of Bauer *et al*¹⁶, and Cassel¹⁷. No experiments have been performed with polarized targets.

The situation with ϕ electroproduction will considerably cleaner than that for the ρ meson (the ρ is very broad whereas the ϕ is narrow; ρ decay suffers interference with the $\pi^+\pi^-$ non-resonant background whereas the comparable problem for ϕ is not significant). However the cross section will be reduced by $\approx \frac{1}{30}$ relative to that of the ρ .

The LAS detector with the CEBAF duty factor will be ideal for making measurements of greatly improved quality and extent. This program will nicely complement the real photon studies (γ, ρ) *etc.* being planned.

2.5 MEASUREMENT OF $R = \sigma_L/\sigma_T$ IN DEEP INELASTIC SCATTERING FROM NUCLEON AND NUCLEAR TARGETS

In the limit $Q^2 \rightarrow \infty$, $R = 0$ due to the quark's possessing $\frac{1}{2}\hbar$ angular momentum; this is born out (within errors) for high energy muon and neutrino scattering. R at SLAC has been shown to be small but non-zero (≈ 0.2), possibly reflecting the influence of primordial transverse momentum, and scattering from di-quarks and pions, *i.e.* 'higher twist' effects. The possible A dependence of R has significant implications for the validity of constituent nucleon models of nuclei. This issue has been discussed by Smith, who within a convolution calculation (the same framework as the calculations which reproduce the magnitude of the high- x EMC effect by binding energies¹⁸) finds essentially *no* nuclear dependence to R (see Figure 2)¹⁹. The recent SLAC experiment E-140 which was proposed as a precision measurement of R in H, D, and several nuclear targets using the 8-GeV spectrometer, was in fact performed between October, 1985 and January, 1986²⁰. However due to difficulties with the machine both the scope and statistics of the run were significantly compromised, and due to mod-

ifications to the linac for SLC operation another such opportunity at SLAC is unlikely if not ruled out completely. Some preliminary data are shown in Figure 3.

The R ratio is an example of how the high duty factor of CEBAF may possibly be of an advantage even in singles experiments. Accurate $L - T$ separation requires measurements at backward angles where the e/π ratio can be extremely low. Electrons could be detected with good efficiency with high average current and in a large-aperture spectrometer at backward angles, which might be precluded with a pulsed linac beam due to pileup.

2.6 π^0 ELECTROPRODUCTION ABOVE THE RESONANCE REGION

As pointed out by Cassel in the 1985 CEBAF workshop, an outstanding problem to date is the comparison of π^0 -photoproduction, and -electroproduction at low- Q^2 . The two should be comparable as $e + p \rightarrow e' + \pi^0 + p$ is dominated by $d\sigma_{\perp}/dt$ (born out by the π^0 angular distribution being a maximum at $\phi = \frac{\pi}{2}$, *i.e.* perpendicular to the scattering plane). The π^0 photoproduction however shows a pronounced minimum at $t \approx 0.5 (GeV/c)^2$ whereas even at the lowest $Q^2 (\leq 0.25 (GeV/c)^2)$ measured in two experiments^{21,22} no dip in the angular distribution was observed, except at $t \rightarrow 0$ (see Figure 4). (At the time, this comparison had been advertized as a test discriminating between dips arising from an impact parameter representation of the scattering amplitude, and dips from the exchanged objects in the t -channel²³.) No further experiments have been forthcoming; the critical measurement will involve measurement of the e' at very low angles, and the proposed high quality π^0 spectrometer.

2.7 DEEP INELASTIC SCATTERING FOR $x > 1$ WITH NUCLEAR TARGETS

It has long been acknowledged that measurements above the quasielastic peak (low energy loss tail) could sensitively isolate details of the nuclear wave function or more exotic features of the nucleon in the nuclear medium, with less ambiguity than for $x < 1$ ²⁴. Figure 5 shows the ratio of νW_2 per nucleon

of a heavy (A) to a light ($B > 4$) nucleus for three calculations: a relativistic wavefunction with $N - N$ correlations (dashed line) where the correlations are responsible for the shoulder at $x \approx 1$ ²⁵; the quark cluster model (solid line) where plateaus are seen between successive integer x due to the differing relative probabilities for finding $6q, 9q \dots$ clusters between nuclei of very different atomic weight²⁶; and a color dielectric model (dot-dashed line) where valence quarks are found to have some long range freedom dependent on Q^2 (see Ref. 27). The working group recognized that while these are low counting rate experiments, neither statistics nor resolution should be compromised by increased solid angle, as in practice such modulations may in fact be very subtle.

2.8 ϕ -ELECTROPRODUCTION IN NUCLEONS AND NUCLEI AS A MEASURE OF THE GLUON STRUCTURE FUNCTION $G(\eta)$

The relevance of leptonproduction of heavy quarks (c, b) in understanding the gluon distribution function in the nucleon was developed in the years following the discovery of J/ψ and Υ and their spectra^{28,29}. The connection is made through the photon-gluon fusion (PGF) diagram (see Fig. 6) where a virtual photon from the scattering muon or electron interacts with a gluon via a *e.g.* $c\bar{c}$ quark line. (Note that the $q\bar{q}$ pair is not a singlet state; and addition gluon is necessary to permit it to be an asymptotic final state.) The topic has attracted considerable interest insofar as one published report claims a 45% increase in the gluon content of ^{56}Fe over than of D_2 in the low- x range, $0.026 < x < 0.085$ (see Ref. 30).

The requirement of production of heavy quarks to uniquely isolate the gluon distribution function results from (i) the interaction being sufficiently pointlike [$c\Delta\tau \approx \hbar c 2\nu/(Q^2 + 4m^2)$] where m is the quark mass and $\nu_{min} = 2m$, and (ii) that *e.g.* ψ production by other mechanisms be sufficiently suppressed. Note that regarding (ii), $q\bar{q}$ production by tunnelling in the string model of hadronization gives $P(q\bar{q}) \propto \exp(-\pi m^2/\kappa)$ where $\kappa \approx 0.2 \text{ GeV}^2$ is the string constant, leading to the ratios $u : d : s : c \approx 1 : 1 : \frac{1}{3} : 10^{-11}$. The question

arises as to whether for CEBAF, where charm production is essentially ruled out at 6 GeV , ϕ production would be reliable to systematically study $G(\eta)$ as a function of A . Whether or not the strange quark is ‘sufficiently heavy’ ($m_s \approx 0.5$, $m_c \approx 1.5$ GeV) is a theoretical issue that needs to be addressed.

2.9 BOSE-EINSTEIN CORRELATIONS IN $\pi^\pm\pi^\pm$ PRODUCTION AND POSSIBLE MEDIUM DEPENDENCE OF HADRONIZATION

Bose-Einstein correlations (Hanbury-Brown/Twiss effect) in like-sign pion production have been measured for virtually every interaction in nuclear and high-energy physics: pp , $p\bar{p}$, pA , AA , e^+e^- , and most recently μp ³¹. The correlation function $R \equiv \rho^{BE}(p_1, p_2)/\rho^{UNCORR}(p_1, p_2)$ at small relative four-momentum differences $(p_{\pi_1} - p_{\pi_2})^2$ reflects the space-time evolution of the emitting source (a superficial discussion, and references are found in Ref. 12). Simple considerations based on the string model of hadronization lead to distance scales over which the flux tube materializes into asymptotic hadrons of several fm in the quark-diquark frame, or ≈ 10 fm in the lab frame, for typical CEBAF and PEP values of Q^2 . These values are well-matched to the size of the nucleus, to be contrasted with hundreds of fermis for the largest Q^2 attained currently in muoproduction. It is thus attractive to look for the effects of the nuclear medium on the string tension κ or flux tube radius.

Outside of relativistic heavy ion collisions which measure source radii of typical nuclear dimensions ($3 - 7 fm$ depending on impact parameter selection), high energy experiments find rather small values on the order of $1 fm$ (*e.g.* $0.7 fm$ in $e^+ + e^- \rightarrow q\bar{q}$ jets at $\sqrt{s} = 29$ GeV ; $0.5 - 0.8$ fm in μp interactions at $E_\mu = 280$ GeV , $1 - 1.5$ fm in pp , $p\bar{p}$, pXe). Little or no discussion has been given so far to the interpretation of HBT for hadrons resulting from string-breaking. We would like to draw some theoretical attention to the problem, and recommend that this program only go ahead when it is better understood what the small values measured in the high energy experiments indicate.

2.10 MEASUREMENT OF STRUCTURE FUNCTIONS F_2^A IN NUCLEI IN COINCIDENCE WITH TARGET SPECTATOR SPECTROSCOPY

That the ratio of structure functions measured within a nuclear target to that of a nucleon, $F_2^A(x)/F_2^N(x)$ was discovered to be non-trivial (the 'EMC effect') encouraged the hope that details of the nucleons' structure and binding within the nucleus might be revealed. While numerous calculations reproduced the shape and magnitude of the effect based on specific dynamics (pions resulting from the binding of the nucleons, QCD rescaling, etc.) a subsequent calculation showed that the high- x depletion could be reproduced if correct shell energies were taken into consideration in the kinematics of a nucleon constituent model¹⁸. No reference is made to what field takes up the momentum from the nucleons which necessarily no longer satisfy the energy-momentum sum rule.

We contend that the situation is principally an experimental one. The role of binding energy could be validated (or invalidated) and isolated from dynamical effects if structure functions in nuclei were measured in coincidence with shell-specific 'tags'. Put differently, what one would like to do is to measure the inelastically scattered electron (which determines F_2), in coincidence with conventional charged particle, n or γ spectroscopy from the target nucleus spectator. This is essentially still an *inclusive* measurement as the observation of the target de-excitation determines the *initial* state of the nucleon before the interaction. Fig. 7 depicts the experiment schematically.

How realistic this technique is depends on two preliminary (experimental) points. The first is the assumption that the struck nucleon and resulting fast hadrons escape the nucleus on the average with minimal final state interaction which would further excite the target spectator. Second, it is important to know systematically how hole states in nuclei decay; of particular interest are the deep-lying holes $(s_{1/2})^{-1}$, $(p_{3/2})^{-1}$, etc.. This calls for a review of and perhaps further work in (p,2p) reactions. The scheme may be most apt for light nuclei (note that the EMC effect is visible even for ${}^4\text{He}$). From an experimental viewpoint we are optimistic that with high current high quality cw beams, conventional Si

and Ge detectors could operate satisfactorily in the proximity of a thin target.

3. Experimental Considerations

The scientific program associated with the x -scaling region cuts across spectrometer hall boundaries. Some experiments ($x > 1$; R-ratio measurement) are singles experiments suited to medium-resolution, medium-solid angle spectrometers; others $[(e, e'\pi^+); (e, e'\pi^0)]$ are best served by two spectrometers, or a spectrometer and a special external detector; and others (ρ, ϕ electroproduction) with LAS.

Our working group made only two specific recommendations concerning the spectrometers and halls. The first is that *adequate open floor area be designed in for presently unspecified specialized experiments that will obviously arise at some point.*

Second, while it is true that this program will utilize all spectrometer halls, it is clear that the Large Acceptance Spectrometer (LAS) is within the special province of this program. *We are very concerned that immediate attention be paid to the problem of instrumenting the forward angles ($3^\circ - 20^\circ$) with magnetic analysis and tracking comparable to or better than that of the wider angles.* How this is to be implemented is a question posing difficult technical choices, but two options are (i) instrumenting a large volume dipole magnet ($2m \times 2m \times 2m$) downstream of LAS, such as the old SLAC streamer chamber magnet, and (ii) significantly extending the LAS sectors themselves in the forward direction, and continuing the sectors in to a small a radius as possible.

4. The Working Group

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Figures

1. (a) The mean charge of hadrons with $x_F > 0$, as a function of x_{Bj} . The QPM prediction is shown for two values of η which lie in between the limiting cases of $\eta = 0$ ($s\bar{s}$ equally probable with $u\bar{u}, d\bar{d}$), and $\eta = -\frac{1}{6}$ ($s\bar{s}$ totally suppressed in the hadronization). (b) The net forward charge for events with $x_{Bj} > 0.2$, as a function of x_F (from Ref. 1).
2. $R_{Fe} = \sigma_L/\sigma_T$ calculated by Smith within a convolution model for (a) various Q^2 as a function of x ; and (b) various x as a function of Q^2 . 'Case 1' (curves not otherwise indicated) and 'Case 2' refer to different prescriptions of taking F_1, F_2 off-shell; the essential point is that no large nuclear difference is predicted from the nucleon case (assumed to be the flat curve $R_N = 0.2$)¹⁹.
3. Preliminary data from E-140. Top: at $x = 0.5$; bottom: at $x = 0.35$. The tic marks on the error bars indicate the expected statistical and total errors when the analysis is complete. 'QCD' = perturbative QCD calculation; 'TM' = target mass correction; 'PT' = primordial transverse momentum.
4. The distribution in $-t$ for π^0 -photoproduction (triangles and squares), and electroproduction (solid dots, Ref. 21). Note that the dip at $-t \approx 0.5 (GeV/c)^2$ in photoproduction has vanished even for $Q^2 \approx 0.2 (GeV/c)^2$ in electroproduction.
5. The ratio of νW_2 between a heavy nucleus (A), and a lighter nucleus ($B > 4$), in the region $x > 1$, calculated within the context of three models. The dashed line is based on a relativistic wavefunction with correlations²⁵; the solid line, the quark cluster model²⁶; and the dot-dash line, a color-dielectric model²⁷.
6. The Photon-Gluon Fusion diagram for the case of muon (or electron) scattering, with charmed (or strange) quark-antiquark pair production.
7. Schematic representation of target spectator spectroscopy after deep inelastic scattering. The baryon and mesons are not observed in coinci-

dence; at sufficiently high momentum transfers the final state with the target nucleon should be minimized.

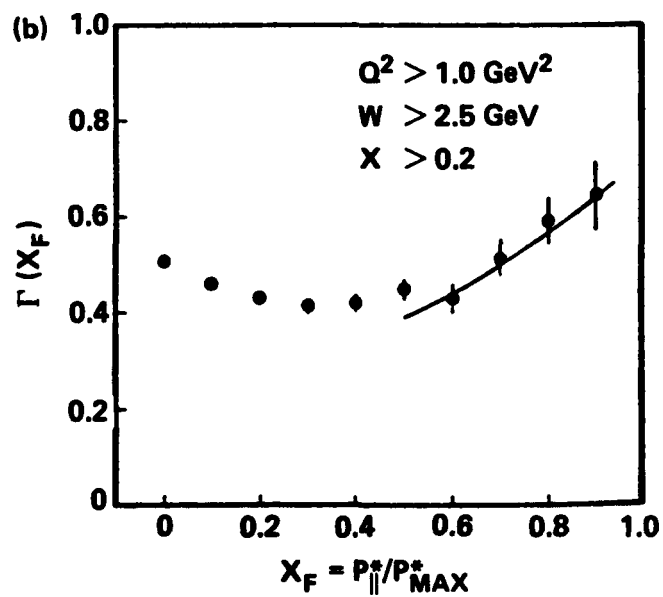
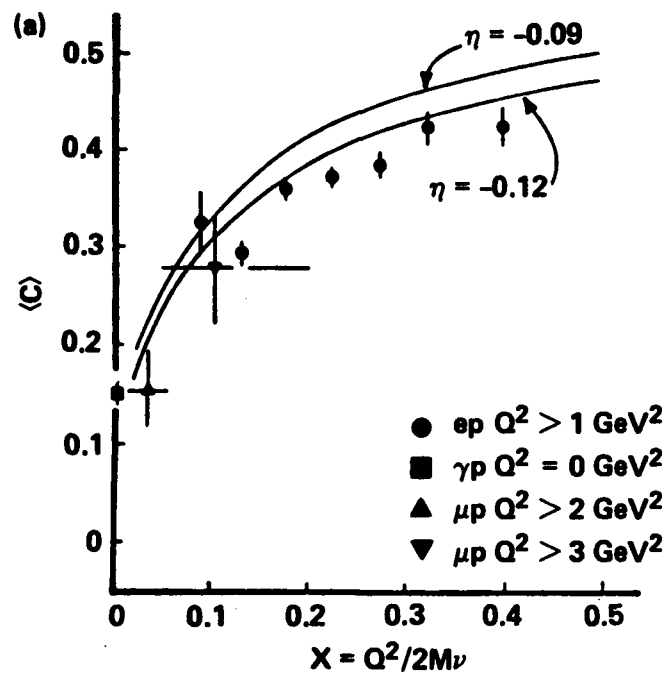


Figure 1

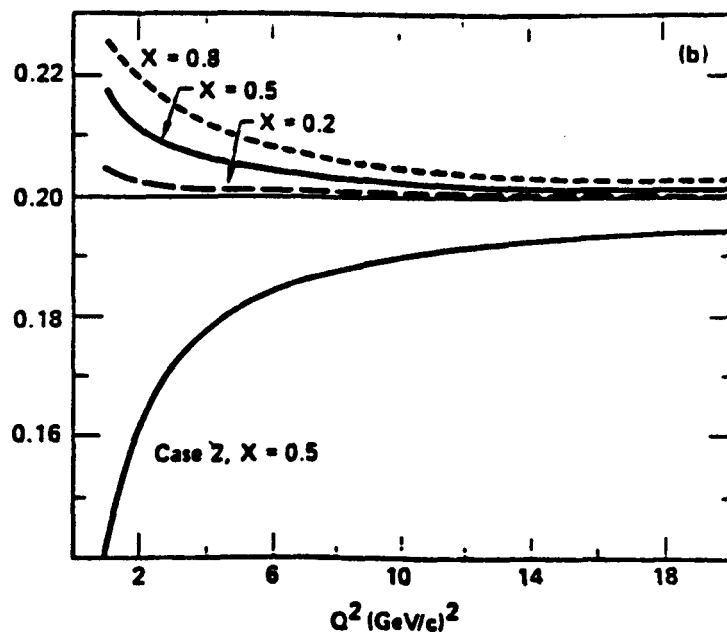
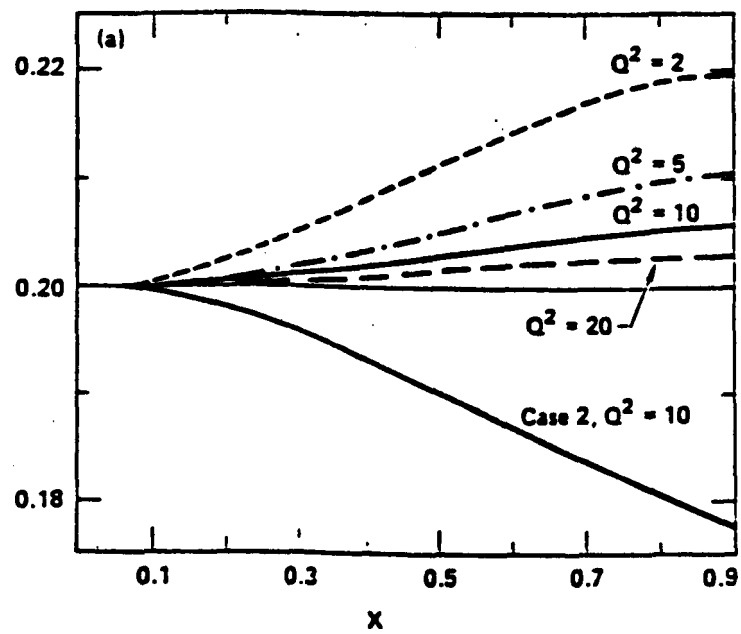


Figure 2

E140 PRELIMINARY RESULTS

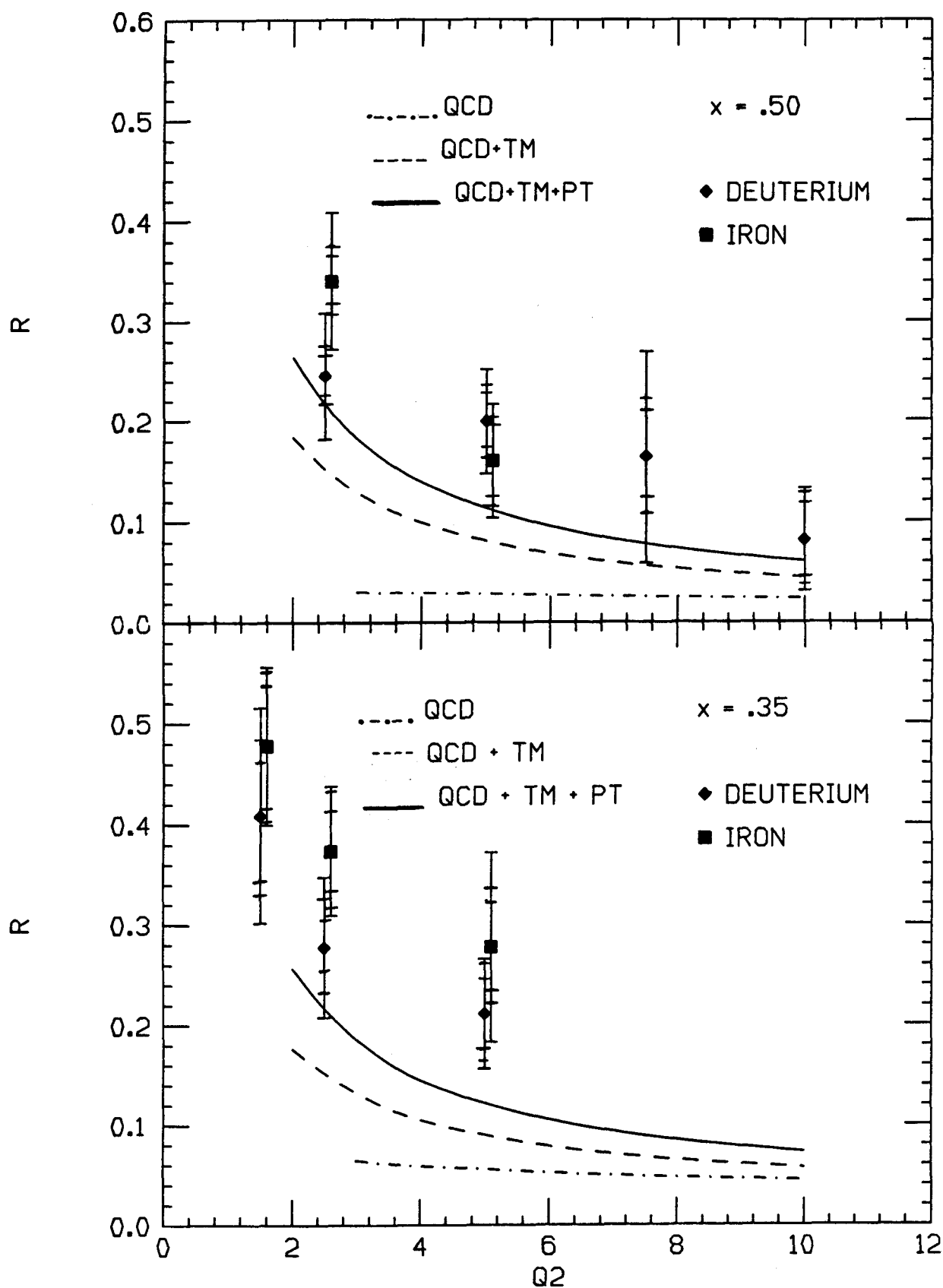


Figure 3

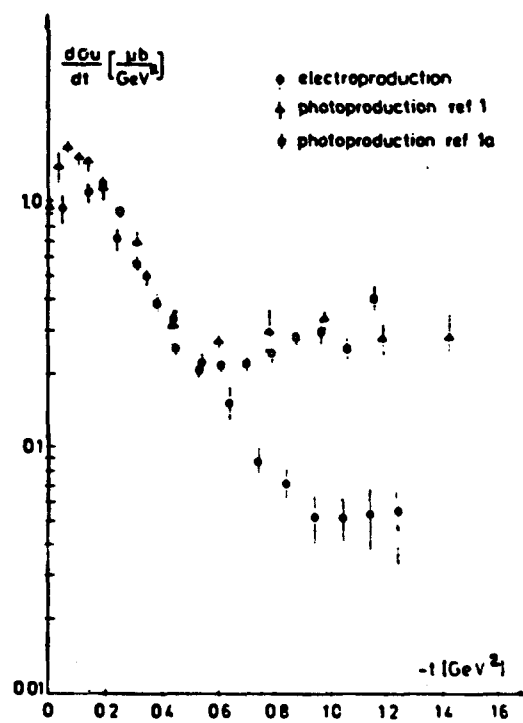


Figure 4

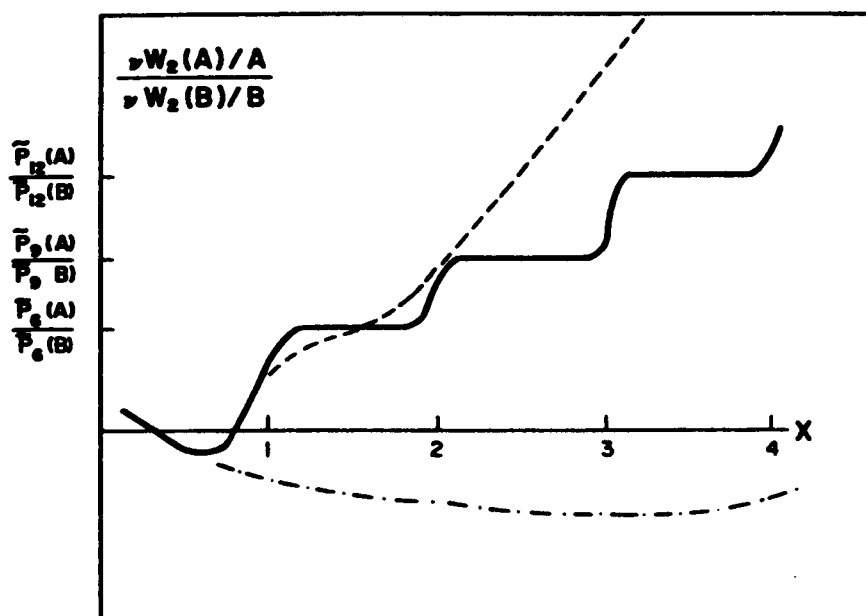


Figure 5

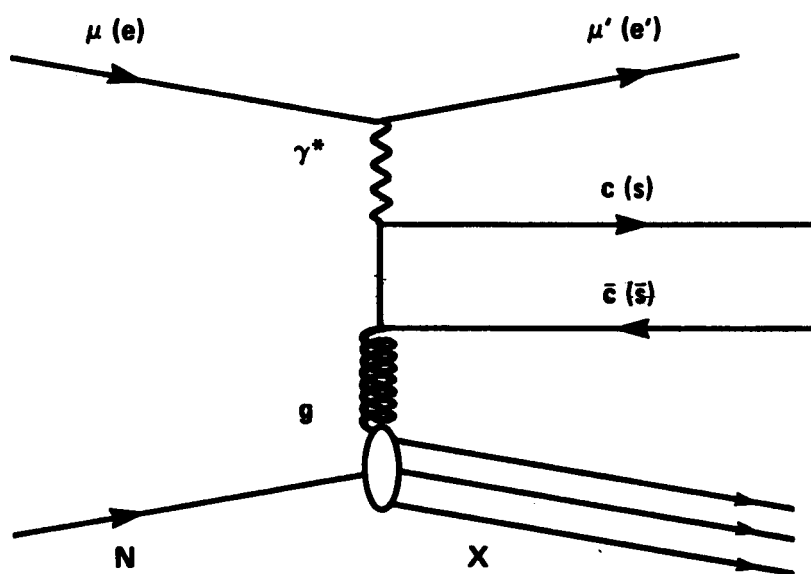


Figure 6

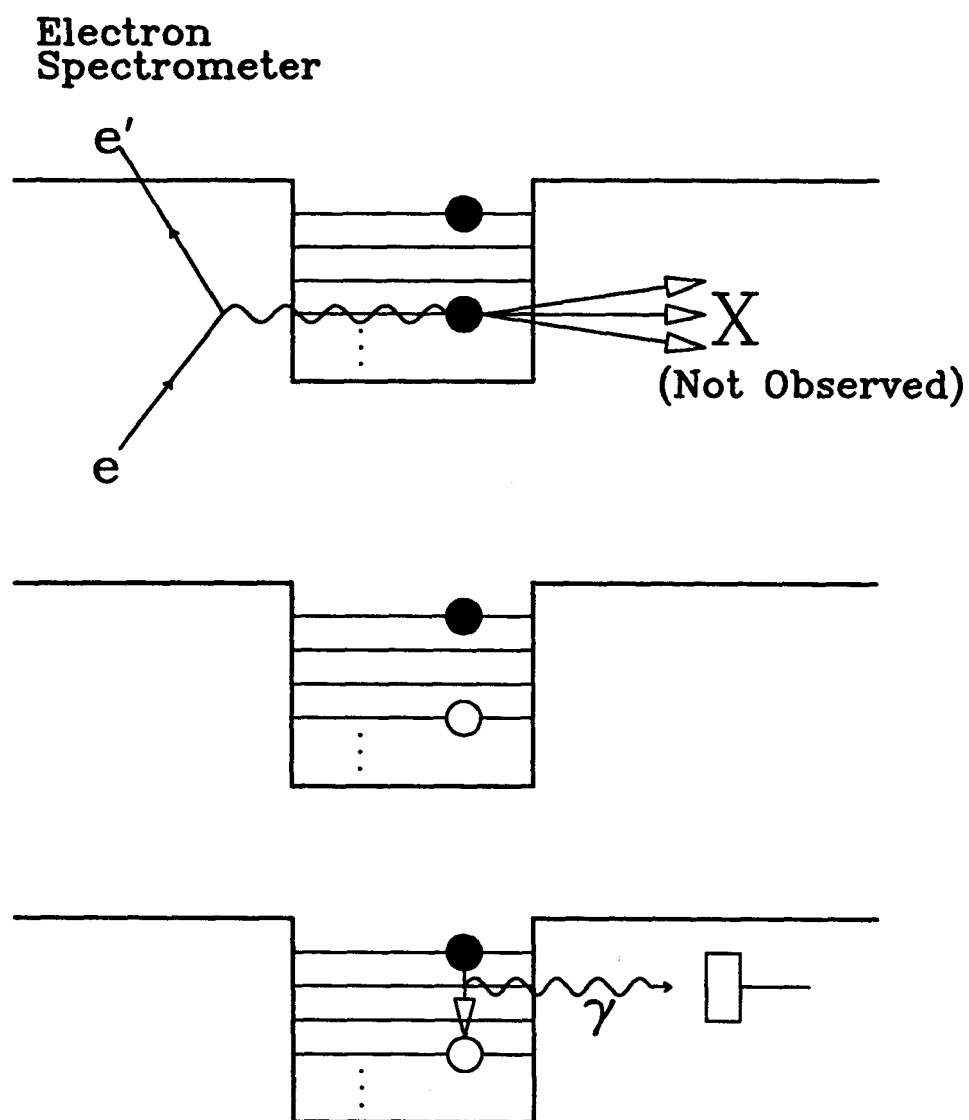


Figure 7

SUMMARY TALK

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It is a pleasure for me to give the last speech in what has been a most stimulating week of talks and discussions. Given the fact that the physics presented is still fresh in everybody's mind (and described on the previous pages of these proceedings), I see little point in doing literally what the title "summary" implies. Rather, I want to share with you some ideas more or less directly related to the many things presented at this workshop.

The external circumstances of this workshop certainly are most positive: it is very likely that construction of CEBAF will be funded in FY 87. The wording of the report of the appropriations committee of Congress is most favourable for both CEBAF and Nuclear Physics as a whole. The signals from the Senate in regards to CEBAF are also most promising. If anything like rational procedure prevails (and if a continuing-resolution situation can be avoided), construction of CEBAF will start in fall of 86.

These good news are the result of many years of work to exhibit the exciting scientific questions that can be studied by Nuclear Physics in general and an electron accelerator in particular, the efforts to produce a coherent long-range plan, and to get the community to support it as a whole. These good prospects for the future also are due to all the people that have advanced the CEBAF project to a state of avantgarde technology that is unanimously accepted by the community. The people who have achieved this are to

be congratulated!

There is a second favourable sign accompanying this work shop: the record participation. (It must have been an unpleasant surprise to those dealing with organizational details, though). This is a proof that the community is realizing that CEBAF is the major new facility for the next few years, and that it realizes that a 4 GeV/CW electron accelerator allows to address a large fraction of the most fascinating question of Nuclear Physics of the Nineties.

CEBAF is a machine that will play a vital role in what we can call a major

CHANGE OF NUCLEAR PHYSICS

Nuclear physics of the past has dealt mainly with nuclei in terms of nucleonic constituents and interactions. Many concepts have been developed to deal with the nucleus as a quantum-mechanical many-body system of strongly interacting constituents. Many models have been developed to describe the complex phenomena occurring in nature. For light nuclei ($A \leq 4$) and nuclear matter, a rather quantitative understanding has been reached, and the theoretical descriptions achieve satisfactory predictive power. Many phenomena, though, are not yet properly understood, and nuclei in terms of nucleons continue to present a challenge to our understanding.

Nuclear physics of the future increasingly recognizes that nucleons have an internal structure, and that nuclei must be understood in terms of their elementary constituents, quarks and gluons. Today, we have at our disposal a theory of the strong interaction, QCD, which holds the promise of a fundamental understanding of nuclei. This theory is simple in the limit of asymptotic energies, but its consequences for low energies (where experiments can be done) are not understood. Its most intriguing property, confinement, and a quantitative derivation of the inter-nucleonic force, remain to be investi-

gated. With QCD we are today at the point where solid state physics (if such existed already) was 100 years ago: While the basic interaction, the Coulomb law, was known, the structure of macroscopic matter was not calculable.

Is QCD the correct theory for the confinement region at low energies? What are its quantitative consequences, and can we observe the new phenomena predicted? These questions can be studied by using the nucleus as a laboratory where partial deconfinement can be achieved.

When trying to address the question of balance between these two focal points, we clearly see a tendency of both experimentalist and theoreticians to move towards a study of quark degrees of freedom of nuclei (this is a purely "experimental" observation). One might consider this to be simply a question of fashion, and an expression of an established tendency to not solve the old problems but to move on to new ones. But I am convinced that it is deeper than that. It is explained by the desire to understand nuclei in terms of the fundamental underlying degrees of freedom. The physics program of CEBAF will thus stand on two legs: The interesting phenomena in nucleonic physics that remain to be studied in better ways using the electro-magnetic probe, and the physics of quarks and mesons that will represent the main theme in the long run.

Proceeding to more concrete considerations, let me make some comments on the

PHYSICS DISCUSSED AT WORKSHOP

In the course of the various talks presented at this workshop, a number of new ideas concerning both theory and experiment have been discussed. They reflect the vitality of the field, and document the shift in interest from nucleon- to quark-dominated aspects of nuclei. In terms of the physics, the workshop clearly has been a success.

If I want to add some critique, I would say that, for

my taste, the talks not always have payed enough attention to the particular characteristics of CEBAF. There was perhaps too much attention payed to the pet ideas of the speakers (a general problem for all of us, including myself) by discussing experiments on inclusive electron scattering, or experiments that would be more appropriate for Bates- or SuperCEBAF-type machines.

The reason for this is obvious: the CEBAF scientific justification and the key experiments are well known. During the past years they have been repeated many times in proposals, committee reports, etc. The 3-star experiments have been elaborated in detail. Today, we simply don't want to discuss them over again - we want to go ahead and do them!

This is fine, but it leaves a somewhat incorrect impression with those participants who are new to the field and not familiar with the past. They could get the idea that much of the physics discussed is at the limit of what is doable, and not entirely thought through. Those who might have obtained an impression of this type should read last years RPAC (Research Program at CEBAF), which does describe the core experiments that will be done during the initial phase of the CEBAF experimental program. This RPAC illustrates why a 4GeV/CW machine is indeed the most exciting nuclear physics facility to be built for the nineties.

Just to remind everybody I give below a short summary of the

BASE PROGRAM FOR CEBAF

It concerns the study of nuclei in the transition region from nucleonic to mesonic and quark degrees of freedom of nuclei, and their understanding in terms of QCD. To achieve this, a spatial resolution of $\Delta x = 1 \text{ fm} \div 0.1 \text{ fm}$ is needed, and this can be achieved with momentum transfers $q = 1.5/\Delta x$. At the same time energy transfers of up to 2 GeV are needed,

such as to go from the nucleon ground state through the resonance up into the structureless region of the response function.

The experiments that form the base program for CEBAF can be put into five categories:

- Single particle densities, measured via $(e,e'N)$. With this type of experiment, we can access the wave function of nucleons in nuclei in the cleanest possible way, we can study deeply bound shells, and large momenta.
- Two-nucleon correlations, measured via $(e,e'2N)$. This type of experiment will allow us to study properties of nucleonic wave functions that go beyond a mean-field description of nuclei, will allow us to look at short-range nucleonic wave functions, and at the shortest range it will tell us about the eventual dissolution of nucleons into objects more appropriately described by n-quark bags.
- Baryon resonances in nuclei, investigated via $(e,e'XY)$ experiments. These experiments will study the properties of the excited nucleon via longitudinal/transverse separations and via a tagging of the states excited by observation of their decay products, they will study the change of these properties due to partial deconfinement in the nuclear medium, and allow the most demanding test of QCD into the strong-interaction limit.
- Hypernuclear spectroscopy, studied via $(e,e'K)$. These experiments will allow us to use the strange quark to probe nuclei, study deeply bound states with the minimum of distortion.
- Form factors of $A \leq 3$, measured by $\vec{A}(\vec{e},e)$. By doing experiments involving both polarized electrons and targets (recoil nuclei), we can separate the different form factors, measure fundamental properties of nucleons, study the S-state interaction between nucleons at short range, and improve our understanding of the S-D transition which is important if we want to finally measure D-state wave functions, calculate meson exchange currents, etc.

These "core" experiments are described in detail in RPAC, and the specific design issues of the hardware needed to do them are discussed in the preceeding summaries of the working group leaders.

Has recent theoretical understanding led to a change of this set of initial experiments? This is an interesting question to discuss since it takes about a decade to progress from concept to realisation of a major accelerator. On the one hand side one could hope that our field is vital enough to produce drastic changes of the physics picture that would change the experimental needs. On the other side, this would complicate very much the design of a facility that is expected to have decades of lifetime.

This question of a change of physics outlook is not entirely an abstract one: in the history of CEBAF there was at one point an (irrelevant) Phys. Rev. Lett. article the new insights of which were misused to undercut the project (before the erratum appeared).

This workshop (or any other source) has not pointed out any development that would suggest that the base-program listed above needs to be changed. The arguments on which this physics rests is based on ideas general enough to remain valid even if specific theories were to undergo drastic changes.

The most relevant progress of particle physics for CEBAF presumably is not to be sought in the development from GUT's to supersymmetries to superstrings in 10 dimensions and beyond. It is easy to predict that experiments at CEBAF will be done in 4 dimensions in agreement with experimental observations of 3.00... space dimensions. The progress of QCD, on the other hand, is of most direct interest to experiments to be done at CEBAF.

From the talks presented at this workshop, we may conclude that the brute-force lattice QCD calculations are not yet close to give quantitative answers. QCD inspired models,

however, are giving predictions for nucleonic properties, and they do begin to deal with the quark structure of nuclei, and the partial deconfinement of quarks. Some calculations already do produce quantitative, and verifiable, results.

Extrapolating over the coming 5 years, we may anticipate that, by the time CEBAF experiments produce first results, QCD based calculations will be ready for a quantitative comparison to data. Given the rich structure of phenomena expected from QCD, we can look forward to a fascinating spectrum of physics to be done at CEBAF. One theoretical development discussed repeatedly at this workshop concerns calculations that predict a

CHANGE OF BOUND-NUCLEON SIZE

(Something I want to comment upon since it touches a pet-topic of my own).

The main experimental evidence cited for justification of a change of N size due to the nuclear medium comes from the lack of strength in the longitudinal quasielastic response function $S_L(\omega)$. The apparent lack of $\sim 30\%$ in $\int S_L(\omega) d\omega$ relative to the sum-rule value Z has led to the suggestion that bound- and free-nucleon size differ by about 20%. This interpretation of experimental data looks naive to me, and ignores well-known complications and other experimental evidence.

The complications result from the fact that the data used concern momentum transfers q of order twice the Fermi momentum, or less. At these transfers Pauli-blocking effects are not negligible, and the final state interaction of recoil nucleons would need to be taken into account. Ignoring them could easily account for 20%. Another complication results from the shape of the experimental $S(\omega)$ at large ω , where S goes to zero in an unphysically fast way; this could

indicate experimental problems (radiative corrections?) that could account for the remaining 10%.

Better evidence on bound-N size can be obtained best from the response function of nuclei as measured at very large momentum transfer, $q = 3-10 \text{ fm}^{-1}$. If such data (which we measured recently for nuclei $A = 4 \dots 197$ at SLAC) is interpreted in terms of γ -scaling, one finds that the radius of the bound nucleon may not differ from the one of the free nucleon by more than a few percent.

Two additional pieces of evidence of the bound-nucleon size have been overlooked:

- Measurements of the M7 form factor of ^{51}V can be interpreted in a rather clean fashion in terms of the $f_{7/2}$ proton radial wave function, and the one-body density matrix element $\psi_{7/2, 7/2}$. Elastic magnetic scattering yields $\psi_{\text{exp}} = 0.93 \pm 0.05$. We know that ψ must be < 1 , and if we assume short-range NN correlations to be present we could expect $\psi \leq 0.9$. The experimental value is compatible with this expectation. It would no longer be compatible if we were to assume a size increase of the bound nucleon of 20%, which would increase ψ_{exp} by 30%. Consequence: the magnetic radius of the bound nucleon is not significantly larger than the one of the free nucleon.

- Measurements of the ^{205}Tl - ^{206}Pb charge density difference can be interpreted in a rather clean way in terms of the 3s proton radial wave function, and the 3s spectroscopic factor s . From elastic charge scattering we find $s = 0.7 \pm 0.07$. Measurements of $(e, e'p)$ and $(d, ^3\text{He})$ on both ^{206}Pb and ^{208}Pb allow us to derive from this spectroscopic factor the occupation of the 3s orbital in ^{208}Pb , which amounts to 0.91 ± 0.09 . We know that this occupation number is ≤ 1 , and we expect it to be ≤ 0.9 if short-range NN correlations are present. Again experiment and expectation agree, and they disagree if we were to assume that the bound proton has a charge radius increased by 20%. In this case the experimental occupation

number (derived mainly from data near 2 fm^{-1}) would be larger by 20%.

