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1. Introduction

The SLAC Linear Collider (SLC) was constructed in the years 1983-1987 for two principal reasons:

1. To develop the accelerator physics and technology that are necessary for the construction of future linear electron-positron colliders.
2. To produce electron-positron collisions at the Z^0 pole and to study the physics of the weak neutral current.

To date, the SLC program has been quite successful at achieving the first goal. The machine has produced and collided high energy electron and positron beams of three-micron transverse size. The problems of operating an open geometry detector in an environment that is more akin to those found in fixed-target experiments than in storage rings have largely been solved. As a physics producing venture, the SLC has been less successful than was originally hoped but more successful than is commonly believed. Some of the results that have been produced by the Mark II experiment with a very modest data sample are competitive with those that have been produced with much larger samples by the four LEP collaborations. At the current time, SLAC is engaged in an ambitious program to upgrade the SLC luminosity and to exploit one of its unique features, a spin polarized electron beam.

These lectures are therefore organized into three sections:

1. A brief description of the SLC;
2. A review of the physics results that have been achieved with the Mark II detector;
3. A description of the SLC's future: the realization and use of a polarized electron beam.

Physics at the SLC*

MORRIS L. SWARTZ

Stanford Linear Accelerator Center
Stanford University, Stanford, California, 94309

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2. The SLAC Linear Collider

2.1. THE IMPORTANCE OF LINEAR COLLIDERS

Circular electron-positron (e^+e^-) storage rings were developed in the 1960's and early 1970's as an inexpensive technique for the production of high energy collisions (the center-of-mass energies were typically a few GeV). They proved to be such extremely important tools in the development of particle physics that two more generations of higher energy machines have been constructed.

Unfortunately, the size and cost of e^+e^- storage rings increase rapidly with the beam energy. This can be illustrated by considering a simple cost model for electron storage rings. The cost C can be expressed as follows,

$$C = \alpha R + \beta \frac{E^4}{R}. \quad (2.1)$$

where: R is the radius of the storage ring; E is the beam energy; and α , β are constants. The first term represents the costs that are proportional to the size of the ring (such as the tunnel, magnets, vacuum chamber, etc). The second term represents the cost (size) of the RF acceleration system that is necessary to replace the energy that is lost to synchrotron radiation. The optimal size of the ring for a given beam energy can be found from the zero of the first derivative dC/dR . The optimal radius and cost, R_{opt} and C_{opt} , scale with the square of the beam energy,

$$R_{opt} = \sqrt{\frac{\beta}{\alpha}} \frac{E^2}{\alpha}, \quad C_{opt} = 2\sqrt{\beta\alpha} E^2.$$

The unpleasant reality of this scaling law can be illustrated by considering the scaling of LEP200 to 1 TeV. The size and cost of a LEP1000 are listed in Table I. The cost is comparable to the entire CERN budget since its inception in 1954. The size is roughly eight times larger than that of the SSC. (If the beams were injected into LEP1000 at CEKN, the 180° interaction region could be located in Zürich!)

Table I

Machine	CM Energy	Length of Tunnel	Cost
LEP200	180 GeV	27 km	FS 1×10^9
LEP1000	1 TeV	833 km	FS 31×10^9
TLC	1 TeV	14 km	FS 3×10^9

Table I also lists a guess at the size and cost of a 1 TeV linear collider (TLC). The size and cost of the linear machine are smaller than the circular one by factors of 60 and 10, respectively. The size is well within the range of existing machines. The cost, while large, is smaller than those of the next generation of hadron colliders. *It appears that linear colliders are the only practical technique for the building of very high energy e^+e^- machines.*

2.2. LINEAR COLLIDERS

Linear colliders differ from circular machines in that the beams are accelerated to collision energy (in one or two linear accelerators), collided, and discarded after only one use.

The luminosity $\mathcal{L}$ of any collider is given by the following expression,

$$\mathcal{L} = \frac{N_+ N_- f}{4\pi\sigma_x\sigma_y}. \quad (2.2)$$

where: N_+ and N_- are the number of positrons and electrons in each bunch; f is the frequency of collision of the bunches; and $4\pi\sigma_x\sigma_y$ is the overlap area of the Gaussian bunches of size σ_x by σ_y .

Circular machines have rather high frequencies of bunch collisions (from 4×10^4 at LEP to $\sim 10^6$ at small storage rings). Unfortunately, non-superconducting linear accelerators are limited to collision frequencies in the range 10^2 to 10^3 . In order to produce comparable luminosity, a linear collider must compensate the low collision frequency with increased bunch population or with reduced beam size. The former

leads to a number of technical difficulties and would require a tremendous amount of RF power. The latter approach is the choice of all linear collider designers. Since the beams are to be discarded after a single use, they can be subjected to extreme perturbations from the transport system or from the other beam. It is therefore quite natural to use very small beams in a linear collider. The SLC has produced beams of $3 \mu\text{m}$ transverse size which is substantially smaller than the LEP design value of $12 \mu\text{m}$ by $300 \mu\text{m}$ (vertical by horizontal size). Future linear colliders are expected to utilize beams that are smaller than those of the SLC by nearly two orders of magnitude.

2.3. TERMINOLOGY

It is clear that the designer of a linear collider must concern himself (herself) with those aspects of the machine design that affect the beam size. The size of a charged particle beam within a magnetic transport system is determined by the focusing strength of the transport system and by the phase space that is occupied by the beam particles.

The concept of the phase space of a beam is an important one. Let us assume that we have a beam that consists of a large number of particles moving principally in the z -direction. We can then use Liouville's theorem to write the following expression,

$$\frac{1}{N} \sum_{\text{beam}} P_y y = \gamma \beta m \frac{1}{N} \sum_{\text{beam}} y' y = \pi \gamma \beta m \varepsilon_y = \text{constant}, \quad (2.3)$$

where: N is the number of particles in the beam; P_y is the momentum of a beam particle along one of the directions that is transverse to the beam direction; y is the displacement of the particle with respect to the beam axis; γ and β are the usual Lorentz quantities; m is the particle mass; and y' is the derivative dy/dz . The sums in equation (2.3) are taken over all particles in the beam. It is customary

to define the *emittance* of the beam (in the y direction) ε_y as follows,

$$\varepsilon_y \equiv \frac{1}{\pi N} \sum_{\text{beam}} y' y. \quad (2.4)$$

Liouville's theorem and the definition of the emittance lead to the following properties of the horizontal and vertical emittances:

1. The horizontal and vertical emittances, ε_x and ε_y , are invariant in conservative fields. Note that magnetic transport systems are composed of conservative fields.
2. The products $\gamma \varepsilon_x$ and $\gamma \varepsilon_y$ are invariant under acceleration (we assume that $\beta \simeq 1$).
3. The emittances ε_x and ε_y are intimately related to the transverse beam sizes. It can be shown that the variances of the (Gaussian) horizontal and vertical beam particle distributions, σ_x and σ_y , are related to the horizontal and vertical emittances by the following expressions,

$$\begin{aligned} \sigma_x^2 &= \varepsilon_x \beta_x(z) \\ \sigma_y^2 &= \varepsilon_y \beta_y(z), \end{aligned} \quad (2.5)$$

where $\beta_x(z)$ and $\beta_y(z)$ are functions which describe the focussing strength of the transport system (a complete description of the β functions formalism can be found in Reference 1).