We therefore have good experimental evidence that the free and bound nucleon have the same form factor. To explain the longitudinal response function cited above, we need more realistic calculations including correlations and FSI. Such a calculation, presented at this workshop by Fantoni, does indeed a good job at explaining the data.

What we need is a better concept to deal with the expected modifications of the nucleon due to the medium. A change occurring during a short-range NN collision only, with the formation of a 6-quark bag, or alike, looks to me like a more fruitful concept in the long range.

Proceeding to other questions raised at this workshop, let me discuss next a point alluded to repeatedly during this meeting, the

QUESTION OF HIGHER DESIRABLE ENERGY OF CEBAF

At present the machine is laid out for 4 GeV, and if the cavities perform better than specified one might hope for 6 GeV. One of the great features of the superconductive RF cavity design: There is a long-term expandability in energy built in, and the machine could ultimately reach 16 GeV if future RF cavity progress keeps up with the past one. The question then is: Do we want a higher energy today?

At this point, one must repeat considerations that have been discussed before, but that apparently still have not been "heard" by everybody. The drive for high energies (momentum transfers) is due to the desire to reach asymptotia, where perturbative QCD is applicable, and life is simple. This regime can be reached in single-arm deep inelastic scattering. This regime cannot be reached for exclusive processes, where much larger q is needed to reach asymptotia. At these transfers the exclusive cross sections

are too small at any conceivable energy. Higher energy thus would not be of any help.

Part of the confusion arises from the erroneous identification of the regime where perturbative QCD is ok, and the regime where the nucleonic response functions become structureless. The latter occurs at $q \gtrsim 2 \text{ GeV}/c$, $\nu \gtrsim 2 \text{ GeV}$, while the former occurs at much larger q, ν . We can, by inclusive and exclusive experiments, cover the region where the nucleon exhibits identifiable excited states, and we can, with 4-6 GeV energy, reach the region where the physics becomes uniform and exhibits no more qualitative change at higher q, ν . Hereby we do cover the entire region of real interest for a quantitative study of QCD in the strong interaction limit. We cannot, at any energy, reach the limit where exclusive processes could be predicted by perturbative calculations.

Turning to more practical matters, I would like to address some questions concerning experimental hardware at CEBAF, in particular

SPECTROMETERS

The different working groups have discussed the hardware requirements needed to carry out specific experiments. I will try to give some general impression on the set of detectors that looks optimum for overall productivity of the facility.

For Endstation A, several types of spectrometers are under discussion: A 4 GeV high resolution device, upgradable to 6 GeV. As the second spectrometer, a 1.2 GeV vertical magnet with $\sim 35 \text{ msr}$ acceptance is planned; alternatively, a 3 GeV, $\sim 11 \text{ msr}$ device is envisaged. These spectrometers would be used mainly for $(e, e'p)$ experiments aiming at initial nucleon momenta up to large values, and $(e, e'K)$ experiments with electron energy of $\sim 2 \text{ GeV}$, low final electron

energy and K-momenta of order 1 GeV/c.

When trying to choose between the 1.2/3 GeV spectrometers, one should keep in mind one scaling law: For realistic spectrometers (of given cost and optimum technology) resolution is proportional to solid angle, and solid angle is inversely proportional to maximum momentum. Maximum performance in one quantity leads to poor performance in the other.

Keeping this scaling in mind, I would clearly favour the choice of a 1.2 GeV spectrometer for the following reasons. The maximum momentum is large enough for most (e,e'p) experiments (a, say, d(e,e'p) experiment to larger q can be performed using low-resolution hardware in ESC). The lower maximum momentum gives the larger Ω clearly needed for (e,e'K). The lower energy gives the shorter spectrometer vital for (e,e'K). The lower energy allows to achieve the 10^{-4} resolution needed for (e,e'p) and (e,e'K). The construction of spectrometers with different maximum momentum (4 GeV, 1.2 GeV) with otherwise optimum performance give maximal flexibility.

This choice would lead to the following setup: 4 and 1.2 GeV/c spectrometers for e' and p in (e,e'p), a few hundred MeV/c tagging (0°) and 1.2 GeV/c spectrometer for e' and K in (e,e'K). Both spectrometers would have $<10^{-4}$ resolution, the tagger $\sim 10^{-3}$. ESA would be optimized for high-resolution work.

For Endstation C, there are 3 spectrometers under discussion. A 5 GeV/c QQD for electrons, and one or two large acceptance detectors (LAD) for hadrons, one rather than two for financial reasons. These detector systems would serve initially for the (e,e'2N) and (e,e'XY) experiments. Their main property would be a large solid angle.

I consider the (e,e'2N) and (e,e'XY) experiments to represent an important fraction of the CEBAF base program. The detectors in ESC will be needed in addition to the 4 π detectors that will be used to mainly study more complex final states at lower luminosity. Consequently, I consider

it necessary to perform (e,e'XY) experiments under optimum conditions, i.e. with discrete detectors of large solid angle. We do need two hadron detectors from the very beginning, and their character needs to be defined together with the rest of the initial experimental equipment of CEBAF.

Two options for the hadron detector are open: 1) non magnetic devices, and 2) magnetic spectrometers.

We need to find out first whether option #1 is a viable one. At NIKHEF tests have been made to determine the response of plastic detectors in direct sight of the target. These tests have been made with a clean enough beam such as to suppress backgrounds due to slit scattering, halo, etc. From the experience at NIKHEF, one should try to extrapolate to CEBAF conditions in order to specify the luminosity achievable. If this luminosity is large enough, non-magnetic devices are preferable, for a number of reasons: they offer a large flexibility for future alternative uses, they serve as neutral-particle detectors as well, and out of plane measurements are more realistic.

Option #2 involves magnetic spectrometers, the only alternative if the luminosity of option #1 is too low. In this case, one should realize that one is talking about 2 hadron spectrometers. These probably would need to have a more closed geometry than the LAD presently discussed given the difficulties with background that excluded option #1.

In either case, one should realize that it does take 2 hadron detectors, not only 1. It probably is also true that a sophisticated non-magnetic detector is as expensive as a magnetic one; money must be set aside for either one.

As a general rule, I would advocate to make ESA and ESC as orthogonal as possible in their performance. ESA should provide the best resolution, ESC the largest solid angles. Mixing the performance goals by combining functions will lead to compromises that hurt in the long range.

Another question repeatedly alluded to during this workshop concerns the need for

OUT OF PLANE MEASUREMENTS

To measure all the structure functions, one obviously has to go out of plane. Realistically, however, nobody will ever measure all of the ~10 structure functions that could be measured. This simply represents too much hassle.

There are at present a number of predictions available for these structure functions, but we do not yet have clear ideas on the relation between a given piece of physics of interest and the optimum structure function that allows to measure it. We will have to analyze in detail many more calculated structure functions in order to develop a qualitative understanding of the importance of a given individual one.

Despite this present lack of qualitative understanding, it seems imperative to me to prepare for out of plane measurements. We have to build this capability right into the CEBAF initial design.

My feeling is that we will not need a continuous variation of the out-of-plane angle ϕ . One value will do. $\phi = 20\div 30^\circ$ looks big enough given the fact that the electron most often will be detected at a small angle.

In order to realize an out-of-plane setup, it seems most realistic to me to move the beam, and leave alone the spectrometers. For several reasons this (unconventional) technique is advantageous:

- The emittance of the beam is infinitely smaller than the solid angle of the spectrometer. Beam transport magnets thus can be orders of magnitude lighter than spectrometers.
- For a superconductive accelerator, the electron polarization can be adjusted at the source such as to have the desired orientation at the target, no matter what deflection the beam undergoes.

- Given the poor quality of beam transport magnets needed, iron-free superconductive coils of high field (small dimension) can be used to move the beam up and down.

I therefore would imagine to take the beam above the scattering plane at the entrance of the Endstation, and then down to the pivot of the spectrometers in the horizontal plane. To recover the outgoing electrons without loss, a second beam dump below the Endstation-floor might be needed.

Concerning the design and construction of these detectors, I would like to add one last consideration relating to

COLLABORATIONS SUBCONTRACTING DETECTORS

Based on construction funding for FY 87, CEBAF will review the plans for experimental equipment in fall 86, and will freeze the design in spring 87. This implies that users who do want to actively participate in the design and construction need to get involved rather quickly.

It is a CEBAF policy to subcontract major pieces of experimental equipment to user groups. These groups thus get an excellent chance to define, via the equipment they build, the physics program of the next decade, and these user groups can get important activity and visibility at their home institution by building the spectrometers and detectors. I therefore encourage the user groups to get organized quickly, and to take on the challenge of building one of the major pieces of experimental hardware!

These remarks bring my summary to an end. This workshop has been most stimulating and enjoyable one, and it has been very profitable in focussing on a number of issues of the CEBAF physics program and experimental equipment. In the name of all the participants I want to thank all those who have contributed to the organization of such a smoothly running meeting!

CHARGE LONGITUDINAL RESPONSE
IN INCLUSIVE (e,e') SCATTERING

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ABSTRACT

The experimental dynamic and static longitudinal structure factors of C^{12} , Ca^{40} and Fe^{56} are investigated with a semiclassical RPA theory, utilizing a semi-phenomenological ph interaction. One obtains a satisfactory account of the experimental data over a wide range of momentum transfer in all the three nuclei, providing that the proton root mean square radius gets increased of about 13% in C^{12} and ~20% in Ca^{40} and Fe^{56} .

INTRODUCTION

As it is well known the charge longitudinal nuclear response has been separated out¹⁻³ in a few light and medium nuclei, over a wide range of momentum transfers. It displays several interesting, and sometimes puzzling, features: indeed, at variance with the universal behaviour of the transverse response, in going from C^{12} to Fe^{56} it displays appreciable changes, which point to the relevance of the nuclear surface.

Moreover, at moderate momenta the longitudinal response turns out to be rather flat or even raising at low energies, revealing the presence of strong collective effects. Finally the integrated response (Coulomb sum rule) even for the largest attained momenta displays a substantial reduction with respect to the free value, thus raising a puzzling problem of missing strength, in a momentum region where correlations are usually believed to be negligible.

As far as collectivity is concerned, the nature of the particle-hole (ph) interaction is different in the two isospin channels: the isovector ($\tau=1$) force is essentially

repulsive and tends to concentrate strength at high energy. On the contrary, the isoscalar ($\tau=0$) ph force is substantially density dependent and becomes strongly attractive on the surface of the nucleus, thus producing an enhancement of the response in the low energy region, particularly pronounced in the lightest nuclei (C^{12}), where the surface plays a major role.

To explore these interesting physical effects, we have evaluated⁴ the charge longitudinal response in the framework of the semiclassical RPA theory introduced by Schuck et al.⁵ The latter allows, in particular, to account for a realistic dependence of the nuclear density upon the radial distance; thus a density dependence of the ph force may yield substantial effects associated with the low-density surface.

THE MODEL

The charge longitudinal response, separated out in inelastic electron scattering, can be expressed in terms of the $\sigma=0$, isoscalar and isovector (semiclassical) polarization propagators as follows

$$R_L(q, \omega) = -4 \int_0^{R_c} dr r^2 \text{Im} \left\{ \prod^{\tau=0} [q, \omega; k_F(r)] + \prod^{\tau=1} [q, \omega; k_F(r)] \right\} \quad (1)$$

where the local Fermi wavenumber, $k_F(r)$, is defined by the equation

$$\frac{\hbar^2 k_F^2(r)}{2m^*(r)} + W(r) = \frac{\hbar^2 (k_F^0)^2}{2m^*(0)} + W(0) \quad (2)$$

and the upper limit of integration R_c is the classical turning point for the single particle motion inside the nucleus, $k_F(R_c) = 0$. In equation (2) the value of $k_F^0 = k_F(0)$ is set by the normalization condition

$$\int d^3r \theta(R_c - r) \rho(r) \approx \int d^3r \theta(R_c - r) \frac{2k_F^3(r)}{3\pi^2} = A \quad (3)$$

(A being the mass number) and $W(r)$ is a suitable, local potential well, which we assumed of the Wood-Saxon type:

$$W_{ws}(r) = \frac{V_0}{1 + e^{[(r-R_0)/a_0]}}. \quad (4)$$

Non-localities of the nuclear mean field have been roughly accounted for by introducing a position dependent effective mass,

$$m^*(r) = m [1 - \alpha \rho_{\text{exp}}(r)], \quad (5)$$

which is lower than m (typically $0.8m$) in the nuclear interior, and smoothly reaches the free value following the experimental nuclear density

$$\rho_{\text{exp}}(r) = \frac{\rho_0}{1 + \exp[(r - R_1)/a_1]}. \quad (6)$$

Actually the effective mass should also depend upon the nucleon momentum and thus, indirectly, upon the momentum transfer q . Indeed we found it more appropriate to assume $m^* = m$ when dealing with the largest q -values. In Table I the parameters of $W_{\text{WS}}(r)$ and α are reported for the nuclei we have analyzed, namely ^{12}C , ^{40}Ca and ^{56}Fe , together with the values ρ_0 , R_1 , a_1 characterizing the experimental proton densities.

The random phase approximation (RPA) for the longitudinal response (1) reads

Table I

	ρ_0 (fm ⁻³)	R_1 (fm)	a_1 (fm)	V_0 (MeV)	R_0 (fm)	a_0 (fm)	α
^{12}C	0.087	2.32	0.42	-52 (-50)	2.59 (2.96)	0.54	2.3
^{40}Ca	0.082	3.60	0.58	-48 (-47)	3.92 (4.36)	0.58	2.5
^{56}Fe	0.068	4.28	0.56	-47 (-46)	4.74 (5.14)	0.62	3.0

Columns 1-3: the parameters of the Fermi distribution (6) which best accounts for the experimental proton density. Columns 4-6: the parameters of the Wood-Saxon well (4) which yield a nuclear matter type density close to (6) for the effective mass (5); in parenthesis are quoted those corresponding to the bare mass. Column 7: the parameter α of the effective mass (5).

$$R_L^{RPA}(q, \omega) = -4 \operatorname{Im} \int_0^{R_c} dr r^2 \Pi^0[q, \omega; k_F(r)] \times \left\{ \frac{1}{1 - V_{ph}^{\tau=0}(q) \Pi^0[q, \omega; k_F(r)]} - \frac{1}{1 - V_{ph}^{\tau=1}(q) \Pi^0[q, \omega; k_F(r)]} \right\} \quad (7)$$

where Π^0 is the free ph polarization propagator and $V_{ph}^{\tau}(q)$ the ph force in a given isospin channel.

Concerning the latter, we have fixed the strength of the ph interaction in the $q \rightarrow 0$ limit through the Landau parameters,

$$\lim_{q \rightarrow 0} V_{ph}^{\sigma, \tau}(q) = \frac{2\pi^2 \hbar^2}{k_F(r) m^*(r)} F_0^{\sigma, \tau}, \quad (8)$$

which, in turn, are related to the compression modulus K and to the symmetry energy coefficient b_{sym} according to the well-known formulas

$$F_0^{0,0} = \frac{3 m^*(r)}{\hbar^2 k_F^2(r)} K - 1 \quad (9a)$$

and

$$F_0^{0,1} = \frac{3 m^*(r)}{\hbar^2 k_F^2(r)} b_{sym} - 1. \quad (9b)$$

This connection entails some density dependence to the ph force, particularly in the isoscalar channel, where the compression modulus of a nucleus like Ca^{40} , according to microscopic calculations⁶, varies from $K \equiv K_V \cong 40$ MeV in the interior to $K \equiv K_S \cong -16$ MeV at the surface. We have smoothly interpolated with the following quadratic density dependence

$$K(\rho) = a \rho^2 + b \rho \quad (10)$$

between these two values of K .

The symmetry energy coefficient is also known to increase somewhat with the density⁷, according to

$$b_{sym}(\rho) = b_{sym}^0 (\rho/\rho_0)^{2/3}, \quad (11)$$

with $b_{sym}^0 \cong 70 \pm 10$ MeV.

Since $\rho = \rho(r) = 2k_F^3(r)/3\pi^2$, a dependence of the ph force upon the radial distance r is obviously implied by the equations (10) and (11).

On the contrary, the above relationship between K , b_{sym} and the ph force is silent about the momentum dependence of

of the interaction itself; the latter is not yet under a firm theoretical control, but for the well-established vanishing of the ph force at very high momenta. Thus we accounted for the q -dependence of V_{ph} with a purely phenomenological "form factor"

$$\Lambda^\tau(q) = \left(\frac{\mu_\tau^2}{q^2 + \mu_\tau^2} \right)^{P_\tau} \quad (12)$$

containing two adjustable parameters for each isospin channel. The resulting ph force reads then

$$\tilde{V}_{ph}^{\sigma,\tau}(q,r) = \Lambda^\tau(q) \frac{2\pi^2 \hbar^2}{k_F(r)m^*(r)} F_0^{\sigma,\tau}[g(r)]. \quad (13)$$

It is interesting to display our interaction (13) as a function of the radial distance r : in agreement with other findings⁸, one sees in Fig.1 that the isovector force is repulsive everywhere in the nucleus, while the isoscalar one, mildly attractive in the nuclear interior, becomes strongly attractive in the surface region.

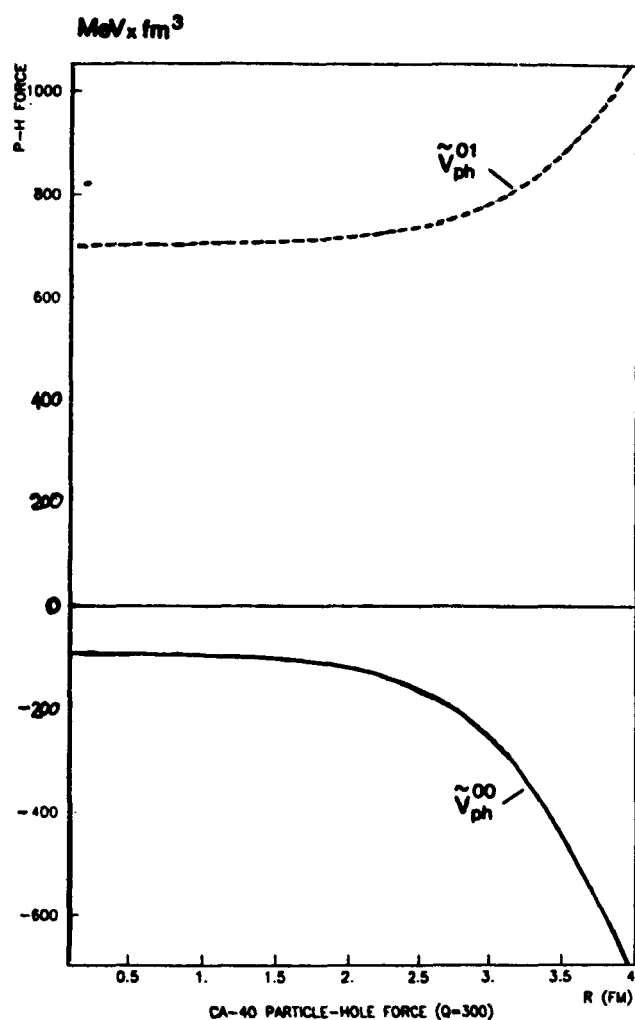


Fig.1 - The isoscalar and isovector ph force as a function of r , at $q = 300$ MeV/c.

RESULTS

Before comparing the longitudinal structure factor (7) with the experimental cross section one has to multiply the former by the electromagnetic proton form factor G_E (which accounts phenomenologically for the size of the proton) and the Darwin-Foldy term (which includes some relativistic effects):

$$S_L(q, \omega) = f^2(q_\mu^2) R_L(q, \omega) \quad (14)$$

where

$$f(q_\mu^2) = \frac{G_E(q_\mu^2)}{1 + (\hbar q_\mu)^2 / 4(mc)^2} \quad (q_\mu^2 = q^2 - \omega^2/c^2) \quad (15)$$

and the usual dipole parametrization for G_E ,

$$G_E(q_\mu^2) = \frac{1}{(1 + q_\mu^2/\lambda^2)^2}, \quad (16)$$

has been assumed.

We have then compared our semiclassical evaluation of $S_L(q, \omega)$ with the available data, measured on C^{12} , Ca^{40} , and Fe^{56} (ref.1,2). As a starting point the free-proton cutoff $\lambda^2 = 18.1 \text{ fm}^{-2}$, which corresponds to a proton root mean square radius (r.m.s.r.)

$$\sqrt{\langle R^2 \rangle_p} = \sqrt{12/\lambda^2} = 0.81 \text{ fm} \quad (17)$$

has been utilized.

With the values of the parameters μ_τ, p_τ of the ph force reported in Table II, we have then obtained a nice reproduction of the shape, as a function of energy, of the experimental cross section over a wide momentum range. However, by keeping the free proton form factor (see eqs.(16) and (17)), there was no way, even with a strongly altered ph force, to reproduce correctly the overall normalization of the data.

On the other hand, we could achieve a satisfactory fit of the experimental cross section by allowing for an increase in the proton size of $\sim 13\%$ in C^{12} and $\sim 20\%$ in Ca^{40} and Fe^{56} . More precisely we employed

$$\sqrt{\langle R^2 \rangle_p} = 0.93 \text{ fm} \quad (\lambda^2 = 14 \text{ fm}^{-2}) \text{ in } C^{12}$$

$$\sqrt{\langle R^2 \rangle}_p = 1.05 \text{ fm} \quad (\lambda^2 = 10.8 \text{ fm}^{-2}) \quad \text{in Ca}^{40}$$

$$\text{and } \sqrt{\langle R^2 \rangle}_p = 1.02 \text{ fm} \quad (\lambda^2 = 11.5 \text{ fm}^{-2}) \quad \text{in Fe}^{56}$$

The results of our calculation are illustrated in Figs. 2, 3 and 4 for C^{12} , Ca^{40} and Fe^{56} , respectively. In each figure the independent particle response (dotted line) as well as the isoscalar and isovector components (dashed lines) of the RPA response are displayed. The effect of the swollen proton is clearly shown in Fig.3b, where the RPA cross section (14), both with the free proton form factor (double-dot-dashed line) and with the modified one (continuous line) are reported.

One should stress that the need for a renormalization of the proton size is independent of the parametrization of the ph interaction, since the latter plays no role at the largest momenta; similar considerations apply to the effective mass.

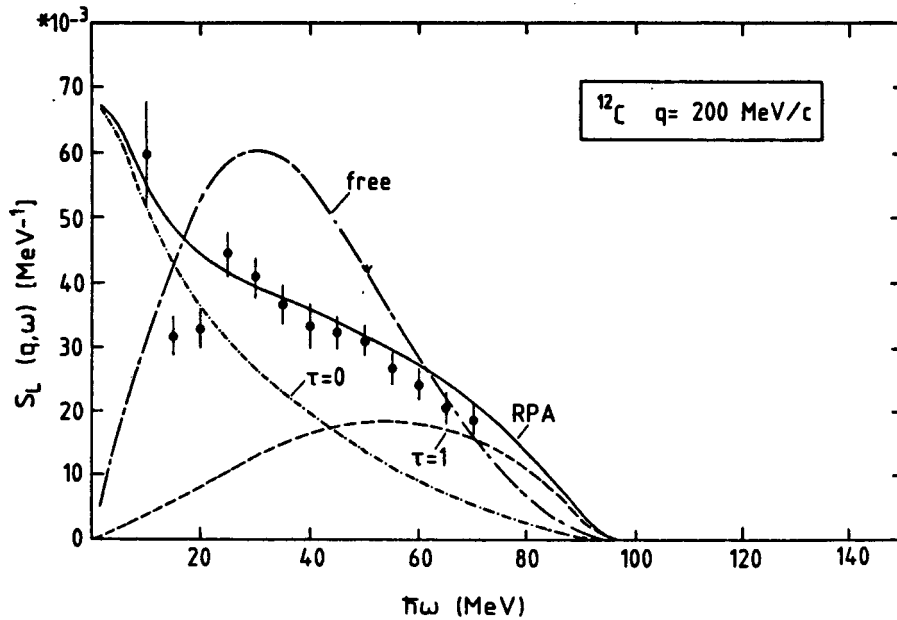


Fig. 2a - RPA charge longitudinal dynamical structure factor in C^{12} (continuous line) at $q = 200 \text{ MeV/c}$. The experimental points are taken from ref.1.

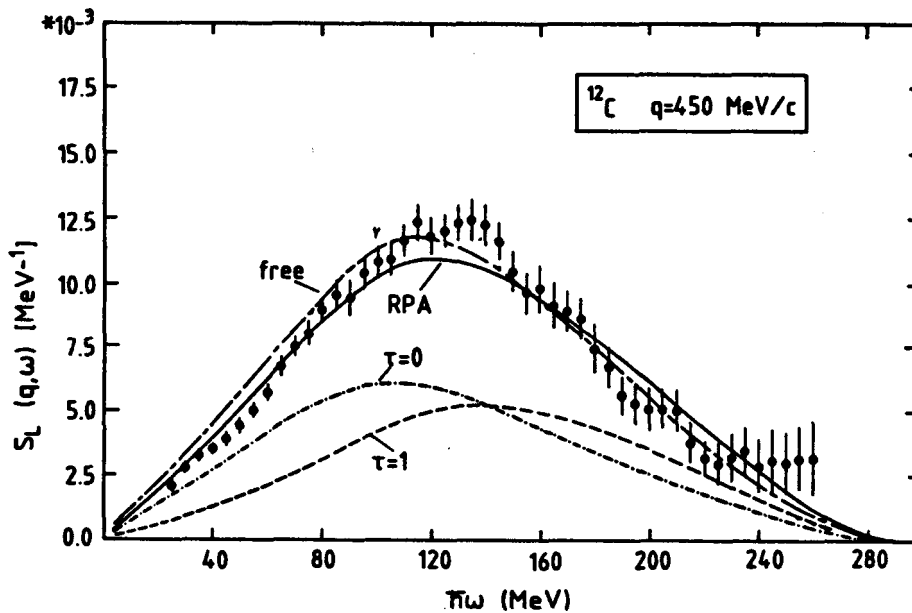


Fig. 2b - The same as in fig.2a, at $q = 450$ MeV/c.

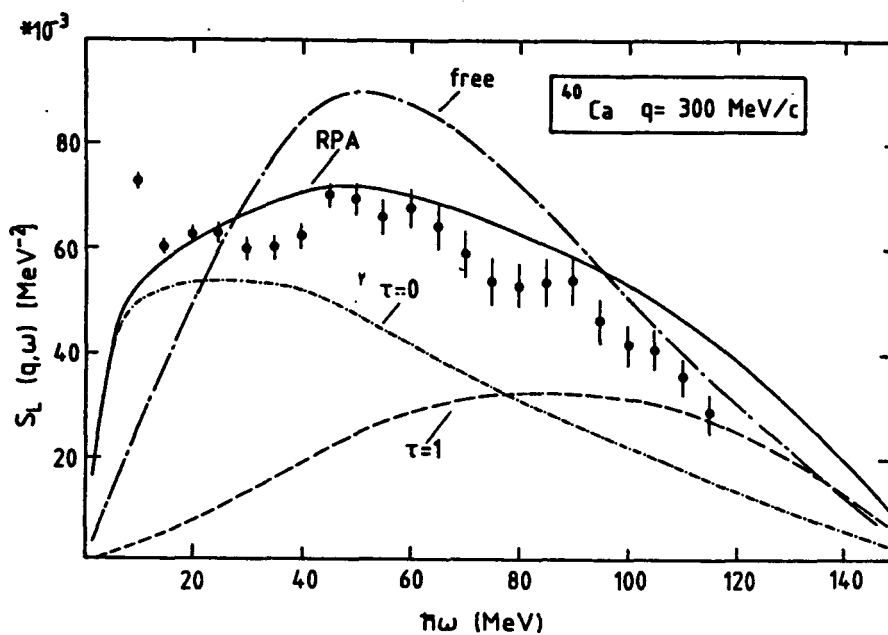


Fig. 3a - RPA charge longitudinal dynamical structure factor in Ca^{40} (continuous line) at $q = 300$ MeV/c. The experimental points are taken from ref.2.

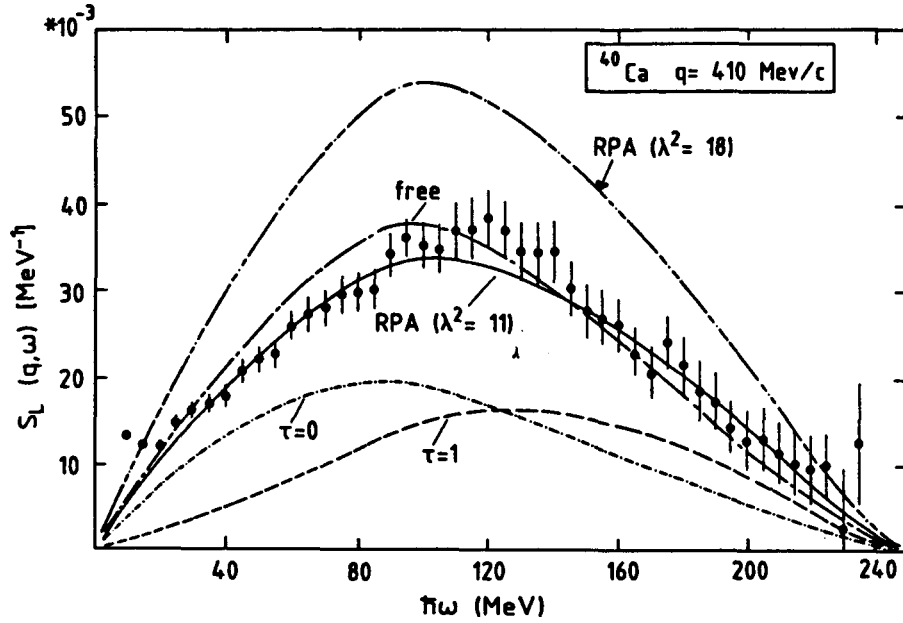


Fig. 3b - The same as in fig.3a, at $q = 410$ MeV/c. The RPA result with the free proton form factor is also shown (double-dot-dashed line).

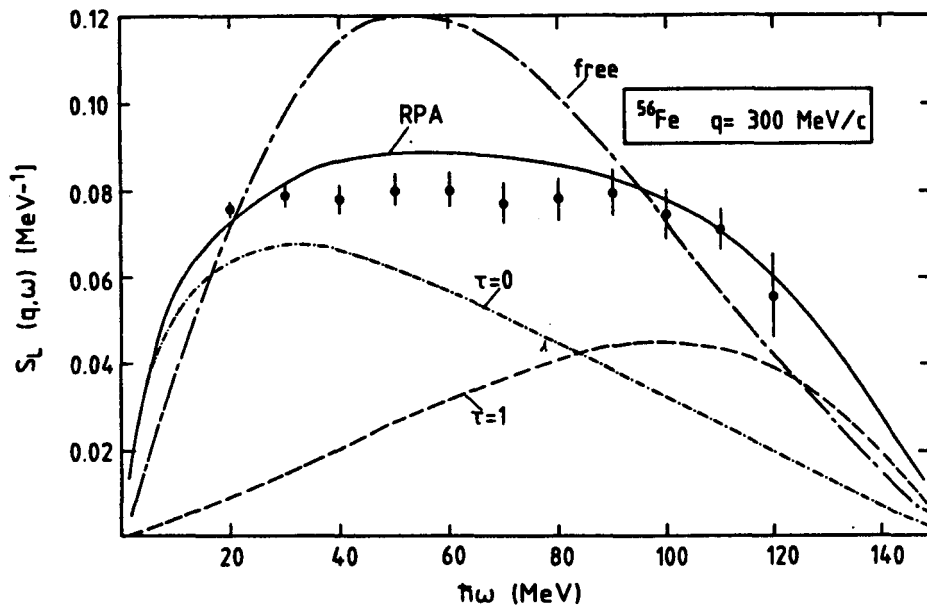


Fig. 4a - RPA charge longitudinal dynamical structure factor in Fe^{56} (continuous line) at $q = 300$ MeV/c. The experimental points are taken from ref.2.

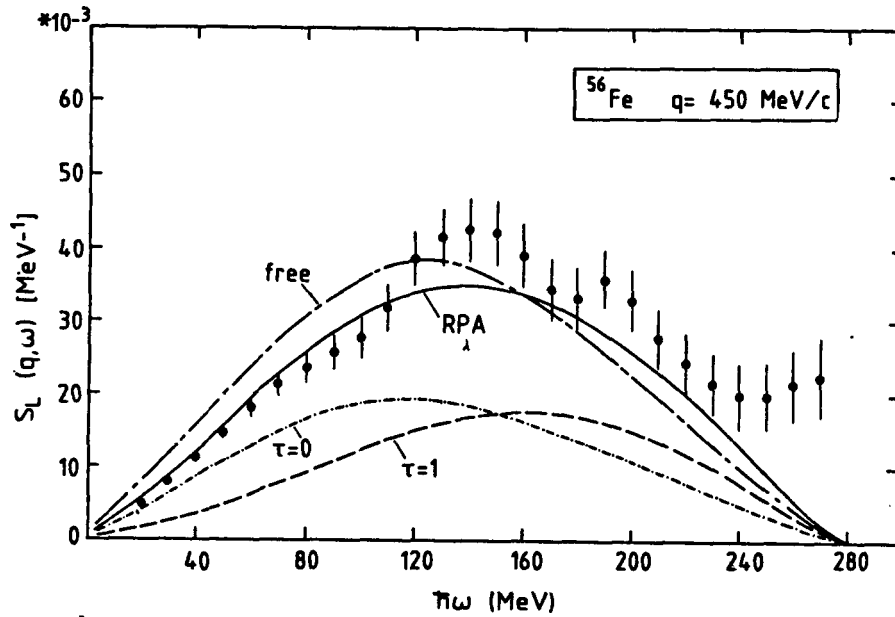


Fig. 4b - The same as in fig.4a, at $q = 450 \text{ MeV/c}$.

Within the same theoretical framework, we have also considered the integrated (over the energy) longitudinal response, usually referred to as "Coulomb sum rule":

$$S_L(q) = \int_0^{\infty} d(\hbar\omega) S_L(q, \omega). \quad (18)$$

In order to compare with the experiment, we have actually limited the integration to the energy range where data are available. For sake of illustration the Fe^{56} RPA sum rule is shown in fig.5, both with the free proton form factor (squares) and with the modified one (black points). Again the latter turns out to be essential in accounting for the experimental strength.

Table II

	μ_0	p_0	μ_1	p_1
C^{12}	430.	2.2	400.	0.95
Ca^{40}	440.	2.15	500.	1.0
Fe^{56}	500.	2.9	400.	0.95

Parameters of the ph force entering into (12). The μ 's are in MeV/c.

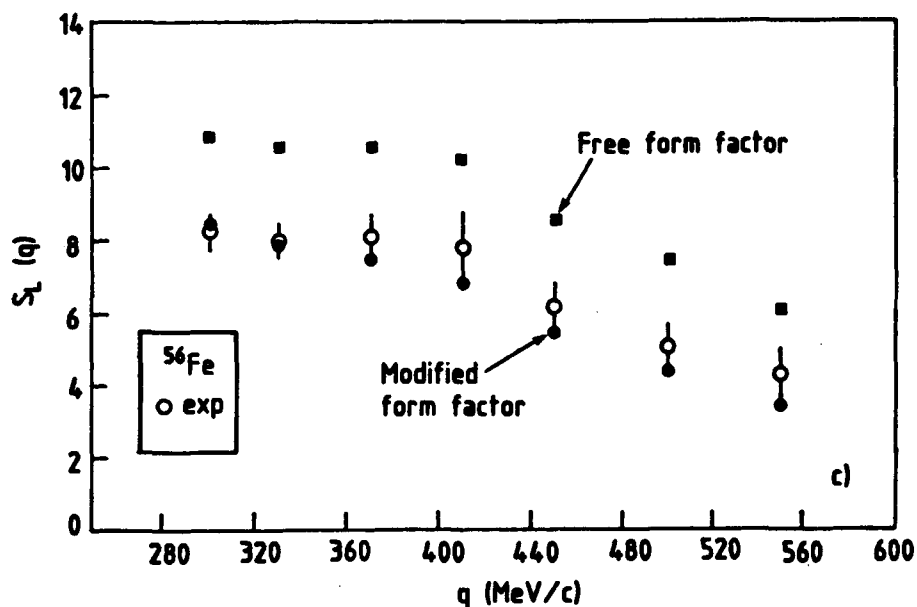


Fig. 5 - The Coulomb sum rule in Fe^{56} .

CONCLUSIONS

We have performed a detailed study of the charge longitudinal response in C^{12} , Ca^{40} and Fe^{56} in the framework of a semiclassical RPA theory, which accounts for finite size effects.

Two points of physical relevance emerge from our analysis:

i) The behaviour of the nuclear response as a function of the energy, at small and intermediate momentum transfers, bears the clear signature of collectivity. Indeed the almost flat shape experimentally observed is the result of a balance between two collectivities acting in opposite directions: the hardening of the isovector response and the softening of the isoscalar one.

As the momentum increases the RPA correlations die out: as a consequence the "coherent" collective nuclear response gradually evolves towards the "incoherent" single particle one. It should be stressed the crucial role played by the nuclear surface in setting the strength of the ph interaction. Indeed on the surface the $\tau = 0$ force is much more attractive than in the interior of the nucleus, thus yielding a low energy collective enhancement of the nuclear response most clearly standing out in C^{12} , and much less pronounced in Ca^{40} and Fe^{56} .

ii) Within the theoretical framework employed here, we have found it possible to account for the magnitude of the experimental cross sections only by allowing for an increase of the proton radius as sizable as 20% in the heavier nuclei. Several conjectures on the physical, microscopic origin of this effect have been recently raised and are complemented, to some extent, by our empirical findings.

In particular the size increase could be ascribed to a change in the QCD mass scale, originated by the modification of the gluon condensate in the nuclear medium⁹. At the level of the proton electromagnetic form factor, this would imply a partial deconfinement of the quarks of a nucleon inside the nucleus. This hypothesis, while fascinating, can hardly be pushed, at present, beyond the stage of a model, owing to our poor control of non-perturbative QCD.

An alternative explanation for the swelling of the proton has been recently proposed in terms of the pionic content of the nucleon isovector electromagnetic form factor¹⁰. In this context, the occurrence of medium modifications of the pionic cloud could lead to a sizable increase of the isovector proton radius. Interestingly, this pionic interpretation of the swelling would have a microscopic origin in two particles - two holes (2p-2h) processes. The latter, in turn, are closely related to the explanation of the missing strength in terms of the depopulation of the single particle orbits, due to 2p-2h correlations.

In conclusion a realistic RPA treatment of the charge longitudinal response, with a suitable ρ - ρ interaction, together with the hypothesis of an increased electromagnetic proton radius, seems to provide a satisfactory description of the experimental data over a large interval of momentum transfers.

However this topic still deserves theoretical and experimental efforts, before a definitive explanation can be reached. We remind, for example, that a relativistic treatment of the proton dynamics inside the nucleus has been proposed¹¹ as capable of accounting for a quenching in the longitudinal response: in this connection the Walecka theory appears well suited to the purpose, but it is not yet clear whether it can actually provide a detailed, quantitative explanation. Analogous considerations apply to the quenching induced by two- (and three-) body correlations, as the ones considered in the framework of a variational calculation¹².

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SUPPRESSION OF THE PHOTONUCLEAR CROSS SECTION AND Δ -HOLE CALCULATIONS

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Calculations based on the Δ -hole formalism¹⁾ was thought to describe reasonably well the experimental data, like e.g. the total photonuclear cross section $\sigma_{\gamma,a}^{\text{exp}}$, except for some underestimation at the low energy side of the Δ -peak²⁾. However, such a comparison between measurement and calculation might be questionable, because the experimental $\sigma_{\gamma,a}^{\text{exp}}(\omega)$, where ω is the incident photon energy, represents the sum of partial cross sections referent to, at least, three competing photoexcitation mechanisms, namely: (1) pion production ($\sigma_{\gamma,a}^{\pi}$), (2) quasi-deuteron excitation ($\sigma_{\gamma,a}^{\text{QD}}$) and (3) the production of e^+e^- pairs ($\sigma_{\gamma,a}^{\text{PP}}$) by γ rays of intermediate energies³⁾ - see diagrams in Fig. 1. So, for $\omega > 140$ MeV, we have

$$\sigma_{\gamma,a}^{\text{exp}}(\omega) = \sigma_{\gamma,a}^{\pi}(\omega) + \sigma_{\gamma,a}^{\text{QD}}(\omega) + \sigma_{\gamma,a}^{\text{PP}}(\omega) \quad (1)$$

It is a well-known and accepted experimental fact that the so-called "quasi-deuteron photoexcitation mechanism", in its modified version or in the approach proposed by Laget, explains quite well the photoabsorption cross sections, experimentally determined in a number of nuclei⁴⁾, for energies above the Giant Dipole Resonance and up to, at least, the photopion threshold (~ 140 MeV). Recently we showed⁵⁾ that, in the fission decay of ^{209}Bi , the QD-mechanism accounts for the nuclear photoexcitation process which drives ^{209}Bi to fission, for energies up to 250 MeV, using the same parameters of $\sigma_{\gamma,a}^{\text{QD}}$ deduced up to ~ 140 MeV - this QD cross section is shown in Fig. 2, along with the total photoabsorption cross section of ^{209}Bi obtained by scaling of $\sigma_{\gamma,a}^{\text{exp}}/A$ (determined for several nuclei, and exhibiting an "universal" behavior²⁾).

In order to get the photopion cross section we must consider, at least, the following (see Eq. 1)

$$\sigma_{\gamma,a}^{\pi}(\omega) = \sigma_{\gamma,a}^{\text{exp}}(\omega) - \sigma_{\gamma,a}^{\text{QD}}(\omega) - \sigma_{\gamma,a}^{\text{PP}}(\omega) \quad (2)$$

Regarding pair production, it is worth noting that recently Kazakov et al.³⁾ investigated the fission of ^{238}U and ^{237}Np induced by back-scattered Compton γ rays over the energy interval 150-710 MeV; according to their data, the cross section for the excitation of a nucleus in pair production (energy transferred by virtual photons - see Fig. 1) is $\sim 0.02\%$ of the total (atomic) photoabsorption cross section. Besides this purely electromagnetic mechanism, a variety of other

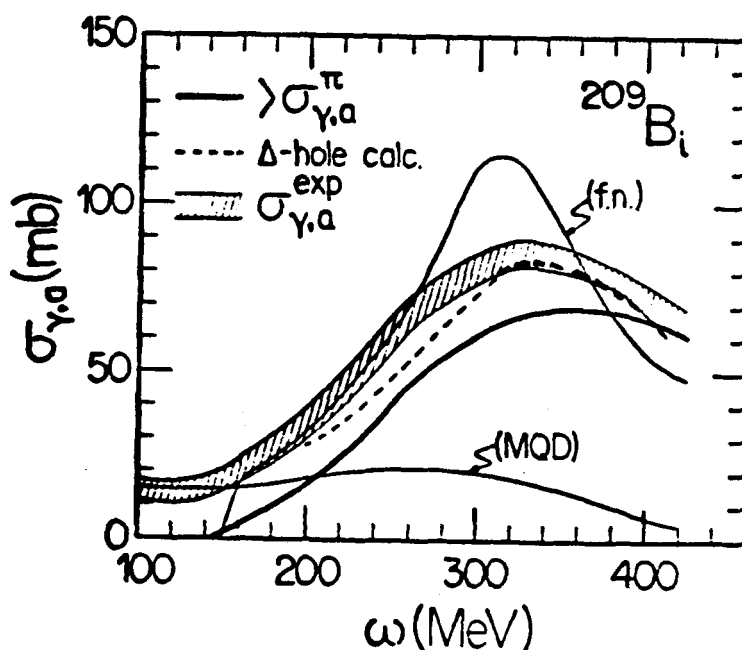
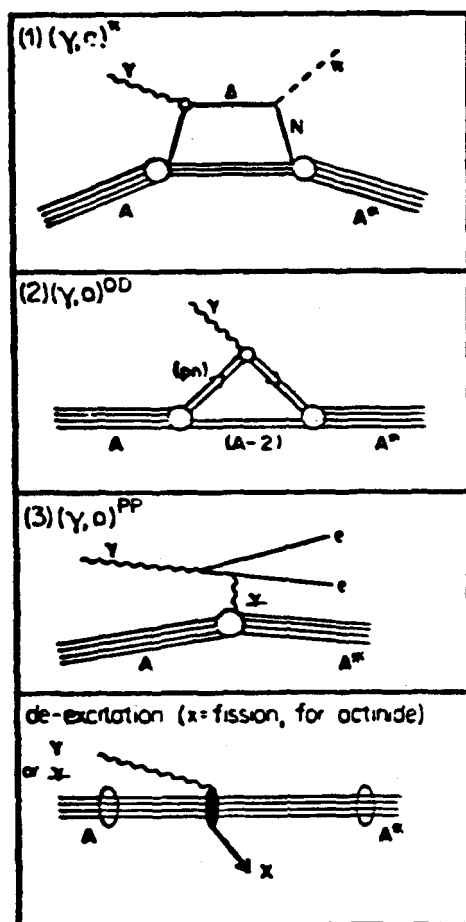
mechanisms should contribute to the total nuclear photoabsorption. However, we decided to estimate an upper limit for $\sigma_{\gamma,a}^{\pi}$ by considering only the QD contribution to $\sigma_{\gamma,a}^{\text{exp}}$; therefore,

$$\sigma_{\gamma,a}^{\pi}(\omega) < \sigma_{\gamma,a}^{\text{exp}}(\omega) - \sigma_{\gamma,a}^{\text{QD}}(\omega) \quad (3);$$

this upper limit is shown in Fig. 2. Now, the disagreement with the Δ -hole calculation (dashed curve in Fig. 2) is much more severe, particularly at the peak of the Δ resonance. From the theoretical point of view such a discrepancy is quite disappointing, for the Δ -hole formalism¹⁾ is, to date, the most complete approach in the treatment of how the nuclear medium modifies $\gamma\text{N}\pi$ processes. We observe a lack of strength (relative to the calculated one) in the entire energy region from m_{π} to the Δ -peak. Now, the suppression of $\sigma_{\gamma,a}^{\pi}$, relative to the free nucleon (curve "f.n." in Fig. 2), is greater than 50%! So, where is the missing strength? This question deserves prompt elucidation.

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GENERATING THE BAG WITH GLUEBALLS

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This research has been done with M.K. Banerjee and T.D. Cohen. We discuss a class of absolutely confining soliton models for baryons. The basic idea is to couple quarks to the glueball in such a way that continuum quark solutions are eliminated. The conjecture that confinement may be the result of quarks creating a cavity in the glueball condensate present in the vacuum, was proposed by Hansson, Johnson and Peterson [1]. Several other authors have also speculated along these lines [2]. Our objective is to form a bridge between the soliton models and the bag models, by producing confinement selfconsistently via interaction of quarks with the 0^{++} glueball, Γ . The coupling of quarks to Γ has a very exotic form. Since the glueball consists entirely of gluons, it is a chiral singlet. In order to preserve the chiral symmetry we have to couple Γ to a quark form which is also a chiral singlet. Restricting ourselves to linear coupling, we arrive at the following term: $i/2 g_\Gamma \Gamma(x) \bar{\psi} \psi$. For concreteness let us consider the following lagrangian:

$$\mathcal{L} = (1+g_\Gamma \Gamma) \frac{i}{2} \bar{\psi} \not{\partial} \psi + g \bar{\psi} (\sigma + i \gamma_5 \vec{\tau} \cdot \vec{\pi}) \psi + \mathcal{L}'(\Gamma) + \mathcal{L}''(\sigma, \vec{\pi}), \quad (1)$$

where g_Γ and g are constants, and σ and $\vec{\pi}$ denote the sigma and pion fields. The glueball lagrangian, $\mathcal{L}'$, consists of a kinetic piece and a quartic potential with a local minimum and a global minimum at $\langle \Gamma \rangle_{\text{vac}}$. The mesonic part, $\mathcal{L}''$, is suitably chosen to generate spontaneous chiral symmetry breaking. As a result, σ acquires a non-zero vacuum expectation value, $\langle \sigma \rangle_{\text{vac}} = -F_\pi$. For details of the lagrangian see refs. 3 and 4.

The operator canonically conjugate to ψ is no longer $\psi^\dagger$, but $(1+g_\Gamma \Gamma) \psi^\dagger$. Accordingly, the quark part of Noether currents acquires the factor $(1+g_\Gamma \Gamma)$. It is convenient to regain the usual forms by transforming $\psi \rightarrow (1+g_\Gamma \Gamma)^{-1/2} \psi$. The quark part of (1) becomes

$$\mathcal{L}_q = \frac{i}{2} \bar{\psi} \not{\partial} \psi + \frac{g \bar{\psi} (\sigma + i \vec{\tau} \cdot \vec{\pi} \gamma_5) \psi}{(1+g_\Gamma \Gamma)}.$$

With the new ψ fields a "denominator" coupling to the glueball has been formed.

There are reasons to believe that gluon-gluon interaction in QCD leads to a glueball condensate in the vacuum [1]. Our central conjecture is that *absolute confinement* implies a very special value for $\langle \Gamma \rangle_{\text{vac}}$, such

that the following condition is satisfied:

$$1 + g_F \langle \Gamma \rangle_{\text{vac}} = 0.$$

One can show easily that if the system is enclosed in a volume V , plane-wave quark states have energies approaching infinity as the $V^{1/3}$. Thus the model has no deconfined quark solutions. However, a soliton multiquark system can alter locally the glueball field from its vacuum value. The quarks can live in that selfconsistently formed cavity (or soft bag). As the quarks move towards the surface of the bag, where $\Gamma \rightarrow \langle \Gamma \rangle_{\text{vac}}$, the potential becomes infinite, preventing the quark from moving farther. Thus the model confines the quarks to a finite volume in space. We should stress that we offer a description of absolute confinement with a *chiral invariant* lagrangian. We also note that the spontaneous symmetry breaking of the vacuum, i.e. $\langle \sigma \rangle_{\text{vac}} \neq 0$, is an indispensable ingredient for absolute confinement in our model.

Fig. 1 shows the numerical solution to (1) with fields of the hedgehog form. We have chosen $g_F = 0.048 \text{ MeV}^{-1}$ and glueball mass = 1400 MeV. For other parameters of $\mathcal{L}'$ and $\mathcal{L}''$ see refs. 3 and 4. We plot the shifted glueball field, $\chi = 1/g_F + \Gamma$. Near the origin, the χ field remains close to the local minimum of the glueball potential. Around $r = 1.5 \text{ fm}$ it drops to its vacuum value, creating a soft surface. At this point the quark fields approach zero very rapidly. The σ field does not depart significantly from its vacuum value. The total energy of the soliton is 1497 MeV. We note that I. Duck has independently found solutions for a model very similar to ours [5].

We will now move to a simplified confining model with quarks and glueballs only [3,6]. Although the model no longer possesses chiral symmetry, hence cannot be a viable candidate for the description of nature, it allows us to concentrate entirely on the role of the glueball. The model is formally obtained from (1) by replacing the chirally invariant coupling to mesons by a mass term:

$$\mathcal{L}_q = i \bar{\psi} \not{\partial} \psi - \frac{\alpha \bar{\psi} \psi}{\chi},$$

where we have chosen $\alpha = (39 \text{ MeV})^2$ and the glueball part of the lagrangian has parameters from ref 3. Fig. 2 shows the solution. It has a total mass of 903 MeV. The character of the solution is reminiscent of the MIT bag. To point it out explicitly, in Table I we show the energy decomposition and some other properties for our solution and for the MIT bag.

There are clear advantages in dealing with a dynamical system of quarks and glueballs. The smooth transition between the central and asymptotic regions removes the bag pathologies and one can attempt to calculate the center of mass corrections, describe the N-N interaction or study the nucleon in the nuclear medium.

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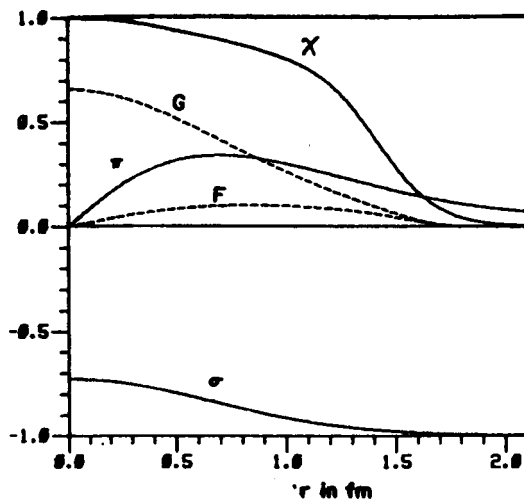


Fig. 1. Solution of the chiral confining model: Quark upper and lower components, G and F (dashed lines), the X field in units of the central value of 42MeV (solid line) and the σ and π fields in units of F_π (dotted line).

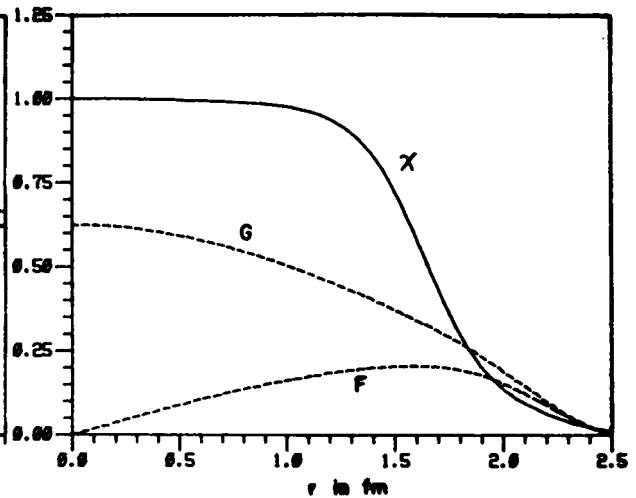


Fig. 2. Solution of the non-chiral confining model: Quark upper and lower components, G and F (dashed lines) and the X field in units of the central value of 32MeV (solid line).

TABLE I. Comparison of energy decomposition and some other properties for the solution to the non-chiral confining model and the MIT bag. E is the baryon mass.

QUANTITY	OUR SOLUTION	MIT
quark eigenvalue/ E	0.248	1/4
glueball kinetic energy/ E	0.089	-
glueball potential energy/ E	0.167	-
$\int d^3r G^2$	0.79	0.74
$g_A(0)$	1.20	1.09
proton magnetic moment $e/2E \mu_p$	3.46	3.30

($e, e'K^+$) and low-lying hypernuclear states using relativistic field theory

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Abstract: We use a relativistic model to describe hypernuclear low-lying excitations via the ($e, e'K^+$) reaction. The transition operator is derived from Feynman diagrams, and the nuclear and hypernuclear wave functions contain Dirac-Hartree orbitals from a relativistic nuclear structure model extended also to hypernuclei. We consider both pseudoscalar (PS) and pseudovector (PV) KNA vertex couplings and discuss the differences between them. We show the effects of the relativistic ingredients--the large relativistic scalar and vector potentials and the four-component spinor wave functions--on predicted cross sections. Shapes and magnitudes of predicted cross sections are significantly altered compared to previous nonrelativistic results, and large differences occur between the PS and PV results.

Recent interest in the electro- and photo-production of a kaon on a nuclear system,¹⁻³ resulting in hypernuclear excitations, is due to the special features of the reaction (which make it interesting for hypernuclear studies) and the attributes of the new generation of cw electron accelerators and high-resolution spectrometers (which will be able to provide, for the first time, the means for studying this reaction). The special features of this method of hypernuclear formation compared with conventional forward-angle (K^- , π^-) formation are large momentum transfers ($q \lesssim 300$ MeV/c) to the nucleus, possible strong excitation of non-normal parity as well as normal parity, high- and low-spin final hypernuclear states, and the reaction mechanism involves only electrons and a K^+ , with minimal distortion- and strong interaction effects.

We shall consider a simple mechanism for the ($e, e'K^+$) reaction (Fig. 1) described in detail in Ref. 3. But, whereas previous

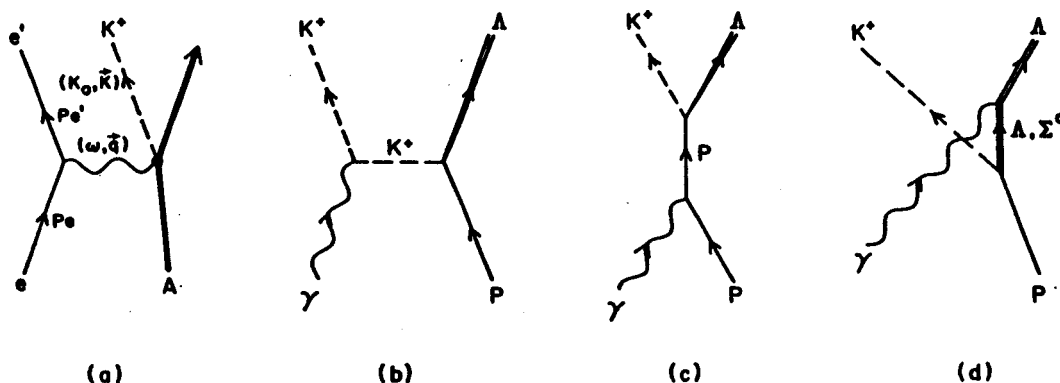


Fig. 1. (a) The nuclear $A(e, e'K^+)A$ reaction. (b)-(d) Feynman diagrams for the elementary reaction $\gamma + p \rightarrow K^+ + A$ for the PS coupling. A contact term is added for the PV coupling.

calculations¹⁻³ use nonrelativistic nuclear structure and reduced operators, we use consistently relativistic field theory in this work. The motivation for such an approach is clear: The process is characterized by important terms involving large four-momentum transfer dependent operators. For such a reaction the use of a non-relativistic formulation is not well-motivated, and the application of the recent relativistic nuclear many-body field theory is a natural theoretical framework to explore. In particular we report results using a relativistic approach for $^{16}\text{O}(e, e'K^+)$ leading to low-lying 0^- and 1^- final hypernuclear states.

The relativistic framework adopted is described by Serot and Walecka.⁴ For the nucleon sector we assume couplings defined for Finite-Hartree QHD, while for the strangeness sector we adopt couplings which are about a factor of 2.5 weaker than the nucleonic case, and are consistent with good fits to experimental spectra.⁵

The following problems are of interest and have been studied: a) the importance of consistently keeping all terms in the transition operator; b) the role of the strong relativistic potentials in affecting the relative importance of the contributions to the reaction mechanism, and to the cross-section as a whole; and c) the effects of the small components in the Dirac-Hartree spinors. An interesting question to explore is the effect of the coupling chosen for describing the KNY vertices. While the differences between the PV and PV couplings are small for positive pion photoproduction in nonrelativistic studies, the case may be completely different when QHD is used.

In Figs. 2 and 3 we present differential cross-sections for the excitation of the $(\Lambda 1s_{1/2}, P1p_{1/2}^{-1})_{0^-1^-}$ low-lying hypernuclear levels in the PS and PV couplings, respectively. Major features associated with the results shown in Fig. 2 are: a) Some previous^{1,2} nonrelativistic calculations have included only the $\vec{\sigma} \cdot \vec{\epsilon}$ term in the operator. This limits significantly the orbital angular momentum transfer L that can contribute to the cross section. A comparison given in Fig. 2 for the two types of calculations (those shown by the solid lines are the full calculations and the dashed-dotted lines represent calculations where the lower components of the Dirac orbitals are set to zero) show large differences between the two cases. Since the lower components (having an opposite parity relative to the upper ones) introduce the extra partial waves, the results indicate that terms should be kept in the transition operator beyond the $\vec{\sigma} \cdot \vec{\epsilon}$ approximation. b) A substantial effect results from the presence of the mean fields in the baryon propagators of the transition operator. This is a result of subtle interference effects between diagrams 1(b) and 1(c), and the changes incurring in 1(c) upon introduction of the relativistic potentials. In Fig. 2 the dashed line represents a calculation without the effect of the potentials in the propagators (but with the complete Dirac spinors included). Note that we use for the intermediate baryons only the nuclear-matter values of the strong potentials, and the actual effects will depend on the degree of surface or interior peaking of the transition. The effects are associated with the 0^- transition because diagram 1(b) does not contribute to the 1^- case.

We have also studied the relative contribution of the various

diagrams. While we leave a detailed description for an extensive report, we note that, in general, all diagrams are important except for 1(d) with a Σ^0 -pole.

In Fig. 3 we present results for the PV coupling. The results are quite different from the PS case and the question naturally arises: "which one is the right thing to use." While the authors of Ref. 4 conclude that the PV coupling is more practical, and the two should be equivalent (as a result of the algebraic transformations from the PS to the PV Lagrangian), we leave this question open for the time being. An interesting place to look for a solution is the $(\gamma, \pi^\pm)$ reaction, where the same formalism as used here is directly applicable and the experimental data exists.

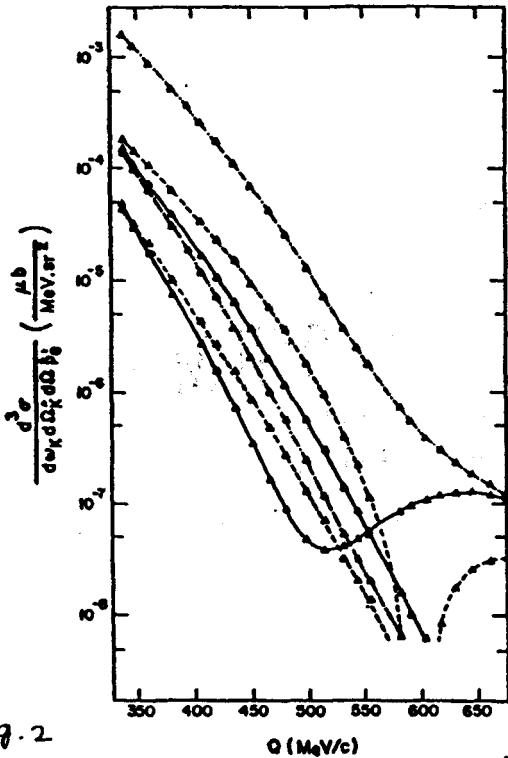


Fig. 2

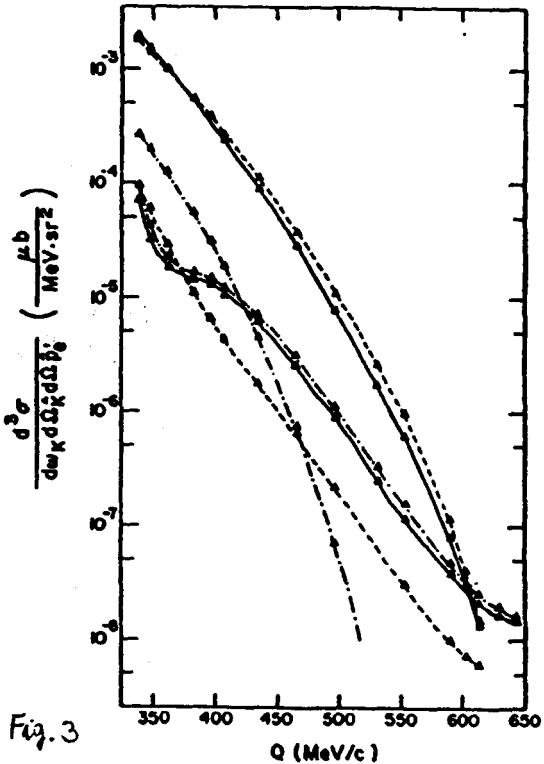


Fig. 3

Cross sections for the $(\Lambda 1s_{1/2}, \text{Plp}_{1/2})_{0^-, 1^-}$ excitations in ^{16}O . The calculations are for $p_e = 2 \text{ GeV/c}$, $p_{e'} = 800 \text{ MeV/c}$, $\theta_{p_{e'}} = 5^\circ$ and $\phi_K = 0$ (incoming electron momentum, outgoing electron momentum and angle, and out-of-the reactions-plane angle, respectively). Full triangles refer to 0^- , open ones to 1^- . The momentum transfer Q is varied by varying the K^+ -angle θ_K^+ . The meaning of the three types of curves is discussed in the text. Fig. 2 is for PS coupling, Fig. 3 for PV.

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DEUTERON POLARIMETERS AT INTERMEDIATE ENERGIES

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1. Introduction

There are fundamental problems for which the measurement of the polarization of recoil deuterons is needed: the separation of the elastic monopole and quadrupole electric form factors of the deuteron in e-d scattering, the complete determination of the $pp \rightarrow d\pi$ and $\pi d \rightarrow \pi d$ scattering matrices and possibly spin transfer experiments in deuteron-nucleus scattering with excitation of giant isoscalar resonances.

In most cases, the interesting range of deuteron energies is 100 - 400 MeV. For example, the - yet to be observed - minimum of the deuteron charge form factor G_C occurs around $q^2 = 20 \text{ fm}^{-2}$ (this is very much model dependent), which corresponds to a deuteron kinetic energy $T_d = q^2/2m_d \approx 200 \text{ MeV}$. To date the highest energy deuteron tensor polarimeters are based on the ${}^3\text{He}(d,p){}^4\text{He}$ reaction at 20 - 25 MeV 1).

2. Measurement of the Analyzing Powers

The analyzing powers for deuteron scattering at intermediate energies were measured at Saturne in experimental conditions such as to approximate a realistic polarimeter. 2)

In inclusive d-A scattering, significant values of iT_{11} ($\sim .4$) are found provided a rough selection on the energy of the scattered deuterons is made by the use of absorber material.

In dp scattering, both final state particles were detected and the elastic scattering was easily identified through angular correlation. At 191 MeV (Fig. 1), the analyzing powers are very much like at 79 MeV 3): the interference between direct dp scattering and neutron exchange creates a characteristic dip in T_{20} around 120°_{cm} . The observables are relatively well understood in terms of 3-body calculations up to 250 MeV 4) and this dip in T_{20} is expected to be present from 80 to 250 MeV. At higher

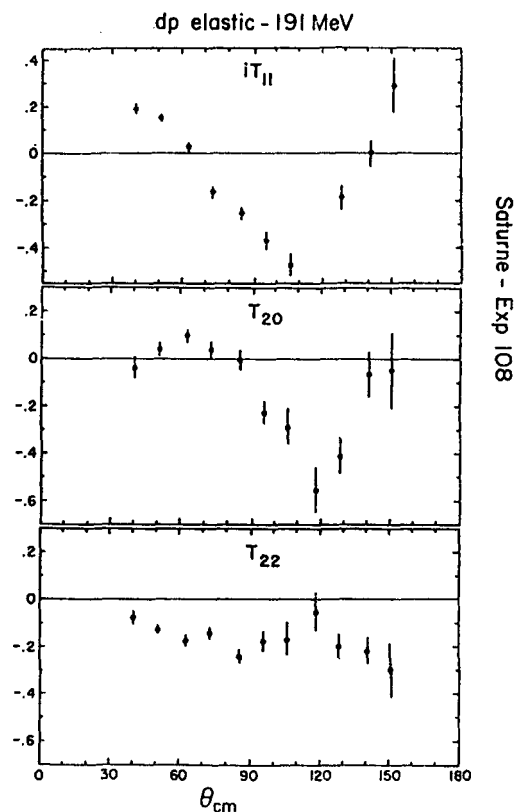


Figure 1.

energies, our measurements at 395 MeV and the excitation function at 180° ⁵⁾ show that the independent 3 nucleons picture is not valid anymore. The dip disappears while $T_{20}(180^\circ_{cm})$ becomes very large and negative.

3. Vector Deuteron Polarimeter

In principle, any proton polarimeter in use at intermediate energies ⁶⁾ may be transformed into a vector deuteron polarimeter, with the only addition of absorber slabs between the last wire chamber and the trigger. Efficiency would be about 1%. Such a polarimeter is to be calibrated at Saturne, with applications to the $pp \rightarrow d\pi$ reaction ⁷⁾. For e-d scattering, the spin transfer from polarized electrons to the deuteron ⁸⁾ ($p_x \sim G_M (G_C + 1/3\eta G_Q)$) could be the only measurable polarization observable above $q^2 \sim 40 \text{ fm}^{-2}$ because of a larger efficiency than in the case of tensor polarization.

A method based on deuteron break-up and subsequent analysis of the proton polarization ⁶⁾ has already been used successfully above $T_d = 400$ MeV, although with no calibration in a d beam of known polarization.

4. Tensor Deuteron Polarimeter

A polarimeter based on dp scattering between 100 and 250 MeV was built at the University of Alberta. Above $T_d = 100$ MeV, the scattered particles energies are large enough to escape a thick target, allowing trajectory determination; below 250 MeV, the same energies are small enough to allow a total energy measurement. Thus, the elastic event is identified by the energy-angle correlation of the scattered particles and by the angular correlation as well. The AHEAD polarimeter (Alberta High Efficiency Analyzer for Deuterons) will be calibrated at Saturne ⁹⁾ and used at Bates for ed scattering ¹⁰⁾. An efficiency of 2×10^{-3} around 150 MeV is expected.

4.1 Specific Advantages of the dp Reaction:

- The analyzing powers show little energy dependence ²⁻⁴⁾.
- T_{21} is expected to be large enough ⁴⁾ for a measurement of t_{21} (in ed scattering, t_{21} is proportional to G_Q).
- T_{20} shows a large angular dependence which will be used to eliminate the need for an absolute measurement of the number of deuterons entering the polarimeter; denoting by c and e the subscripts relating to the use of the polarimeter during the calibration (with unpolarized beam) and during the double scattering experiment, the number of counts at angle θ_i is in each case proportional to the efficiency ϵ :

$$\left. \begin{aligned} N_{i,c} &= C_c \varepsilon(\theta_i) \\ N_{i,e} &= C_e \varepsilon(\theta_i)(1 + t_{20} T_{20}(\theta_i)) \end{aligned} \right\} + N_{i,e} = k N_{i,c} (1 + t_{20} T_{20}(\theta_i))$$

A fit of all counts to the last relation yields the (irrelevant) normalization constant k and the observable t_{20} . Δt_{20} is roughly inversely proportional to $\max |T_{20}(\theta_i) - T_{20}(\theta_j)|$. In our case, T_{20} varies from +0.1 to -0.55. In practice, one has to eliminate, or measure precisely, any instrumental asymmetries (ϕ dependence of ε) which, when expressed in terms of a Fourier expansion, give additive contributions to t_{20} .

4.2 Other Polarimeters

A new version of the low energy polarimeter was built at ETHZ 11): divided into (θ, ϕ) sectors, it also allows the determination of all components of the deuteron polarization. Above 250 MeV, there are, to our knowledge, no definite plans for building tensor polarimeters, but three possible ways to proceed are: a) degrade the deuteron energy to fall within the range of an AHEAD-like polarimeter (this has its limitations); b) use dp scattering at 180° ⁵⁾: T_{20} is very large, but the efficiency small; c) charge exchange (d,pp) ¹²⁾, which needs experimental investigation.

4.3 Polarimeter vs Polarized Target in ed Scattering

	Performed Experiments			Objectives			
	External Target		Internal Target	External Target		Polarized Internal Target	
	Unpolarized D ₂ O Bates 13)	Polarized ND ₃ Bonn-ELSA 14)	Atomic Beam Polarized Novosibirsk 15)	Unpolarized LD ₂ Bates 10)	Polarized ND ₃ Bonn-ELSA 14)	Storage Cell Optical Pumping + Spin Exchange Aladdin 16)	Goal for Bates 17) PSR 17)
Beam intensity	30 μ A	.33 nA	300 mA	50 μ A	7 nA	200 mA	80 mA
Target density (at/cm ²)	4.2×10^{21}	9×10^{22}	10^{11}	2.9×10^{23}	9×10^{22}	10^{14}	4×10^{15}
Luminescence (cm ⁻² s ⁻¹)	7.9×10^{35}	1.9×10^{32}	1.9×10^{29}	9.1×10^{37}	3.9×10^{33}	1.3×10^{32}	2×10^{33}
e solid angle (msr)	20	5	120	20	5	150	150
Losses	.5	-	-	.3	-	-	-
Polarimeter efficiency	10^{-4}	-	-	2×10^{-3}	-	-	-
Analyzing power t_{20}	- 0.8	-	-	- 0.4	-	-	-
Target polarization $A/\sqrt{2}$	-	.12	.57	-	.21	.6	.6
Figure of merit $F = L \Delta t_{20}^2$	5×10^{29}	1.4×10^{28}	7.3×10^{27}	1.8×10^{32}	8.7×10^{29}	6.8×10^{30}	1.1×10^{32}

At given incident electron energy and momentum transfer, the beam time necessary to achieve a statistical accuracy Δt_{20} is inversely proportional to the figure of merit F which appears in the above table. The new tensor polarimeter AHEAD, together with the use of a thick LD_2 target and the energy upgrade at Bates, appears to be in the near future the most efficient way to proceed. A polarized target experiment would fill the gap $2 < q < 3.7 \text{ fm}^{-1}$ not accessible to the low and intermediate energy polarimeters and could achieve smaller systematic errors if a precise method to measure the alignment of a polarized atomic target is available.

This work is the result of a collaboration between DPhN/ME and LNS at Saclay, and the University of Alberta. The author gratefully acknowledges support from the University of Alberta.

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QUARK-INTERCHANGE EFFECTS IN REACTIONS INVOLVING TWO-NUCLEON SYSTEMS

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ABSTRACT

Assuming that two-nucleon wave functions are antisymmetrized at the quark level, we expect that, for a reaction involving the two-nucleon system, there are reaction mechanisms which involve quark interchange [between the two nucleons] and cannot easily be rephrased in terms of only hadronic degrees of freedom. In this paper, we wish to demonstrate the importance of incorporating explicitly reaction mechanisms involving quark-interchange by considering specific numerical results related to nucleon-nucleon scattering and pion production via nucleon-nucleon scattering.

INTRODUCTION

Recent lattice QCD calculations¹ suggest possible existence of both the deconfinement and chiral-symmetry-restoration phase transitions. Thus, low-lying hadrons, as ground states of systems of quarks and gluons, may be described by an effective lagrangian with built-in spontaneous symmetry breaking, in a way similar to how we model superconductivity in QED. Once we describe low-lying hadrons as ground states with confinement and chiral-symmetry breaking built in via spontaneous symmetry breaking, we may assume, as suggested by the asymptotically free nature of QCD, that perturbative treatment of residual interactions is adequate. This implies that effects arising from one-gluon exchange [OGE] are more important than those due to multi-gluon exchanges. It also implies that degrees of freedom required to simulate chiral symmetry breaking, namely, (σ, π) Goldstone bosons, couple directly to quarks. Folding of such coupling

yields, among others, a πNN coupling so that the point tail of NN interactions follows. I believe such picture for low-lying hadrons provides us a reasonable starting point for investigating electroweak or nuclear reactions involving nuclei.

In the problem of nucleon-nucleon [NN] interactions, search for possible hidden-color dibaryons may present us a unique opportunity for establishing the role played by color, an important feature possessed by quarks but not by hadrons. Unfortunately, it is a very difficult task to formulate and make reliable predictions on the location and width of any such dibaryon, owing to possible model dependence as well as potential important coupling among the NN, $N\Delta$, $\Delta\Delta$, and hidden-color channels. For this reason, it may be useful to focus some attention on possible symmetry-violating effects associated with NN interactions, including isospin-symmetry violations or parity nonconservation. Symmetry violation studies in the NN problem may be unique for a couple of reasons: First, symmetry-violating effects are small and can be treated as a perturbation while ordinary [symmetry-conserving] NN interactions can be chosen as an empirical one such as given by a phase-shift analysis of scattering data or in a phenomenological potential. In this way, the potential important coupling among the NN, $N\Delta$, $\Delta\Delta$, and hidden-color channels can be avoided to some extent in making a theoretical prediction. Second, it is commonly accepted that symmetry-violating effects at the quark level in general give rise to effects of similar magnitude at the hadron level and predictions of this kind are often not very sensitive to which specific quark model has been used.

Using the rationale described earlier which links a certain version² of chiral bag models³ to the hadron world, we expect that reaction mechanisms illustrated in Fig. 1 constitute an important class of diagrams which may be of numerical significance in the NN problem. Here we note that one-pion exchange [OPE] or single

sigma-exchange [SSE] between two nucleon clusters must also be there following this line of reasoning. The remaining important feature is presumably inclusion of confining forces via formation of flux tubes, a characteristic common to several quark models including bag models⁴ and the flux-tube model⁵. I believe major uncertainty in this program has to do with treatment of nonlinearities inherent in any such model. Justification for neglecting nonlinearities can only come from a brute-force approach such as in a lattice QCD calculation.¹ I suspect that a clean picture concerning the role of quarks in nuclear physics may be obtained primarily in connection with symmetry violation studies.

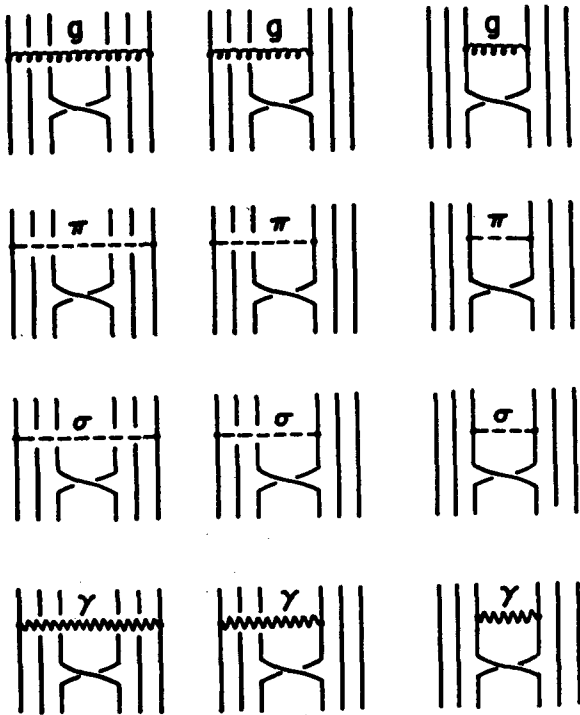


Fig. 1. Reaction mechanisms involving quark interchange for nucleon-nucleon scattering. Here we have assumed a chiral bag model in which both confinement and chiral symmetry breaking are built in via spontaneous symmetry breaking.

Analogously, if we assume that pion emission takes place effectively at the quark level and that quark-quark interactions are

described by one-gluon exchange [OGE] and one-pion exchange [OPE], then we are led to consider two different classes of diagrams for pion production via nucleon-nucleon scattering, i.e. $N + N \rightarrow N + N + \pi$. [Here we shall not consider explicitly single sigma-exchange, which can be incorporated in a similar way.] Typical examples are given in Figs. 2 and 3. Diagrams illustrated in Fig. 2 can be organized using

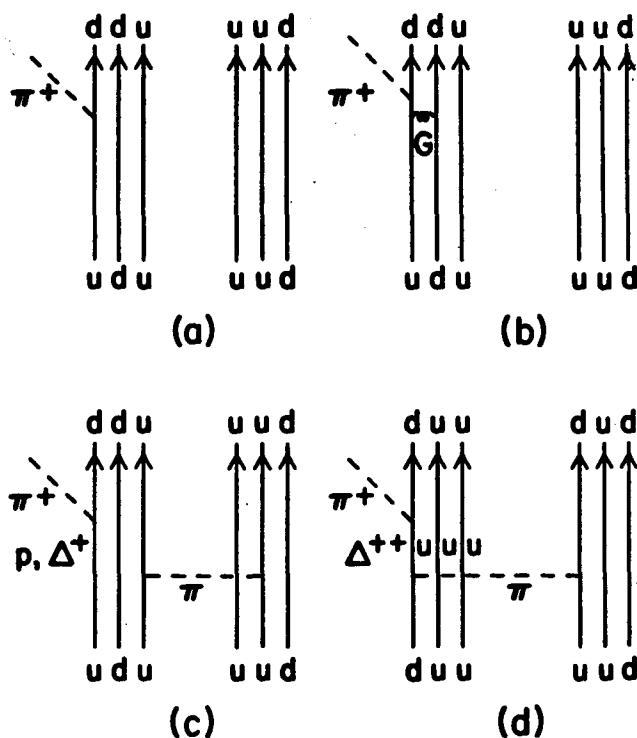


Fig. 2

organization scheme using ONM and TNM appears to be a natural one, although we do expect that the form factors, or the q^2 -dependence of the vertex functions, are no longer free parameters in the quark-model formulation. On the other hand, diagrams illustrated in Fig. 3 must

hadrons as the basic degrees of freedom, owing to the strong color-correlation requirement that hadrons are colorless. To argue further along this line, it is then an issue of great importance to decide which hadrons need to be included in the kinematic region of interest. For instance, Fig. 2(d) calls for inclusion of the (3,3) resonance [$\Delta(1232)$] in the intermediate state. Nevertheless, we should not anticipate anything dramatic for the first-class diagrams so that the standard

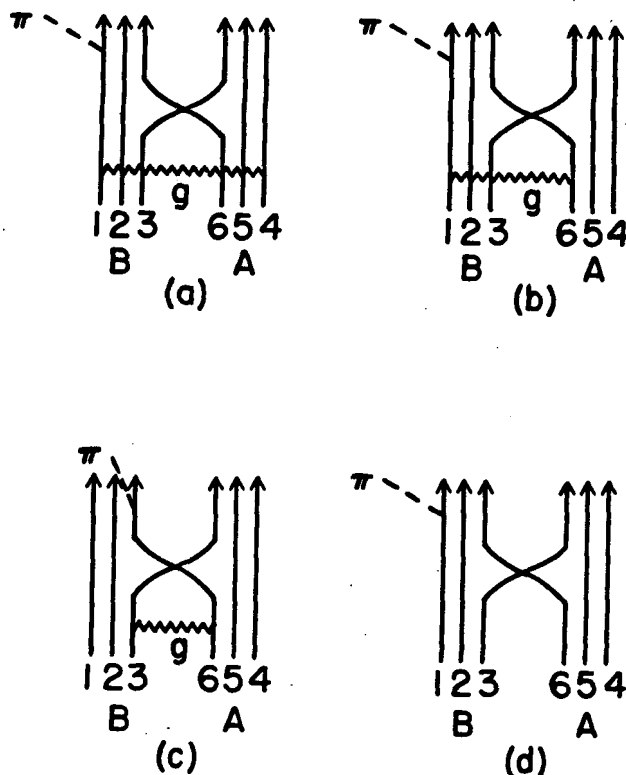


Fig. 3

In what follows, I wish to report some numerical results related to nucleon-nucleon scattering and pion production via nucleon-nucleon scattering, which demonstrate numerical importance of quark-interchange effects.