4. The variances of the angular distributions of the beam particles, $\sigma_{x'}$ and $\sigma_{y'}$, are also related to the emittances and the β functions,

$$\begin{aligned} \sigma_{x'}^2 &= \varepsilon_x / \beta_x(z) \\ \sigma_{y'}^2 &= \varepsilon_y / \beta_y(z). \end{aligned} \quad (2.6)$$

2.4. LINEAR ACCELERATORS

The main element(s) of a linear collider is a linear accelerator. The first linear accelerators were Cockroft-Walton and Van der Graaf accelerators. They consist of a linear drift space across which a large voltage difference V is maintained. This generates a strong axial electric field which is used to accelerate charged particles to kinetic energies that are equal to the voltage V . Unfortunately, it is not possible to maintain an arbitrarily large voltage across the accelerating structure. As the voltage is increased, one inevitably exceeds the dielectric strength of the insulators being used and discharges to ground occur. The maximum voltage that is typically obtained is 10-20 MV.

This limitation was overcome by the development of the traveling wave accelerator in the late 1940's.^[2] The traveling wave accelerator is based upon the observation that the TM_{01} mode of an electromagnetic field in cylindrical waveguide has a longitudinal electric field. The electric field is oriented along the axis of the cylinder which is the direction of propagation of the electromagnetic field (this is quite different from the case of a freely propagating em field). Charged particles can therefore be accelerated by a moving pulse of RF power and there is no need to produce a huge voltage along the entire length of the accelerator. Unfortunately, the phase and group velocities of the TM_{01} mode conspire to complicate the design of the accelerator. The phase velocity of the cylindrical waveguide is larger than the speed of light. A bunch of charged particles would see a longitudinal electric field of constantly changing sign and no acceleration would take place. The solution to this problem is to *load* the cylinder with annular disks. The phase velocity can be adjusted to be equal to the speed of light.

The group velocity of a radio frequency pulse of electromagnetic energy in a cylindrical accelerating structure is much less than the speed of light (in the SLC it is 1%-2% of c). A bunch of electrons cannot travel along with an electromagnetic pulse of energy. The solution to this problem is to feed-in RF power at short intervals along the length of the accelerator. In the SLC, this is done at intervals

of 3 m.

2.5. A DESCRIPTION OF THE SLC

A layout of the SLAC Linear Collider is shown in Figure 1. A cycle of the machine begins when one bunch of positrons and two bunches of electrons are extracted from the damping rings and are accelerated down the linac structure. As the after the positron bunch and the first electron bunch pass the 2/3 point of the linac (the 3 km linac is composed of 30 sectors), a pulsed kicker magnet diverts the second electron bunch onto the positron production target to make more positrons. The positrons are returned to the beginning of sector 1 by a long return line. The electron gun at the front end of the machine fires to produce two electron pulses which are coaccelerated with the positrons to 1.15 GeV in the first linac sector. The beams are then injected into their respective damping rings.

The positron bunch and the first electron bunch continue to be accelerated to ~ 46.5 GeV at the end of the linac. The bunches pass through a large dipole magnet which sends them into the north (electron) and south (positron) machine arcs. The beams each lose approximately 1.0 GeV as they traverse the 1.5 km arcs. They then enter the final focus regions which cause them to be demagnified to sizes of a few microns at the interaction point. After the collision, the beams are ejected into beam dumps by kicker magnets.

The following sections describe each of the machine subsystems in more detail.

The Linac

The actual linac has been modified substantially for use in the SLC. The klystron power supplies which provide the s-band RF power (2860 MHz) have been upgraded from 20 MW devices to 67 MW devices. The energy upgrade along with the implementation of a pulse compression technique have increased the maximum energy of the linac from 20 GeV to more than 50 GeV.

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The increased energy has required that the focussing strength of the of the quadrupole lattice be increased. Improved quadrupole magnets are placed at 12 m intervals along the machine.

As we have already mentioned, the production of very small beams is a critical design feature of any linear collider. This requires that the emittance of the beam be kept as small as is possible. Unfortunately, there are several effects that can increase the emittance of a beam as it is accelerated in a linac. Collectively, they are known as wakefield effects. They fall into two categories:

1. Transverse wakefield effects are caused by the interaction of a bunch with its own image fields or with the image fields of other bunches. Within a single bunch, the image fields of the head of the bunch can affect the transverse positions of the particles in the tail of the bunch. The ensuing rotation of the bunch causes an effective increase in the transverse emittance of the beam. The solutions to this problem are to make the bunches as short as possible and to steer the beam as close to the axis of the accelerator structure as is possible (the effect vanishes on the accelerator axis).
2. Longitudinal wakefields effects are caused by the intrabunch electrostatic fields. Fields from the bunch head tend to decelerate the particles in the bunch tail. This causes an increase in the energy spread of the beam as it is accelerated in the linac. This problem is minimized by making the beam bunches as long as is possible.

In order to control wakefield effects, a number of changes to the linac and its mode of operation have been implemented:

1. The linac is operated with a short (optimized) bunch length of 2 mm.
2. Beam position monitors and corrector magnets have been installed at intervals of 12 m along the linac. The system can control the trajectory of a single beam bunch to ± 100 μm of the accelerator axis.
3. The linac RF phases are adjusted to introduce and to remove an energy

spread as the beam is transported down the linac. This causes the beam to decompose into filaments of different momenta which follow different orbits. The transverse wakefield effects are reduced by this technique (which is called BNS damping).

The Electron Source

The electron source consists of several components. An electron gun produces two 2-ns pulses of up to 2×10^{11} electrons from a hot cathode. The 175 kV electron pulses are separated in time by 61 ns.

A system of three RF bunchers is then used to reduce the bunch length from ~ 20 cm at 175 kV to 2 mm at 40 MeV. The bunchers make use of the non-relativistic velocity of the electrons that are emitted from the gun. A long wavelength axial electric field is used to accelerate the tail of the bunch and to decelerate the head of the bunch. This velocity dispersion decreases the bunch length until the increasing energy causes the velocity to saturate at c along the entire bunch.

Finally, an accelerator section is used to increase the beam energy to 200 MeV at the entrance to the first linac sector. The electron bunches are then coaccelerated with positrons returning from the positron production target to 1.15 GeV for injection into the damping rings.

Damping Rings

The electron and positron bunches that are produced by the respective sources have invariant emittances that are too large for the high luminosity operation of a linear collider. It is therefore necessary to make use of a phenomenon that doesn't conserve energy to reduce the beam emittances.

The emission of synchrotron radiation is particularly useful for this purpose. Let us consider a bunch of electrons circulating in a storage ring. Let the z -axis define the instantaneous direction of a particle that is traveling along the central orbit of the bunch. Most of the remaining particles have some momentum components that are perpendicular to the central orbit (let x and y be the horizontal

and vertical directions in the perpendicular plane). As the average beam particle passes through the machine arcs, it radiates photons that are collinear with its instantaneous direction of motion. The momenta along the three directions, p_x , p_y , and p_z , are reduced by the emission process. This lost energy is replaced in an RF accelerating cavity. *Note, however, that the energy is replaced along the z direction only.*

In the vertical direction (which is orthogonal to the bending plane of the arcs), the transverse momentum of the beam is reduced without affecting the spatial distribution of the particles. *The vertical emittance, $\gamma\epsilon_y$, therefore becomes smaller.*

In the horizontal direction, the transverse momentum components are also reduced. Unfortunately, the particle trajectory moves horizontally when the particle energy changes (the radius of the orbit becomes smaller when energy is lost). Thus, we have two competing effects: one that reduces the transverse momentum components, and one that increases the spatial distribution of the beam. The machine lattice can be designed to enhance the damping effect and to reduce the horizontal emittance, $\gamma\epsilon_x$, at a cost in longitudinal energy spread.

The SLC has two small storage rings that are designed to reduce the vertical and horizontal emittances of the electron and positron bunches. The north (electron) damping ring is shown in Figure 2. The ring is designed to reduce the emittance of a particle bunch with a characteristic $1/e$ time of 3 ms. Since the positron bunches are produced with larger emittances, they must be stored for two machine cycles (16 ms) to be sufficiently damped.