ISOSPIN SYMMETRY VIOLATIONS IN NUCLEON-NUCLEON SCATTERING AT LOW ENERGIES

Nucleon-nucleon [NN] scattering at low energies, say < 5 MeV, is dominated by the $^1S_0$ scattering, which can be characterized by the scattering length and effective range parameters $\{a, r\}$ in the well-known effective-range theory. Given a list of QI reaction mechanisms as in Fig. 1, we know how to evaluate each of them from lessons in treating a number of similar problems. For instance, we

be treated at the quark level, owing to quark interchange [QI] between the two interacting nucleons. This is where the formulation becomes highly nontrivial. In view of its complexities, we wish to refer an interested reader to ref. 6. for a detailed description of the technical procedure in treating these second-class diagrams [and the first-class diagrams as well].

need to do folding on electron-electron interactions in order to obtain the force between two atoms. The methodology is often referred to as the Born-Oppenheimer approximation. We also need to do folding on nucleon-nucleon interactions in order to find the interaction between two complex nuclei. The organization scheme has been formalized as the "resonating-group" [RG]⁷ or "generator-coordinate" [GC]⁸ method in nuclear physics. It has also been a subject of extensive investigations⁹ as to how the relativistic nature of the constituents can be treated properly, especially in connection with certain atomic-physics problems where the Dirac-particle nature of electrons may be relevant. In all cases, the cluster-cluster interaction in a specific configuration can be expressed in terms of the constituent-constituent interaction evaluated between the various initial and final wave functions of the constituents. To illustrate this point, we quote¹⁰, in the limiting case with $m_d = m_u$,

$$\begin{aligned} & \langle p(\uparrow)p(\uparrow) | \epsilon_{14}^g \hat{P}^{\tau\sigma X} | p(\uparrow)p(\uparrow) \rangle \\ &= \frac{1}{324} \eta^2 \left\{ 72f_{11}^g + 52f_{33}^g + 10f_{44}^g \right\}. \end{aligned} \quad (1)$$

Here η is the overlap integral defined by

$$\eta \equiv \langle Q_B(\uparrow) | Q_A(\uparrow) \rangle, \quad (2)$$

where $Q_A(\uparrow)$ [$Q_B(\uparrow)$] denotes a $d(\uparrow)$ quark sitting inside cluster A [cluster B]. Note that we use Q to label either a u or d quark. In addition, the various QQ amplitudes are specified as follows:

$$f_1^g = \langle Q_A(+)|Q_B(+)|\varepsilon_{14}|Q_A(+)|Q_B(+)\rangle, \quad (3a)$$

$$f_3^g = \langle Q_A(+)|Q_B(+)|\varepsilon_{14}|Q_A(+)|Q_B(+)\rangle, \quad (3b)$$

$$f_4^g = \langle Q_A(+)|Q_B(+)|\varepsilon_{14}|Q_A(+)|Q_B(+)\rangle. \quad (3c)$$

In a independent-particle model where quarks are treated as Dirac particles, these amplitudes are determined by¹⁰

$$\begin{aligned} f_1^{gab}(\underline{r}) = & \int d^3r_1 d^3r_4 \int \frac{d^3q}{(2\pi)^3} \exp\{i \underline{q} \cdot (\underline{r}^1 - \underline{r}^4)\} \cdot \\ & \cdot \Psi_a^\dagger(\underline{r}_A^1) S^\dagger(\underline{v}) \gamma_4 \gamma_\mu S(\underline{v}) \Psi_a(\underline{r}_A^1) \cdot \\ & \cdot g^2 i D_{\mu\nu}(\underline{q}) \cdot \Psi_b^\dagger(\underline{r}_B^4) S^\dagger(-\underline{v}) \gamma_4 \gamma_\nu S(-\underline{v}) \Psi_b(\underline{r}_B^4). \end{aligned} \quad (4)$$

Here $S(\underline{v})\Psi(\underline{r}_A^1)$ denotes the quark wave function in a cluster of velocity $\underline{v}$ as seen in the CM frame so that both the standard Lorentz transformation $S(\underline{v})$ for a Dirac spinor and a transformation on the relative coordinate $\underline{r}_A^1 [\equiv \underline{r}^1 - \underline{r}_A^1]$ are implied.¹¹ Note that $iD_{\mu\nu}(\underline{q})$ is the gluon propagator and $g^2/(4\pi) [\equiv \alpha_S]$ is the QCD coupling constant. We also note that $\underline{r}_A^1 [\underline{r}_B^4]$ is the coordinate of the first [fourth] quark expressed to the center of the A [B] cluster.

To study isospin symmetry violations, we must take into account the difference between ϕ_u and ϕ_d so that Eqs. (1)-(4) becomes considerably more complicated. Although we may obtain easily the n - n matrix elements from the corresponding p - p matrix elements by

exchanging u and d , we need to work a little harder to obtain the various n - p matrix elements when the d quark is treated differently from the u quark. The final formulae are elucidated in ref. 10, to which I wish to refer an interested reader for the details. An important aspect is that, although there are a large number of formulae similar to Eq. (1) which are relevant for isospin symmetry violation studies, it is a straightforward task to evaluate six-dimensional integrals, such as those given by Eq. (4) with the integration over d^3q trivially carried out, using a Monte-Carlo technique.¹²

For numerical results to be described below, we use the quark wave function which is of the form suggested by the MIT bag model⁴ but expressed with respect to the "center" of the three-quark [baryon] cluster. We assume that, at a given instant, the two baryon clusters move relative to each other with some instantaneous velocity. Recoil effects have been incorporated in the way suggested in the literature.¹¹ We have sorted out numerically all possible QI terms corresponding to diagrams illustrated by Fig. 1. Pauli blocking effects as described by the "norm" and "kinetic-energy" kernels in the resonating-group equation are also incorporated in a way described in ref. 10.

The input parameters in this calculation are obtained primarily from an overall fit to the observed mass splittings of low-lying baryons.¹³ For a nucleon radius R of 0.8 fm, we have

$$\begin{aligned} \alpha^*_S &= 1.97, & \alpha^*_\pi(r=0) &= 0.310, \\ m_u &= 7.12 \text{ MeV}, & m_d &= 12.88 \text{ MeV}. \end{aligned} \tag{5}$$

Note that we insist on reproducing the well-known Yukawa pionic tail for two nucleons in separation by choosing a linear interpolation to obtain $\alpha^*_\pi(r)$. The mass of σ is fixed to be 520 MeV, a value used in

the phenomenological Bonn potential.¹⁴ The σ_{NN} couplings are first calculated from the given σ_{QQ} coupling using the same quark wave function but later rescaled to the Bonn value [$g_{opp} = 5.650$]¹⁴ or a value [$g_{opp} = 4.540$] which reproduces the slope of the $^1S_0$ phase shift at zero energy and so is used in calculating the low-energy scattering parameters.

To obtain the scattering length a for low-energy NN scattering, we adopt the well-known effective-range formula,

$$k \cot \delta = -\frac{1}{a} + \frac{1}{2} r k^2. \quad (6)$$

We determine the phase shift δ at two different CM kinetic energies [chosen to be 1 MeV and 2 MeV] so that Eq. (6) can be inverted to obtain both a and r . As indicated earlier, the isospin-symmetry violating potential is treated as a perturbation. This has been accomplished by two different means: In the first method, we compute QI diagrams twice, first with $m_d \neq m_u$ and $m(\pi^+) \neq m(\pi^0)$ and then with $m_d = m_u$ and $m(\pi^+) = m(\pi^0)$, and then obtain the perturbing potential by taking the differences. In the second method, QI diagrams are evaluated only once but we take the differences $V_{nn} - V_{pp}$ and $V_{np}(I=1) - V_{pp}$ as our perturbing potentials in the n - n and $I=1$ n - p channels. Both methods must yield numerical identical results, as they should, so that we use the first method to eliminate any error in all the derived formulae similar to Eq. (1) but with $m_d \neq m_u$ as well as in the computer code. Thus, the second method can be safely adopted so that, with the Reid $^1S_0$ soft-core potential as the input and with Coulomb force explicitly subtracted, our calculated a_{pp} is fixed to be - 17.0324 fm.

In Table I immediately below, I list the results on the scattering length parameters:

TABLE I: The scattering length parameters in the various $^1S_0$ channels. The experimental values for a_{pp} and a_{np} are from ref. 15 while the quoted a_{nn} is from Gabioud *et al.*¹⁶ Note that the Coulomb force in the p-p channel is already subtracted.

	R = 0.8 fm	R = 1.0 fm	experiment
a_{pp}	-17.03 fm	-17.03 fm	-17.1 ± 0.2 fm
a_{np}	-21.16 fm	-21.60 fm	-23.75 ± 0.01 fm
a_{nn}	-18.22 fm	-19.21 fm	-18.45 ± 0.45 fm

It is clear that the prediction for $R = 0.8$ fm compares very well with the experimental result, especially if we add the additional contribution -1.1 ± 0.4 fm arising from the simultaneous $(\pi\gamma)$ two-boson exchange diagram.¹⁷ For $R = 0.8$ fm [1.0 fm], QI diagrams altogether [with Pauli blocking included] contribute -0.11 fm [-0.75 fm] to $|a_{np}| - |a_{nn}|$ and -0.21 fm [-1.23 fm] to $|a_{pp}| - |a_{nn}|$. As a well-known result¹⁸, the remaining contribution to $|a_{np}| - |a_{nn}|$ as listed in Table I comes from "direct" one-pion exchange with $m(\pi^+) \neq m(\pi^0)$. On the other hand, "direct" single sigma-exchange [between the two nucleon clusters] is responsible for most of the remaining contribution in the listed result on $|a_{pp}| - |a_{nn}|$.

In passing, we wish to note that, with $g_{opp} = 5.650$ [the Bonn value]¹⁴ and $R = 0.8$ fm and all remaining parameters fixed as described earlier, the same simple one-boson-exchange model [at the quark level] reproduces very well the $^1S_0$ phase shift for the CM kinetic energy from 10 MeV to 150 MeV [a kinematic region of interest in most of the phenomenological potential]. We refer an interested reader to ref. 10 for further details.

It is clear that the results summarized in Table I are very encouraging. It should be useful to investigate how nucleons behave in nuclear medium, especially in light of the EMC effect¹⁹ which suggests that the nucleon structure gets modified in nuclear matter. Several competing interpretations²⁰ have been proposed for this effect, including an increase in nucleon size or the relevant scale. Our intention here in this paper is not to advocate one picture over the others but, instead, is to point out possible existence of other EMC-type effects in nuclear physics. One important candidate which I wish to address is the well-known Nolen-Schiffer anomaly²¹ that appears consistently for the binding energy difference in mirror nuclei. Taking the ^{41}Ca - ^{41}Sc mirror pair as an illustrative example, Negele²² has shown for some time that Coulomb effects and their manifestations as many-body effects cannot account for the entire ^{41}Ca - ^{41}Sc binding energy difference and, in addition, he has also shown that presence of a charge-symmetry-breaking force of the size that may give rise to $a_{nn} - a_{pp} = -0.8$ fm is capable of removing the anomaly completely.

In light of the listed results in Table I, I suspect that Negele's suggestion should be taken seriously as a simple way to understand Nolen-Schiffer anomalies. Allowing the nucleon size to increase by about 10% in nuclear medium, some authors argues that the EMC effects can be understood. For an isolated nucleon radius R_N of 0.8 fm, a 10% increase in size has at least three important consequences, viz.:

- (1) The neutron-proton mass difference in nuclear medium is larger than the free-space value, i.e. 1.293 MeV, by 70 keV. This effect works against us in terms of understanding the Nolen-Schiffer anomaly.
- (2) The prediction on $a_{nn} - a_{pp}$ which is - 1.2 fm in free space will increase in magnitude to - 1.5 fm in nuclear matter. Both the free-space and nuclear-matter results are more than adequate to resolve the Nolen-Schiffer anomaly.

(3) The delta-nucleon mass difference in nuclear medium is smaller than the free-space value, i.e. 294 MeV, by 27 MeV.

Note that we have obtained (1) and (3) making use of the procedure described in ref. 11 while (2) follows from Table I given earlier. If we combine Items (1) and (2) above with Negele's suggestion, we find that, as long as the increase in nucleon size is not much worse than 10%, there is not a serious problem in understanding the Nolen-Schiffer anomaly.

An important reason which I mention Item (3) is due to a recent experiment of O'Connell et al.²³, in which the Δ -peak position in deep inelastic electron scattering off complex nuclei is seen to move downward relative to that in electron-proton scattering by about 10 MeV. Recent calculations²⁴ basing on the nucleon-plus-delta-only picture yield an upward shift of about 20 MeV. Combining the result of a conventional many-body calculation with Item (3) above, we find that a net downward shift of about 10 MeV comes out in a simple way.

Besides all these interesting observations, there are yet other important lessons to learn. Firstly, the proposed increase in nucleon size in nuclear matter is better within 10% in order to avoid to do fine-tuning in explaining both the Nolen-Schiffer anomaly²¹ and the experiment of O'Connell et al.²³ Secondly, before arguing that explicit incorporation of the quark degree of freedom is essential for the typical nuclear physics phenomena addressed in the present paper, we must ask whether a very different model such as the Skyrmion picture²⁵ may also provide a simple interpretation of the otherwise-mysterious experimental results ranging from differences in the $^1S_0$ nucleon-nucleon scattering length parameters^{15,16}, Nolen-Schiffer anomalies²¹, and the experiment of O'Connell et al.²³ I wish to leave this important topic to an interested reader to think about. It is clear that the naive model used in the present paper does tie together all the experimental results in a simple manner and more quantitative nuclear structure calculations appear to be an important next step.

To summarize this section, we have investigated potential important role played by interactions involving quark interchange in the determination of the scattering length parameters a_{pp} , a_{np} and a_{nn} relevant for low-energy p-p, n-p, and n-n $^1S_0$ scattering, respectively. The quark-quark interactions being considered include one-gluon exchange, one-pion exchange, single sigma-exchange, and one-photon exchange. We have found that, for a nucleon radius R of 0.8 fm with the up-down quark mass difference as determined by an overall fit to the observed baryon mass splittings, the predicted charge-symmetry-breaking effect via the difference $|a_{pp}| - |a_{nn}|$ is in accord with the recent experimental result of Gabioud *et al.*¹⁶ while the charge-dependent effect as reflected by $|a_{np}| - |a_{nn}|$ may also be understood. We note that the model may also provide a simple scenario for understanding Nolen-Schiffer anomalies²¹ in complex nuclei.

ROLE OF QUARK INTERCHANGE IN PION PRODUCTION VIA NUCLEON-NUCLEON SCATTERING

As already explained earlier in the introduction, reaction mechanisms for $N + N \rightarrow N + N + \pi$ may then be grouped into several distinct categories, viz.: (i) one-nucleon-mechanism [ONM] or the nucleon-only impulse approximation [NOIA] where a pion is emitted from one of the two nucleon clusters; (ii) two-nucleon-mechanisms [TNM] or meson-exchange currents [MEC] where pion emission takes place while the two nucleon clusters exchange a virtual pion, and here we include isobar[Δ] currents, ρ - π currents, and pair currents; (iii) reaction mechanisms involving quark interchange [QI] with or without one-boson exchange between the two nucleon clusters. Describing quarks as Dirac particles moving in a harmonic-oscillator-type confining potential, we have been able to simplify partial-wave transition amplitudes due to QI reaction mechanisms from 15-dimensional integrals into 6-dimensional ones.⁶ These 6-dimensional integrals are then carried out by a Quasi Monte-Carlo method.¹² We note that, unlike the work of Miller and Kisslinger²⁶ where ONM and TNM are yet to be incorporated

on the same footing, we have in fact taken into account all categories of reaction mechanisms in this calculation. Thus, although our results provide a quantitative support of their observation²⁶ concerning potential importance of quark effects, it is clear that our conclusion has been obtained in a considerably more consistent manner. However, we do not know at this point how to make a precise link to the approach^{27,28} in which one attempts to solve the Faddeev-like three-body equations which treat the NN and NN π systems as coupled channels, since, by treating quarks explicitly, we have to deal with a (6+1)-body problem. Of course, the sophisticated Faddeev-type approach can also be an approximation to reality since inclusion of certain form factors may spoil the exact-solvability of the problem.

Before presenting our major numerical results, we wish to also emphasize that the classification mentioned here is a phenomenological one since it is known that inclusion of pion distortions is of numerical importance and the pion-deuteron optical potential, as chosen by us to describe pion distortions, already contains the resonance behavior due to isobars [$\Delta(1232)$] and others. In practice, we choose the pion-deuteron optical potential which provides the best fit [in terms of the χ^2] the data at several incident pion energies. Thus, to avoid double counting with contributions from isobar currents, we do not include pion distortions in the ONM or NOIA contributions. In this way, we are able to carry out a parameter-free calculation for $p + p \rightarrow d + \pi^+$. Our results support the major conclusion reached by Miller and Kisslinger²⁶ that, for a nucleon radius of greater than 0.8 fm, quark effects, which are identified in our formalism as contributions arising from reaction mechanisms involving quark interchange, play an important role even near threshold [where their argument regarding the large momentum transfer remains valid]. Our study of the $n + p \rightarrow p + p + \pi^-$ reaction using the π -d optical potential to describe pion distortions also points to numerical importance of quark effects.

We show in Figs. 4 and 5 the predicted cross section and analyzing power, respectively, for the reactions $p + p \rightarrow d + \pi^+$ and $n + p \rightarrow p + p + \pi^-$. In making these predictions, we have used an energy-dependent isobar width obtained experimentally,²⁹

$$\Gamma_{\Delta} = (0.35/6) \cdot (q^{*3}/m_{\pi}^2) \cdot [(M^* + M)^2 - m_{\pi}^2] \cdot M^{*-2},$$

$$\text{with } q^* = q_{\pi} \cdot \left(1 - \frac{E_{\pi}}{2(M+E_{\pi})}\right).$$

The nucleon radius parameter R , the strong coupling α_S^* , and the effective pion-nucleon coupling α_{π}^* are chosen as follows¹³:

$$R = 0.57154 \text{ fm [corresponding to } R_{MIT} = 0.8 \text{ fm]},$$

$$\alpha_S^* = 1.97,$$

$$\alpha_{\pi}^*(r = 0) = 0.31 \text{ and } \alpha_{\pi}^*(r = \infty) = 3.572.$$

Note that the above nucleon radius parameter in our Gaussian-type wave function corresponds to a bag radius of 0.8 fm in the MIT model.⁴ It is important to keep in mind that all the parameters in our calculation are determined by other means and none of them have been fine-tuned to best fit the observed data related to the reaction $p + p \rightarrow d + \pi^+$.

In Fig. 4, the calculated cross section for the pion-absorption reaction $\pi^+ + d \rightarrow p + p$ and the analyzing power A_y for the reaction $p + p \rightarrow d + \pi^+$ at $T_p = 180$ MeV are shown respectively as a function of $\eta [= q_{\pi}/m_{\pi}]$ and as a function of the pion scattering angle θ_{CM} [in the CM frame]. The experimental data points are obtained from ref. 30. Here we compare the prediction calculated with quark-interchange [Q.I.] reaction mechanisms with that calculated without Q.I. [NO Q.I.]. The p - p partial waves included in the calculations are $^3P_1$, $^1S_0$, and $^1D_2$. Although importance of Q.I. varies with the kinematic

region, this result alone already confirm the observation made by Miller and Kisslinger²⁶ that effects caused by short-distance quark-interchange mechanisms need to be considered even near threshold.

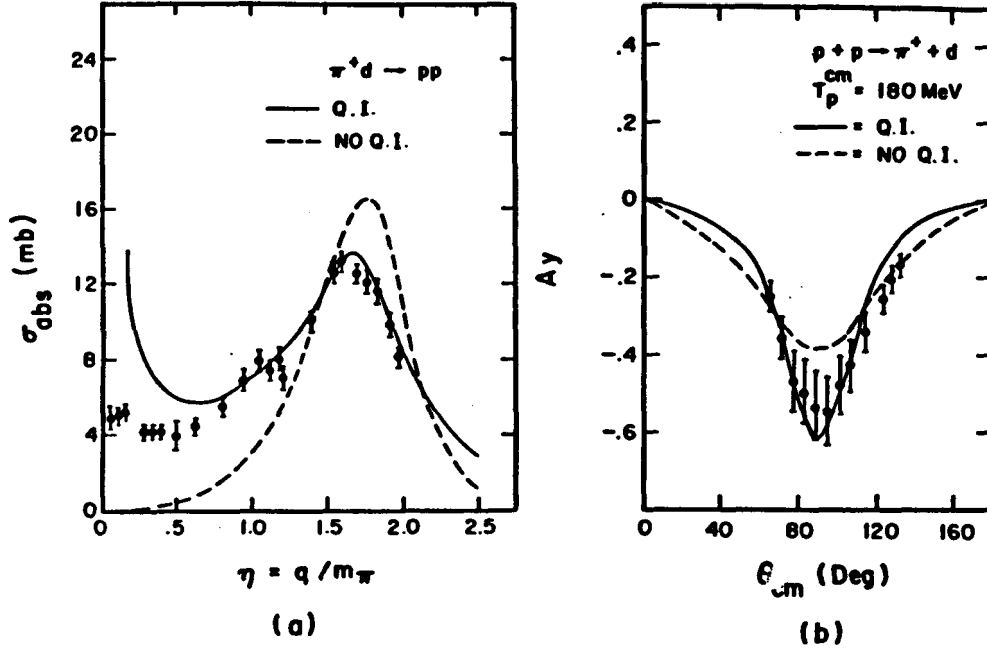


Fig. 4

In Fig. 5, the predicted double differential cross section $d^2\sigma/dE_\pi d\Omega_\pi$ and analyzing power A_y for the reaction $n + p \rightarrow p + p + \pi^-$ at the p-p excitation energy $E_x = 4$ MeV and the pion emission angle $\theta_\pi = 30^\circ$ are both shown as the initial proton energy T_p^{cm} . Specifically, we have included the following channels in the calculations: $^3P_0 \rightarrow ^1S_0$, $^1S_0 \rightarrow ^3P_0$, $^3S_1 + ^3D_1 \rightarrow ^3P_1$, $^3D_2 \rightarrow ^3P_2$ [all for S-wave pions], and $^3S_1 + ^3D_1 \rightarrow ^1S_0$, $^3P_1 \rightarrow ^3P_0$, $^3P_J \rightarrow ^3P_1$, $^3P_J \rightarrow ^3P_2$ [all for P-wave pions] for ONM and TNM; $^3P_0 \rightarrow ^1S_0$ and $^3S_1 + ^3D_1 \rightarrow ^1S_0$

for the various quark-interchange [Q.I.] reaction mechanisms. Once again, we see that effects due to quark-interchange reaction mechanisms become of numerical importance and should be included even near the threshold. Note, however, that the optical potential for the π -(pp) system may differ significantly from the π -d optical potential. Results from a further investigation of this aspect, as reported in ref. 6, suggest that our predictions on the reaction $n + p \rightarrow p + p + \pi^-$ are not modified severely by adoption of a rather different π -(pp) optical potential.

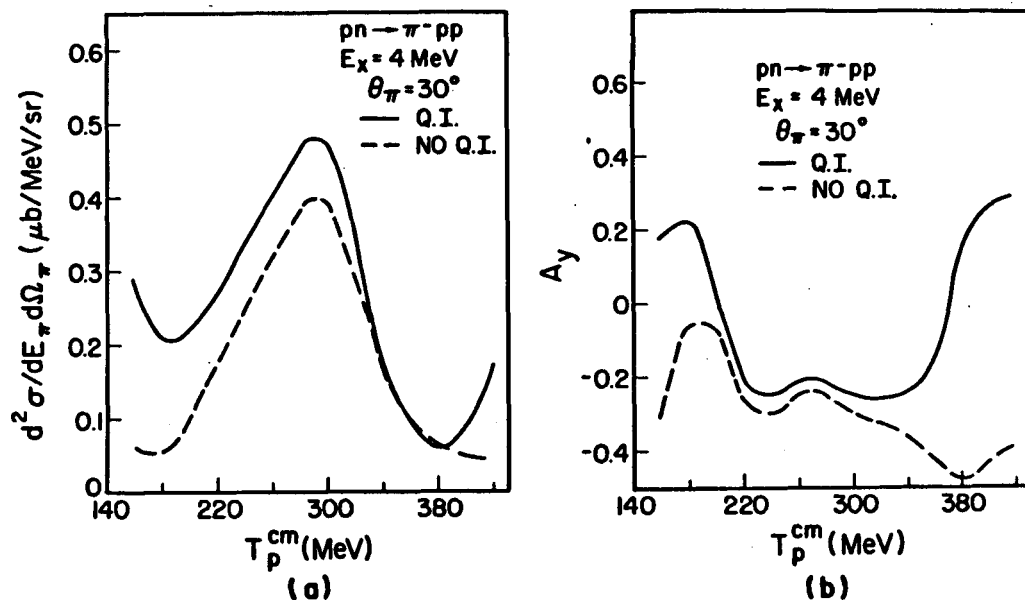


Fig. 5

To summarize this section, we have investigated pion production via nucleon-nucleon scattering, i.e. $N + N \rightarrow N + N + \pi$, assuming that pion emission takes place effectively at the quark level and that

quark-quark interactions are described to a sufficient approximation by one-gluon exchange plus effective one-pion exchange. The six-quark wave function is totally antisymmetrized. As a result, reaction mechanisms may be grouped into several categories: (1) one-nucleon mechanism [ONM] or the nucleon-only impulse approximation [NOIA], (2) two-nucleon mechanisms [TNM] or meson-exchange currents [MEC], and (3) reaction mechanisms involving quark interchange [QI]. Our numerical predictions for the reactions $p + p \rightarrow d + \pi^+$ and $n + p \rightarrow p + p + \pi^-$ indicate that, for a nucleon radius of greater than 0.8 fermi [as measured in the MIT bag model], reaction mechanisms involving quark interchange are as important as the conventional one-nucleon or two-nucleon reaction mechanisms even near the pion production threshold. This specific result provides a quantitative support of the conclusion reached by Miller and Kisslinger²⁶ in the hybrid quark-baryon model concerning the necessity of incorporating quark effects in describing pion production or absorption reactions.

CONCLUDING REMARK

Assuming that two-nucleon wave functions are antisymmetrized at the quark level, we expect that, for a reaction involving the two-nucleon system, there are reaction mechanisms which involve quark interchange between the two nucleons and cannot easily be rephrased in terms of only hadronic degrees of freedom. In this contribution, I have reported some numerical results related to nucleon-nucleon scattering and pion production via nucleon-nucleon scattering, which demonstrate numerical importance of quark-interchange effects. I believe that these results shed some light on the importance of formulating certain nuclear physics problems at the quark level.

ACKNOWLEDGEMENT

This work is supported in part by the US Department of Energy and in part by the National Science Foundation.

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A Ratio Technique to Measure the Quadrupole
Contribution to the $N \rightarrow \Delta$ Transition Using the
 $(e, e' \gamma)$ and $(e, e' \pi^0)$ Reactions

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ABSTRACT

A technique is proposed for measuring the quadrupole contribution to the $N \rightarrow \Delta$ transition. The method involves simultaneous measurement at 2 or more angles of the reaction $p(e, e' p \gamma)$. This is shown to minimize the effects of theoretical and experimental uncertainties on the analysis.

The spherically symmetric quark model predicts that the $N(938) \rightarrow \Delta(1232)$ transition is a pure M1 transition. Techniques previously used for measuring C_2 and E_2 in the $N \rightarrow \Delta$ transition include (γ, π) ,¹ (e, e') ,² and the exclusive reactions $(e, e' \pi^+)$ and $(e, e' \pi^0)$ ³ ($(e, e' \gamma)$ has never been observed). These measurements have generally suffered from large systematic errors (>5%) and analysis problems in separating resonant from non-resonant amplitudes.³ Any new measurement must minimize theoretical uncertainties in the non-resonant background and must find ways to limit systematic error to a much higher degree than has been done previously.

It is possible to measure both $(e, e' \gamma)$ and $(e, e' \pi^0)$ by doing an $(e, e' p \gamma)$ experiment.³ The advantage to measuring both reactions is that the non-resonant backgrounds of the 2 processes are of a different nature. This provides a constraint on the analysis used to extract the resonant amplitudes from the data. The problem in performing an $(e, e' \gamma)$

measurement is in distinguishing photons from $\Delta \rightarrow p + \gamma$ (.6% decay Prob.) from $\Delta \rightarrow p + \pi^0 \rightarrow p + \gamma + \gamma$ (66%). The photon energy difference between the 2 decays is typically only 10 MeV out of 250 MeV. However, by measuring the direction and momentum of the recoil proton it is possible to distinguish $\Delta \rightarrow p + \gamma$ from $\Delta \rightarrow p + \pi^0$. The solid curve in Figure 1 shows how the proton momentum from $\Delta \rightarrow p + \gamma$ varies with photon angle ($E_0 = 750$ MeV, $\theta_e = 30^\circ$, $\omega = 386$ MeV). The dashed curve shows the proton momentum from $\Delta \rightarrow p + \pi^0$, where the proton angle is fixed from the $\Delta \rightarrow p + \gamma$ kinematics. The gap between the 2 decays is typically $\Delta p/p = 5\%$, or $\Delta E/E = 10\%$ at 200 MeV. Even at CEBAF energies with $E_0 = 4$ GeV the momentum gap is still approximately 5%. A proton spectrometer of only moderate resolution could be used to resolve the 2 processes. The main purpose of the gamma detector would be to reduce accidental coincidences between the electron and proton spectrometers. Lead glass is a good candidate for the gamma detector because it has good timing and little sensitivity to low energy backgrounds.

Figure 2 shows the expected photon angular distribution for $\Delta \rightarrow p + \gamma$ for varying ratios of C2/M1. It is readily apparent that the sensitivity to C2 is not large. Because of this low sensitivity any precision measurement of the angular distribution will always be limited by systematic error when the points of the angular distribution are measured sequentially. What must be done is to measure 2 or more angles simultaneously and then form ratios of yields. By forming a ratio much of the systematic error cancels out and accuracy is limited largely by statistics. A typical ratio to be measured in these kinematics is the ratio of yields at photon angles of 120° and 200° . Figure 3 shows how this ratio varies as a function of the C2 content of the transition. At these angles the sensitivity to C2/M1 is approximately equal to the precision of the yield ratio. For example, to measure C2/M1 to the 1% level the total error in the yield ratio cannot exceed 1%. This ratio technique has been recently applied to the 2^+ state at 4.4 MeV in ^{12}C .⁴ In this transition a transverse-longitudinal interference term as small as $E2/C2 = 7\%$ was successfully observed.

Count rate estimates for $p(e,e'\gamma)$ (assuming typical spectrometer acceptances, beam currents and target thicknesses) gives true rates of greater than 1 per minute,³ making the proposed technique quite attractive. With the 3-fold coincidence and 100% duty factor, the accidental rate should be very small. This experiment could run at high beam currents, limited only by target heating effects.

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Fig. 1

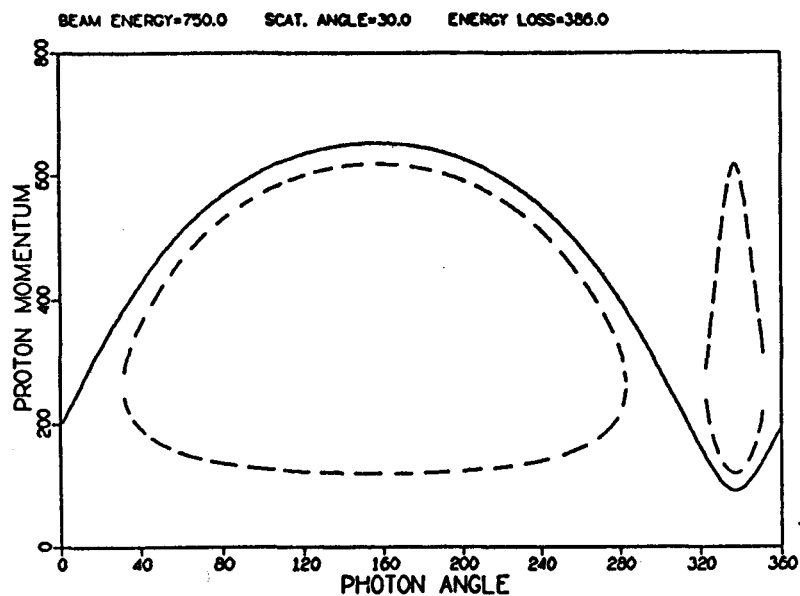


Fig. 2

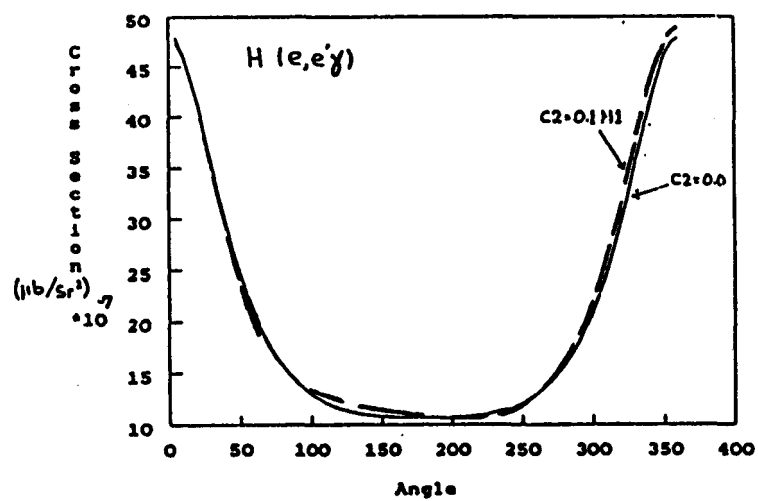
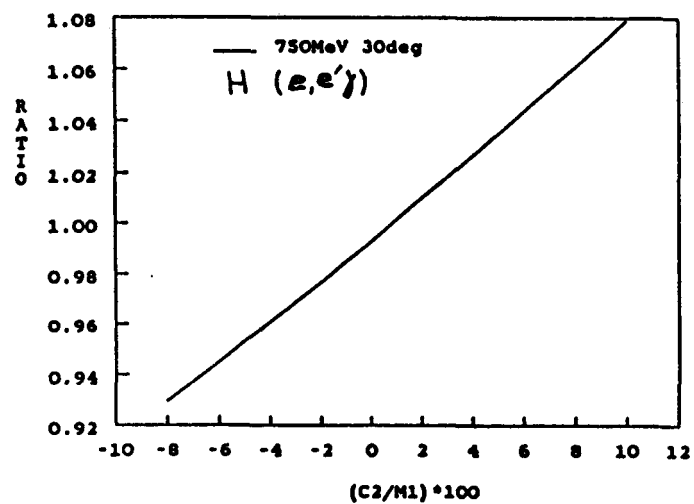


Fig. 3



Contributed Paper 1986 CEBAF Summer Workshop

Structure Functions and
Quark Distributions in Nuclei

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ABSTRACT

A relativistic Fermi Motion model for nuclear structure functions is presented. Comparison with experimental data, and predictions for the quasi-elastic ($x \sim 1$) and $x > 1$ regions are given. The same model allows us to obtain quark distributions in several nuclei, which can be used in many high energy nuclear reactions.

If we know the probability $G_{a/A}(y, \vec{k}_T)$ of finding a constituent of type "a" in a nucleus "A" with fractional (light-cone) momentum y and transverse momentum $\vec{k}_T$, then scattering interactions of that nucleus may be reliably predicted for sufficiently large q^2 , where q is the momentum transferred. This is due to the fact that final-state interactions (higher twist) fall-off as inverse powers of q^2 .

In reference (1), a model for $G_{a/A}(y, \vec{k}_T)$ was constructed when "a" = nucleon.

$$G_{a/A}(y, \vec{k}_T) = \frac{cy^h(1-y)^g}{\left[k_T^2 + M_O^2(y)\right]^2 \left[k_T^2 + M_O^2(y) + \delta^2\right]^{g-1}} \quad (1)$$

where $M_O^2(y) = (1-y)M_N^2 + yM_\alpha^2 - y(1-y)M_A^2$

$$g = 6A-7 \quad (2)$$

$$\delta^2 = .2A$$

$$M_N = \text{nucleon mass}$$

$$M_A = \text{mass of nucleus "A"}$$

$$M_\alpha = \text{mass of nucleus "A-1"}$$

h varies from nucleus to nucleus and can be determined by electron scattering. In Table 1, h is given for some nuclei. C is a normalization constant which must be fixed by the conservation of momentum

$$\sum_a \int dy d^2k_T y G_{a/A}(y, \vec{k}_T) = 1 \quad (3)$$

The nucleus inelastic structure function, $F_{2A}^I(X_A, q^2)$ where $X_A = -q^2/2M_A\nu$, may then be calculated from Figure 1, which gives

$$F_{2A}^I(X_A, q^2) = \sum_a \int_0^1 dy \int_{-\infty}^{+\infty} dk_1 \int_{-\infty}^{+\infty} dk_2 G_{a/A}(y, k_1, k_2) F_{2a}(-q^2/2k \cdot q) \quad (4)$$

where the constituent (nucleon) structure function is F_{2a} . $a =$ proton, neutron, .. Here $k_T^2 = k_1^2 + k_2^2$ and the region of integration is restricted to

$$k_T^2 + M_0^2(y) - \frac{Q^2}{x_A} y(1-y) + Q^2(1-y) + 2k_T \sqrt{Q^2} (1-y) < 0 \quad (5)$$

$$\frac{Q^2}{2k \cdot q} = x_A / (y - (2k_T x_A / (\sqrt{Q^2}))) \quad (6)$$

We define $Q^2 = -q^2$ and $Z = 2x_A / (1 + \text{SQRT}(1 + \frac{M_A^2 x_A^2}{Q^2}))$.

Near $x_A \sim 1/A$, there will be a quasi-elastic contribution, which may be calculated from Figure 2, with the result

$$F_{2A}^Q(X_A, q^2) = \sum_a \int_{Y_1}^{Y_2} y dy \pi G_{a/A}(y, k_T^2 = K^2) F_a^2(q^2) \left[\frac{Q^2}{x_A} \frac{1-y}{1-Z} \right] \quad (7)$$

where $F_a(q^2)$ is the nucleon form factor. The total structure function for a nucleus is

$$F_{2A}(X_A, q^2) = F_{2A}^Q(X_A, q^2) + F_{2A}^I(X_A, q^2) \quad (8)$$

K^2, Y_1, Y_2 are given by

$$K^2 = \frac{1}{1-Z} \left\{ (Y-Z) \left[-\frac{Q^2}{Z} (1-Y) + M_A^2(1-Y) - M_\alpha^2 \right] - M_N^2(1-Y) \right\} \quad (9)$$

$$Y_1 = C_0 \mp \text{SQRT} \left\{ C_0^2 - Z \left(1 - \frac{M_\alpha^2}{\frac{Q^2}{Z} + M_A^2} \right) - \frac{M_N^2}{\frac{Q^2}{Z} + M_A^2} \right\} \quad (10)$$

where $C_0 = \frac{1}{2}(1+Z-(M_\alpha^2-M_N^2)/(\frac{Q^2}{Z}+M_A^2))$. As an example, we give in Table 2 the prediction for $(\tilde{F}_2)_{\text{carbon}}/(\tilde{F}_2)_{\text{helium}}$ versus the variable $x_p =$

$(X_A)A$. This is a continuation of the "EMC" effect into $X_p > 1$. $\tilde{F}_2 \equiv \frac{1}{A} F_2$.

Finally, having an estimate for $G_{a/A}(y, \vec{k}_T)$ allows one to compute directly the probability of finding a parton in a nucleus with a fraction X_A of the total nucleus' momentum. This probability is the convolution of the probability of finding a nucleon in a nucleus with fractional momentum y , with the probability of finding a parton in a nucleon with fractional momentum X_A/y . For $N=Z = A/2$, the up or down valence momentum distribution is

$$X_A G_{uV/A}(X_A, q^2) = \int_{X_A}^1 dy G_{a/A}(y) \frac{X_A}{y} \left[U_V\left(\frac{X_A}{y}\right) + d_V\left(\frac{X_A}{y}\right) \right] \quad (10)$$

correct, in the large q^2 limit. Here we have defined

$$G_{a/A}(y) = \int d^2k_T G_{a/A}(y, \vec{k}_T) \quad (11)$$

$zU_V(z)$ and $zd_V(z)$ are the up and down valence distributions in a proton. The sea in a nucleus is

$$X_A S(X_A, q^2) = \sum_a \int_{X_A}^1 dy G_{a/A}(y) \frac{X_A}{y} S\left(\frac{X_A}{y}\right) \quad (12)$$

where $zS(z)$ is the sea in a nucleon. If we use the parton distributions of Buras, reference (2), we obtain Tables 3, 4 for $Q^2 = 7.86 \text{ GeV}^2/c^2$. Knowing $X_A G_{p/A}(X_A, q^2)$ for valence quarks, gluons and sea in a nucleus allows one to extend to nucleus interactions the same QCD techniques used in hadron interactions.

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TABLE AND FIGURE CAPTIONS

Table 1. h for selected nuclei.

Table 2. Prediction of the "EMC effect" for $X_p > 1$. $Q^2 = 1.8 \text{ GeV}^2/c^2$.

Table 3. Parton distributions in deuterium

Table 4. and Helium. ($Q^2 = 7.86 \text{ GeV}^2/c^2$)

Figure 1. The inelastic structure function of a nucleus, for high q^2 .

Figure 2. The quasi-elastic structure function of a nucleus, for high q^2 .