The Positron Source

As was described at the beginning of this section, a positron bunch and two electron bunches are extracted from the damping rings and accelerated in the linac. As they pass the 19th sector of the machine, the trailing electron bunch is diverted onto the positron production target by a pulsed kicker magnet and a short beam transport system. The electrons produce electromagnetic showers in the 6 radiation length target. Positrons between the energies of 2 MeV and 20 MeV are captured

by a solenoidal focusing system and are accelerated to 200 MeV. The system is designed to capture 2 positrons for each incident electron. The positrons are then returned to the front end of sector 1 for acceleration and storage in the south damping ring. Including losses enroute, the system is designed to store 1 positron in the damping ring for each electron striking the positron production target.

The Arcs

After the positron and leading electron bunches are accelerated to full energy (which is approximately 1 GeV larger than the interaction point energy), they are transported to the final focus systems in 1.5 km, S-shaped arcs. In our discussion of radiation damping, we noted that particles traversing a magnetic field lose energy along three coordinates. Since we have no RF system to replace the lost energy, the emittance of a beam bunch that is transported through a large arc must increase. The horizontal emittance is also increased by the horizontal displacement caused by the energy change (this is called dispersion).

The increase in the horizontal emittance can be minimized by keeping the beam strongly focused as it is transported through the arcs. In order to do this, the SLC arcs consist of three-pole, combined-function magnets. The fields produced by these magnets have strong dipole, quadrupole, and sextapole components. Each arc is constructed of 460 such magnets arranged to alternately focus and defocus in each plane. The arcs are designed to be achromatic to second order and to be capable of transporting a beam with an energy spread of 0.5%.

Final Focus

The last 150 m of each (electron and positron) beamline is called the final focus. Each final focus system is a transport system that consists of 8 bending magnets, 26 quadrupole lenses, 8 sextapole magnets, and a number of correction and monitoring devices. These systems are designed to demagnify the $250\ \mu\text{m} \times 30\ \mu\text{m}$ beams that leave the arcs to spots of $2\ \mu\text{m} \times 2\ \mu\text{m}$ at the interaction point. The final focus systems are designed to cancel all geometrical and chromatic aberrations to second order.

Beam Monitoring

The beams that are stored in storage rings have stable orbits and energies. After stored electron and positron beams are brought into collision, they will generally remain in collision for some macroscopic time. Unfortunately, this convenient behavior is not necessarily true in a linear collider.

The problem of energy stability in the linac is confronted with feedback systems and by making the arcs and final foci fairly achromatic (they can transport momenta over a range $\Delta p/p = 0.5\%$). Residual energy drifts are measured on each pulse by spectrometers that are placed in the extraction lines. These will be described in chapter 4.

Linear colliders must rely heavily on sensors and feedback systems to control the orbits of the beams. The SLC makes use of several techniques to bring the beams into collision:

1. A system of beam position monitors is used to measure the positions and directions of the beams at the interaction point. These devices measure the beam centroid position by comparing the beam induced signals in pickup loops that are placed on either side of the vacuum chamber. They are capable of steering the beams to within $20 \mu\text{m}$ of each other.
2. The phenomenon of beam-beam deflection provides the single most important technique for establishing collisions.^[3] The fields of each beam deflect the other in a manner that depends upon their transverse sizes, the distance of closest approach, and number of particles in the bunches. The deflection angle θ of an infinitely narrow beam by a target beam of finite size is given by the following expression.

$$\theta = \frac{-2r_e N_e}{\gamma} \cdot \frac{1 - e^{-\Delta^2/2\sigma^2}}{\Delta}, \quad (2.7)$$

where: r_e is the classical radius of the electron; γ is gamma factor of the deflected beam; N_e is the number of particles in the target bunch; Δ is

the (signed) miss distance of the beams; and σ is the size of the target beam. Note that the deflection angle has maxima at the miss distances, $\Delta \simeq \pm 1.6\sigma$. As the beams are moved closer together, the deflection angle becomes smaller. It passes through zero when the beams collide and changes sign as the original beam positions are interchanged. Using the system of beam position monitors to measure the deflection angles (which are of the order of $100 \mu\text{rad}$), it is straightforward to target the beams to within a small fraction of a beam size.

3. The strong magnetic fields that are associated with the beam-beam deflection process (up to 100 T) also produce a large quantity of synchrotron radiation (10^6 to 10^{10} photons of energy larger than 20 MeV). This radiation, called *beamsstrahlung*, is separated from the electron beam by a large bending magnet in the final focus. Since there is a large background of lower energy photons (typically 2 MeV) from the focusing elements of the beam transport system, the *beamsstrahlung* photons are converted into e^+e^- pairs with a radiator and detected in a cerenkov counter.
4. The transverse profiles of the electron and positron beams are measured by the devices called *wire scanners*. These devices work by passing the beams through a very fine carbon filament (the smallest has a $4 \mu\text{m}$ diameter) and by detecting scattered radiation. The wire scanner used in the interaction point produces *bremsstrahlung* photons that are detected by the *beamsstrahlung* monitor.

2.6. SLC PERFORMANCE

The performance of the SLC has not yet approached the level that was intended when the machine was designed. The technical capabilities of the machine were recently assessed by a committee of experts.^[4] They conclude that it is possible to improve the performance of the SLC to produce 10^4 to 10^5 Z events per year. The correctness of this assessment will be established in the next several

years. In chapter 6, we shall see that if this level of performance is achieved, the implementation of the polarized electron beam should provide an interesting and unique test of the Standard Model.

3. Physics with the Mark II Detector

The Mark II detector was originally constructed to study e^+e^- collisions at the PEP storage ring. It was upgraded for use at the SLC by the replacement of the tracking system. The detector has operated at the SLC in 1989 and 1990. Beginning in 1991 it will be replaced by the new SLD detector.

3.1. THE MARK II DETECTOR

A schematic diagram of the Mark II detector is shown in Figure 3. The detector consists of a system to reconstruct the tracks of charged particles, a calorimeter to measure the energies of charged and neutral particles, and a system to identify penetrating charged particles (which are presumed to be muons).

The cylindrically symmetric tracking system consists of three distinct devices. A three-layer silicon strip detector (SSD) occupies the region from a radius of 2.5 cm from the beam axis to a radius of 5.0 cm from the beam axis. Each of the three measurements is made in the azimuthal direction with a precision of approximately 10 μm . The SSD is followed by a twelve-layer high pressure drift chamber microvertex detector (DCVD). The DCVD occupies the region between the radii 5.0 cm and 17 cm. It is capable of measuring tracks in the azimuthal direction with a precision of approximately 40 μm per measurement. The DCVD is followed by a large cylindrical drift chamber that spans the region between the radii 19 cm and 147 cm. The sense wires of the chamber are organized into 6-wire vector cells. There are 12 layers of vector cells which provide 72 measurements of each track. Alternate layers are oriented parallel to the beam axis or are tilted by $\pm 3^\circ$ with respect to the axis (the axial coordinate is provided by this small-angle stereo arrangement). The typical precision of each measurement is 145 μm .

The tracking system is capable of reconstructing tracks in the region of polar angle $|\cos \theta| < 0.8$ with an efficiency of 99%. The momentum resolution of the combined system is given by the following expression, $\sigma_p/p \sim 0.002p$, where p is the momentum in GeV. The large drift chamber has the capability to measure the ionization energy loss in the gas with a precision of roughly 7%.