Table 1

<u>Nucleus</u>	<u>h</u>
D	5
He	4
Be	2.5
C	3
Al	2
Fe	0
Au	-0.5
U	-1.5

Table 2

$\underline{x}_p$	$\underline{R} = \frac{1}{12} (F_2)_{\text{Carbon}} / \frac{1}{4} (F_2)_{\text{Helium}}$
.7	.949
.8	.952
.9	.978
1.0	.933
1.1	.871
1.2	.779
1.4	.586
1.6	.432
1.8	.319
2.0	.243
2.2	.195

Table 3

Deuterium ($Q^2 = 7.86$)

<u>X</u>	<u>VAL</u>	<u>SEA</u>	<u>GLUE</u>	<u>CHARM</u>
0	0	3.579	13.224	0.493
.05	0.631	1.854	6.471	0.188
.1	0.835	0.931	3.063	0.068
.15	0.910	0.449	1.390	0.024
.2	0.912	0.209	0.605	0.008
.25	0.868	0.093	0.251	.236E-2
.3	0.795	0.039	0.099	.678E-3
.35	0.705	0.016	0.037	.181E-3
.4	0.605	0.006	0.013	.448E-4
.45	0.506	.207E-2	.413E-2	.103E-4
.5	0.409	.679E-3	.124E-2	.219E-5
.55	0.320	.205E-3	.345E-3	< 10 ⁻⁶
.6	0.240	.574E-4	.889E-4	"
.65	0.173	.149E-4	.215E-4	"
.7	0.117	.368E-5	.498E-5	"
.75	0.074	< 10 ⁻⁶	.114E-5	"
.8	0.043	"	< 10 ⁻⁶	"
1.0	.122E-2	"	"	"
1.25	.971E-5	"	"	"
1.50	.490E-7	"	"	"
1.75	.245E-10	"	"	"

Table 4

Helium ($Q^2 = 7.86$)

<u>X</u>	<u>VAL</u>	<u>SEA</u>	<u>GLUE</u>	<u>CHARM</u>
0	0	7.276	26.883	1.003
.05	1.310	3.674	12.800	0.368
.1	1.717	1.798	5.894	0.129
.15	1.860	0.850	2.617	0.044
.2	1.849	0.387	1.117	0.014
.25	1.743	0.169	0.457	0.004
.3	1.579	0.071	0.178	.123E-2
.35	1.382	0.028	0.066	.336E-3
.4	1.172	0.011	0.023	.863E-4
.45	0.964	.385E-2	.777E-2	.209E-4
.5	0.768	.131E-2	.244E-2	.476E-5
.55	0.590	.416E-3	.718E-3	.102E-5
.6	0.435	.124E-3	.199E-3	< 10 ⁻⁶
.65	0.307	.350E-4	.519E-4	"
.7	0.205	.925E-5	.128E-4	"
.75	0.129	.232E-5	.299E-5	"
.8	0.075	< 10 ⁻⁶	< 10 ⁻⁶	
1.0	.348E-2	"	"	"
1.25	.143E-4	"	"	"
1.5	.262E-7	"	"	"

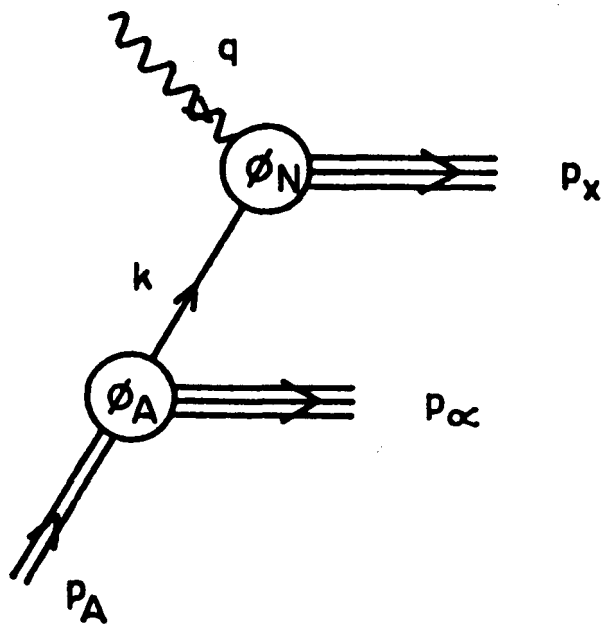


FIG. 1

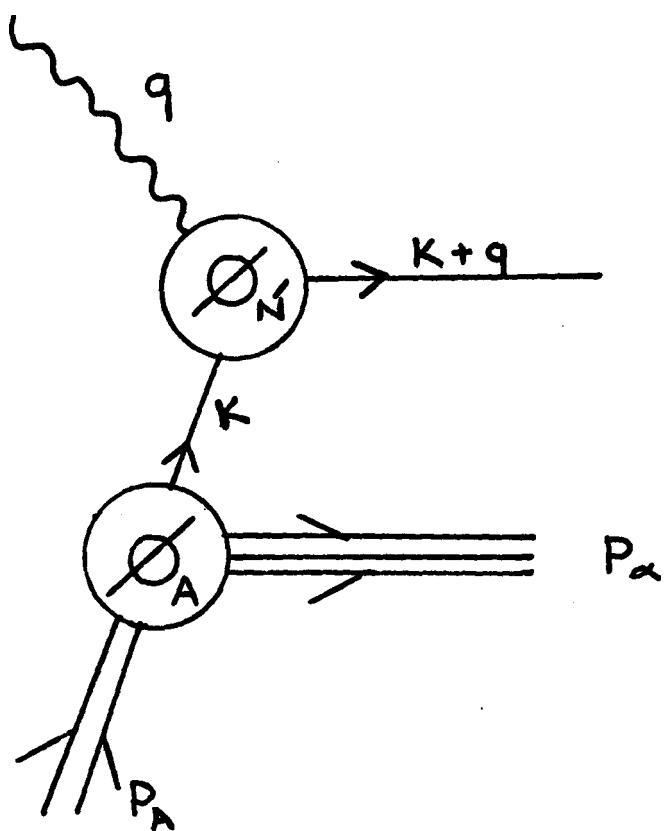


FIG. 2

THREE-BODY MODEL OF THE ${}^6\text{Li} \rightarrow p + (n\alpha)$ SPECTRAL FUNCTION*

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The ${}^6\text{Li} \rightarrow p + (n\alpha)$ vertex has been investigated within the framework of three-body models of ${}^6\text{Li}$, complementing earlier work on the ${}^6\text{Li} \rightarrow \alpha + d$ vertex.¹ The key element in the examination of the ${}^6\text{Li} \rightarrow p + (n\alpha)$ vertex is the overlap integral of the ${}^6\text{Li}$ ground state with a state that describes a proton moving freely relative to the $(n\alpha)$ center-of-mass while the $(n\alpha)$ pair interact. The strength of the three-body model in describing this vertex lies in the fact that the dynamics in the initial and final states are fully consistent. Furthermore, the predictions of the model can be tested by means of kinematically complete coincidence experiments, such as ${}^6\text{Li}(p,2p)n\alpha$ and ${}^6\text{Li}(e,e'p)n\alpha$, when the incident projectile has sufficient energy ($\gtrsim 500$ MeV) to assure that the ejected proton leaves the nucleus with $\gtrsim 100$ MeV energy relative to the $(n\alpha)$ system so that the rescattering between the ejected proton and the $(n\alpha)$ system can be neglected.

The ${}^6\text{Li}$ wave functions used in this work are those of Lehman, Rai, and Ghovanlou.² The interactions chosen to generate the ${}^6\text{Li}$ wave functions ($J^\pi, T=1^+, 0$) are the ${}^3S_1$ - ${}^3D_1$ np and the $S_{1/2}$, $P_{1/2}$, and $P_{3/2}$ n α . The $S_{1/2}$ n α interaction makes manifest the repulsion due to the

Pauli exclusion between a nucleon and the composite alpha particle. It is represented two different ways: 1.) A purely repulsive interaction; and 2.) An attractive interaction with the forbidden $n\alpha$ bound state completely removed.³ The $(n\alpha)$ scattering wave function that is required in the overlap integral consists of a plane-wave term plus the scattered-wave term containing the $S_{1/2}$, $P_{1/2}$, and $P_{3/2}$ components.

In this contribution, we suggest that the ${}^6\text{Li}(e,e'p)n\alpha$ reaction is an ideal means of testing the three-body predictions for the ${}^6\text{Li} \rightarrow p + (n\alpha)$ vertex. Specifically, it might be possible to learn about the representation of the underlying $S_{1/2}$ $(n\alpha)$ interaction. The obvious advantage of the $(e,e'p)$ reaction compared with the $(p,2p)$ is the absence of projectile rescattering effects. Although the magnitudes of the $(e,e'p)$ cross sections are smaller than $(p,2p)$, with 100% duty-factor machines which have relatively high currents, this aspect will be somewhat alleviated.

With the following Feynman diagram for the ${}^6\text{Li}(e,e'p)n\alpha$ reaction in the laboratory:

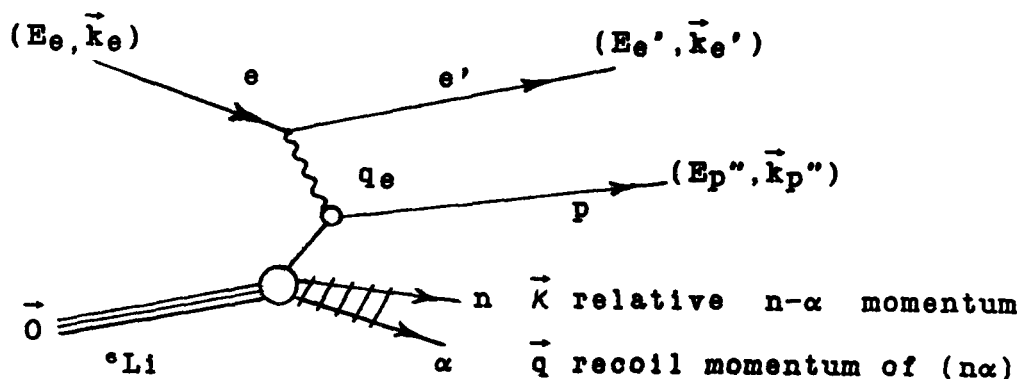


Fig. 1. Feynman diagram for ${}^6\text{Li}(e,e'p)n\alpha$

the four-fold coincidence cross-section becomes

$$\frac{d^4\sigma}{dE_e' d\Omega_e' d\Omega_p'' dE_K} = \left[\frac{k_p'' E_p''}{|1 - (E_p''/E_q)(\vec{k}_p'' \cdot \vec{q}/(k_p'')^2)|} \right] \times$$

$$\left[\frac{d\sigma}{d\Omega_e' \text{ lab}} \right]_{\text{(w.r.f.)}} \left[\mu_{n\alpha K} V(K, q) \right]$$

, (1)

where $E_K = K^2/2\mu_{n\alpha}$ is the relative energy of the $(n\alpha)$ pair, the first bracket is a kinematical factor, the e-p 'cross-section' is in the laboratory frame without recoil factor, and the last bracket contains the physics of the ${}^6\text{Li} \rightarrow p + (n\alpha)$ vertex. $V(K, q)$ is the square of the ${}^6\text{Li} \rightarrow p + (n\alpha)$ overlap integral, summed and averaged over magnetic substates and integrated over the angular direction of the relative-momentum vector of the $(n\alpha)$ pair.

The ${}^6\text{Li} \rightarrow p + (n\alpha)$ spectral function is defined by the factor in the third bracket in (1). It gives the energy spectrum of the $(n\alpha)$ pair by holding q fixed and varying E_K . Such a spectrum, below the threshold for ${}^6\text{Li} \rightarrow p + (d \text{ } {}^3\text{H})$, shows a prominent peak due to the $P_{3/2}$ $(n\alpha)$ resonance. On the other hand, integration of the spectral function with respect to E_K over the $P_{3/2}$ -resonance peak, e.g., $0 \leq E_K \leq 2 \text{ MeV}$, yields the ${}^6\text{Li} \rightarrow p + (n\alpha)$ momentum distribution at the resonance.

At $q=0$, the momentum distribution possesses a non-zero minimum, that arises from the destructive interference between the plane-wave and the $S_{1/2}$ scattered-wave terms of the $(n\alpha)$ scattering state.⁴ Figure 2 shows an example of one of our spectral function predictions.

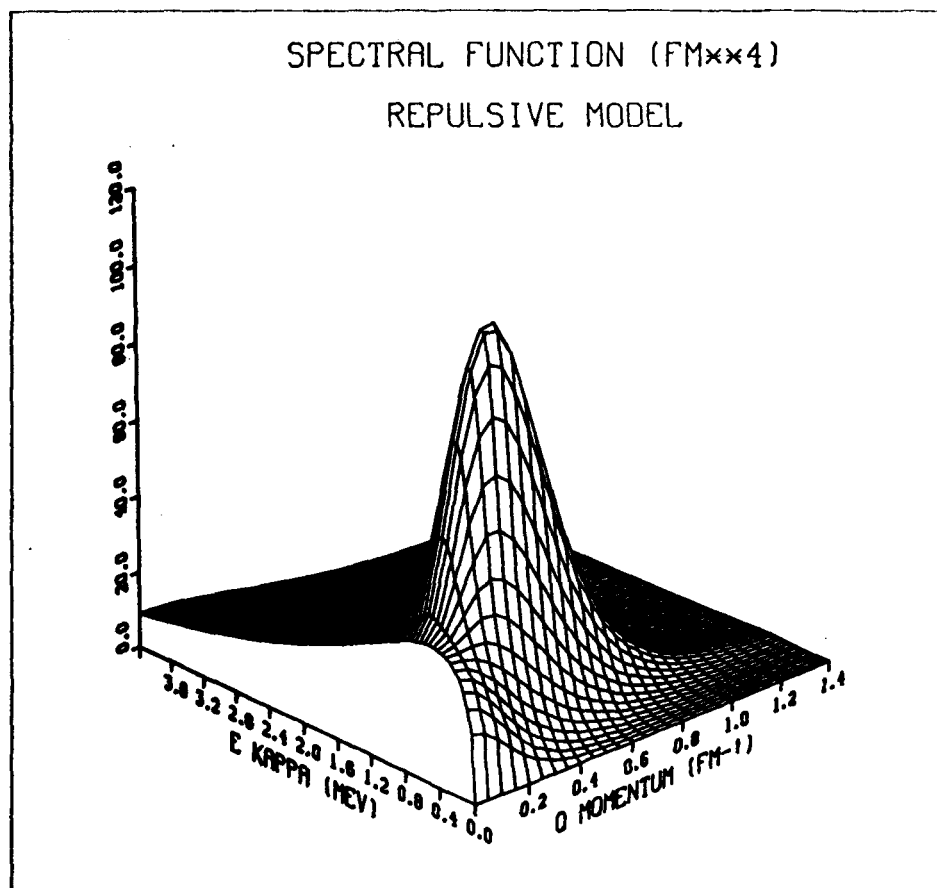


Figure 2. ${}^6\text{Li} \rightarrow p + (n\alpha)$ spectral function.
(Model: Repulsive $S_{1/2}$ and 4% np tensor)

* Supported in part by the U.S. Department of Energy
Grant No. DE-FG05-86-ER40270.

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A DWIA Study of the (γ, K^+) Reaction

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ABSTRACT

The effects of distortion and of a complete photoproduction operator on the calculated differential cross sections for the (γ, K^+) reaction are examined. Distortion effects are found to be important while inclusion of terms other than $\vec{t} \cdot \hat{e}$ leads to a significantly richer spectrum.

The use of tagged photons to explore hypernuclear structure holds fascinating possibilities. Most of our knowledge of hypernuclear spectra comes from the (K^-, π^-) reaction, a reaction which does not tend to excite spin-flip transitions. A better understanding of the Λ -nucleus interaction will require more complete spectra. The (γ, K^+) reaction can generate such spectra since both spin-flip and non-flip modes are excited. Previous studies have focused on the spin-flip excitation induced by this reaction by employing an effective $\vec{\sigma} \cdot \hat{\epsilon}$ form of the photoproduction operator^(1,2). These earlier studies also neglected distortion effects. A more realistic calculation of this reaction will be useful for planning the characteristics of proposed experiments.

At the energies of greatest interest, around $k_{\text{LAB}} = 2$ GeV/c, the kaon momentum in the KN center-of-momentum frame is ≈ 0.7 GeV/c which is above the threshold for pion production. Such inelasticities lead to a KN total cross section which exceeds 20 mb in this kinematic region⁽³⁾ and consequently the usual arguments for negligible kaon distortions are questionable. We have therefore constructed a first order optical potential of $t\rho$ form using Martin's parameterization of the K^+N amplitudes⁽⁴⁾. This has been used in a coordinate space DWIA calculation with a $\gamma + p \rightarrow K^+ + \Lambda$ amplitude of CGLN form

$$F = F_1 \vec{\sigma} \cdot \hat{\epsilon} + i F_2 \vec{\sigma} \cdot \hat{q}_c \vec{\sigma} \cdot \hat{\epsilon} \times \hat{k}_c + F_3 \vec{\sigma} \cdot \hat{k}_c \hat{q}_c \cdot \hat{\epsilon} + F_4 \vec{\sigma} \cdot \hat{q}_c \hat{q}_c \cdot \hat{\epsilon} \quad (1)$$

where k_c (q_c) is the CM photon (kaon) asymptotic momentum. The F_i are shown at zero degrees in figure 1. They are obtained from the microscopic model of reference 5 by fitting the extant photokaon data and are qualitatively similar to the amplitudes obtained using other models^(6,7). All such models, which employ lowest order Feynman diagrams or Breit-Wigner type terms for the exchanged resonances are in qualitative agreement with the data from threshold through $k_{\text{LAB}} = 1.4$ GeV/c. At the higher energies of interest our amplitudes (and all others) are obtained by extrapolation and must be regarded as

speculative.

The important amplitudes at forward angles are F_1 and F_2 . At zero degrees, spin-flip hypernuclear excitations are proportional to $F_1 - F_2$ while at all angles non-flip transitions are proportional to F_2 . Since $F_2 \approx F_1/2$ at $k_{\text{LAB}} = 2 \text{ GeV/c}$, we expect the spin-flip and non-flip transitions to be comparably excited.

With a zero range approximation the cross section becomes

$$\sigma(\theta) = \rho_{1f} \left(\frac{1}{2(2J_i+1)} \right) \sum_{\lambda M_i M_f} |\langle J_f M_f; K^+(\vec{q}) | H | J_i M_i; \gamma(\vec{k}, \lambda) \rangle|^2 \quad (2)$$

where the notation is that of reference 8. The transition operator is

$$H = \frac{4\pi s_c^{1/2}}{\sqrt{M_p M_\Lambda}} F \quad (3)$$

where s_c is the total γp energy in the γp center-of-momentum frame. The momenta in eqn.(1) are converted to gradients in the γ -nucleus (ACM) center-of-momentum frame, according to

$$\begin{aligned} \vec{\sigma} \cdot \hat{\epsilon} \vec{\sigma} \cdot \vec{q}_c &\rightarrow (-1/q_{\text{ACM}}) \vec{\sigma} \cdot \hat{\epsilon} (\vec{\sigma} \cdot \vec{\nabla}_q) \\ \hat{q}_c \cdot \hat{\epsilon} \vec{\sigma} \cdot \vec{k}_c &\rightarrow (-1/q_c k_{\text{ACM}}) (\vec{\nabla}_q \cdot \hat{\epsilon}) (\vec{\sigma} \cdot \vec{\nabla}_k) \\ \vec{\sigma} \cdot \hat{q}_c \hat{q}_c \cdot \hat{\epsilon} &\rightarrow (1/q_{\text{ACM}}) (\vec{\sigma} \cdot \vec{q}_{\text{ACM}}) (\hat{q}_c \cdot \hat{\epsilon}) \\ &= (1/q_{\text{ACM}}) (\vec{\sigma} \cdot (\vec{k}_{\text{ACM}} + \vec{p}_i - \vec{p}_f)) (\hat{q}_c \cdot \hat{\epsilon}) \\ &= (-1/q_{\text{ACM}} q_c) (\vec{\sigma} \cdot \vec{\nabla}_k + \vec{\sigma} \cdot \vec{\nabla}_i - \vec{\sigma} \cdot \vec{\nabla}_f) \hat{\epsilon} \cdot \vec{\nabla}_q \end{aligned} \quad (4)$$

where ∇_i and ∇_f are gradients on the initial and final baryon wave functions. Calculations were performed for $^{12}\text{C}(\gamma, K^+)_{\Lambda}^{12}\text{B}$ and $^{208}\text{Pb}(\gamma, K^+)_{\Lambda}^{208}\text{Tl}$ to study the effects of distortion and of the various operators for extremes in mass number. Oscillator wave functions were employed throughout. For ^{12}C , $b = 1.7$ fm and for ^{208}Pb , $b = 2.4$ fm. The mass excess for $_{\Lambda}^{12}\text{B}$ was taken as $180 \text{ MeV}/c^2$ and $150 \text{ MeV}/c^2$ for $_{\Lambda}^{208}\text{Tl}$.

Figures 2 and 3 show predictions for the $(p_{3/2}^{-1} p_{3/2})^J$ multiplet in $_{\Lambda}^{12}\text{B}$ and for the $(h_{11/2}^{-1} s_{1/2})^{J=6}$ state in $_{\Lambda}^{208}\text{Tl}$ at $k_{\text{LAB}} = 1.84 \text{ GeV}/c$. Distortion effects are clearly significant. For the stretched configurations shown in figure 2 the distortion is a relatively angle independent effect amounting to a factor of ≈ 1.5 for $A = 12$ and ≈ 5 for $A = 208$. However, other states, such as the 1^+ in $_{\Lambda}^{12}\text{B}$ have no such simple suppression factor; at the very smallest angle the distorted wave cross section is even enhanced relative to the plane wave one.

The significance of operators other than $\vec{\sigma} \cdot \hat{e}$ are also shown in these figures. The operators have a minor effect on the angular distribution of the $J = 3$ state in $_{\Lambda}^{12}\text{B}$ and a negligible effect for the stretched state in $_{\Lambda}^{208}\text{Tl}$. However, as expected, the F_2 operator gives rise to significant non-flip transition such as the 2^+ in $_{\Lambda}^{12}\text{B}$. Notice that where the 2^+ cross section peaks at 10° it is half as large as the dominant 3^+ . The angular signatures of these states are sufficiently unique to enable them to be distinguished by proposed spectrometers⁽⁹⁾.

Of course, a pure particle-hole representation of these states is useful only for a first orientation. When calculations are performed using more reliable configuration mixed wave functions we find an additional suppression by a factor of ≈ 2 for most $_{\Lambda}^{12}\text{B}$ states. Details will be published elsewhere.

We conclude that accounting for distortion effects will be essential in any attempt to extract hypernuclear information from measured (γ, K) cross

sections. At $k_{\text{LAB}} = 1.8 \text{ GeV}/c$ the suppressions are typically a factor of 1.5 for light nuclei and a factor of five for heavy nuclei. Such effects will be crucial, for example, in determining the systematics of spin excitation suppression for hypernuclei. In addition, we predict a richer spectrum of excited states than a pure $\hat{\sigma} \cdot \hat{e}$ form of the amplitude would suggest, a richness which is potentially useful for extracting details of the ΛN interaction.

This work was supported in part by a grant from Research Corporation and the National Science Foundation under grants Phy-8317084 and Phy-8516195.

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Figure 1 - The forward angle real parts of the CGLN amplitudes.

Figure 2 - Differential cross sections for the stretched states of the considered $^{12}_{\Lambda} \text{B}$ and $^{208}_{\Lambda} \text{Tl}$ configurations. Curves plotted are with Coulomb distortion only (short dashed curves), full optical distortion but $\hat{\sigma} \cdot \hat{e}$ only (long-short dashed curves) and full distortion with all amplitudes (solid curves).

Figure 3 - Differential cross sections for the $J = 0, 1, 2$ states in $^{12}_{\Lambda}B(p_{3/2}^{-1} p_{3/2})$. The 1^+ cross section is shown with Coulomb distortions only (dashed curve). All others are calculated with full optical distortion.

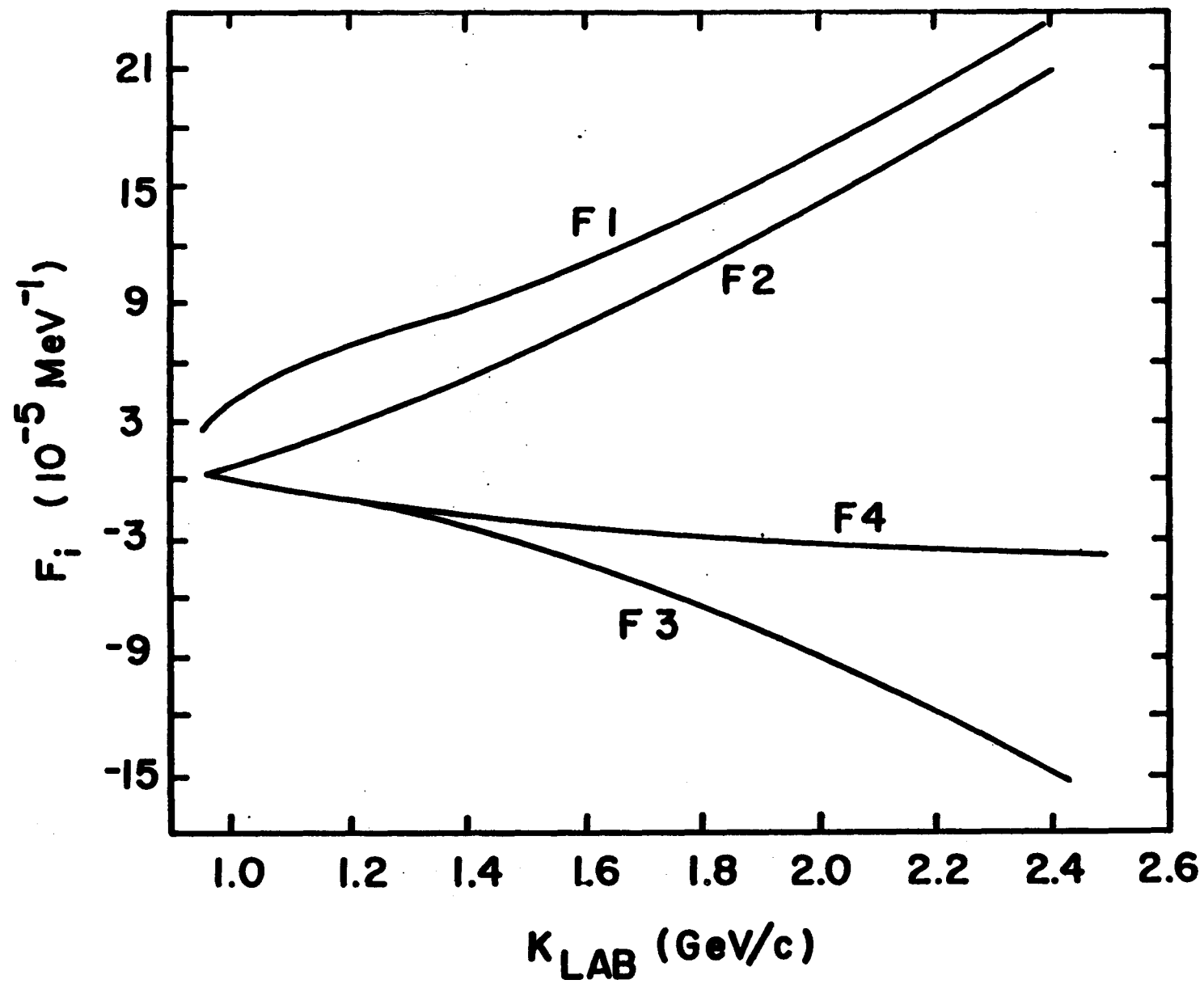


fig 1

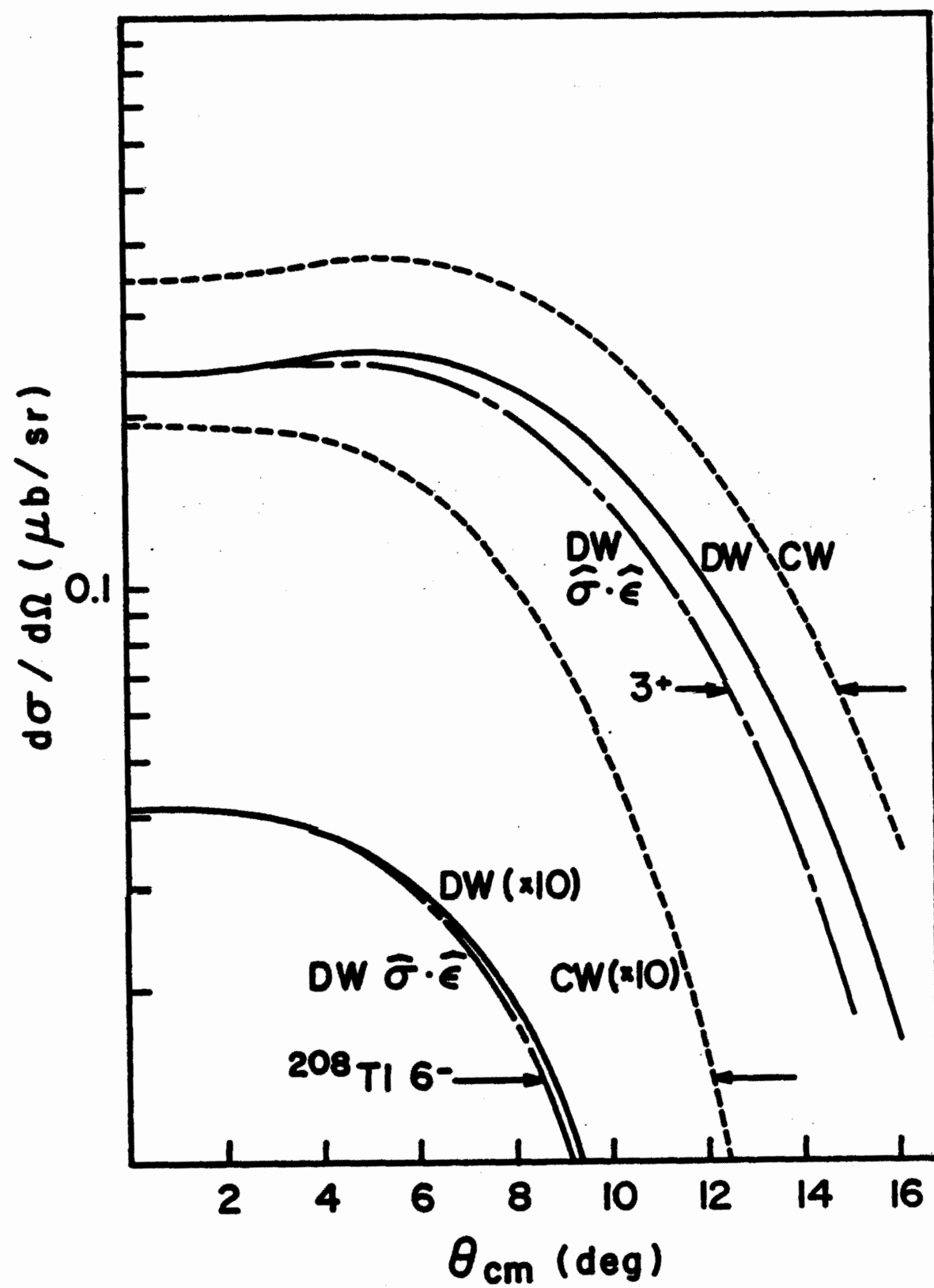


fig 2

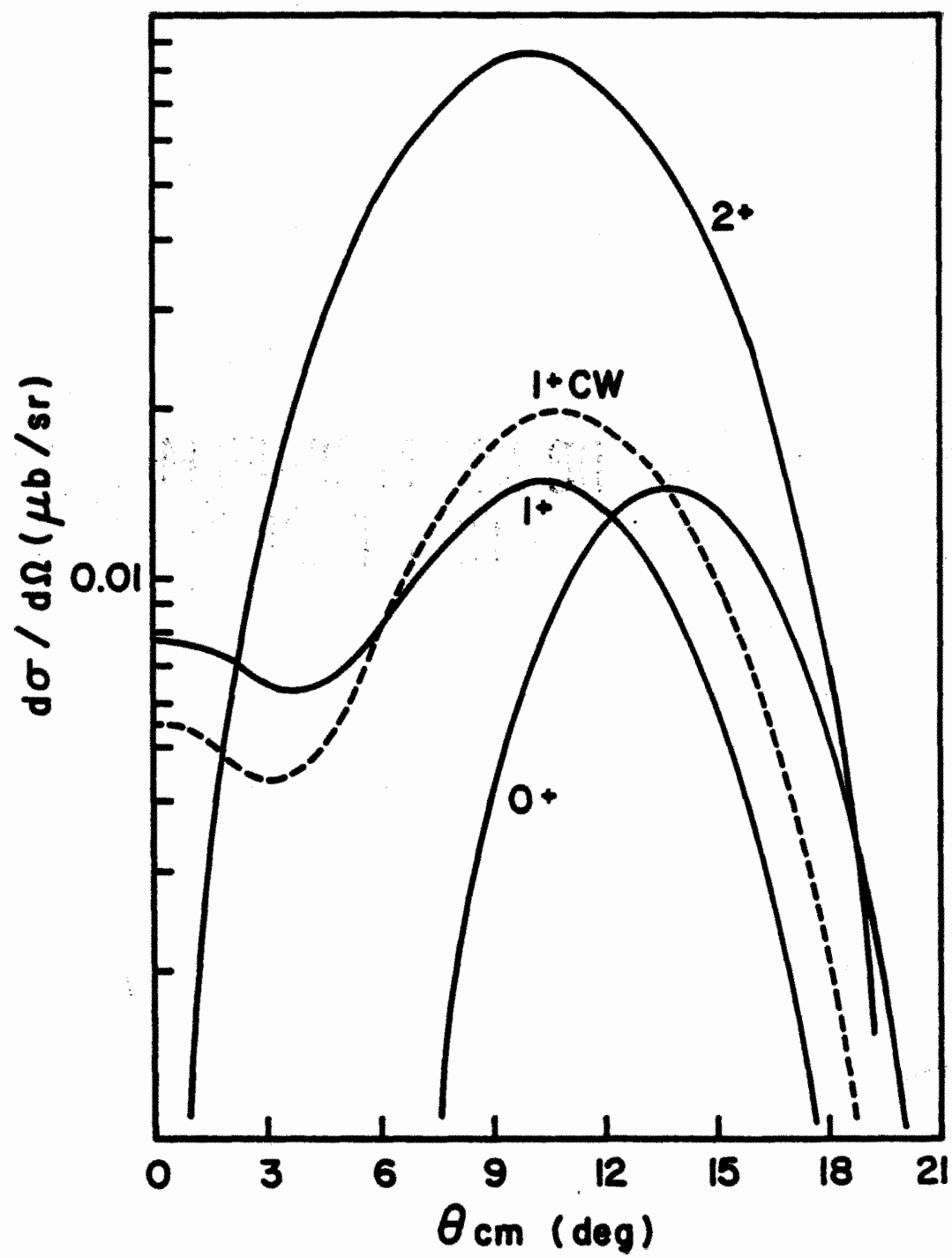


fig 3

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TOWARDS RELATIVISTIC ELECTRO-MAGNETIC BARYON COUPLINGS MOTIVATED BY QCD

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1. INTRODUCTION

Quantumchromodynamics (QCD) is generally supposed to be the field theory of strong interactions. Thus this theory should also be able to understand all properties of all hadrons. Since most hadrons known up till now have inner structures where confining forces dominate^{/1/}, there is presently no reliable mathematical framework to calculate hadron data starting from the QCD-Lagrangian. Of course, the numerical evaluation of lattice gauge field theory can finally arrive at this goal but in the meantime one must be content with a semiphenomenological description of the hadrons which includes as many qualitative features of QCD as possible.

Since in this way one can at least develop an effective theory it is advisable to use models which are as simple as possible. In this respect the naive quark model has proven to be a very successful starting point. Especially the "Soft QCD-Model" elaborated by Isgur and collaborators^{/2/, /3/} has demonstrated how one can understand even more detailed hadron properties by using simple tools like an effective hamiltonian with instantaneous forces between quarks.

The success of these calculations also for hadrons composed of the light u-, d- and s-quarks encouraged us to investigate in more detail the "relativization" of the QCD-motivated quark model. Relativistic effects will not only lead to spin splittings of hadronic levels but also change the properties of the couplings of hadrons with other hadrons and with electromagnetic and weak currents. Here the photon-hadron interactions are especially suited for a thorough research: They have a better theoretical foundation than purely hadron vertices where in the quark model approach only first steps for a QCD-inspired calculation have been done^{/4/}. On the other hand, there is considerably more experimental information on photonic, than on weak couplings of hadrons. We can use the rich data of photoproduction and electron scattering experiments to compare theoretical results with the empirical

information on photoexcitation amplitudes and form factors of hadrons.

In this paper we deal with the electromagnetic properties of nonstrange baryons. In a first part we explain the general theoretical ideas which lead to a quasi-relativistic interaction hamiltonian between quarks and the electromagnetic fields expressed in a formal series in powers of $1/c$ and we give explicit formulae for the $(1/c)^2$ and $(1/c)^3$ -terms. The second part is devoted to an application of these results to the Isgur-Karl model for baryons. We compare our theoretical findings with the experimental helicity amplitudes for the photoexcitations of baryon resonances and their dependence of the momentum transfer.

2. A QUASI-RELATIVISTIC PHOTON-QUARK INTERACTION UP TO ORDER $(v/c)^2$

The local form of the fundamental electromagnetic interaction

$$(1) \quad e \bar{\psi} Q \gamma^\mu \psi A_\mu$$

where ψ denotes the Dirac spinor of the light quarks and Q is the charge matrix in flavour-space, is only valid for the pointlike "current" quarks. It must be modified in various respects to describe the photonic interaction of "constituent" quarks, of those quarks which are the building blocks of the quark model.

a) The single quark interaction hamiltonian

As a first step one replaces the pointlike relativistic coupling (1) by applying a Foldy-Wouthuysen transformation to the Dirac spinors. Retaining terms up to order $\frac{1}{c^2}$, one finds from textbooks^{/5/} a hamiltonian acting on the Pauli spinors of the quarks:

$$(2) \quad H_{ee} = H_{ee}^{(1)} + H_{ee}^{(2)} + H_{ee}^{(3)}$$

with

$$(2a) \quad H_{ee}^{(1)} = -\frac{eQ}{2mc} (2\vec{p} \cdot \vec{A} + \vec{\sigma} \cdot \vec{B})$$

$$(2b) \quad H_{ee}^{(2)} = \frac{eQ}{2mc} \frac{\omega}{4mc} (\vec{\sigma} \cdot \vec{B} - 2i \vec{\sigma} \cdot (\vec{A} \times \vec{p}))$$

$$(2c) \quad H_{ee}^{(3)} = \frac{eQ}{2mc} \frac{1}{2mc^2} \left\{ (2\vec{p} \cdot \vec{A} + \vec{\sigma} \cdot \vec{B}, \frac{\vec{p}^2}{2m}) \right\}$$

These expressions are matrices in spin space - via the Pauli matrix $\vec{\sigma}$ - and in flavour space - via the charge matrix Q . In (2c) an anticommutator occurs. The three terms (2a-2c) illustrate what is meant by orders $\frac{1}{c}$, $\frac{1}{c^2}$, and $\frac{1}{c^3}$. Physically $1/c$ is a short term for following ratios:

$$(a) \text{ the quark velocity } \frac{v}{c} = \frac{p}{mc}$$

$$(b) \text{ photon momentum/mc } \frac{k}{mc} = \frac{\lambda_c}{\lambda}$$

$$(c) \text{ photon energy/mc } \frac{\omega}{mc}$$

((b) and (c) coincide for real photons).

If (a) determines $1/c$ the increasing powers of $1/c$ correspond to an increasing non-locality of the hamiltonian. We are aware of the fact that for hadrons build up by light quarks the first of these parameters is not necessarily small. But our calculations show that there is a high degree of mutual compensation of different terms so that the employed expansion appears to be justified finally.

If one remembers the only loose connection between current and constituent quarks the hamiltonian (2) may appear as a very special ansatz. But if one analyses the helicity content of (2a-c) one finds that only three of the four terms of the general Melosh ansatz^{/6/} for a single-quark photon interaction are present; the exotic D-term does not occur. Indeed if one considers a helicity $\lambda = +1$ photon moving along the negative z-direction: $\vec{A} = \frac{1}{\sqrt{2}} (\vec{e}_x + i\vec{e}_y) e^{-i(kz+\omega t)}$ the hamiltonian (2) - being a 2x2 spin matrix - can be written in the form

$$(3) \quad H_{ee}^{\lambda=+1} = A^+ 1 + B^+ \sigma^+ + C^+ \sigma^z + D^{++} \sigma^-$$

where $\sigma^{\pm} = \sigma^x \pm i\sigma^y$. Under a rotation around the z-axis by an angle φ the four amplitudes of (3) behave like

$$(3a) \quad A^+ \sim C^+ \sim e^{+i\varphi}$$

$$(3b) \quad B^2 \sim 1$$

$$(3c) \quad D^{++} \sim e^{+2i\varphi}$$

corresponding to the helicities +1, 0 and +2 respectively. The convection term of (2a) is an A-amplitude, while its magnetic part behaves like a B-amplitude as does the first term of (2b). Therefore, in the pure non-relativistic quark model only the amplitudes A and B contribute contradicting strongly the experimental facts^{/28/}. The second part of (2b) is the prototype of a spin-orbit term, due to the relation $(\vec{\sigma} \times \vec{p})^+ = i(\sigma^z p^+ - \sigma^+ p^z)$ it contributes to the B- and the C-amplitudes. The 3rd order hamiltonian (2c) gives new contributions to A and B^{/7/}. Only the helicity 2 term D does not show up in (2).

Therefore, the ansatz (2) is not a general one even if we allow for arbitrary reduced matrix elements when analysing transition amplitudes. (2) represents only the first, trivial step of an approach to a relativistic photon- constituent quark interaction.

b) The interaction dependent photon hamiltonian

The Foldy-Wouthuysen transformation employed to derive the hamiltonian (2) starts from free quark spinors. However, the strong interactions between the quarks will produce new pieces in the effective photon-quark hamiltonian which we call the "interaction dependent terms". In principle, these contributions should be calculated by integrating out the gluon fields from the complete QCD + photon-Lagrangian. Thereby one would obtain at the same time a model hamiltonian for the strong quark-quark forces. Such an ambitious program cannot be carried through at present. Therefore, one has to recourse to a more simple method where general considerations are combined with plausible approximations gained by a careful look to the experimental data. In view of the last point the successful study of the charmonium system provides a very strong help. Using the idea of the seminal paper of DeRujula, Georgi and Glashow from 1975^{/8/} the one-gluon exchange potential should be a good starting point for the quark-quark interaction. This should be especially true for its momentum dependent parts which decrease stronger at larger distances. By the minimal substitution of $\vec{p}$ by $\vec{p} - \frac{e}{c}\vec{A}$ exactly these terms produce the wanted interaction dependent part of the photon coupling. Therefore, we are led to start from the wellknown Fermi-Breit form of the quark potentials. But the derivation of this force by reducing the field theoretic expression for one-gluon-exchange to a quasi-relativistic potential^{/9/} encounters an ambiguity problem. Its origin is immediately clear by looking at the expression for the gluon propagator

$$(4) \quad \frac{1}{q^2} = \frac{1}{q_0^2 - \vec{q}^2}$$

For the potential one needs the Fourier transform with respect to the spatial part $\vec{q}$ of the momentum transfer. But the energy transfer q_0 is not fixed by $\vec{q}$ unless one chooses a special Lorentz system. Usually one works in the center of mass system of the two quarks where q_0 vanishes. But for our goal - the coupling of the electromagnetic field to the three quark system we need the quark-quark potential in any Lorentz frame.

To solve this ambiguity we used the results of the long lasting discussion on the possibility of a relativistic theory for interacting particles without fields^{/10/}. This discussion started about 25 years ago and in spite of a no-go theorem^{/11/} it leads to concrete results which may get new attention in view of the attempts to formulate a relativization of the constituent quark model^{/12/}. We based our calculations especially on an approach started in 1961 by L.L. Foldy^{/13/} and worked out in detail in 1975 by Krajcik and Foldy^{/14/}. Parts of the results of these lengthy calculations have also been obtained by other methods^{/15/}. A thorough analysis of^{/14/} from various points of view and a series of applications have been given in^{/16/}.