The tracking system is surrounded by barrel and endcap calorimeters. The barrel calorimeter is a lead-liquid argon device that covers the region of polar angle $|\cos \theta| < 0.7$. Since it is only 14 radiation lengths thick, the energy resolution is given by the following expression,

$$\frac{\sigma E}{E} = \begin{cases} 0.12 \cdot \text{GeV}^{1/2} / \sqrt{E} & E < 15 \text{ GeV} \\ 0.15 \cdot \text{GeV}^{1/2} / \sqrt{E} & 15 \text{ GeV} < E < 50 \text{ GeV} \end{cases}$$

The resolution of the barrel calorimeter in azimuthal and polar angle is approximately 3.5 mRad. The endcap calorimeter is a lead and proportional wire chamber device that covers region of polar angle $0.7 < |\cos \theta| < 0.98$. It has a thickness of 20 radiation lengths and is sampled each 0.5 radiation length of thickness. The energy resolution is given by the following expression,

$$\frac{\sigma E}{E} = 0.20 \cdot \text{GeV}^{1/2} / \sqrt{E}.$$

The muon system is a sandwich structure composed of iron of the flux return of the solenoid and an triangular proportional tubes. The system covers 60% of the solid angle and has a total thickness of 1.05 m. The probability that a hadron be misidentified as a muon is less than 1%.

The luminosity is determined from the number of small angle Bhabha scattering events that are detected in two different detectors. The minisam consists of a pair of cylindrical tungsten-scintillator calorimeters that detect both legs of Bhabha events in the region of polar angle between 15 and 25 milliradians. Each calorimeter is segmented into four quadrants. The small angle monitor (SAM) is comprised of a pair of lead-proportional wire chamber calorimeters that detect both legs of events in the region of polar angle between 50 and 160 milliradians.

4. The Lineshape of the Z^0 Boson

The measurement of the Z^0 lineshape provides a great deal of information about the Standard Model. In order to fully appreciate the relevance of the mass measurement, we must discuss the Standard Model briefly.

4.1. PARAMETERS OF THE STANDARD MODEL

The minimal Standard Model contains some 21 empirical parameters. They are listed in Table II with their approximate values. The dynamics of electroweak physics are determined (at tree level) by three of the parameters: the $SU(2)$ coupling constant (g), the $U(1)$ coupling constant (g'), and the vacuum expectation value of the Higgs field ($\langle\phi\rangle$). The complete specification of the electroweak sector of the Standard Model requires that all three parameters be *precisely* known. The values of these quantities are extracted from the measurement of three related quantities: the electromagnetic fine structure constant (α), the Fermi coupling constant (G_F), and the mass of the Z^0 boson (M_Z). The current values of these quantities are listed in Table III.

The value of α is extracted from a very precise measurement of the anomalous magnetic moment of the electron.^[5] The value of G_F is derived from the measured value of the muon lifetime.^[6] The first precise measurements of the Z^0 mass have been made quite recently at the Tevatron collider, the SLC, and at LEP. The quoted value is determined mostly by the remarkable measurements of the LEP collaborations.^[6] Although M_Z is determined with far less accuracy than are α and G_F , it is expected to remain the most well-determined Standard parameter for the foreseeable future. It is clear that the measurement of a fourth physical quantity should overconstrain the determination of the electroweak parameters. We should therefore be able to *test* the electroweak sector of the Standard Model.

Unfortunately, the expression given in Table III that relates M_Z to g , g' , and $\langle\phi\rangle$ is valid only at tree-level. Since M_Z is measured at a substantially larger energy scale than are α and G_F , we must include virtual electroweak corrections

Table II

Parameter	Description	Approximate Value
g_s	$SU(3)$ coupling constant	1.3 @ 34 GeV
g	$SU(2)$ coupling constant	0.63
g'	$U(1)$ coupling constant	0.35
$\langle\phi\rangle$	VEV of the Higgs field	174 GeV
M_H	Higgs boson mass	?
m_{ν_e}	electron neutrino mass	< 12 eV
m_{ν_μ}	muon neutrino mass	< 0.25 MeV
m_{ν_τ}	tau neutrino mass	< 35 MeV
m_e	electron mass	0.511 MeV
m_μ	muon mass	106 MeV
m_τ	tau mass	1.78 GeV
m_u	up-quark mass	5.6 MeV
m_d	down-quark mass	9.9 MeV
m_s	strange-quark mass	199 MeV
m_c	charm-quark mass	1.35 GeV
m_b	bottom-quark mass	5 GeV
m_t	top-quark mass	?
$\sin\theta_{12}$	K-M Matrix parameter	0.217-0.223
$\sin\theta_{23}$	K-M Matrix parameter	0.030-0.062
$\sin\theta_{13}$	K-M Matrix parameter	0.003-0.010
$\sin\delta$	K-M Matrix parameter	?

in order to extract accurate values for the electroweak parameters. In principle, this requires a knowledge of all of the parameters listed in Table II. In practice, a dispersion relation is used to determine the dominant correction (due to low mass fermion loops) from the low energy e^+e^- total cross section. The largest remaining corrections depend upon the top quark mass (strongly) and the Higgs boson mass (weakly). A reasonably precise test of the Standard Model therefore requires at least two more experimental measurements (ideally a measurement of m_t would be

Table III

The current values of the physical parameters that determine the determine the electroweak sector of the Standard Model.

Quantity	EW Parameters	Current Value	Precision (PPM)
α	$\frac{1}{4\pi} \frac{e^2}{g^2 + g'^2}$	$[137.0359895(6)]^{-1}$	0.015
G_F	$\frac{1}{(\sqrt{2})^2 \sqrt{s}}$	$1.16637(2) \times 10^{-5} \text{ GeV}^{-2}$	17
M_Z	$\sqrt{\frac{g^2 + g'^2}{2}} (\phi)$	91.16(3) GeV	320

one of them).

4.2. MASS AND WIDTH OF THE Z^0

We have already discussed the importance of a high precision measurement of the mass of the Z^0 . The width of the Z^0 has a tree-level dependence upon the parameters of the Standard Model and the particle content of the theory. The total width is the sum of the partial widths for the decay into each fermion-antifermion final state.

$$\Gamma_Z = \sum_f \Gamma_{ff} = \frac{G_F M_Z^3}{24\pi\sqrt{2}} \sum_f C_f (v_f^2 + a_f^2), \quad (4.1)$$

where: Γ_{ff} is the partial width for the decay $Z^0 \rightarrow f\bar{f}$; v_f and a_f are the vector and axial vector coupling constants.

$$\begin{aligned} v_f &= \tau_3^f - 4Q_f \frac{g^2}{g^2 + g'^2}; \\ a_f &= -\tau_3^f \end{aligned} \quad (4.2)$$

τ_3^f is twice the third component of the fermion weak isospin; Q_f is the fermion charge; and the constant C_f is defined as

$$C_f = \begin{cases} 1 + \frac{3\alpha}{4\pi} Q_f^2 & \text{for leptons} \\ 3 \cdot \left[1 + \frac{3\alpha}{4\pi} Q_f^2 + \frac{\alpha_s}{\pi} \right] & \text{for quarks.} \end{cases}$$

Note that the expression of each partial width in terms of M_Z has the advantage that the m_{top} and m_{Higgs} dependences are minimized.

The partial widths for a generation of quarks and leptons are listed in Table IV. The last line shows the expected total width for three lepton flavors and five quark flavors. A small phase space suppression factor is included for the $b\bar{b}$ final state.

Table IV

Final State	Γ_{ff}
$\nu\bar{\nu}$	166 MeV
l^+l^-	83 MeV
$u\bar{u}$	297 MeV
$d\bar{d}$	383 MeV
2.75 Generations	2.481 GeV

The actual measurement of M_Z and Γ_Z is made by measuring the cross section for the process $e^+e^- \rightarrow Z^0 \rightarrow f\bar{f}$ for a number of center-of-mass energies about the Z^0 pole. The theoretical Z lineshape is then fit to the measured cross section points to extract the desired parameters. The theoretical lineshape has been the subject of much analysis.^[7] It can be shown that the tree-level lineshape for the process $e^+e^- \rightarrow Z^0 \rightarrow f\bar{f}$ is well-approximated by a relativistic Breit-Wigner form,

$$\sigma_f^0(s) = \frac{12\pi}{M_Z^2} \cdot \frac{s\Gamma_{ee}\Gamma_{ff}}{(s - M_Z^2)^2 + \Gamma_Z^2 s^2/M_Z^2}. \quad (4.3)$$

Equation (4.3) does not apply to the process $e^+e^- \rightarrow e^+e^-$ which occurs via both s-channel and t-channel subprocesses.

The electron and positron radiate real photons rather copiously in a hard collision. The lineshape is strongly affected by the initial state radiation. This effect can be treated in a Drell-Yan-like formalism by introducing an electron structure function. The electron structure function $D(x, s)$ is defined as the probability that

an electron (positron) radiates a fraction $1 - x$ of its initial energy during the collision (of cm energy $\sqrt{s}$). The radiatively corrected cross section can then be written as,

$$\sigma_f(s) = \int dx_1 dx_2 D(x_1, s) D(x_2, s) \sigma_f^0(\hat{s} = x_1 x_2 s), \quad (4.4)$$

where x_1 and x_2 the electron and positron energy fractions. The leading term of the electron structure function has the form,

$$D(x, s) \simeq \frac{\beta}{2} (1 - x)^{\frac{\beta}{2} - 1}, \quad (4.5)$$

where the dimensionless constant β is the effective number of radiation lengths for the process,

$$\beta \equiv \frac{2\alpha}{\pi} \left[\ln\left(\frac{s}{m_e^2}\right) - 1 \right] \simeq 0.11.$$

The effect of the convolution described in equation (4.4) is to reduce the peak cross section by $\sim 25\%$ and to shift the peak of the cross section by roughly 120 MeV from the pole position.

It is convenient to write the radiatively corrected cross section in a form that is close to the underlying Breit-Wigner form,

$$\sigma_f(s) = \frac{12\pi}{M_Z^2} \cdot \frac{s\Gamma_{ee}\Gamma_{ff}}{(s - M_Z^2)^2 + \Gamma_Z^2 s^2/M_Z^2} \cdot [1 + \delta_{RC}(s)], \quad (4.6)$$

where the effects of the radiative corrections are contained in $\delta_{RC}(s)$. Using equation (4.1), we can express all of the quantities that appear in equation (4.6) in terms of a single parameter, M_Z . Note that this choice of parameters minimizes the sensitivity of the lineshape to higher-order terms in m_{top} and m_{higgs} .

Equation (4.6) is the basis for the measurement of a number of Z resonance parameters. The Mark II analysis was performed with several sets of constraints:

1. All resonance parameters are constrained to their Standard Model values. In this case, the only free parameter is M_Z . The measurement was performed by summing all of the final states except the electron-positron final states.

2. The visible partial widths are constrained to their Standard Model values and the invisible width is allowed to vary as a free parameter. The total width Γ_Z is decomposed into visible and invisible portions,

$$\begin{aligned} \Gamma_Z &= \sum \Gamma_{q\bar{q}} + 3\Gamma_{\ell^+\ell^-} + 3\Gamma_{\nu\nu} \\ &= \Gamma_{vis} + \Gamma_{inv}, \end{aligned} \quad (4.7)$$

where the visible width Γ_{vis} contains all hadronic final states and all charged lepton pairs, and Γ_{inv} contains the neutrino decays and any additional unobserved particles. All of the final states except the electron pairs are used to perform the measurement. The data are therefore fit to a function of two parameters (M_Z and Γ_{inv}).

$$\sigma_f(s) = \frac{12\pi}{M_Z^2} \cdot \frac{s\Gamma_{ee}\Gamma_{ff}}{(s - M_Z^2)^2 + (\Gamma_{vis} + \Gamma_{inv})^2 s^2/M_Z^2} \cdot [1 + \delta_{RC}(s)], \quad (4.8)$$

3. The resonance parameters of the total hadronic cross section are not constrained to their Standard Model values. The hadronic cross section is described by the model-independent form,

$$\sigma_{had}(s) = \frac{12\pi}{M_Z^2} \cdot \frac{s\Gamma_{ee}\Gamma_{had}^0}{(s - M_Z^2)^2 + \Gamma_Z^2 s^2/M_Z^2} \cdot [1 + \delta_{RC}(s)], \quad (4.9)$$

where the free parameters are: M_Z , Γ_Z , and the tree-level hadronic peak cross section σ_{had}^0 . The Standard Model prediction for the tree-level peak cross section is,

$$\sigma_{had}^0 = \frac{12\pi}{M_Z^2} \cdot \frac{\Gamma_{ee}\Gamma_{had}}{\Gamma_Z^2} \simeq 41.5 \text{ nb}^{-1}. \quad (4.10)$$

Scanning Theory

A hadron collider gives the experimenter a free energy scan. The hadron structure functions are quite broad in that reasonable quark-quark luminosity is produced over a large range of energies. The electron structure functions have an

integrable singularity at $x = 1$. Most of the e^+e^- luminosity is produced near the nominal value of $\sqrt{s}$. The experimenter can therefore choose the most efficient energy scan to optimize the measurement he/she wishes to measure. Note that an optimal scanning strategy requires some *a priori* knowledge of the parameters that one desires to measure. In the earliest runs of the SLC, the Z^0 mass was not well known and it was necessary to search for an enhancement in the event rate. Once M_Z became somewhat constrained, it was possible to choose very efficient operating points. The presence of the Standard Model as a predictor of widths and couplings made this task much easier.

Let us consider a hypothetical scan of N energy-luminosity points:

$$E_b = E_1, E_2, \dots, E_N$$

$$\int \mathcal{L} dt = L_1, L_2, \dots, L_N.$$

We assume that a cross section σ_i is measured at each point,

$$\sigma_{\text{measured}} = \sigma_1, \sigma_2, \dots, \sigma_N.$$

The M parameters a_j ($j = 1, M$) of our theoretical lineshape $\sigma(E)$ can be extracted from a χ^2 fit to the measured points. The quantity χ^2 is defined as,

$$\chi^2 \equiv \sum_{i=1}^N \frac{[\sigma_i - \sigma(E_i)]^2}{(\delta\sigma_i)^2}, \quad (4.11)$$

where $\delta\sigma_i$ is the error on the i^{th} measurement.

The best estimate of the parameters ($\bar{a}_j$) is the one that minimizes χ^2 . The parameter errors are found from a Taylor expansion of χ^2 about the minimum

value,

$$\chi^2 = \chi^2(\bar{a}) + \frac{1}{2} \sum_{j,k=1}^M \frac{\partial^2 \chi^2}{\partial a_j \partial a_k} (a_j - \bar{a}_j)(a_k - \bar{a}_k) \quad (4.12)$$

$$= \chi^2(\bar{a}) + \sum_{j,k=1}^M (C^{-1})_{jk} (a_j - \bar{a}_j)(a_k - \bar{a}_k)$$

where the matrix C^{-1} is the inverse of the parameter covariance matrix. The error hyperellipsoid is determined by changing χ^2 by one unit about the minimum value. It is straightforward to show that the parameter errors are given by the diagonal elements of the covariance matrix C ,

$$(\delta a_j)^2 = C_{jj}. \quad (4.13)$$

Averaging equation (4.12) over many experiments, the inverse matrix can be expressed in the following form,

$$(C^{-1})_{jk} = \sum_{i=1}^N \frac{1}{(\delta\sigma_i)^2} \cdot \left[\frac{\partial \sigma}{\partial a_j}(E_i) \right] \cdot \left[\frac{\partial \sigma}{\partial a_k}(E_i) \right]. \quad (4.14)$$