The idea of Foldy's method^{/11/} can be described briefly as follows: One tries to formulate a relativistic theory for a many particle system approximating the fully fledged field theory in two important aspects

- creation and annihilation of particles are neglected, the number of particles is fixed,
- for the description of the interaction no fields are used, the retardation effects are taken into account step by step by using an expansion in powers of $1/c$.

To carry this program through one starts with a naive relativistic hamiltonian for the n-particle system

$$(5) \quad H = \sum_{i=1}^n c \sqrt{(m_i c)^2 + \vec{p}_i^2} + U^{(0)}(\vec{r}_{ij})$$

where only the kinetic part has a relativistic shape. The potential $U^{(0)}$ describes instantaneous forces depending only on the relative spatial coordinates ($\vec{r}_{ij} = \vec{r}_i - \vec{r}_j$). To get a truly relativistic theory from (5) one replaces $U^{(0)}$ by a potential

$$(6) \quad U(\vec{r}_{ij}, \vec{p}_i, \vec{s}_i)$$

which can also depend on momenta and spins of the quarks so that the hamiltonian (5) acts as a generator of the Poincaré group. The basic result of ^{/13/} and ^{/14/} can be stated as follows: By an expansion in powers of $(\frac{1}{c})^2$ a potential U can be constructed so that the hamiltonian H forms together with the total momentum $\vec{P}$, the total angular momentum $\vec{J}$ and the boost operator $\vec{K}$ a representation of the Poincaré algebra. Thereby, $\vec{P}$ and $\vec{J}$ keep their simple additive free form while for the boost $\vec{K}$ one has to include terms depending on the interaction. In this way one finds that the ambiguity mentioned in connection with (4) can be solved by the following simple recipe: Let in a general Lorentz form $\vec{p}_1, \vec{p}_2$ respectively $\vec{p}_1', \vec{p}_2'$ be the momenta of two quarks before and after one gluon exchange, so that

$$(7) \quad \vec{p}_1' = \vec{p}_1 - \vec{q} \quad ; \quad \vec{p}_2' = \vec{p}_2 + \vec{q}$$

Then the energy transfer q_0^2 must be understood according to the symmetric rule

$$(8) \quad q_0^2 \rightarrow \frac{1}{c^2} (E_1' - E_1)(E_2 - E_2') \approx \frac{1}{4m_1 m_2 c^2} (\vec{p}_1'^2 - \vec{p}_1^2)(\vec{p}_2^2 - \vec{p}_2'^2)$$

where an approximation of order $(\frac{1}{c})^2$ has been used in the last step. Instead of (8) several other procedures have been used in the literature ^{/29/}. But only our recipe seems to be consistent with the Poincaré algebra. To calculate the potential U up to this order one has to use the Feynman rule result for the one-gluon exchange amplitude, employ the kinematic relations (7) and (8), do an expansion of all kinematic quantities and free Dirac spinors up to order $(\frac{1}{c})^2$ and then Fourier transform with respect to $\vec{q}$. The explicit results will be given in a forthcoming publication. The second order parts contain a lot of momentum dependences like

$$\sum_{i < j} \frac{1}{4m_i m_j c^2} \vec{p}_i V(|\vec{r}_i - \vec{r}_j|) \vec{p}_j$$

where $V(r)$ denotes e.g. the static potential: $V(r) = -\frac{2}{3} \frac{\alpha_s}{r}$.

By applying the minimal substitution one finally obtains the wanted interaction dependent parts of the photonic hamiltonian. We quote the results by distinguishing between terms which follow when using $q_0=0$ in (4) and those which one finds from (8) for equal quark masses ^{/17/}

$$\begin{aligned}
 H_{ee}^{(3)}(q_0=0) = \sum_{i < j} \frac{e Q_j}{4 m^2 c^3} & \left[\{ 2 \vec{p}_i \cdot \vec{A}_j, V(r_{ij}) \} + \right. \\
 (9a) \quad & \left. + \frac{V'(r_{ij})}{r_{ij}} (\vec{\sigma}_j + 2 \vec{\sigma}_i) \cdot \vec{r}_{ij} \times \vec{A}_j \right]
 \end{aligned}$$

$$\begin{aligned}
 H_{ee}^{(3)}(q_0 \neq 0) = \sum_{i < j} \frac{e Q_j}{4 m^2 c^3} & \left[\{ \vec{p}_i \cdot \vec{A}_j, V(r_{ij}) \} + \right. \\
 (9b) \quad & + 2 \vec{A}_j \cdot \vec{r}_{ij} \frac{V'(r_{ij})}{r_{ij}} \vec{r}_{ij} \cdot \vec{p}_i - \\
 & \left. - i \left(3 \frac{V'(r_{ij})}{r_{ij}} + V''(r_{ij}) \right) \vec{r}_{ij} \cdot \vec{A}_j \right]
 \end{aligned}$$

Both terms are naturally of order $(\frac{1}{c})^3$ and thus have to be considered together with (2c). The physical significance of (9a, b) can be described as follows: The first terms in the square brackets of (9a) and (9b) are consequences of the well-known Darwin-Breit term resulting from an unretarded retarded vector meson exchange^{/18/}

$$\frac{1}{\vec{q}^2} \gamma_{(i)}^\mu \gamma_{(j)}^\nu \rightarrow V(r_{ij}) + \frac{1}{c^2} \vec{p}_i V(r_{ij}) \vec{p}_j$$

The second term of (9a) is a consequence of a spin orbit interaction between the quarks, while the second and third terms of (9b) have no simple physical interpretation.

To simulate the effects of confining forces we did analogous calculations for the exchange of a scalar gluon. The results for the energy-transfer part are identical with (9b), while typical differences occur for the term analogous to (9a).

While the hamiltonians (2a-c) have a one-quark structure and give rise to processes like those of fig. 1 all terms in (9a, b) depend on the coordinates of two quarks. The resulting contributions to the photo-excitation of a three quark system are

illustrated in fig. 2. This two-quark structure also leads to new amplitudes if one analyses the helicity content of (9a, b). Most parts of these two expressions are independent from spin and then proportional to the unit matrix in spin space. Therefore, they contribute to the generalized A-type amplitude, cp. eq.(3). The only spin dependence stems from the spin orbit terms in (9a), they lead to new B- and C-type amplitudes. To be more specific we consider again a three quark system and let the photon couple to the third quark. Then in (9a) expressions of the form

$$(\vec{\sigma}_1 \times \vec{\tau}_{12}) \cdot \vec{A}(\vec{r}_{12}) \quad ; \quad (\vec{\sigma}_2 \times \vec{\tau}_{23}) \cdot \vec{A}(\vec{r}_{23})$$

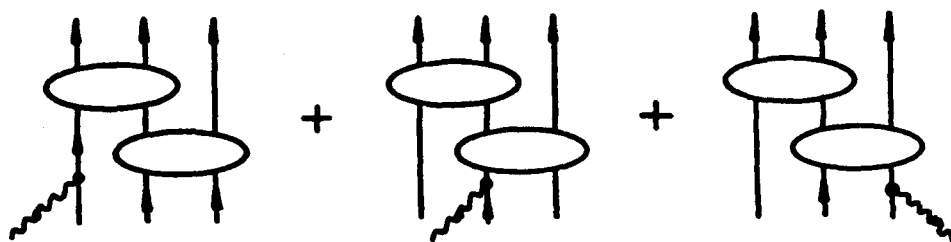


Fig. 1. One particle photon interaction in a three quark system
(The bubbles indicate the effects of the strong interaction)

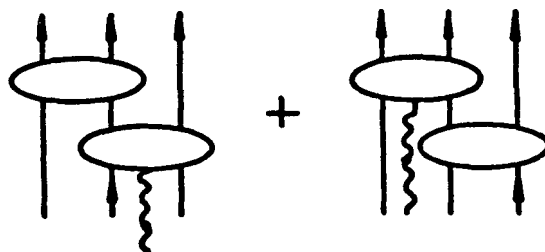


Fig. 2. Examples for the interaction dependent two particle photonic couplings

occur. Introducing matrices

$$\vec{\sigma}_2 = \frac{1}{2} (\vec{\sigma}_1 + \vec{\sigma}_2) , \quad \vec{\sigma}_3 = \frac{1}{2} (\vec{\sigma}_1 - \vec{\sigma}_2)$$

using the notation of Isgur-Karl^{/19/} for the 3-quark system and observing again $(\vec{\sigma} \times \vec{\xi})^z = i(\sigma^z \xi^+ - \sigma^+ \xi^z)$ one finds the following structure for (9a, b)

$$(10) \quad A \mathbb{1} + B_2^z \sigma_2^+ + C_2^+ \sigma_2^z + B_3^z \sigma_3^+ + C_3^+ \sigma_3^z$$

which can be employed when calculating matrixelements for baryon (=3-quark) wave functions.

c) The Center of Mass-corrections

The relativistic definition of the center of mass coordinates is a long standing problem. For an early reference see^{/20/} where various physical possibilities are discussed. The basic property of a Center of Mass should be that it moves uniformly also for an interacting system of particles. This condition will be automatically fulfilled if the generators of the Poincaré group especially the hamiltonian H and the boostoperator $\vec{K}$ can be written in the same form as for a single particle, where the center of mass coordinate $\vec{R}$, the total momentum $\vec{P}$ and the total angular momentum $\vec{J}$ play the role of the particle quantities $(\vec{r}, \vec{p}, \vec{j})$. This goal was in principle achieved already in 1961 by Foldy^{/13/} and later on worked out in detail by various authors^{/14/}, ^{/15/}. The essential result is that one can arrive at this "one particle form" by applying a suitable unitary transformation to the hamiltonian (5), (6) and all other generators, so that the following relations hold

$$(11) \quad H' = e^{i\Phi} H e^{-i\Phi} = \sqrt{h^2 + c^2 \vec{P}^2}$$

$$(12) \quad \vec{K}' = e^{i\Phi} \vec{K} e^{-i\Phi} = \frac{1}{2c^2} \{H', \vec{R}\} - \frac{\vec{S} \times \vec{P}}{H' + h} - t \vec{P}$$

Here $\hbar$ depends only on the internal coordinates and plays the role of a mass operator for the composite system and the total spin $\vec{S}$ is defined by

$$\vec{S} = \vec{J} - \vec{R} \times \vec{P}$$

The second term in (12) takes account of the Wigner rotation and leads to the Thomas precession of the spin of an accelerated body^{/21/}. The mathematical details of this transformation are worked out in^{/14/} and^{/22/}. The relation between the particle coordinates $(\vec{r}_i, \vec{p}_i, \vec{s}_i)$ on one hand and the CMS-quantities $\vec{R}, \vec{J}$ and the "relative", internal coordinate $(\vec{x}_i, \vec{q}_i, \vec{\sigma}_i)$ on the other hand are rather complicated. Starting from the transformation $\vec{r}_i = e^{-i\phi} (\vec{x}_i + \vec{R}) e^{i\phi}$ one gets in the $\frac{1}{c^2}$ - approximation the following expression

$$(13) \quad \vec{r}_i = \vec{x}_i + \vec{R} - \frac{1}{2c^2} \left[\vec{x}_i \cdot \frac{\vec{P}}{M} \left(\frac{\vec{q}_i}{m_i} + \frac{\vec{P}}{2M} \right) + h.c. \right] - \frac{\vec{\sigma}_i \times \vec{P}}{2m_i M c^2} + \text{many other terms}$$

(Here M denotes the total mass: $M = \sum m_i$)

Because of the Wigner rotation term in (12) also the spin variables undergo a transformation

$$(14) \quad \vec{S}_i = \vec{\sigma}_i - \frac{1}{2m_i M c^2} \vec{\sigma}_i \times (\vec{q}_i \times \vec{P}) + \text{further terms}$$

Both matrices $\vec{s}_i$ and $\vec{\sigma}_i$ obey the usual angular momentum commutation relations.

Because of (11) the eigenfunctions of

$$H'^2 = c^2 \vec{P}^2 + \hbar^2$$

can be factorized in the usual way known from non-relativistic quantum mechanics

$$(15) \quad \Psi(\vec{R}, \vec{q}_i, \dots) = e^{i\vec{P}_0 \cdot \vec{R}} \varphi(\vec{q}_i, \dots)$$

where $\vec{P}_0$ is an eigenvalue of $\vec{P}$. All kinematical relativistic effects like Lorentz construction or Thomas precession are thus taken into account by the non-trivial formulae (13) and (14) for the internal coordinates.

Until now all calculations described in this subsection have been without the electromagnetic field. If one applies the minimal substitution to the generating function Φ of the unitary transformation (11) one gets an electromagnetic function

Φ_{e1} and one has to transform the hamiltonian by

$$\mathcal{U} = e^{-i(\Phi + \Phi_e)}$$

Since Φ_{e1} contains the time dependent vector potential $\vec{A}$ the correct rule for the new hamiltonian is

$$(16) \quad H' = \mathcal{U}^{-1} \left(H - i \frac{\partial}{\partial t} \right) \mathcal{U}$$

Developing in powers of $\frac{1}{c}$ one obtains for the terms up to third order

$$(17) \quad H'_{ee} = H_{ee} + i c \omega \Phi_{ee}^{(3)} + i [\Phi_{ee}^{(2)}, H_{ee}^{(1)}] + i [\Phi_{ee}^{(1)}, H_{ee}^{(0)}]$$

where H_{e1} is the sum of (2a, b, c) and (9a, b) and $H^{(0)}$ denotes the hamiltonoperator for the non relativistic quark model. By working out the commutators one arrives at a very complicated expression which is written out in ref. ^{/23/}.

For our further discussion it is necessary to describe the helicity structure of (17). H_{e1} and H'_{e1} are matrices in the spin space of the three quarks and therefore, are of a 8x8 type. But both operators are sums over one-body and two-body spin matrices. A careful counting shows that at most $1 + 3 \times 3 + 2 \times 3 \times 3 = 28$ independent matrices can occur. In H_{e1} only 7 occur (cp. (2a-c) and (10)). By the transformation (17) one generates 7 new terms. The result can be written in the form

$$(18) \quad \begin{aligned} H'_{ee} = & A^+ \mathbb{1} + B^2 \sigma_3^+ + C^+ \sigma_3^2 + D^{++} \sigma_3^- + \\ & + B_\lambda^2 \sigma_\lambda^+ + C_\lambda^+ \sigma_\lambda^2 + D_\lambda^{++} \sigma_\lambda^- + \\ & + B_\vartheta^2 \sigma_\vartheta^+ + C_\vartheta^+ \sigma_\vartheta^2 + D_\vartheta^{++} \sigma_\vartheta^- + \\ & + C_{3\lambda}^+ \sigma_3^+ \sigma_\lambda^- + \\ & + C_{3\vartheta}^+ \sigma_3^+ \sigma_\vartheta^- + \\ & + E_{3\lambda}^- \sigma_3^+ \sigma_\lambda^+ + E_{3\vartheta}^- \sigma_3^+ \sigma_\vartheta^+ \end{aligned}$$

Here we have singled out again the third quark. The first line of (18) contains single-quark interactions while all other terms depend on two quarks. We have written down only those spin matrices which have non vanishing coefficients. One observes that the exotic D-amplitudes occur which are generated e.g. by the following interaction

$$(19) \quad \frac{-ieQ_3}{8Mm^2c^3} 2\vec{q}_3 \cdot \vec{A}_3 \sum_{i=1}^3 (\vec{\sigma}_i \times \vec{q}_i) \cdot \vec{k}$$

which contain - among others - the following helicity contributions

$$(20) \quad \frac{-eQ_3}{8Mm^2c^3} 2k \left[\overset{D^{++}}{\downarrow} q_3^+ q_3^+ \sigma_3^- + \overset{D_\lambda^{++}}{\downarrow} q_3^+ (q_1^+ + q_2^+) \sigma_\lambda^- + \overset{D_s^{++}}{\downarrow} q_3^+ (q_1^+ - q_2^+) \sigma_s^- \right]$$

Their physical origin is a recoil correction which is indicated by the total mass M in the denominator. In a similar way the CMS-transformation induces new E-type amplitudes which are due to an expression of the form

$$(21) \quad \frac{-ieQ_3}{8Mm^2c^3} (\vec{\sigma}_3 \cdot \vec{B}) \sum_{i=1}^3 (\vec{\sigma}_i \times \vec{q}_i) \cdot \vec{k}$$

The symbolic formula (18) represents our main result. Its explicit content will be given in ^{/23/}. The consequences for the electromagnetic transition form factors and the electromagnetic excitations of nucleon resonances are described in the next section.

3. RELATIVISTIC EFFECTS IN THE PHOTO-AND ELECTROEXCITATION OF NUCLEON ISOBARS

For a quantitative study of the relativistic corrections to the photon hadron amplitudes we applied the general formalism described in section 2. We have in mind an exploratory work and shall use - as explained below - a rather primitive inter-quark potential to calculate the interaction dependent parts of the electromagnetic hamiltonian. In spite of this simplification our results exhibit a typical structure built up by many partly compensating terms.

The relativistic corrections have to be applied both to the hadronic wave functions and to the photon interaction. With respect to the first point we have used the nucleon wave function with its admixtures caused by the hyperfine interaction^{/24/}

$$(22) \quad |N(938)\rangle = 0.93 |56, 0^+\rangle_0 - 0.29 |56, 0^+\rangle_2 - 0.23 |70, 0^+\rangle_2$$

Here we used the notation $|n, L^P\rangle_N$ where n refers to the SU(6) representation; L and P are the orbital angular momentum and the parity of the 3-quark system respectively; N denotes the oscillator quantum number. This configuration mixing is a relativistic effect of order $(\frac{1}{c})^2$, therefore, we write generally for the hadron wave functions

$$|\psi\rangle = |\psi^{(0)}\rangle + |\psi^{(2)}\rangle$$

If we use an analogous notation for the different terms of the electromagnetic interaction amplitudes up to order $(\frac{1}{c})^3$

$$(23) \quad \begin{aligned} \langle B | H_{ee} | A \rangle &= \langle B^{(0)} | H_{ee}^{(1)} | A^{(0)} \rangle + \\ &+ \langle B^{(0)} | H_{ee}^{(1)} | A^{(2)} \rangle + \langle B^{(2)} | H_{ee}^{(1)} | A^{(0)} \rangle \\ &+ \langle B^{(0)} | (H_{ee}^{(2)} + H_{ee}^{(3)}) | A^{(0)} \rangle \end{aligned}$$

In this equation the first line gives the non-relativistic amplitudes and the second line the corrections due to the wave function mixing. The third line contains the corrections to second and third order in $1/c$ which follow from equations (2a-c), (9a-c) and (11).

a) Description of the model and the interquark potential

For the evaluation of the matrix elements in eq. (23) we employed the oscillator wave functions of the Isgur-Karl model with their parameter and mixing angles^{/19/}. Especially we take their values for the constituent mass of the u- and d-quark and for the oscillator constant: $m_u = m_d = 350$ MeV, $\alpha = 300$ MeV.

The complete model contains three more constants which enter in the quark-quark potential

$$(24) \quad V_{qq}(r) = \frac{2}{3} \left(-\frac{\alpha_s}{r} + br + c \right)$$

The strong fine structure constant α_s , the string tension b and the parameter c should be fixed in our context by the hadronic mass spectrum. But in the Isgur-Karl calculations for baryons they enter only in an indirect way. We need these constants explicitly for the interaction dependent parts of H_{e1} .

Unfortunately the constants are not fixed by the spectroscopy and one finds quite different values in the literature^{/1/, /25/, /26/, /27/} where e.g. the constant c varies between -1 and $+0.2$. Therefore, we studied in detail how a change of these parameters alters the electromagnetic amplitudes for hadrons. Moreover, we tried an overall fit to the following data:

- The magnetic form factors of the nucleons up to $Q^2 \approx 2$ GeV²,
- the photon excitation amplitudes obtained by the various multipole analyses of the photoproduction data^{/28/},
- the results of various analyses of the electroproduction data^{/29/}.

Our numeric investigation showed only a weak dependence on α_s but a strong sensitivity to the string constant b . To fit the data we could allow only a very small value for $|b| \leq 0.01$ GeV². In our final numerical analysis we put $\alpha_s = 0$ and $b = 0$ and found as the best value for the only remaining constant

$$c = -0.165 \text{ GeV}$$

Thus we calculated the interaction dependent photon-quark coupling with a constant potential, which can only be a coarse approximation. That it works might be due to the fact that we need only averages of the correct potential over a spatial region where the wave functions are significantly different from zero. We shall see that already such a single interquark potential leads via formulae (9a, b) and (11) to very non-trivial properties of the $(1/c)^3$ terms.

b) General description of the relativistic contribution
to the electromagnetic amplitudes

The photoexcitation of a nucleon isobar N^* can be described by the commonly used photon helicity amplitudes

$$(25) \quad A_\lambda(Q^2) = \langle N^*, j_z = \lambda | H_{\text{em}} | N, j_z = \lambda-1 \rangle$$

We consider only transversely polarized photons and let them move in the z-direction. The helicity λ can then take only the values $\lambda = \frac{1}{2}$ and $\frac{3}{2}$. (24) describes real photons ($Q^2 = 0$) and virtual photons ($Q^2 > 0$)^{/31/2}. Each A_λ can be further analyzed in terms of the 14 different amplitudes occurring in eq. (18).

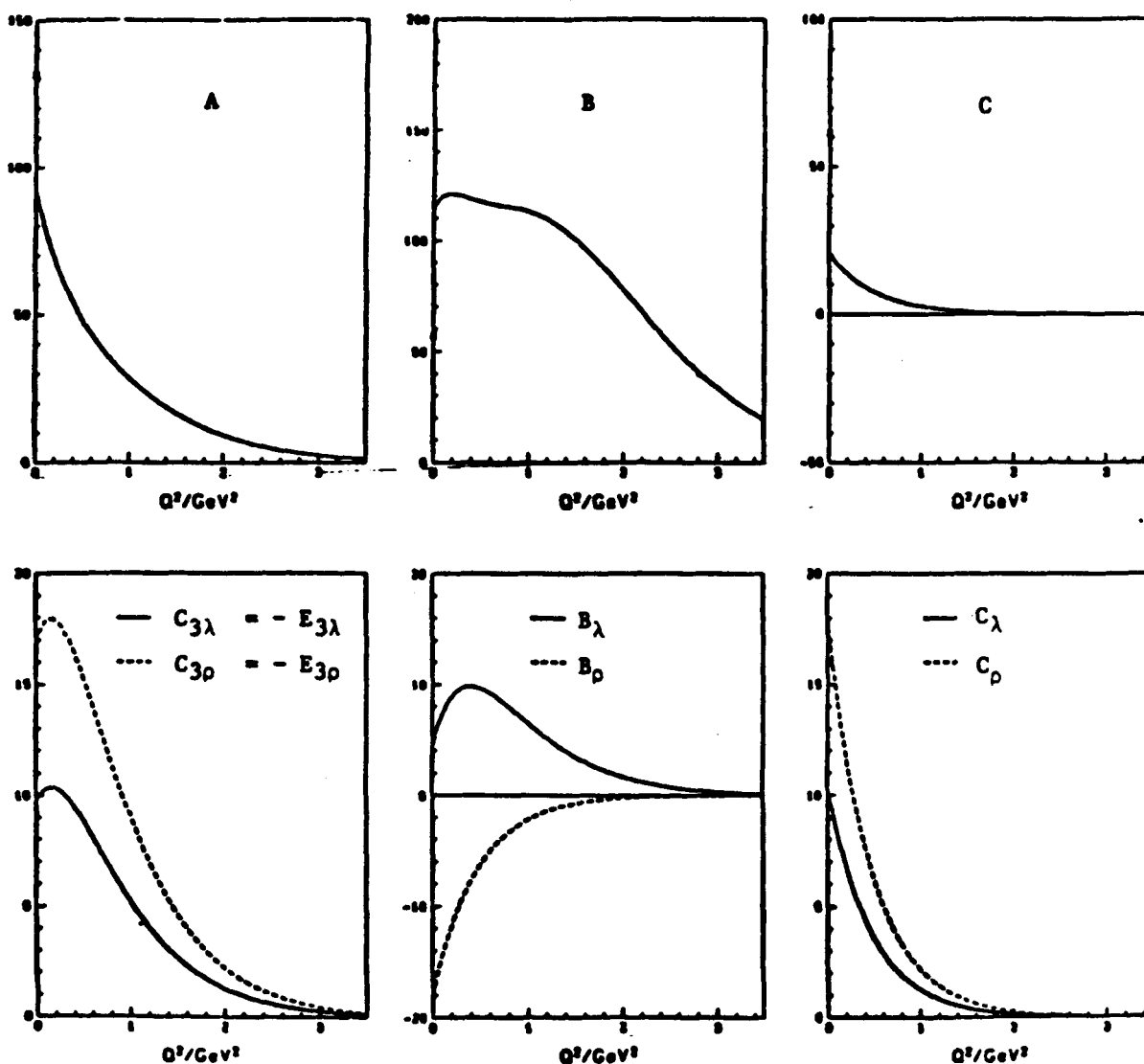


Fig. 3 The reduced matrix elements for the $[70, 1^-]_1$ multiplet in units $[10^{-3} \cdot \text{GeV}^{-1/2}]$

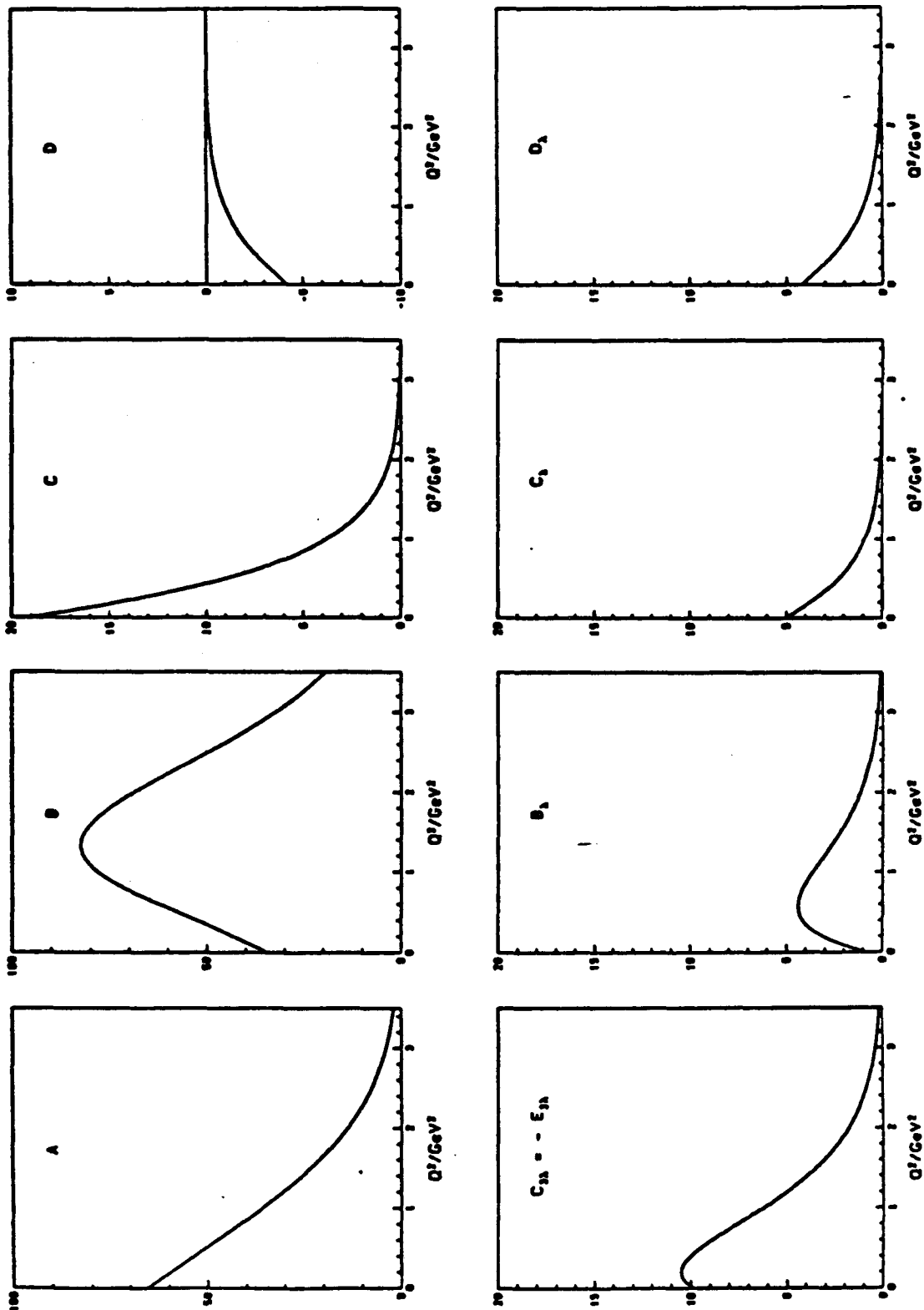


Fig. 4 The reduced matrix elements for the $[56, 2^+]_2$ -multiplet in units $[10^{-3} \cdot \text{GeV}^{1/2}]$

In accordance with the Wigner-Eckart theorem the spin- and isospin dependence of the A_λ can be factorized off and a "reduced matrix element" contains the dynamics. The latter can be characterized by the SU(6)-multiplet to which the N^* belongs. Ref.^{/23/} gives the tables for the generalized Clebsch-Gordan coefficients.

To illustrate our results we present in fig. 3 and 4 the reduced matrix elements for the excitation of the P-wave $[70, L^P = 1^-]_1$ - and the D-wave $[56, L^P = 2^+]_2$ - multiplet. Fig. 3 shows that apart from the D-amplitudes all terms of eq. (18) contribute. Since the D's carry two units of orbital helicity they can be excited only in multiplets with an orbital angular momentum $L \geq 2$. In fact as fig. 4 illustrates D-amplitudes are present in the $[56, L^P = 2^+]_2$.

The leading amplitudes are naturally of B-type (magnetic) and A-type (convection current). The third single quark transition amplitude C is smaller roughly by a factor 5. All other amplitudes have a purely relativistic origin but their values can be as large as the C-amplitude. See e.g. $B_\rho, C_\rho, C_{3\lambda}$ and $E_{3\lambda}$ in fig. 3. In general, the size of the relativistic amplitudes amount to 10% of the non-relativistic ones.

The leading amplitudes A and B have also relativistic contributions, which are quantitatively displayed in fig. 5. Its first column shows the total effects of the different orders in $\frac{1}{c}$. Evidently the second order is rather small while the third order contributes appreciably. The second column gives the 2nd order alone. Here one observes for the B amplitude that the effect of the Foldy-Wouthuysen transformation (FW) which is small by itself is further reduced by the center of mass corrections (CM) having an opposite sign. In the third order similar compensations take place between the three terms: the FW-, the CM- and "potential"-part which are caused by the interaction dependent hamiltonian (9). This is shown in the 3rd and 4th column of fig. 5. The different contributions of order $(1/c)^1$ can be as large as the non-relativistic amplitudes but after summing up all pieces we are down to about 20% of the non-relativistic amplitudes.

Summarizing this subsection we have verified our frequent statement that the relativistic corrections up to order $(1/c)^3$ are originating from different sources and have appreciable magnitudes. But the contributions of different origins - FW, CMS, potential - collaborate in such a way that their total size is less than 20% of the non-relativistic amplitudes. Thus the general features of the non-relativistic quark model are not spoiled.

c) A sample comparison with the data

With our theory we have calculated the nucleon form factors and the transverse photon transition amplitudes for all isobars with oscillator quantum number $N = 0, 1, 2$ and four momentum transfers below $\sqrt{3}$ GeV. Since we are using oscillator wave functions our results have a Gaussian Q^2 -dependence and therefore, we cannot trust them for large Q^2 .

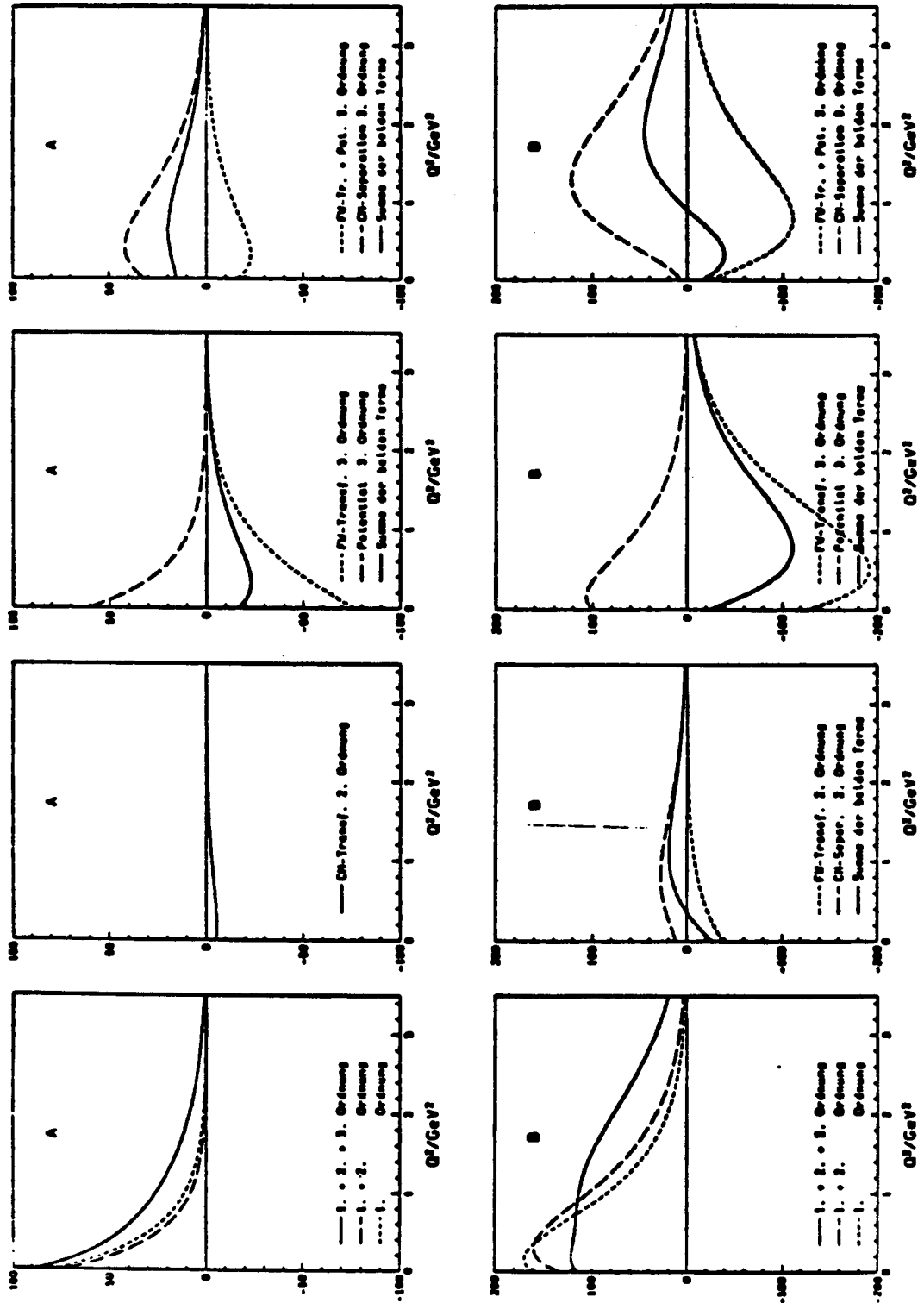


Fig. 5 The different contributions to the reduced matrix elements A, B for the $[70, 1^-]_1$ -multiplet in units $[10^{-3} \cdot \text{GeV}^{-1/2}]$.

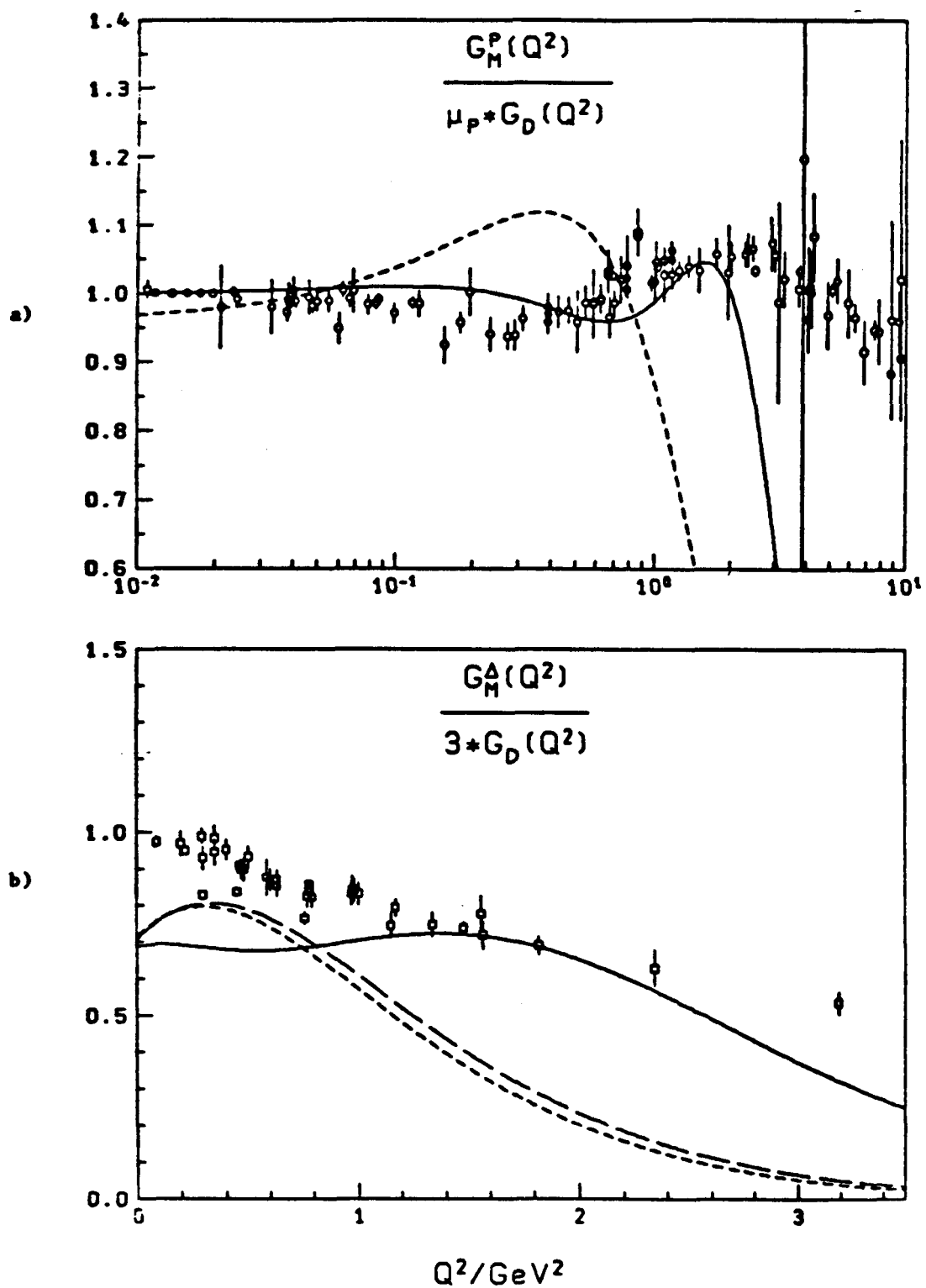


Fig.6 Magnetic form factors for nucleon and Δ -isobar comparing the order in $1/c$.
 ----- 1. Order; ——— 1.+2.+3. Order; - - - - 1.+2. Order

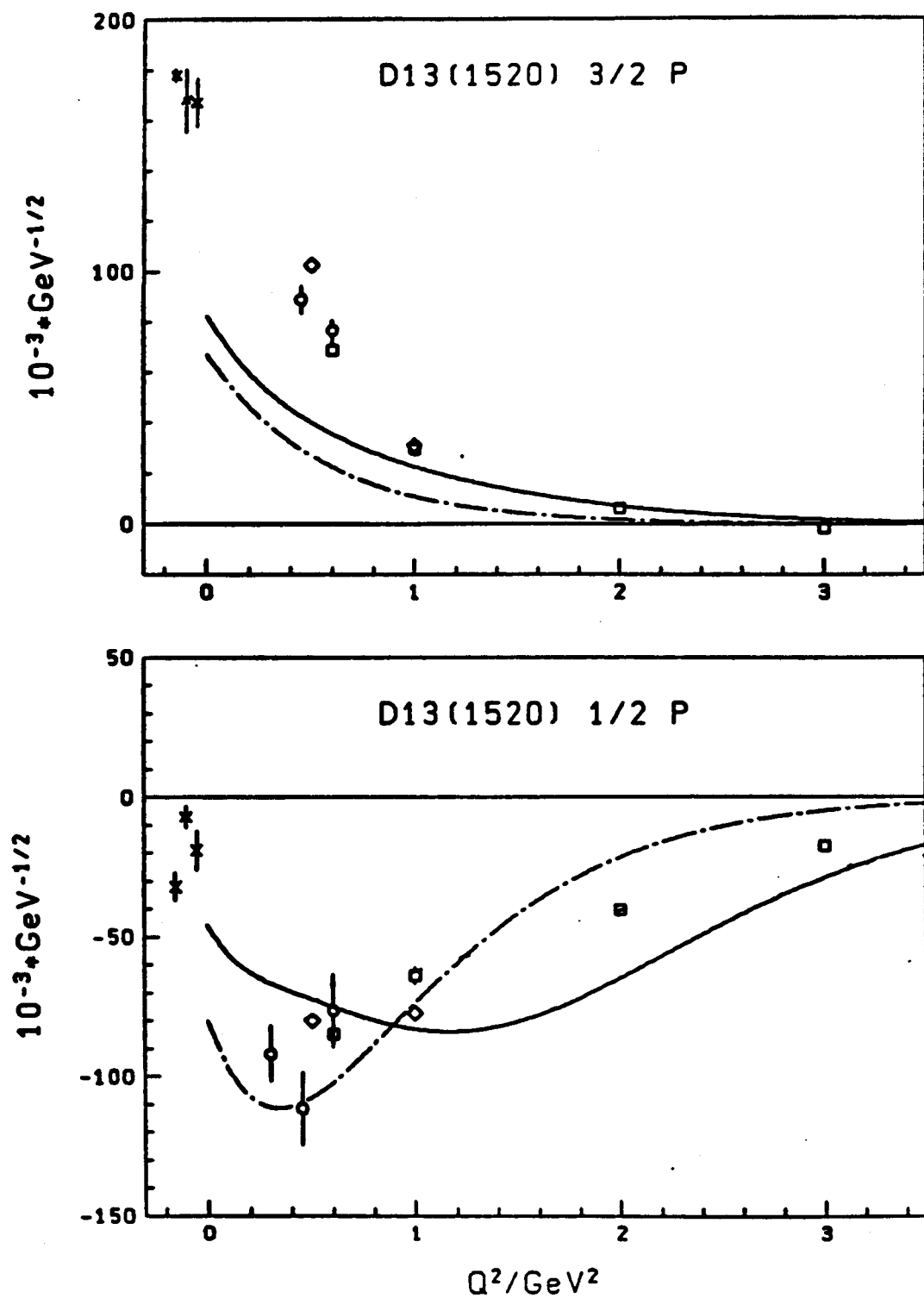


Fig. 7 Transition amplitudes for the $D_{13}(1520)$ isobar.