Although equation (4.14) is quite general, it is useful to express the cross section errors in terms of the luminosity and the theoretical cross section. Ignoring the statistical errors on the luminosity measurements*, we can express the cross section errors as $(\delta\sigma_i)^2 = \sigma(E_i)/L_i$. Equation (4.14) can then be written as,

$$(C^{-1})_{jk} = \sum_{i=1}^N \frac{L_i}{\sigma(E_i)} \cdot \frac{\partial \sigma}{\partial a_j}(E_i) \cdot \frac{\partial \sigma}{\partial a_k}(E_i) = \sum_{i=1}^N L_i \cdot S(E_i, a_j) \cdot S(E_i, a_k), \quad (4.15)$$

where we define the so-called *sensitivity function* $S(E, a_j)$ as

$$S(E, a_j) \equiv \frac{1}{\sqrt{\sigma(E)}} \cdot \frac{\partial \sigma}{\partial a_j}(E). \quad (4.16)$$

If the lineshape is a function of a single parameter or if the off-diagonal elements of the inverse matrix C^{-1} are small, the parameter errors have a particularly simple

* This assumption is quite valid for the measurement of non-resonant cross sections.

form,

$$(\delta a_j)^{-2} \simeq \sum_{i=1}^N L_i \cdot \left[S(E_i, a_j) \right]^2. \quad (4.17)$$

Equation (4.17) implies that the error δa_j is minimized when the integrated luminosity is concentrated in regions of scan energy where $|S(E, a_j)|$ is large. Note that $|S(E, a_j)|$ is large where the derivative $|\partial\sigma/\partial a_j|$ is large and where the cross section is small.

The correlations between the parameters are described by the off-diagonal elements of the matrices C^{-1} and C (the error ellipsoid is unrotated if they vanish). The presence of non-zero correlation always increases a parameter error beyond the value given in equation (4.17).[†] It is clearly important to minimize the off-diagonal elements by our choice of the scan point luminosities.

Equations (4.15) and (4.13) predict the complete parameter error matrix in terms of the theoretical lineshape and the scan point luminosities. Note that it is assumed that χ^2 is well-defined ($N > M$) and that a sufficient number of events is collected at each point that the errors are Gaussian.

Since any cross section measurement has an associated normalization uncertainty, it is important to consider the sensitivity of the final result to systematic shifts in the measured cross sections. Expanding the theoretical cross section in parameter space about the best estimates $\bar{a}_j$, it is straightforward to derive the average shift in a parameter Δa_j caused by shifts in the measured cross sections $\Delta\sigma_i$.

$$\langle \Delta a_j \rangle = \sum_{k=1}^M C_{jk} \cdot \sum_{i=1}^N L_i \cdot \frac{\Delta\sigma_i}{\sigma_i} \cdot \frac{\partial\sigma}{\partial a_k}(E_i). \quad (4.18)$$

It is clear that we would like to choose the energies and luminosities to minimize the parameter errors and the correlations between the parameters. We can be guided in this task by examining the energy dependence of the functions $S(E, a_j)$.

[†] The presence of non-zero correlation allows the error associated one parameter to leak into the error associated with another parameter.

As an example of the usefulness of the sensitivity functions, let us consider the measurement of the model-independent parameters of the hadronic cross section. For simplicity, we assume that values of M_Z , Γ_Z , and $\sigma_{had}^0(M_Z^2)$ are 91 GeV, 2.5 GeV, and 40 nb, respectively. The sensitivity functions for M_Z , Γ_Z , and $\sigma_{had}^0(M_Z^2)$ are plotted in Figures 4-6 as functions of $E - M_Z$. The maximum sensitivity to M_Z occurs at the scan energies -0.8 GeV and +1.0 GeV about the pole. Note that there is little sensitivity to Γ_Z at these points. The maximum sensitivity to Γ_Z occurs at points that are approximately ± 2 GeV about the pole. If we choose our energy-luminosity points symmetrically about the pole, the sum of the products $S(E_i, M_Z) \cdot S(E_i, \Gamma_Z)$ will tend to cancel since $S(E, M_Z)$ is odd about the pole and $S(E, \Gamma_Z)$ is even about the pole. The maximum sensitivity to σ_{had}^0 occurs at the pole. The same odd-even effect that cancels the M_Z - Γ_Z correlation will cancel the M_Z - σ_{had}^0 correlation. The Γ_Z - σ_{had}^0 correlation cannot be cancelled by a choice of scan energies. However, it is not intrinsically large since $S(E, \Gamma_Z)$ is small in the energy region where $S(E, \sigma_{had}^0)$ is large.

In general, a scan strategy that is based upon equations (4.15) and (4.13) is a problem in linear programming. The scan planner must decide how important various parameters are and what constraints must be satisfied. Nevertheless, fairly simple considerations lead to the conclusion that a minimal Z -pole scan should include points at 0, ± 1 , and ± 2 GeV about the pole.

4.3. THE EXPERIMENTAL ANALYSIS

In order to appreciate the selection criteria that must be applied to the data, we must first discuss the signatures and the relative rates of various processes that occur in an electron-positron collider.

The Electron-Positron Environment

Unlike the situation with hadron colliders, the most copious processes in a high energy e^+e^- collider are also the most interesting ones. The signatures and relative sizes of the various processes are indicated in Table V. The most serious

background to Z^0 production is due to the various two-photon processes. The two-photon background is rather trivial to remove from the data sample (a total energy cut is sufficient to suppress it by several orders of magnitude).

Table V

Event Type	Signature	$\sigma(\sqrt{s} = M_Z)$
$e^+e^- \rightarrow Z^0 \rightarrow \text{hadrons}$	2-3 jets ≥ 20 charged tracks	~ 30 nb
$e^+e^- \rightarrow e^+e^-$ (small angle)	45 GeV clusters in small angle <i>tagger</i>	$\sim 50\text{-}200$ nb (dep on acceptance)
$e^+e^- \rightarrow e^+e^-\ell^+\ell^-$ $e^+e^- \rightarrow e^+e^-h^+h^-$	Transversely balanced low energy track pairs	$\sim 7\text{-}8$ nb (dep on acceptance)
$e^+e^- \rightarrow Z^0 \rightarrow \mu^+\mu^-$	back-to-back high energy tracks	~ 1.5 nb
$e^+e^- \rightarrow Z^0 \rightarrow \tau^+\tau^-$	acolinear track pairs 1-3 combinations	~ 1.5 nb

Event Selection

The Mark II Z^0 mass and width measurement^[6] was performed with a data sample that corresponds to a total integrated luminosity of 19.3 nb^{-1} that was collected at 10 center-of-mass energies. The hadronic final states were selected with the following criteria:

1. Each event is required to contain three or more charged tracks. Each charged track must originate from within a cylindrical volume of 1 cm radius and 6 cm length that is centered upon the nominal interaction point. The momenta of all reconstructed tracks must be larger than 110 MeV/c and the reconstructed polar angle must fall within the region $|\cos\theta| < 0.92$.

2. The visible energy (track momenta and/or calorimeter energy) that is observed in each of the forward and backward hemispheres must be larger than 5% of the center-of-mass energy. The criterion suppresses beam-gas and

two-photon events.

3. Any events which are also identified as $\tau^+\tau^-$ pairs are removed from the sample.

A total of 450 events passed the selection criteria. The detection efficiency for hadronic events is $95.3 \pm 0.6\%$. The residual background contamination is at level of a few parts in 10^3 (mostly from $\tau^+\tau^-$ events).

Leptonic events were selected with the following criteria:

1. Each event is required to contain between two and six charged tracks.
2. The polar angle of the thrust axis must be contained within the region $|\cos\theta_{\text{thrust}}| < 0.65$.
3. The energy measured in the calorimeters must be less than 80% of the center-of-mass energy. (This criterion eliminates e^+e^- pairs.)
4. Each event must satisfy either of the following:

- (a) If the momenta of each of two tracks are larger than 50% of the beam momentum, the event is categorized as a muon pair.
- (b) If the event fails the previous criterion and has a total visible energy that is larger than 10% of the center-of-mass energy, the event is categorized as a $\tau^+\tau^-$ pair.

A total of 30 events passed the lepton selection criteria. The efficiencies of the selection criteria are $99 \pm 1\%$ and $96 \pm 1\%$ for muon and tau pairs produced within the fiducial region $|\cos\theta| < 0.65$, respectively.

Luminosity Measurement

The luminosity at each scan point is inferred from the measured rate of small angle (15-160 milliradian) Bhabha scattering. In the small angle region, this process is dominated by t-channel exchange of photons and is independent of the

parameters of the Z^0 system. The tree-level differential cross section has the form,

$$\frac{d\sigma_{lum}}{d\theta} \simeq \frac{4\pi\alpha^2}{s} \cdot \frac{1}{\theta^3}, \quad (4.19)$$

where the scattering angle θ is assumed to be small. An accurate determination of the luminosity requires that the radiative corrections be included in equation (4.19). Nevertheless, equation (4.19) does illustrate one of the difficulties in the measurement of the luminosity. The measured cross section σ_{lum}^{meas} is a sensitive function of the angular acceptance of the detector edges,

$$\sigma_{lum}^{meas} \simeq \frac{2\pi\alpha^2}{s} \left(\frac{1}{\theta_1^2} - \frac{1}{\theta_2^2} \right), \quad (4.20)$$

where θ_1 and θ_2 are the angles of the inner and outer detector edges.

It is clear from equation (4.20) that the very small angle luminosity monitor, the minisam, has a much larger rate than does the larger angle SAM (the counting rate of the former is six times larger than the latter). However, the larger angle device has good spatial (angular) resolution whereas the minisam has no segmentation in polar angle. The systematic error that is associated with the detector acceptance is substantially smaller for the SAM than for the minisam. The strategy that was used to determine the luminosity of each scan point was therefore to use the minisam to determine the relative luminosities of different scan points and to use the SAM to determine the overall normalization.

Each SAM event was required to satisfy the following selection criteria:

1. The measured energies of the scattered electron and positron were required to be larger than 40% of the beam energy.
2. The measured scattering angles had to satisfy one of the following criteria:
 - (a) Both scattering angles were larger than 65 milliradians. Events in this category were assigned a weight of 1.0.

- (b) One of the scattering angles was larger than 60 milliradians and the second was larger than 65 milliradians. Events in this category were assigned a weight of 0.5.

The weighting procedure reduced the sensitivity of the result to detector misalignments and to radiative corrections. A total of 485 events satisfied the selection criteria.

The theoretical cross section for accepted SAM events is given by the following expression,

$$\sigma_S(E_{cm}) = 25.2 \text{ nb} \cdot \left[\frac{91.1 \text{ GeV}}{E_{cm}} \right]^2,$$

where E_{cm} is the center-of-mass energy. The systematic error due to uncertainties on the detector resolution and position is 2%. The systematic error due to uncertainties on the radiative corrections are taken to be 2%. The combined systematic error is therefore 3%.

The total integrated luminosity for the 10 scan points is therefore evaluated to be,

$$\int \mathcal{L} dt = 19.3 \pm 0.9 \text{ nb}^{-1},$$

where the 5% error includes both statistical and systematic effects.

Each minisam event was required to satisfy the following selection criteria:

1. The measured energies of the scattered electron and positron were required to be larger than 25 GeV in diagonally opposite quadrants.
2. The timing of the minisam signals was required to be consistent with that expected for a scattering process.

A total of 4299 minisam events were recorded during the energy scan. The theoretical cross section for accepted minisam events is given by the following

expression,

$$\sigma_M(E_{cm}) = 230 \text{ nb} \cdot \left[\frac{91.1 \text{ GeV}}{E_{cm}} \right]^2.$$

The luminosity of each point, L_i , is given by the following expression,

$$L_i = \frac{(N_S^i + N_M^i)/(\sigma_S(E_i) + \epsilon_M^i \sigma_M(E_i))}{\sum_i (N_S^i + N_M^i)/(\sigma_S(E_i) + \epsilon_M^i \sigma_M(E_i))} \cdot \int \mathcal{L} dt, \quad (4.21)$$

where: N_M^i and N_S^i are the number of SAM and minisam events recorded at energy E_i , respectively; and where ϵ_M^i is the minisam efficiency for scan point i (the minisam was sensitive to radiation background during some runs).

Experimental Results

In order to improve the statistical accuracy of the mass and width measurements, the $q\bar{q}$ final states are combined with the leptonic final states ($\mu^+\mu^-$ and $\tau^+\tau^-$ within the region $|\cos\theta| < 0.65$) to calculate the cross section at each scan point. The average energy, integrated luminosity, and measured cross section are listed in Table VI.

The measured cross sections and the results of the one, two, and three parameter fits are shown in Figure 7. The parameter estimates for the three fit hypotheses are listed in Table VII.

The measured value of M_Z agrees well with the results of the LEP measurements. The number of neutrino species is consistent with the expected number of 3. Including the systematic errors, the upper limit on the number of light neutrinos is 3.9 with 95% confidence.

Systematic Errors

The errors listed in Table VII include systematic uncertainties on the cross section normalization and on the energy scale of the SLAC. The various resonance parameters vary in their sensitivity to the energy scale and normalization uncertainties.

Table VI

Average energy, integrated luminosity, number of events, MiniSAM efficiency and σ_Z for each energy scan point. The luminosity for each scan point is given by $\text{Lum} = (N_S + N_M)/\sigma_L$, where $\sigma_L = \sigma_S + \epsilon_M \sigma_M$. The given error is the statistical error on N_S and N_M only; there are additional statistical errors on σ_L due to the scaling errors on σ_S and σ_M .

Scan Point	$\langle E \rangle$ (GeV)	N_S	N_M	ϵ_M	Lum. (nb ⁻¹)	Z Decays			σ_Z (nb)
						Had.	Lep.	Tot.	
3	89.24	24	166	0.99	0.68±0.05	3	0	3	4.5 ^{+4.5} _{-2.5}
5	89.98	36	174	0.99	0.76±0.05	8	2	10	13.5 ^{+6.0} _{-4.3}
10	90.35	116	617	1.00	2.61±0.10	60	2	62	24.8 ^{+3.8} _{-3.3}
2	90.74	54	266	0.96	1.21±0.07	33	3	36	31.7 ^{+6.8} _{-5.5}
7	91.06	170	923	0.99	4.08±0.12	114	6	120	31.6 ^{+3.4} _{-3.1}
8	91.43	164	879	0.91	4.12±0.13	108	6	114	29.8 ^{+3.3} _{-2.9}
4	91.50	53	275	0.99	1.23±0.07	33	6	39	34.3 ^{+7.0} _{-5.7}
1	92.16	31	105	0.97	0.54±0.05	11	0	11	21.5 ^{+9.2} _{-6.6}
9	92.22	128	680	0.98	3.05±0.11	67	4	71	24.3 ^{+3.4} _{-3.0}
6	92.96	39	214	0.98	1.00±0.07	13	1	14	14.6 ^{+5.4} _{-4.0}
Totals		815	4299		19.3±0.9	450	30	480	

Table VII

Z resonance parameters. The three fits are described in the text.

Fit	m_Z GeV/c ²	N_ν	Γ GeV	σ_0 nb
1	91.14±0.12	-	-	-
2	91.14±0.12	2.8±0.6	-	-
3	91.14±0.12	-	2.42 ^{+0.45} _{-0.35}	45±4

The determination of M_Z depends completely on the accelerator energy scale. The energies of the SLC beams are measured by a pair of energy spectrometers after they have collided. A schematic diagram of the north (electron) energy spectrometer is shown in Figure 8. The beam is focused by a quadrupole doublet to a point at the detector plane. The beam passes through a small horizontal bend dipole magnet, a large vertical bend dipole magnet, and a second small horizontal bend magnet. The passage of the beam through the horizontal bend magnets produces flat distributions of synchrotron radiation which are detected by a phosphor screen detector. The separation of the flat distributions is proportional to the beam energy. The spectrometers have sufficient resolution to determine the center-of-mass energy (and the value of M_Z) to ± 40 MeV.