— Our results - · - Non-relativistic model with configuration mixing

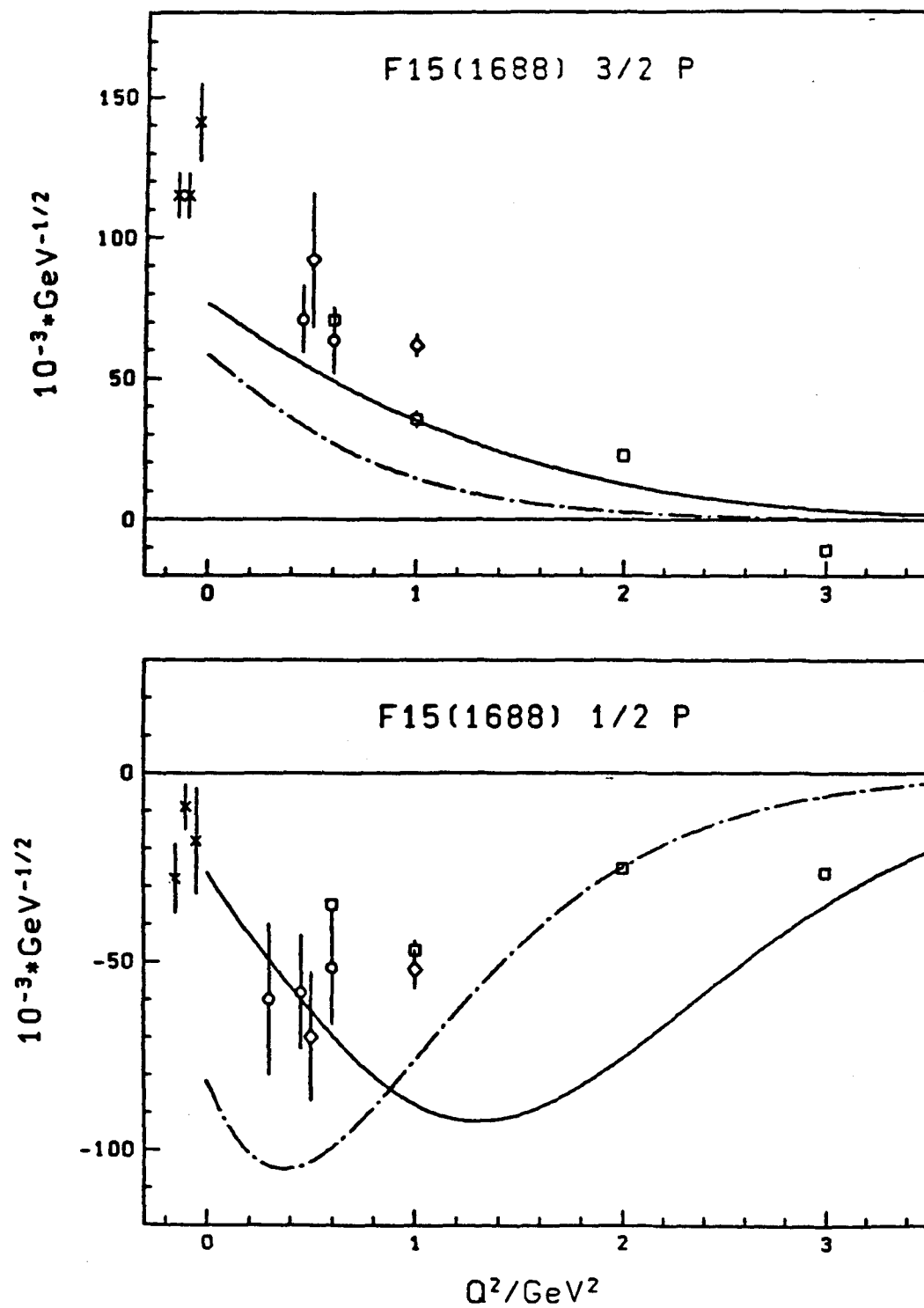


Fig. 8 Transition amplitudes for the $F_{15}(1688)$ isobar.

— Our results - · - · - Non-relativistic model with configuration mixing

We illustrate our results with the magnetic form factor of the proton (fig.6a) and the magnetic transition form factor for the $\Delta(1232)$ (fig. 6b) and two prominent isobars of the 2nd and 3rd resonance region (fig. 7,8).

The proton form factor (fig.6a) is rather well reproduced up to $Q^2 \approx 2 \text{ GeV}^2$. It should be stressed that for this good result the hyperfine configuration mixing (see eq.(22)) was needed. The figure shows the importance of the relativistic corrections already for Q^2 -values between 0.1 and 1 GeV^2 . The results for the Δ -transition form factor (fig. 6b) are good for Q^2 -values between 1 and 2.4 GeV^2 if the $(1/c)^3$ terms are induced. But there remains still the notorious discrepancy for the photoproduction amplitude ($Q^2=0$) which is known since the advent of the SU(6) or the non-relativistic quark model^{/30/}. The relativization gives no new insights into this old problem at $Q^2=0$ but improves the Q^2 -dependence considerably.

For the higher isobars $D_{13}(1520)$ and $F_{15}(1690)$ there is a marked difference between the two helicities (figs.7,8). For $\lambda = \frac{3}{2}$ the relativistic contributions increase the amplitudes systematically where the effect is larger for the heavier isobar. The experimental data lie above the theoretical curves especially for the D_{13} and for small Q^2 -values. But we should keep in mind that the experimental amplitudes given in the figures have been extracted from the raw data by assuming the $SU(6)_W$ single quark structure for the transition matrix elements. For helicity $\lambda = \frac{1}{2}$ the relativistic effects change their sign at about $Q^2 \approx 1 \text{ GeV}^2$ thus reproducing the data for lower Q^2 better than for larger values. The qualitative difference between the two helicities are formally based on the fact that the important B-type amplitude only contributes for $\lambda = \frac{1}{2}$.

A complete representation and a more detailed discussion of all our results will be published elsewhere^{/23/}.

4. OUTLOOK

Our approach to a relativistic photon-hadron interaction clearly demonstrates the importance of relativistic terms. They improve the understanding of the experimental data especially with respect to the momentum transfer dependence, but the relativistic effects remain small enough so that the non-relativistic quark model retains its merits even for the light quarks.

There are certainly many points where our approach should be improved. We mention only two aspects: From a technical point of view the oscillator wave functions should be replaced by more realistic ones. But more importantly, the crude interquark potential should be improved. Due to the importance of the interaction dependent terms in the photon quark hamiltonian we even may learn something about the quark-quark potentials from form factors.

Our result on the presence of non-single quark terms which amount up to 20% has important consequences for the analysis of experimental data: The new types of ampli-

tudes should be taken into account from the beginning.

Finally, our approach should also be worked out for longitudinal photons, which is presently under investigation.

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NUCLEAR EFFECTS IN DEEP-INELASTIC LEPTON-NUCLEUS SCATTERING

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INTRODUCTION

Deep-inelastic lepton-nuclear scattering (DILNS) experiments¹⁻³ have focused the attention to many important aspects of nuclear and particle physics. The main question raised by these experiments is: Where is the boundary between the interpretation in terms of quarks and that in terms of effective nucleons and mesons degrees of freedom? The deviation between the free nucleon structure function (SF), $F_2^N(x)$, and the nuclear SF, $F_2^A(x)$, [the European Muon Collaboration (EMC) effect] has led to many theoretical models which explain the EMC effect by assuming: The nucleon's quark structure changes inside a nucleus, the appearance of multiquark configurations, the enhancement of the "sea", etc.⁴

We consider the DILNS within the conventional nuclear theory taking into account the effects of nuclear binding and two-body short range correlations. We show that the EMC effect is mainly due to nuclear binding. It is also shown that the DILNS data³ for ^{12}C in the $x > 1$ region are well reproduced when two-body short range correlations are taken into account; x is the Bjorken variable for a free nucleon (see refs. 5-7 for details). We add that these nuclear effects were introduced in some papers⁸⁻¹¹ through parameters fitted to the data. In our work we carry out the calculations on the basis of known nuclear data with no free parameters.

We assume that a nucleus is the bound system of nucleonic and mesonic fields and at asymptotically large $-q^2 \gg m_N^2$ the virtual photon is scattered from the nuclear constituents incoherently. We adopt the

independent particle model description of the nucleus. We also take the bound nucleon's SF to coincide with that of the free nucleon.

THE EMC EFFECT

We consider the deep-inelastic scattering (DIS) from single nucleons occupying the nuclear shell characterized by the separation energies, ϵ_λ (negative) and the single particle wave functions ϕ_λ . Writing the tensor $W_{\mu\nu}$, which describes the hadronic part of the DIS cross section, in terms of single-nucleon electromagnetic currents, we find⁵⁻⁷ that

$$F_2^A(x) = \int_x^\infty dz f^A(z) F_2^N(x/z) \quad (1)$$

$$\begin{aligned} f^A(z) &= \sum_\lambda \int \frac{d\vec{p}}{(2\pi)^3} \delta(z - \frac{p_\lambda q}{m_N q_0}) \left| \phi_\lambda(\vec{p}) \right|^2 \\ &= \sum_\lambda \int_{|m_N(z-1) - \epsilon_\lambda|}^\infty dp \frac{p m_N}{(2\pi)^2} \left| \phi_\lambda(\vec{p}) \right|^2, \end{aligned} \quad (2)$$

where $\phi_\lambda(\vec{p})$ is the Fourier transform of $\phi_\lambda(\vec{r})$, m_N is the nucleon mass and p_λ and q are the four vector momenta of the nucleon and the virtual photon, respectively. The sum in (2) is over all occupied single particle states. The physics of the binding energy effect is in taking $p_\lambda = (m + \epsilon_\lambda, \vec{p})$ since the off-mass shell nuclear nucleon has energy ϵ_λ and momentum $\vec{p}$, with scaling variable $x_\lambda = -q^2/(2p_\lambda q)$. Taking the z axis in the direction of $\vec{q}$ we have $p_\lambda q = (m + \epsilon_\lambda)q_0 - p_z q_z$ with $q_z/q_0 = 1$ in the scaling limit. The well-known Fermi smearing effect is obtained from (1) and (2) using $\epsilon_\lambda = 0$. Figure 1 compares the results of our calculation for the ratio $R(x) = F_2^A(x)/[A F_2^N(x)]$ with the data of ref. 2. We have used harmonic oscillator wave functions (HOWF) for ϕ_λ with $\hbar\omega = 45/A^{1/3} - 25/A^{2/3}$, which nicely reproduce the measured root mean square radii of light and heavy nuclei. The values of ϵ_λ are taken from a compilation¹² of experimental data for the $(e, e'p)$ reaction. Note that the average separation energy is $\langle \epsilon_\lambda \rangle \approx -30$ MeV. We take the nucleons in the deuteron to be free and use the

parameterization of ref. 13) for $F_2^N(x)$. We note that the low x behaviour of $R(x)$ is due to nuclear binding and the rise of $R(x)$ for $x > 0.6$ is due to Fermi motion.

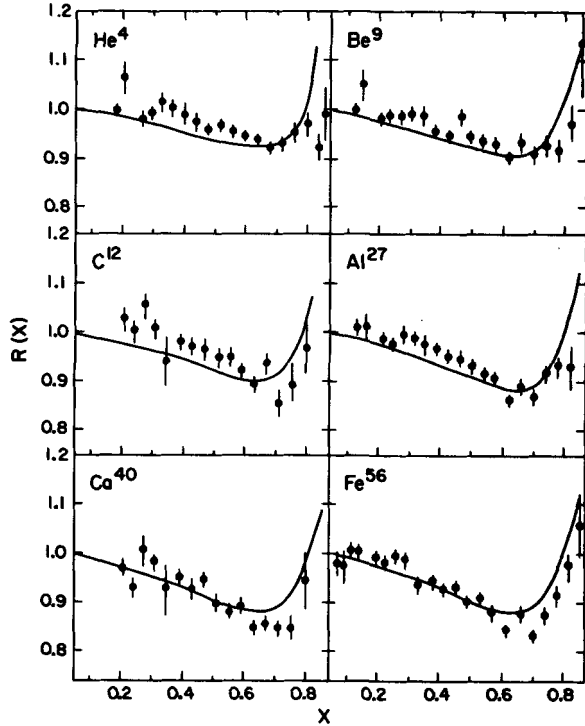


FIG. 1. $R(x) = F_2^A / (AF_2^N)$ as a function of x for various nuclei. The data are taken from Ref. 2. The curves are the results including nuclear binding. See text for details.

THE $x > 1$ REGION

For sufficiently large x , ϵ_λ can be neglected so that $f^A(z)$ can be written in terms of $N(p)$, the momentum distribution of nucleons in the nucleus. We thus write

$$f^A(z) = f_{mf}^A(z) + f_c^A(z) \quad (3)$$

where the mean field contribution $f_{mf}^A(z)$ was calculated above using HOWF and experimental values of ϵ_λ and

$$f_c^A(z) \approx \int_{p_z}^{\infty} \frac{dp}{(2\pi)^2} \frac{p m_N}{4\pi} \frac{(2\pi)^3}{4\pi} N_c(p) \quad (4)$$

is due to short-range correlation. We adopt the normalization $\int_0^\infty p^2 N(p) dp = A$. For the high momentum tail, $N_c(p)$, we use the results of Brueckner theory of finite nuclei with realistic two-body interaction¹⁴, i.e.,

$$N_c(p) = \begin{cases} 0 & \text{for } p < p_c \\ N_c(p_c) \exp [-\beta(p-p_c)] & \text{for } p > p_c, \end{cases} \quad (5)$$

with $\beta = 1.5 \text{ fm}^{-1}$, $p_c = 2 \text{ fm}^{-1}$ and $N_c(p_c) = 0.021A \text{ fm}^{-3}$. In (3), $p_z = \max [p_c, |m_N(z-1) - \langle \epsilon_\lambda \rangle|]$. We note that $f^A(z)$ is normalized according to

$$\int_0^\infty f^A(z) dz = A. \quad (6)$$

The results of our calculation for ^{12}C is compared with the experimental data³ in fig. 2. This clearly demonstrates the important effect of short-range correlation. Note that due to two-body short-range correlations, the DILNS cross section for $x > 1$ is enhanced by arbitrary order of magnitude. Since the magnitude of $N_c(p)$ in (5) is proportional to A , the results for other nuclei are very similar to those shown in Fig. 2.

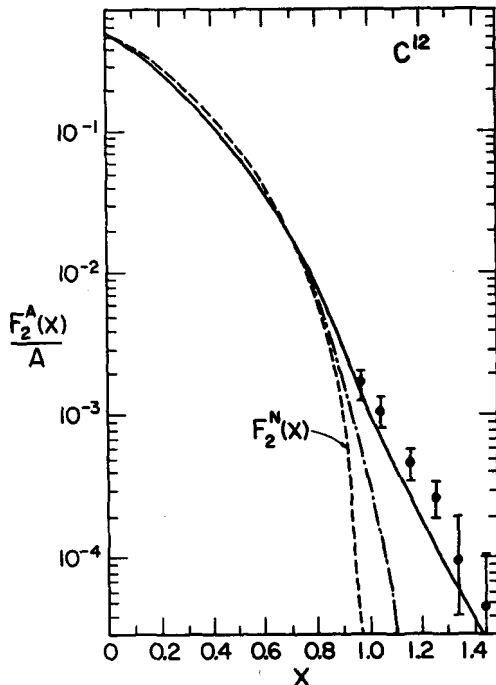


FIG. 2. Plot of structure functions as a function of x . Data are from the Dubna experiment (Ref. 3). The solid line corresponds to our full result, the dash-dotted line to that obtained when the correlation part of the distribution function $f_c^A(z)$ is neglected.

CONCLUSIONS

Before closing, we now mention very briefly other possible effects which should be taken into account in a refinement of the calculations presented above. Using (1) and (2) we have that

$$\frac{1}{A} \int_0^A dx F_2^A(x) - \frac{1}{A} \int_0^A dz z f^A(z) = 1 + \frac{\langle \epsilon_\lambda \rangle}{m_N} \approx 0.97. \quad (7)$$

This violation of the momentum sum rule can be restored by including the contributions of mesons⁵ and the simultaneous scattering of the electron from two nucleons¹⁵. These effects lead to an increase of $R(x)$ for $x \leq 0.3$, improving the agreement with data. The relatively small effects of the deuteron SF and the $Q^2 = -q^2$ dependence of the neutron and proton SF should be included. Due to the smearing effect of the deuteron wave function, one has that $F_2^D(x)/(2F_2^N(x)) > 1$ for $x > 0.6$. Also the problems related to off-mass shell analytic continuation of the scattering amplitude, final state interaction and relativistic corrections to the nuclear wave function should be investigated. However, these effects are not expected to change our main conclusion concerning the significance of nuclear binding, Fermi motion and short-range correlations. Our considerations seem to indicate the surprising stability of the nucleon inside the nucleus. Therefore more accurate data, particularly in the small x region (≤ 0.3), will help to clarify the possible role of pure quark effects.

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DELTA ELECTRO- AND PHOTO-PRODUCTION IN NUCLEI

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As the lowest and most prominent baryonic excitation the delta plays a central role in medium energy nuclear physics. Reactions involving momentum transfers greater than a few fm^{-1} or energies of more than a few hundred MeV cannot be consistently described without taking the delta into account. At momentum transfers in the medium energy range the higher resonances are much smaller, strongly overlapping, and less well described, so that their role in nuclear physics has so far been negligible.

Among the various probes applied to the study of the delta in nuclei, the photon plays a special role since it illuminates the entire nuclear volume, and its couplings are well understood. The incident real or virtual photon can be set to define the delta formation. There are various combinations of final state particles which are sensitive to delta propagation or absorption. For example, the detection of two nucleons with invariant mass near 2070 MeV ($= M_\Delta + M_N$) is related to the absorptive part of the $N\text{-}\Delta$ interaction of the type $\Delta + N \rightarrow N + N$. A series of measurements of the reaction (γ, pp) at delta energies were made at Tokyo (Ho-84), and a study of the reaction ${}^3\text{He}(\gamma, pp)X$ is underway at Saclay (Au-86). This type experiment with real or virtual photons will certainly be studied at CEBAF.

Another example is that of pion production in which a single pion, or a pion and nucleon are detected at kinematics which define the delta final state invariant mass, momentum, or production angle. For neutral pion production the delta, as well as the higher resonances, dominate the cross section. Thus neutral pion production will be an attractive area of study. A few measurements

have been performed in recent years, for example ${}^4\text{He}(\gamma, \pi^0){}^4\text{He}$ at Bates (Ti-84), which strongly supports the delta hole model (Ko-84). However, technical difficulties make it difficult to detect neutral pions efficiently, or with good energy or momentum resolution. Thus, it will be difficult to measure them as part of multi-particle coincidence experiments, except perhaps in conjunction with a very large angular acceptance detector.

On the other hand charged pions can be measured efficiently, with precise determination of momentum and energy, using standard general purpose magnetic spectrometers. However, the charged pion photo or electro production amplitude has the added complication of a large non-resonant background which interferes with the resonant amplitudes, and is dominant for all resonance energies higher than the delta.

The present study concentrates on the potential of doing charged pion photo or electroproduction experiments at CEBAF as a means of studying delta production and propagation in a nucleus. In the 1986 and 1985 CEBAF Summer Study Proceedings, I have reviewed some of the recent and current real photon charged pion photoproduction experiments which are aimed at studying delta dynamics in nuclei (St-86). All of these experiments, can be generalized to virtual photons and higher momentum transfers available at CEBAF. Several potential CEBAF experiments are discussed there. A summary is given here, and the reader is referred to the above reports for more details.

Single arm charged pion photoproduction at delta energies has been measured for nuclei $A \leq 16$ at several laboratories, mostly with bremsstrahlung or a continuous energy virtual photon distribution due to undetected electrons. Currently, the most interesting results are from the reaction ${}^{14}\text{N}(\gamma, \pi^+){}^{14}\text{C}_{\text{gs}}$ (Te-86, Se-86), and ${}^{13}\text{C}(\gamma, \pi^-){}^{13}\text{N}_{\text{gs}}$ (St-84, Ko-86). These cases are interesting because the specific

nuclear structure suppresses the non-resonant amplitudes such that both cases are dominated by delta production in the transition amplitude. Both experiments deviate from expectations of the most current DWIA (Wi-86) and delta-hole calculations (Su-86) at photon energies near the peak of the free nucleon delta. This is an indication that the delta may still not be treated properly in nuclei. An example of the results for the ^{14}N case is shown in Figure 1.

An obvious extension of these experiments is to measure the reaction $A(e, e'\pi)A$, where now σ_L , σ_T , σ_I and σ_p are all present. At $\theta_\pi = 0$ deg. relative to the momentum transfer only σ_T and σ_L are non-zero. The most favorable case to study is $^3\text{He}(e, e'\pi^+)^3\text{H}$, because the cross sections are very high. Using theoretical cross section estimates of (Ti-81) we find that coincidence counting rates are greater than 10^3 per hr. for electrons scattered at forward angles ($\epsilon \sim .9$), and greater than 10^2 per hour for large angle scattering ($\epsilon \sim .4$). For targets such as ^{14}N or ^{13}C the counting rates will be an order of magnitude lower due to small cross sections.

Less data exists for the $(\gamma, p\pi^-)$ experiments because of duty factor and energy constraints of current electron accelerators. The only series of measurements to date were done on ^2H at Saclay several years ago (Ar-78). By fixing the momentum and energy of the outgoing particles they were able to minimize the normally dominant quasi-free production and concentrate on reactions in which both nucleons share the energy of the incident photon. An example of the limited results they were able to obtain, shown in Figure 2, can be explained by a two pion mechanism in which one of the pions is off shell and reabsorbed on the other nucleon. This mechanism is related to the absorptive part of the Δ -N interaction.

The generalization of this reaction is $A(e, e'p\pi^-)X$. Several possible examples of this type experiment, which use the full range

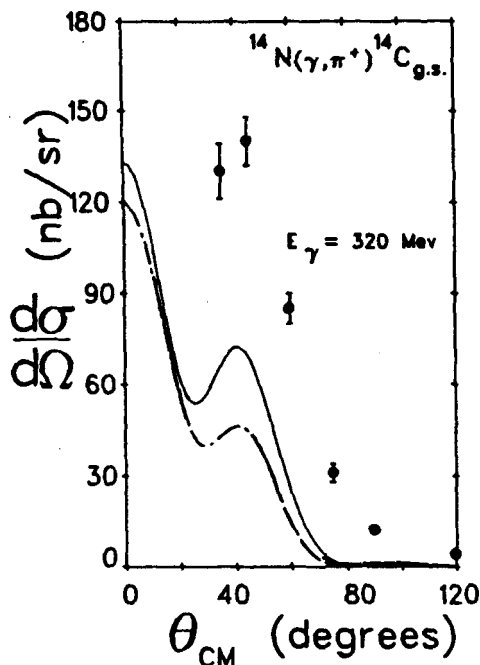
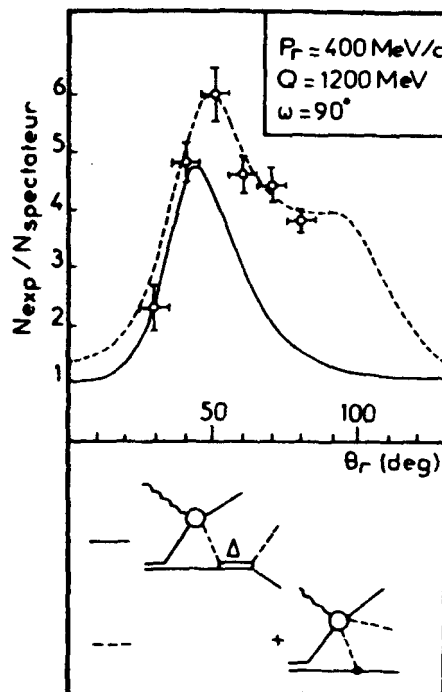


Figure 1 Cross section for the reaction $^{14}\text{N}(\gamma, \pi^+)^{14}\text{C}_{\text{gs}}$ at $E_\gamma = 320$ MeV. Experimental data are from (Te-86). Hatched curve is due to a calculation of (Wi-86) using the standard DWIA theory. The solid curve is due to (Wh-86), using an improved delta amplitude and distorted pion wave function.

Figure 2 Cross sections for the reaction $^2\text{H}(\gamma, \text{pp}_r)\pi^-$ at high recoil momentum P_r as a function of θ_r , measured at Saclay (Ar-78). The deviations from the quasifree model can be explained (La-81) by single pion rescattering and a meson exchange mechanism shown in the diagrams under the figure.



of the CEBAF energy and duty factor are discussed in (St-85). A simple Fermi gas model, together with realistic spectrometer parameters give quasi-free counting rates near 10^3 per hr. The most interesting experiments may be at kinematics which strongly deviate from the quasi-free ones, in which more than one of the nucleons participate in the formation and decay of the delta, and quark degrees of freedom may be important. These will have much lower counting rates. Due to the strong interaction of pions in the nucleus at delta energies, this type experiment may be feasible only with very light nuclei such as ^2H , ^3He and ^4He .

By making simulations based upon the solid angle, momentum acceptance, and resolution properties of the Medium Resolution Spectrometers (MRS) proposed for Experimental Hall C (Bu-86), it is found that there are feasible experiments of the type $(e, e'\pi)$ and $(e, e'p\pi)$ which can be done with good counting rates. However, since most of the simulations are based on the naive Fermi gas model with no final state interactions, realistic theoretical calculations are needed to get better estimates, and to determine if the experiments can actually yield the desired information.

An important requirement of the MRS is that in all cases it will be necessary to detect at least one hadron at forward angles, and it will be useful for one of the hadron spectrometers to be able to accommodate the backgrounds due to a high electron beam current.

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SOME THEORETICAL AND EXPERIMENTAL ASPECTS
OF THE DEUTERON PROBLEM

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The essential problems of the intermediate energy nuclear physics appear even in the investigation of the simplest nucleus- deuteron. We shall review here some of our results and list the most interesting theoretical and experimental problems arising in the investigation of the deuteron electromagnetic structure. All comments given below represent our proper point of view on the role of different degrees of freedom in deuteron and the importance and priority of different experiments at low and intermediate energy. Let us emphasize that we shall not discuss the six quark deuteron structure since we believe that the description of the present situation in elastic and quasi-elastic ed-scattering needs no quark degrees of freedom.

I. Theory.

I.1. Some results.

The main ideas of our approach to the deuteron electromagnetic form factors (DEMFF) and the essential results are summarized in our report ^{/1/}. Here we shall list some new results concerning several aspects of the deuteron problem.

Let us begin with the results obtained in the frame of an approach to relativistic impulse approximation (RIA) with the relativistic electromagnetic current operator and nonrelativistic deuteron wave functions (DWF)^{/2/}. In the case of Paris potential this approach gives the structure functions $A(Q^2)$ in accord with the experimental data; as to $B(Q^2)$ it differs from the data just by the generally accepted meson exchange current contribution. The relativistic corrections are positive as well as for the quadrupole moment Q_d ($Q_d = 0,2794$ in nonrelativistic and $Q_d = 0,2856$ in relativistic case). Note that in nonrelativistic limit the equations^{/2/} are the same as well known Jackson-Maximon equations for DEMFF in the impulse approximation.

On the base of the whole set of experimental data on the elastic ed-scattering at low $Q^2 \leq 1 \text{ fm}^{-2}$ (see /3/ for references) we have performed the direct calculation of the root mean square radius of the deuteron charge distribution without any dynamical model. The radius occurs to be $\langle r_{ch}^2 \rangle^{1/2} = r_c = 2.11 \pm 0.01 \text{ fm}$. It is impossible to make the dispersion smaller because of the present accuracy of experimental data. It seems us that the more accurate values of r_c obtained by various extracting schemes (e.g. $r_c = 2.116 \pm 0.006 \text{ fm}$ in /4,5/) can not be regarding as following directly from the experiment. The real dispersion of experimental r_c covers almost all theoretical results obtained with the use of realistic and even model NN - interaction potentials. The accuracy of present experiments does not provides a test for theoretical predictions so that the unique dynamical model of deuteron can not be choosen. Let us note in addition that the problem of accuracy of r_c is very important for the theoretical investigation of corrections to Lamb shift in deuterium due to finite nucleus size.

In recent years a number of papers (see e.g. /6-8/) concern a class of unitary equivalent transformations of radial S-, D-deuteron wave functions $u(r)$, $w(r)$ at small distances ($r < 1 \text{ fm}$): $u, w \rightarrow \bar{u}, \bar{w}$ and the effect of these transformations on DEMFF $G_{id}(Q^2) \rightarrow \bar{G}_{id}(Q^2)$, $i = c, M, Q$. On the base of Paris DWF we have performed the full analysis of the /6-8/ transformation action on the whole set of observables in ed-scattering. Although we could not prove the general statement about the DEMFF invariance ($G_{id} = \bar{G}_{id}$) (such a statement had been done by some authors), we confirm the weak sensitivity of all of DEMFF to the transformations /6-8/ in the wide region of Q^2 . In the case of deuteron magnetic (μ_D) and quadrupole moments the situation becomes quite different. It is possible /9/ that the D-state weight p_D can occur to be a nonobservable parameter so that at present the study of the μ_D value at $p_D \rightarrow \bar{p}_D$ seems to have no sense and the only possibility is to analyse Q_d (as to r_c - see above). For the set of transformations /8/ the respective variation of Q_d is much greater than the accuracy of the experimental determination of its value. For Paris potential ($Q_d = 0.279 \text{ fm}^2$) the transformations 6, 7 of the table I of /8/ give $Q_d = 0.267 - 0.275 \text{ fm}^2$, i.e. make the agreement of theoretical and experimental results

even worse. Two conclusions follow from this consideration. First, the Q_d invariance requirement eliminates some of $u, w \rightarrow \bar{u}, \bar{w}$ transformations with the radius $r = 0.8 - 1$ fm. Second, since there are some doubts^{/10/} about the correctness of the extracting of the value $Q_d = 0.286 \pm 0.0015$ fm² from the radio-frequency experiments on molecular deuterium, it is desirable to perform some new measurements of Q_d directly by different nuclear experiments. Let us point in this connection that there is a preliminary information^{/11/} announcing the new value of $Q_d = 0.282 \pm \pm 0.019$ fm² obtained by the analysis of subcoulomb scattering of polarized deuterons in the reaction $\vec{d} + {}^{208}\text{Pb} \rightarrow \vec{d} + {}^{208}\text{Pb}$. We shall more discuss Q_d in section II.

In nonrelativistic limit our approach^{/12/} gives in the frame of RIA the following constructive formulae for the s-state DWF $u(p)$:

$$u(p) = r^{-1} \Delta(2p^2/m), \quad (1)$$

$$\Delta(k) = \frac{\sin \delta(k)}{k} \left(\frac{k}{k_0} \right)^{\frac{\delta(k)}{\pi}} \exp \left\{ \frac{k+k_0}{\pi} \int_0^\infty \frac{dk' [\delta(k') - \delta(k)]}{(k'-k)(k'+k_0)} \right\}. \quad (2)$$

In (2) $\delta(k)$ is the ${}^3S_1$ neutron-proton scattering phase. The normalization constant r in (1) is determined by the condition $\int_0^\infty u^2(p) dp = p_s$. Here p_s is the s-state weight. The representation (1) is obtained by comparison of nonrelativistic limit^{/13/} with the standard Jackson-Maximon construction^{/14/} for charge DFF in p-space. It is known that for the present experimental and theoretical situation in NN-scattering the permanent appearance of new sets of phase shifts and the refinement of the previous ones is specific. This fact is the consequence not only of progress in the obtaining of new experimental data on NN-interaction but also of the fundamental ambiguities in the phase analysis itself. At present there are 6-7 different realistic ${}^3S_1$ -phase shifts in the large energy region. So (1) gives at once a set of realistic nonrelativistic DWF^{/15/} which are difficult to obtain by any other way. As the new phase analysis appears the simple equations (1), (2) can be used for the immediate derivation of new DWF without use of cumbersome construction of realistic NN-potentials. Let us note that in our approach^{/12/} the relativisation (full or even to first order in $(v/c)^2$) of the

nonrelativistic eq.(17) of the paper^{/13/} breaks down immediately the factorisation typical for Jackson-Maximon equation, so it becomes impossible to define and obtain any representation for the relativistic DWF. So the deuteron treatise in the frame of our relativistic approach enables us to conclude that the numerous relativisation ansatz's for DWF can not be fully adequate.

In the frame of the approach^{/12/} we have also calculated the relativistic corrections (RC) to the S-wave part of nonrelativistic deuteron magnetic moment $\mu_d^S = \mu_n + \mu_p$. This problem is rather of interest. In fact, the true value of the D-state weight in the deuteron is unknown, it differs from 4% to 7%. A number of present experiments on NN-scattering demonstrates a permanent tendency towards systematic decrease of the tensorial component of NN-forces, i.e. towards the lowering of the D-wave weight in deuteron: $p_d = 7.53\% \rightarrow 7\% \rightarrow 6\% \rightarrow ?$ For the most popular among realistic NN-potentials, the Paris potential, $p_d = 5.77\%$. For the Bonn potential the value of p_d is extremely low: $p_d = 4.27\%$. It is possible that the value of p_d shall be even lower in future and the limit is unknown to-day. Moreover, there is a possibility to pass (via the Dyson transformation) from any "pure" variant to a linear combination of pseudoscalar and pseudovector variants of interaction with an arbitrary mixing parameter (the unitary freedom parameter) so that the p_d value itself can occur to be an additional parameter one needs for calculations, with no physical sense^{/9/}. That's why one needs first to analyse the RC to μ_d^S in the frame of the RIA to estimate the role of auxiliary degrees of freedom in deuteron and (or) the relativistic enhancement of S-wave deuteron magnetic moment to fit theoretical and experimental values.

The RC δ to μ_d^S is of the form

$$\mu_d^S(\text{rel}) = (1 + \delta) \mu_d^S(\text{nonrel}) \quad (3)$$

$$\delta = \frac{\int_0^\infty \Delta^*(E) \sqrt{E} dE}{(2M + \sqrt{2M(2M+E)})} \bigg/ \int_0^\infty \Delta^*(E) \frac{dE}{\sqrt{E}}, \quad (4)$$

and $\Delta(E)$ is given by eq.(2). The value of δ was calculated in ^{/16/} for all experimental $^3S_1$ -phase shifts and is equal to $(4 - 4.5) \cdot 10^{-3}$, so that the correction to the deuteron magnetic moment is

$$\delta\mu_d^s = (0.35 - 0.40) \cdot 10^{-2} \text{ n.m.} \quad (5)$$

Let us do some comment. 1. The value of $\delta\mu_d^s$ depends on the np-system dynamics rather weakly. 2. $\delta\mu_d^s > 0$, so that in the case of any "familiar" D-state weight the theoretical value of μ_d shifts towards the experimental one. For Paris potential the correction (5) reduces the discrepancy approximately twice. For Bonn potential the situation is quite different. 3. RC to μ_d in S-wave deuteron may be of considerable value (e.g. our correction (5) is much greater than the correction $\delta\mu_d^s = 0.25 \cdot 10^{-2}$ in^{/17/}). Let us note in this connection that the RC to S-state deuteron has been considered as far back as in 1947 in the classical paper^{/18/}. Let us emphasize in conclusion that, to our opinion, this problem needs more investigation.

In discussion of standard approaches to the RIA it is important to note that there are^{/19-20/} quite different points of view on the contributions of different degrees of freedom (the off mass shell (MS) effect for the nucleon-spectator or for the interacting nucleon) to DEMFF. The "asymptotic" analysis of these problems can be easily performed^{/21/} by use of model relativistic (R) WF proposed in^{/22/}:

$$\Psi(p, P) \sim \exp [a(p^2 - 2\lambda(pu)^2)] \quad , \quad (6)$$

here $u_\mu = P_\mu / M_d$ is the velocity of the system as a whole, λ - an arbitrary parameter. The choose of RWF in the form (6) guarantees automatically that all relativistic effects of "contraction" and "recoil" and a part of retardation effects are taken into account. That is why the RWF of this kind are usually exploited in the calculation of hadron-hadron high energy scattering. In nonrelativistic limit the WF (6) becomes a gaussoid and any realistic DWF can be approximated by a linear combination of gaussoids. In the relativistic case the linear combination of the functions (6) can be a good approximation to any known RDWF. Let us consider now the one-loop RIA integral for the scalar S-wave deuteron in the general case when both of nucleons are out off MS with the npd-vertex function $V(p, P)$ corresponding to (6):

$$V(p,P)/[(\frac{1}{2}P + p)^2 - M^2][(\frac{1}{2}P - p)^2 - M^2] \sim \exp[a(p^2 - 2\lambda(pu)^2)]. \quad (7)$$

For the structure function CDEFF asymptotics we obtain immediately the result well known from the automodel approach and from the euristical relativisation schemes

$$D_c(Q^2) \sim (\frac{1}{\lambda} a) Q^{-2}, \quad Q^2 \rightarrow \infty \quad (8)$$

Let us consider now the one-loop integral in the case when the nucleon-spectator is on the MS. In this case the asymptotics becomes

$$D_c(Q^2) \sim (\frac{1}{\sqrt{\lambda} a}) Q^{-2}, \quad Q^2 \rightarrow \infty \quad (9)$$

i.e. only the preasymptotical multiplier is changed. One can show that the on-MS-limits of the positive and negative energy nucleon-spectator states give the same contribution in (9). Note that the calculation in^{/19/} has been done neglecting the negative energy state contribution. Let us consider now the third case when the interacting nucleon is on MS and the nucleon-spectator is off MS. The calculations show that in this case the contribution to the asymptotics tends to zero

$$D_c(Q^2) \sim (\frac{1}{\sqrt{\lambda} a}) Q^{-2} \exp[-aQ^2(\frac{13}{4}\lambda - 1)], \quad Q^2 \rightarrow \infty \quad (10)$$

The comparison of asymptotics (8), (9), (10) leads to the conclusion that the on-MS-limit of nucleon-spectator does not lead to the loss of generality (contrary to the result of^{/20/}).

Let us note in conclusion that in the case of the vector deuteron we have faster decreasing than in scalar case (8):

$$D_c(Q^2) \sim (\frac{1}{\lambda a}) Q^{-4}, \quad Q^2 \rightarrow \infty$$

It means that the correct relativisation restores the sufficiently slow decreasing regime of DEMFF similar to realistic nonrelativistic regime. This fact can be useful for the description of high Q^2 experimental data in the frame of standard RIA.

I.2. Selected topics.

The progress in the following domains is desirable.

1. The decrease of arbitrariness in the calculation of elastic and nonelastic DEMFF due to possible variation of the nucleon EMFF (this concerns particularly the neutron charge FF

$G_{En}(Q^2)$). At present this arbitrariness is of very high degree especially in the region of mean and large momentum transfers.

2. The investigation of effects in γNN -vertex in the case of far off MS effect for the interacting nucleon. These effects can radically change the value of DEMFF in the frame of RIA as can be seen from the preliminary analysis^{/23/}.

3. The clarification of the nature of RC in the obtained variants of the RIA. Why the different approaches to the DEMFF theory give even different signs of RC? Is it possible to formulate any deuteron concepts common for relativistic and nonrelativistic cases?

4. The relativisation of meson exchange currents (MEC) theory of deuteron. This theory at present is essentially nonrelativistic. This problem is particularly important for the intermediate energy region because at low energy one has the well known low energy theorems for pions and vector dominance model and at high energy one has quantum chromodynamics.

II. Experiment. What is to be measured and what for?

It should be likely to perform the following experiments (with nonpolarized as well as with polarized particles).

II.1. Elastic electron-deuteron scattering.

1. The detailed measurements of the longitudinal part of $A(Q^2)$ of the elastic ed-scattering cross section at small $Q^2 \leq 1 \text{ fm}^{-2}$, particularly at $Q^2 \leq 0.044 \text{ fm}^{-2}$. The information on $A(Q^2 \rightarrow 0)$ is essential for reliable extraction of two the most important quantities: the root mean square radius of the deuteron charge distribution $\langle r_{ch}^2 \rangle^{1/2}$ and the slope of the neutron CFF $G'_{En}(0)$. The experimental situation is comparatively unfortunate for the deuteron: in the cases of ep- and $e^3\text{He}$ - scattering there is a much greater number of experimental points than in the case of ed-scattering.

2. The accurate measurement of the magnetic (M) DFF $G_{Md}(Q^2)$ at small $Q^2 \leq 1 \text{ fm}^{-2}$ at the same Q^2 as $A(Q^2)$. The existing experimental data on $G_{Md}(Q^2 \leq 1 \text{ fm}^{-2})$, few in number, were obtained long ago and are of low accuracy. The importance of such a measurement is caused by the following facts. For the extraction of $G_{En}(Q^2)$ from $A(Q^2)$ one needs first of all to extract correctly the MDDFF contribution $2/3 \eta G_{Md}^2$. At present for this purpose the scaling $G_{Md}/2\mu_d = G_{Cd}$ is used, the experimental values of G_{Md} in appropriate points of Q^2 being unknown.

The scaling itself does not follow from theory. Our experience is that the value of $G'_{\text{En}}(0)$ is rather sensitive to the method of extracting of G_{Md} from A . That is why one needs to measure G_{Md} at $Q^2 \leq 1 \text{ fm}^{-2}$ and $A(Q^2 \leq 1 \text{ fm}^{-2})$ to extract correctly $G'_{\text{En}}(0)$.

3. The measurement of G_{Md} at high Q^2 . It is likely to increase twice the interval $0 < Q^2 \leq 28.1 \text{ fm}^{-2}$ (s.f./24/) up to $Q^2 \sim 50-60 \text{ fm}^{-2}$ to localise "fine structure" of $G_{\text{Md}}(Q^2)$ (the zero position, the altitude and the position of secondary maximum). This program, aside from the obvious causes, is interesting because of the possibility to obtain from the reactions on nonpolarized particles ($\text{ed} \rightarrow \text{ed}$, $\Theta_e = 180^\circ$) the same information one has to derive from the reaction on polarized ones (the fine structure G_{Cd} in $\text{ed} \rightarrow \text{ed}$ reaction). The G_{Md} measurement should be rather fruitful for this purpose because one shall expect the progress (in Q^2) of polarization experiments to be very slow.

4. The detailed measurements of the polarization tensor $T_{20}(Q^2)$ of the recoil deuterons in the $\text{ed} \rightarrow \text{ed}$ reaction at small $Q^2 \leq 1 \text{ fm}^{-2}$. Since $T_{20} \sim G_{\text{Cd}}$, such measurements extrapolation to $Q^2 = 0$ gives one more independent value Q_d through a nuclear experiment (see the discussion in I.1). The results of recent $T_{20}(Q^2)$ measurement /25/ can not be used for this purpose.

5. The $T_{20}(Q^2)$ measurement in the second nontrivial region $Q^2 > 15 \text{ fm}^{-2}$ (the comments are given below). It is likely to measure T_{20} at the same points as A .

6. The recoil deuteron polarization vector $P(Q^2)$ measurement on polarized electron beam $\text{ed} \rightarrow \text{ed}$ under the same condition as in the preceding case. The experiments 5,6 are important to define $G_{\text{Cd}}(Q^2)$ in the region $Q^2 \gtrsim 16 \text{ fm}^{-2}$ where one expects to see the fine structure (the presence of zero) in $G_{\text{Cd}}(Q^2)$ on the base of traditional approaches. It is however hardly probable that the deformation of the standard fine structure (the absence of zero) in $G_{\text{Cd}}(Q^2)$ shall unambiguously indicate that the six-quark (6q) component in deuteron is essential. The smoothing off the minimum in certain region of Q^2 (e.g. near $Q^2 = 16 \text{ fm}^{-2}$) can be essentially a two-particle effect (the arbitrary choosen NN-potential; in terms of potential description it is sufficient for this purpose to choose the Graz-II potential). To-day the main trend in phase analysis is to reduce the repulsion in NN-interaction at small distance, the zero of $G_{\text{Cd}}(Q^2)$ shifting to the right

from $Q^2 = 16 \text{ fm}^{-2}$ to $Q^2 = 23 \text{ fm}^{-2}$. In the extreme case of weak attraction at small distance (and this case can not be excluded!) the zero in G_{Cd} disappears. The usefulness of the mentioned experiments for $G_{Md}(Q^2)$ fine structure (associated with $Q^2 \sim 2 (\text{GeV}/c)^2$) is even greater. It is not clear at present what experiment shall be the crucial test for the $6q$ -contribution to the deuteron structure in intermediate and high Q^2 regions.

II.2. The deuteron electrodisintegration $ed \rightarrow e'np$.

1. The measurement of the cross section $d^2\sigma/dE'_e d\Omega_e$ as a function of Q^2 at $\theta_e = 180^\circ$ in the neighbourhood of quasielastic peak $E'_e \simeq (E'_e)_p$. At present there is no systematic data of such a kind. 2. The same cross section measurement near the deuteron disintegration threshold $E'_e \leq (E'_e)_{th}$ at $Q^2 > 18 \text{ fm}^{-2}$. The experiments 1,2, if performed, would provide a new information on the relative contribution of meson and nucleon degrees of freedom to the deuteron structure. Our opinion is that the generally used conclusion about the dominance of the isovector ($2N$)-MEC contribution to the cross section of 2 is to be analysed in details. 3. The measurement of the nucleon polarization vector $\vec{P}_n$ in the disintegration of the nonpolarized deuteron by the nonpolarized electrons $d(e, e'N)N$ in the region between the quasielastic peak and the deuteron disintegration threshold $(E'_e)_p < E'_e < (E'_e)_{th}$. To our surprise we found no experimental data (even on protons) although the quantity $\vec{P}(ed \rightarrow e'NN)$ itself is discussed from 1959. The theoretical analysis of $\vec{P}_n(ed \rightarrow e'np)$ structure shows that this quantity mixes up the following four problems: 1) The choose of the most adequate np-scattering phases. 2) The nucleon degrees of freedom contribution to deuteron. 3) The possible existence of dibaryon resonances (at present the base for the investigation of this problem is arbitrarily narrowed and limited mainly by the measurement of the proton polarization vector in deuteron photodisintegration). 4) The $G_{En}(Q^2)$ extraction. 4. The measurement of the recoil neutron polarization vector $\vec{P}_n$ in the electrodisintegration of nonpolarized deuteron by polarized electrons $d(\vec{e}, e'\vec{n})p$ at $E'_e \simeq (E'_e)_p$. The $\vec{P}_n$ is almost deuteron structure independent so that the $\vec{P}_n$ measurement is a convenient method to obtain $G_{En}(Q^2)$. The

The measurements in the neighbourhood of $Q^2 = 18 \text{ fm}^{-2}$ are of the particular importance; it is interesting to establish the validity of the well known relation $G_{En} = -\mu_n \tau G_{Ep}$, G_{En} being maximal at the mentioned Q^2 ($G_{En} \approx 0.1$).

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RECEIVED 1984 08
1984 08

Secondary Beams Possibilities at CEBAF

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INTRODUCTION

CEBAF's high-power primary electron beam is a potential source for creating secondary beams. CEBAF's 100 per cent duty factor and commitment to build large acceptance detector systems make secondary beams particularly interesting. Early in the project, work was begun on producing secondary tagged photon beams. This work has progressed very favorably. (1,2,3,4) This paper addresses the possibilities for other secondary beams.

MOTIVATION

Multi-GeV nucleon and particle physics can be investigated by any particle with appropriate energy. Electrons see only the quarks, not the gluons. Secondary beams of muons have the same features as the electrons with the additional advantage of less background due to radiation effects. Secondary beams of pions, kaons and neutrons see both quarks and gluons, with various quarks in their own makeup. These particles frequently induce reactions complementary or similar to those induced by electrons. To understand nuclear and particle physics, all of these beams as well as others are required.

The use of the same detector and target with more than one incident type of particle can help to remove experimental uncertainties. For example, CEBAF is planning a large acceptance detector to be used in conjunction with a tagged photon beam of 10^7 photons per second. Hadronic beams with rates on the order of 10^5 particles per second are expected to produce approximately the same triggering rate in the detector system as the photon beam.

There are several questions to be addressed for secondary beams at CEBAF:

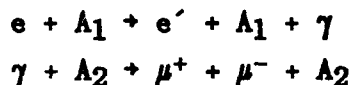
- o Beam intensity

- o Beam purity
- o Beam quality
- o Similar beams available at other facilities
- o Physics usefulness.

In the following sections these questions are discussed for several different secondary particles. One of the questions can be answered in general. In that all of the beams considered are secondary, the beam purity at CEBAF will be similar to that obtained at other accelerator facilities. Electrostatic and/or rf separators, differential production and/or absorption, time of flight, decay in flight, etc., are used to enhance the purity.

MUON BEAM

A primary electron beam dominantly produces muons via a two-step process:



The target nuclei, A_1 and A_2 , provide the requisite coulomb field.

If the electron beam is longitudinally polarized, then the bremsstrahlung will be polarized. Likewise the muons will also be polarized. Figure 1 qualitatively indicates the polarization P for the photons.

100% Longitudinally Polarized Electrons Incident

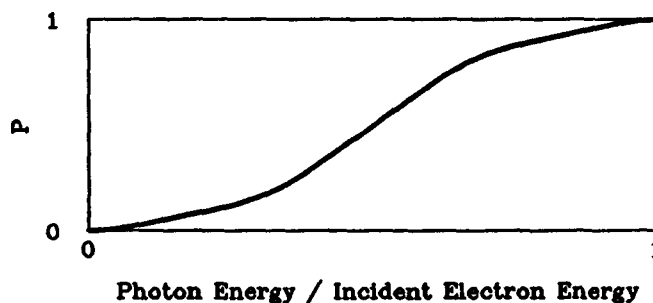


Figure 1: Qualitative Bremsstrahlung Polarization

A similar curve gives the muon polarization.

The production target for muons is typically tungsten for 3-5 radiation lengths, followed by 15-35 radiation lengths of other material. Most of the higher-energy muons are produced in the first 3-5 r.l. of tungsten. The extra thickness of production target is to eliminate the electrons and to some extent the pions. Three to five r.l. of W translates to 1-1.7 cm in space.

CEBAF's high-quality electron beam can have a transverse extent of 0.1-0.5 mm and a very narrow angular divergence. The bremsstrahlung photons are at angles on the order of $1/\gamma_e$ where $\gamma_e = E_e/m_e$, m_e is the electrons mass and E_e is the electrons incident energy. The muons will be produced in a volume roughly 0.1-0.5 mm in diameter and 15 mm long.

In the several GeV/c region the muons will mostly be inside a cone of less than $\pm 1^\circ$. The size of the cone comes from $1/\gamma_\mu$ of the muons and multiple scattering in the tungsten. Figure 2 shows the muon production geometry.

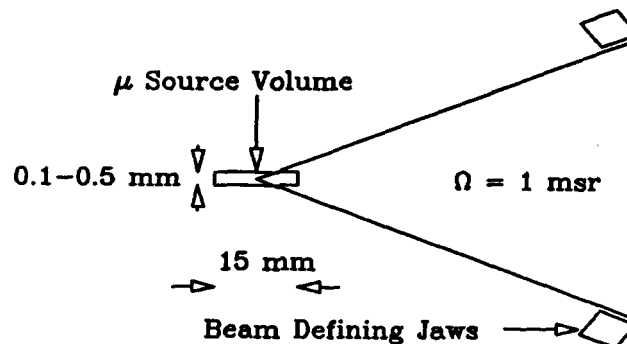


Figure 2: Muon Production

The small production volume and solid angle for the muon beam translates into a small transverse emittance of $\epsilon_T = 1 \times 10^{-5}$ m rad. This is much better than could be achieved in this energy region for muons coming from pion decay. Small targets can also be used if the beam is refocused to the size of the production region.

Beam purity is expected to be quite high as the 20-40 radiation length production target will degrade any electrons and absorb most pions if the last 15-35 radiation lengths are made of appropriate material. Secondary absorbers along the beam line can also be used.

Muon beam rates have been calculated according to the methods of Refs. 5 and 6 and are shown in Figures 3 and 4 for two conditions.

Muon beam rates are from 10^7 to 10^8 muons/sec for energies from 0.5 to 5 GeV for a 6 GeV incident electron beam. The rates are from 10^7 to 2×10^8 muons/sec for energies from 0.5 to 9 GeV for a 10 GeV incident electron beam. For comparison the several hundred GeV muon beams at Fermi and CERN are typically run at the several times 10^7 muons/sec rate. At these rates the muons in the beam can be individually detected so that muons which have lost significant energy or are outside the multiple scattering cone can in effect be anti-tagged in the event trigger. By combining the beam counter with a momentum dispersed beam, energy resolution can be significantly enhanced.

Muon Beam

$$\Delta p/p = \pm 5\%, \epsilon_T = 1 \times 10^{-5} \text{ m rad}$$

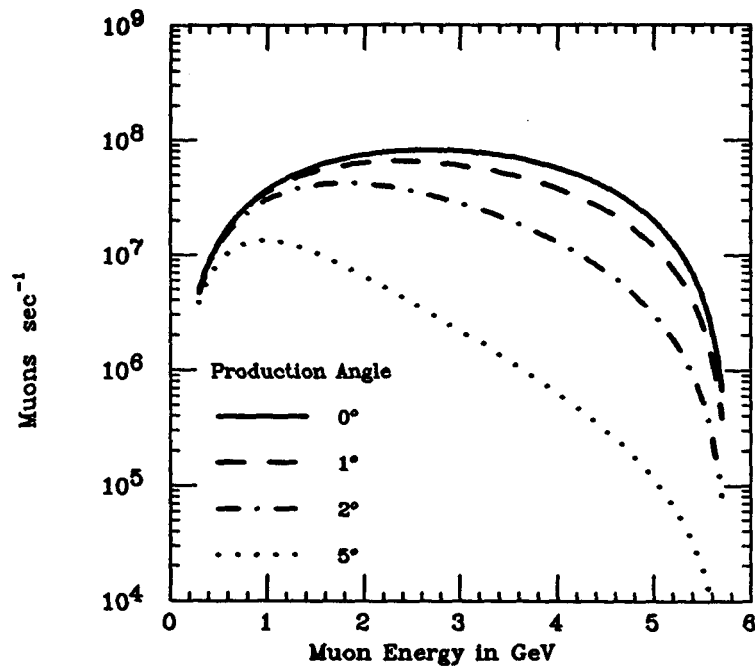


Figure 3: $E_0=6$ GeV, $I=200 \mu\text{A}$, Target=40 rl W, Source $\Omega=1$ mr.

Muon Beam

$$\Delta p/p = \pm 5\%, \epsilon_T = 1 \times 10^{-5} \text{ m rad}$$

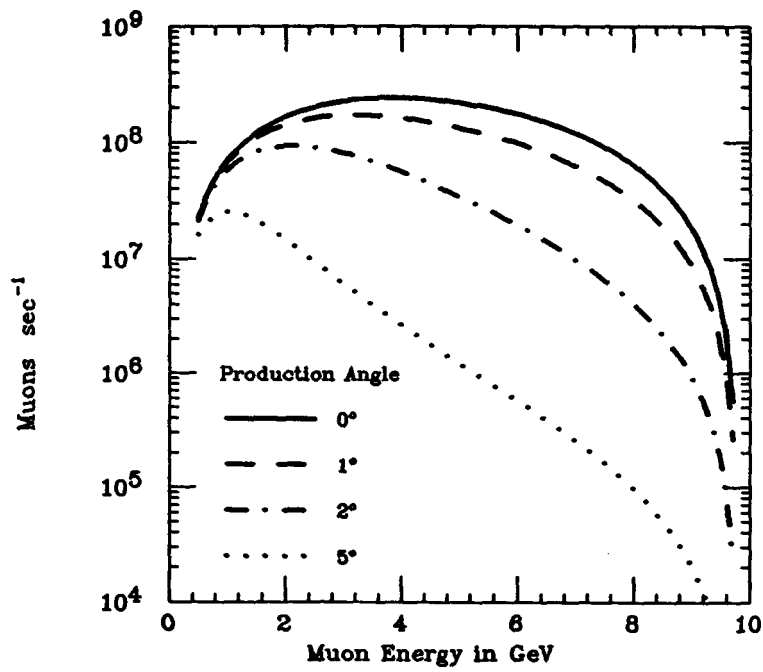


Figure 4: $E_0=10$ GeV, $I=200 \mu\text{A}$, Target=40 rl W, Source $\Omega=1$ mr.

SECONDARY PION AND KAON BEAM INTENSITIES

The secondary pion and kaon beam rates induced by multi-GeV electrons can be calculated according to the equation below, obtained from Roger Gearhart of SLAC. The equation was empirically fit to the SLAC data.

$$\frac{\text{yield}}{\text{incident electron-sr-GeV/c}} = Z \cdot A \cdot X^2 \cdot e^{-\left[BX + \frac{q^2}{P_{Ti}^2} \right]}$$

where: E_o is the incident electron energy in GeV

θ is the secondary production angle in radians

$Z = 1$ for $0.3 \leq B_o$

4.5-5 for $1.0 \leq B_o$

$$A = a_i \left(\frac{E_o}{18} \right)^{\frac{3}{2}}$$

$$X = \frac{P_{\text{secondary}}}{E_o}$$

$$B = b_i \left(\frac{E_o}{18} \right)^{\frac{1}{4}}$$

$$q = E_o \cdot X \cdot \theta$$

The i subscripted coefficients are:

	$\pi^{\pm}$	K^+	K^-
a_i	0.084	0.0048	0.0034
b_i	7.9	8.1	9.3
P_{Ti}	0.40	0.57	0.58

These rates were measured using a beryllium target. By using sandwiches of higher Z and A material with the beryllium, particular components of the secondary beam production can be enhanced. (5)

In calculating the beam intensities at the target, the length of the beam line has been taken into account for particle decay. Shown in Figures 5 and 6 are estimated secondary hadronic beam rates for π^+ , π^- , K^+ , K^- particles for CEBAF. At the price of making a very short beam line, higher beam fluxes at low momenta could be achieved than indicated in the figures. Plotted on the graphs are the secondary beam rates obtainable at Brookhaven⁽⁷⁾ which has a duty factor of 40%. The rates for LAMPF are not included as most

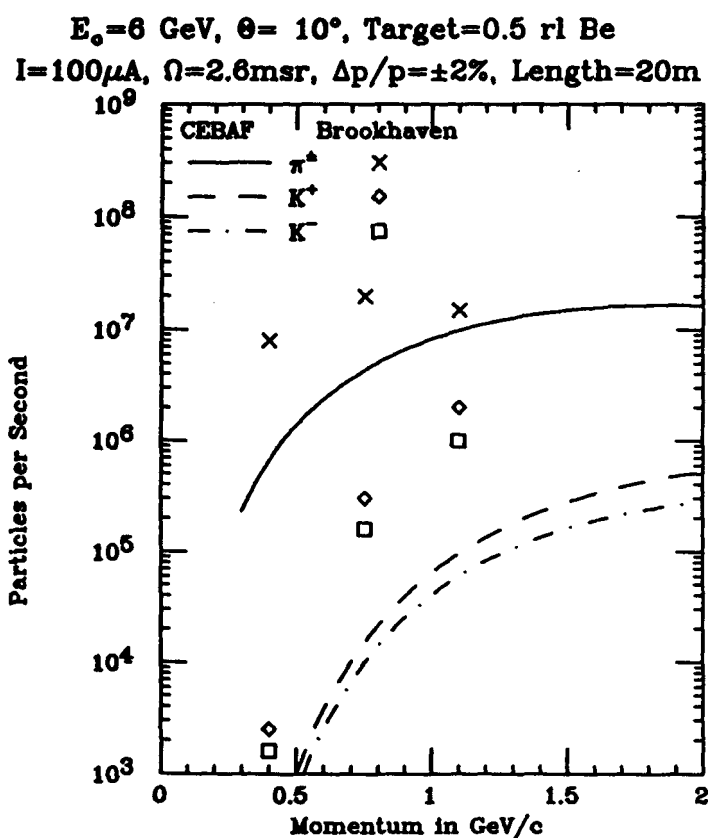


Figure 5. Low momentum charged secondary beams.

$E_0 = 6 \text{ GeV}$, $\theta = 2^\circ$, Target = 0.5 rl Be
 $I = 100 \mu\text{A}$, $\Omega = 0.3 \text{ msr}$, $\Delta p/p = \pm 3\%$, Length = 40m

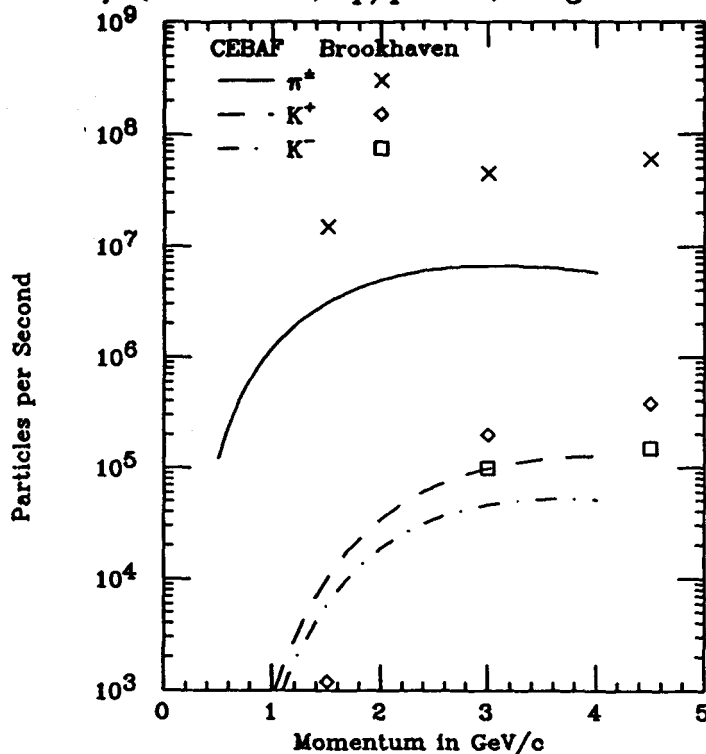


Figure 6. High momentum secondary beams.

secondary beams at LAMPF fall below 0.5 GeV/c. The only other facilities in the world that can produce these beams at reasonable intensities are CERN, Serpukhov, Fermi and SLAC.

The conditions assumed for CEBAF parameters are believed to be reasonable. Note that it is possible to achieve 200 microamps of beam current, which would double the CEBAF projected rates. The solid angles and momentum bytes were taken to be the same as those of separated secondary beams at Brookhaven such that a direct comparison could be made between CEBAF and Brookhaven rates. It should be noted that in several years Brookhaven expects an increase in intensity by approximately a factor of 4 when their booster project is completed.

NEUTRAL SECONDARY HADRONIC BEAMS

Shown in Figure 7 is a sample of the electron-induced production of K_L^0 mesons from a beryllium target. The results are taken from Ref. 8. CEBAF would produce these particles with sufficient fluxes to make them useful for certain classes of experiments. Neutrons are also produced. Clearly proton accelerators can make more neutrons, but the total intensity may not be critical in some cases. As indicated in the section on motivation, detector availability or the CW character of the beam may be more important.

K_S^0 are also be produced. The question of their rates as compared to proton machines is being investigated. It may be that using a primary electron beam rather than a proton beam is more useful so as to keep the K_S^0 /neutron ratio as high as possible.

K_L^0 Yield from 1.75 r.l. Be Target, $E_0=10\text{GeV}$

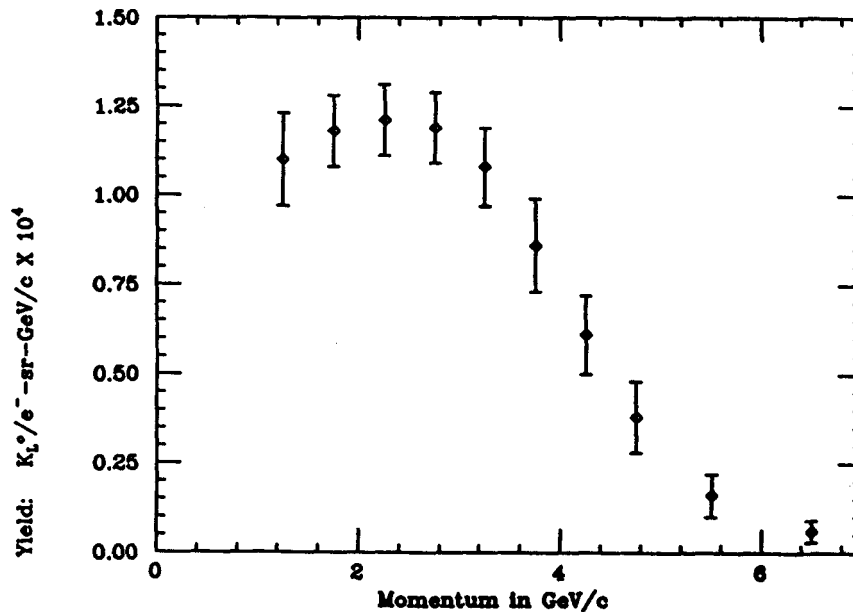


Figure 7. Secondary neutral kaon yield at $\theta=2^\circ$, Ref. 8.

SUMMARY AND COMPARISONS

CEBAF will be capable of producing multi-GeV/c secondary beams of muons, pions, kaons and neutrons. While the rates in some cases are somewhat below those currently available at other labs, the rates for use in combination with large solid-angle detectors are quite high, or even too high if used at maximum. In some instances the ability to do a particular experiment with both an incident photon and incident secondary beams may be particularly worthwhile from the point of view of the physics derived. In other instances the 100% duty factor may be critical.

While higher-rate secondary beams may be available at two or maybe three other places in the world, most of these beams are not 100% duty factor and will not have detectors of the character envisioned at CEBAF. Another important point for U.S. physicists is that the secondary beams at Brookhaven and at Fermi are currently booked two to three years ahead.

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2. R. R. Whitney, "Workshop Report: Tagged Photons--Low Current Electron Beams and Large Acceptance- 4π Detectors," Proceedings of the CEBAF 1984 Summer Workshop, p. 445.
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mt-WHI-AH3

SEARCH FOR DIBARYON RESONANCES IN THE REACTION $\gamma d \rightarrow pp\pi^-$.

W.J. Schwille, Bonn

Contribution to the Working Group on Excited Baryons,
CEBAF Summer Study 1986.

The quest of the existence of bound two-baryon states besides the deuteron started very early in the 1960's, when Oakes (1) suspected strange dibaryon companions of the deuteron from SU_3 and Dyson (2) gave a mass formula, for non-strange two baryon states also, in the frame work of SU_6 . Within the development of QCD the interest in dibaryon states has grown considerably, looking not for such molecule like states but rather for $6q$ -states, or more generally for "exotic" quark (and antiquark) objects with baryon number two. The possible existence of such objects is of some interest for understanding the confinement hypothesis. In nuclear physics it might give important contributions to the description of the nucleon-nucleon interaction and the structure of nuclei. The latter point of view is as well one of the reasons for the current interest in the EMC effect. Compared to the many and contradictory explanations of the EMC effect a unmistakable experimental identification of $6q$ -dibaryon states would be an unequivocal signature. In this short contribution, I will not give a summary of the very embarrassing story of dibaryon searches (for this I refer to summaries by M.P. Locher (3) and K.K. Seth (4), which also are somewhat contradictory in their conclusions), but I will rather report the observation of a narrow enhancement in the pp mass distribution from $\gamma d \rightarrow pp\pi^-$ at Bonn.

Our experiment was done at the Bonn 2.5 GeV synchrotron using a bremsstrahl photon beam with 520 MeV endpoint energy. The kinematics of the reaction were determined by measuring the two protons with time-of-flight hodoscopes (fig.1). The obtained overall time resolution was 250 ps (σ). On calculating the proton momenta, energy loss and different path lengths were taken into account. By comparison with Monte Carlo calculations we find a momentum resolution $\Delta p/p$ better than 1.5% (σ), resulting in an invariant mass resolution $\sigma(M_{pp})$ of $2.4 \text{ MeV}/c^2$. The protons were accepted and identified in the momentum ranges $290 \text{ MeV}/c < p_A < 700 \text{ MeV}/c$ and $335 \text{ MeV}/c < p_B < 600 \text{ MeV}/c$. The total sample consists of 315700 pp-events. In the invariant mass distribution one observes a resonance-like structure at $M_{pp} = 2014 \text{ MeV}/c^2$ ($\pm 2 \text{ MeV}/c^2$ systematic uncertainty) with a width of $\Gamma \leq 5 \text{ MeV}/c^2$, being comparable to the experimental resolution. Different checks on subsamples of the data have been performed showing that this enhancement is not due to a special group of counters or selected kinematics and does not depend on a single data file. The resulting M_{pp} -distribution from all events is shown in fig.2. To further enrich the sample with nonspectator events a cut in the polar angle θ_π^* giving the direction of the pion in the two proton rest frame has been applied. Fig.3 shows the result for $-.6 < \cos \theta_\pi^* < .6$. The spectrum was fitted to a

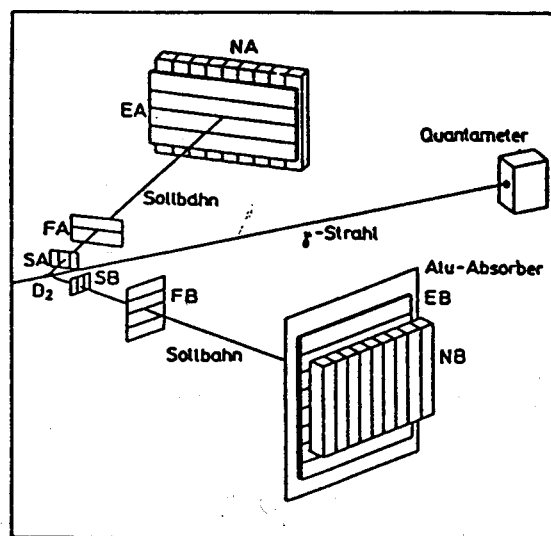


Fig.1 - Experimental set-up

Breit-Wigner resonance on top of a 4th order polynomial background. Applying standard techniques and taking into account the variance of the background contribution we determine a significance of 4.7 standard deviations. A detailed description of this experiment can be found in ref. (5), here I will only present the experimental results. The total data set of this experiment was generally very clean, so we restricted our analysis to events with single hits in the End Hodoscopes EA - EB, losing only 2.7% of the triggers. Due to their short distances from the target, the hit analysis for the start Hodoscopes was more complicated, but in most cases (~60%) SA and SB had single unique hits, too. In addition to the analysis of all events, these "SAB"-events were analysed

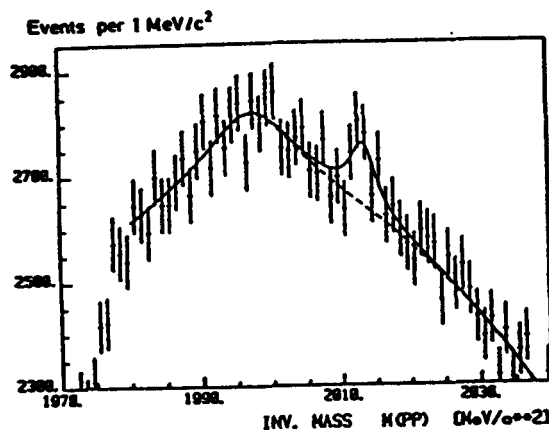


Fig.2 - The pp mass spectrum from $\gamma d \rightarrow pp\pi^-$ (all events)

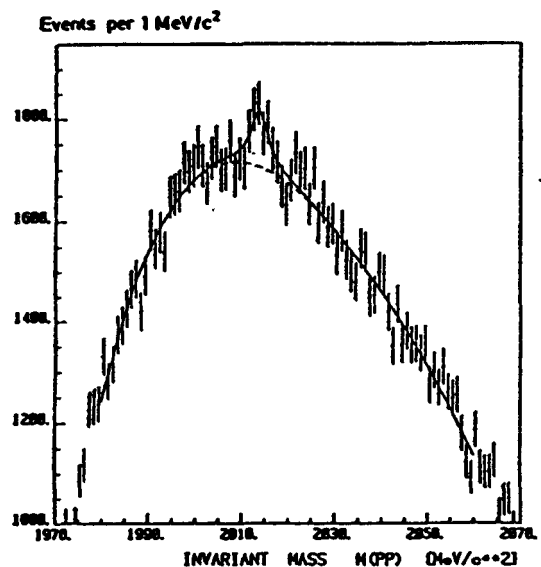


Fig.3 - The pp mass distribution for events which are selected by the cut $-0.6 < \cos\theta^* < 0.6$

separately, because here the accuracy in the determination of the absolute m_{pp} - values is somewhat better. In this analysis, two additional enhancements have been observed, but a with lower statistical significance of 3.5 and 2.7 standard deviations respectively (Fig.4). In order to avoid the discontinuities due to the arbitrary choice of the binning in usual histogramming, we additionally used the method of the k^{th} nearest neighbour (KNN) (6). A result of this analysis for the total data set as well as for "SAB"-events is shown in Fig.5. We will continue our investigations on these candidates with the SAPHIR detector at ELSA, and I hope to encourage the CEBAF physics community to consider the dibaryon search in the physics program. This should be done with tagged photons in a large-solid-angle spectrometer as well as in a missing mass experiment at the high intensity electron beam.

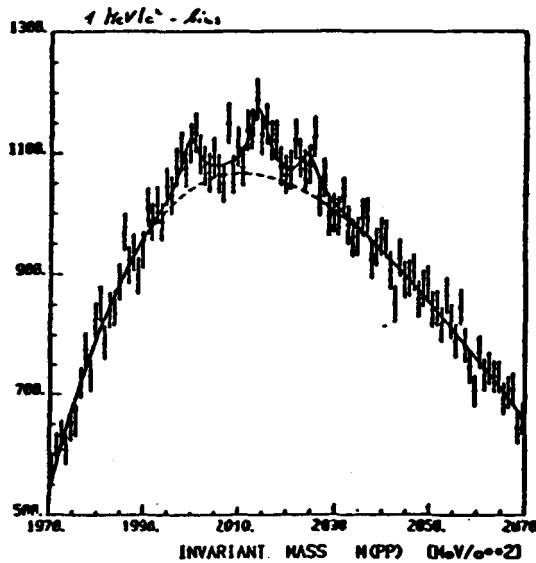


Fig. 4: M_{pp} distribution for "SAB"-events.

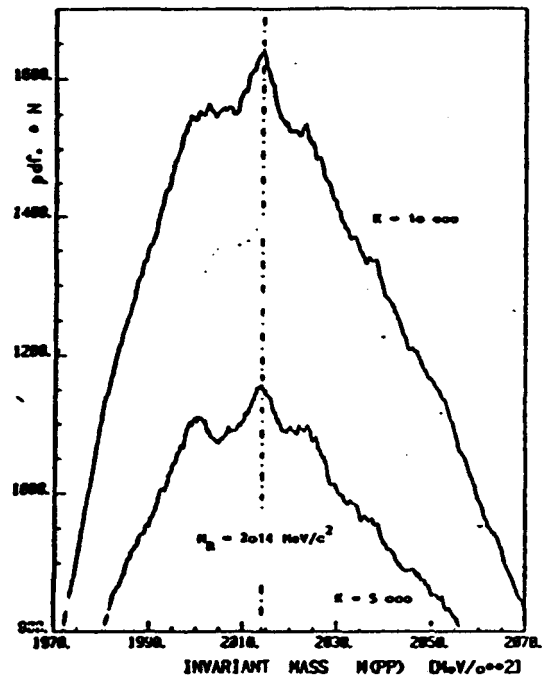


Fig.5: KNN-plot for the total data set ($K = 10000$) and "SAB"-events ($K = 5000$).

References

- (1) R.J. Oakes, Phys.Rev. 131, 2239 (1963)
- (2) F.J. Dyson et al., Phys.Rev.Lett. 13,26, p. 815 (1964)
- (3) M.P. Locher et al., Adv. in Nuc. Phys., Vol.17, 47 (1986)
- (4) K.K. Seth: "The Current Status of Dibaryon Searches", invited talk, Bad Honnef 1984, Springer Lecture Notes in Physics, Vol. 234.
- (5) W. Ruhm et al., to be published in Nuc.Phys. A 180
B. Bock et al., to be published in Nuc.Phys. A 180
- (6) J.H. Freedman, in Data Analysis Techniques for High Energy Particle Physics, CERN-74-23

Appendix A

PROGRAM FOR THE CEBAF 1986 SUMMER WORKSHOP

JUNE 23-27, 1986

Continuous Electron Beam Accelerator Facility
12070 Jefferson Avenue, Newport News, VA 23606

Monday, June 23 *(All events at Christopher Newport College until 4:45 p.m.; at 4:45 p.m. return to CEBAF by bus.)*

8:00	Registration begins
9:00-9:05	Opening remarks - R. Minehart (UVa)
9:05-9:10	Greeting by John E. Anderson, President, Christopher Newport College
9:10-9:30	Welcome and Overview by the Director - H. Grunder (CEBAF)
9:30-9:50	The CEBAF Scientific Program and Organization of the Workshop - F. Gross (CEBAF)
9:50-10:10	The Accelerator and its Operation - C. Leemann (CEBAF)

COFFEE

10:40-11:00	Engineering, Cost and Schedule - A. Chargin (CEBAF)
11:00-11:30	Experimental Equipment for CEBAF - J. Mougey (CEBAF)
11:30-12:30	From the Standard Model to Superstrings - P. Langacker (Penn)

LUNCH (at CNC)

2:00-3:00	Current Technology of Particle Physics Detectors - T. Ludlam (BNL)

Organization of the Working Groups

3:00-3:15	Single Nucleon Emission - A. Gerard (Saclay)
3:15-3:30	Two-Body Correlations - J. Lightbody (NBS)
3:30-3:45	Excited Baryons - N.C. Mukhopadhyay (RPI)

COFFEE

4:15-4:30	Hypernuclear Physics - E. Hungerford (Houston)
4:30-4:45	Studies in the x-Scaling Region - K. van Bibber (LLNL)
4:45-5:00	<i>Buses return to CEBAF</i>
5:00-6:00	Organizational Meetings of the Working Groups (at CEBAF)
6:00-7:00	<i>Social Hour (at CEBAF)</i>

Tuesday, June 24 (All events at CNC until 3 p.m.; at 3 p.m. return to CEBAF.)

8:30-9:15 New Physics from the (e,e'n) Reaction - S. Boffi (Pavia)
9:15-10:00 Dynamic Structure Functions in Nuclear Matter -
S. Fantoni (Pisa)

COFFEE

10:30-11:30 Quark Models of Nucleons in Nuclei - A. Thomas (Adelaide)
11:30-12:30 Nuclear Physics with Hyperons - P. Barnes (CMU)

LUNCH (at CNC)

2:00-3:00 Hadron Production in Charged Lepton Scattering -
F. Janata (DESY)
3:00-3:15 Buses return to CEBAF
3:15-6:00 Working Group Meetings [at CEBAF]

Wednesday, June 25 (Morning events at CNC; return to CEBAF at 12:30.)

8:30-9:00 Review of the Bates Experimental Program -
W. Turchinets (MIT)
9:00-9:30 Nuclear Physics at PEP: First Test of Future Plans -
K. van Bibber (LLNL)
9:30-10:15 Modifications of the Nucleon in the Nuclear Medium -
A. Gerard (Saclay)

COFFEE

10:45-11:45 Recent Progress in Non-Perturbative QCD - H. Pirner (CERN)
11:45-12:30 The ASTERIX Detector at LEAR and its Future Development -
U. Gastaldi (CERN)
12:30-12:45 Buses return to CEBAF

LUNCH (at CEBAF)

2:00-3:00 User Group Meeting
3:00-5:00 Working Group Meetings

BANQUET (at Cascades in Williamsburg; buses depart CEBAF at 5:45)

Thursday, June 26 (Morning events at CNC; return to CEBAF at 12:30.)

- 8:30-9:00 Review of Nuclear Physics Program at SLAC -
 R. Arnold (SLAC/AU)
- 9:00-9:30 Present Status and Experimental Program for the Electron
 Stretcher Ring (ELSA) - W. Schwille (Bonn)
- 9:30-10:15 Multiparticle Final States from Pion-Induced Reactions -
 P. Roos (Maryland)

COFFEE

- 10:45-11:45 Electron Scattering from Nucleons and Nuclei in the Skyrme
 Model - D. Riska (Helsinki)
- 11:45-12:30 The Proton Form Factor in the Time-Like Region -
 J. Duclos (Saclay)
- 12:30-12:45 Buses return to CEBAF

LUNCH (at CEBAF)

- 2:00-3:00 Comments on the Planned CEBAF Experimental Program -
 J. Duclos (Saclay)
- 3:00-6:00 Working Group Meetings

Friday, June 27 (All events at CEBAF)

Working Group Reports

- 8:30-9:00 Single Nucleon Emission - A. Gerard (Saclay)
- 9:00-9:30 Two-Body Correlations - J. Lightbody (NBS)
- 9:30-10:00 Excited Baryons - N. C. Mukhopadhyay (RPI)

COFFEE

- 10:30-11:00 Hypernuclear Physics - E. Hungerford (Houston)
- 11:00-11:30 Studies in the x-Scaling Region - K. van Bibber (LLNL)
- 11:30-12:30 Workshop Summary - I. Sick (Basel)

LUNCH

SEE REVERSE SIDE FOR WORKING GROUP SCHEDULE

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REVISED 6/18/86

CEBAF 1986 SUMMER WORKSHOP

WORKING GROUP SCHEDULE AND MEETING ROOMS

The 5 working groups are:

- A - Single Nucleon Emission (A. Gerard)
- B - Two-Body Correlations (J. Lightbody)
- C - Excited Baryons (N.C. Mukhopadhyay)
- D - Hypernuclear Physics (E. Hungerford)
- E - Studies in the x-Scaling Region (K. van Bibber)

Monday	5:00-6:00	Working Group A	Room 53
		Working Group B	Room 55
		Working Group C	Room 47
		Working Group D	Small Conference Room
		Working Group E	Room 40
	All Day	HUGS	Room 41
Tuesday	3:00-4:00	Working Group A	Room 53
		Working Group C	Room 55
		Working Group D	Small Conference Room
		HUGS	Room 41
	4:00-5:00	Working Group B	Room 53
		Working Group C	Room Room 55
		Working Group E	Small Conference Room
		HUGS	Room 41
	5:00-6:00	Working Group B	Room 53
		Working Group D	Room 55
		Working Group E	Small Conference Room
		HUGS	Room 41
Wednesday	3:00-4:00	Working Group A	Room 53
		Working Group C	Room 55
		Working Group E	Small Conference Room
		HUGS	Room 41
	4:00-5:00	Working Group B	Room 53
		Working Group C	Room 55
		Working Group D	Small Conference Room
		HUGS	Room 41
Thursday	3:00-4:00	Working Group A	Room 53
		Working Group D	Small Conference Room
		Working Group E	Room 55
		HUGS	Room 41
	4:00-5:00	Working Group B	Room 53
		Working Group D	Small Conference Room
		Working Group C	Room 55
		HUGS	Room 41
	5:00-6:00	Working Group B	Room 53
		Working Group A	Room 55
		Working Group E	Small Conference Room
		HUGS	Room 41

Appendix B

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