The model-independent determinations of M_Z are completely insensitive to the normalization uncertainty. The model constrained determinations of M_Z have a slight sensitivity to the normalization uncertainty. These uncertainties are typically a few MeV or less (even with the model constraints, most of the M_Z information is derived from the resonance shape).

The peak cross section and the invisible width are strongly affected by normalization uncertainty. This can be seen from an inspection of equation (4.8). The invisible width enters the cross section as a component of the total width. The influence of the total width is maximized when the center-of-mass energy is $s = M_Z^2$. The effect of the normalization uncertainty $\delta\sigma$ upon the invisible width is approximately,

$$\begin{aligned}\delta\Gamma_{inv} &\simeq 1.5 \text{ GeV} \cdot \left(\frac{\delta\sigma}{\sigma}\right) = 75 \text{ MeV} \\ \delta N_{\nu} &\simeq 9 \cdot \left(\frac{\delta\sigma}{\sigma}\right) = 0.45.\end{aligned}$$

The measurement of Γ_Z depends almost entirely upon the measurement of the resonance shape. It is therefore insensitive to the absolute energy and normalization errors. It is sensitive to point-to-point errors in the energy and luminosity. These are much smaller than the absolute errors.

The effects of the theoretical uncertainties on the Z lineshape upon the value of the extracted parameters are small as compared with the energy scale and normalization uncertainties.

4.4. MASS AND WIDTH OF THE W

The measurement of the W boson mass and width will become possible in the second phase of LEP operation. The installation of superconducting RF cavities will permit the beam energy to be increased to a value above the threshold for the process $e^+e^- \rightarrow W^+W^-$. The low event rates at the W pair threshold will make essential to optimize the scan to measure the threshold shape.

High Energy e^+e^- Cross Sections

The tree-level expression for the W -pair cross section is somewhat complex.^[9] The inclusion of initial state radiation (as in equation (4.4)) and finite widths for the final state W bosons involves a four dimensional convolution of the tree-level expression. We therefore choose to present only the result of a Monte Carlo integration. The cross section for the process $e^+e^- \rightarrow W^+W^-$ is plotted in Figure 9 as a function of $E_b - M_W$ where E_b is the single beam energy. The mass and width of the W are assumed to be 80 GeV and 2.1 GeV, respectively. Note that three curves are plotted: the dashed curve is the basic tree-level cross section; the dashed-dotted curve is the cross section including the effect of initial state radiation; and the solid curve is the cross section including initial state radiation and the effect of a finite W width. The inclusion of initial state radiation reduces the size of the cross section. The finite W width produces non-zero cross section at energies below the nominal threshold at $E_b = M_W$.

The basic $e^+e^- \rightarrow f\bar{f}$ cross section for five quark and three lepton flavors increases from about 7 units of R at center-of-mass energies below the Z^0 pole to 10 units of R at energies above the Z^0 pole.* At $\sqrt{s} = 160$ GeV, the tree-level cross

* The unit of R is the cross section for $e^+e^- \rightarrow \gamma^* \rightarrow \mu^+\mu^-$. Numerically, the cross section has the value $\sigma_R = 86.8 \text{ nb-GeV}^2/s$.

section is approximately 34 pb. Unfortunately, the initial state radiative corrections increase this number enormously. Although the photon structure functions decrease greatly as x is decreased from 1, the Z pole is sufficiently large that the convolution given in equation (4.4) is several times larger than the tree-level cross section. The process $e^+e^- \rightarrow \gamma Z^0$ therefore dominates the visible cross section at W -pair threshold. Using equation (4.4), we estimate the size of the visible cross section to be ~ 150 pb at $\sqrt{s} = 160$ GeV.

$e^+e^- \rightarrow W^+W^-$ Threshold Scan

There are several different techniques that can be used to measure the W mass at LEP II. It is possible to extract M_W from the measured distributions of jet masses or lepton energies. These methods are described in Reference 10. The technique that we'll discuss here is the measurement of the threshold behavior of the W pair cross section.

It is clear that the W mass can be extracted from the $slep$ in the cross section that is shown in Figure 9. Since there is a large background from ordinary processes, it is necessary to apply selection criteria to the data to improve the signal-to-noise ratio. The background processes produce mostly two- and three jet hadronic events or lepton pair events that are often highly boosted along the beam direction. The visible energy of the background is often small as compared with $\sqrt{s}$. The W -pair events appear most often as four-jet events ($\sim 44\%$ of W -pairs) or as an energetic lepton and two jets ($\sim 44\%$ of W -pairs). The authors of Reference 10 have studied a number of selection criteria to reduce the background cross section to less than ~ 1 pb while retaining $\sim 75\%$ of the four-jet and $\sim 45\%$ of the lepton+two-jet events (we assume that τ leptons cannot be used and that one third of the remaining events are eliminated by the isolation cut used to suppress heavy flavor events). Assuming that the residual background is due to the large $\sqrt{s}$ continuum, the measured cross section would have the following form,

$$\sigma_{meas}(E_b) = \varepsilon \sigma_{WW}(E_b) + \frac{B}{(2E_b)^2}, \quad (4.22)$$

where: ε is the efficiency to identify a W -pair event ($\varepsilon \simeq 0.53$); $\sigma_{WW}(E_b)$ is the cross section plotted in Figure 9; and B is a constant that represents the residual background (which presumably scales as $1/s$).

Sensitivity Functions

We can analyze the M_W and Γ_W sensitivity of a cross section scan of the W pair threshold by using the scanning theory that was discussed in the last section. Numerically differentiating the measured cross section (as defined in equation (4.22)), it is straightforward to calculate the sensitivity functions for M_W , Γ_W , and the background constant B . For the purpose of this exercise, we assume that $B = 1 \text{ pb} \cdot (2M_W)^2$ or that the background cross section is 1 pb at W -pair threshold.

The sensitivity function $S(E_b, M_W)$ is plotted in Figure 10 as a function of $\epsilon_b = E_b - M_W$. Note that the maximum sensitivity occurs at $\epsilon_b \simeq 0.5 \text{ GeV}$.

The sensitivity function $S(E_b, \Gamma_W)$ is shown in Figure 11 as a function of ϵ_b . As one would expect, it peaks just below the nominal threshold ($\epsilon_b = -1 \text{ GeV}$) where the width-induced *tail* in the cross section is largest. The function $S(E_b, \Gamma_W)$ decreases rapidly as E_b is increased. It passes through zero near $\epsilon_b = 1 \text{ GeV}$ and plateaus above $\epsilon_b = 3 \text{ GeV}$. The sensitivity in the plateau region is due to the reduction in the cross section caused by the finite width (see Figure 9). The maximum value of $|S(E_b, \Gamma_W)|$ is smaller than the maximum value of the mass sensitivity function by a factor of three. A good measurement of Γ_W will clearly require a substantial commitment of luminosity to a point of very small cross section. Note that the product $S(E_b, M_W) \cdot S(E_b, \Gamma_W)$ is an odd function about the point $\epsilon_b = 1 \text{ GeV}$. In principle, the M_W - Γ_W correlation can be cancelled by measuring the cross section on both sides of this point. The functions $S(E_b, M_W)$ and $S(E_b, \Gamma_W)$ are not large in the region $\epsilon_b > 1 \text{ GeV}$. The cancellation of the correlation therefore requires a substantial commitment of luminosity to a relatively insensitive region.

The function $S(E_b, B)$ is plotted as a function of ϵ_b in Figure 12. As one would

expect, the background sensitivity is largest at small beam energy and decreases dramatically as E_b increases through the W pair threshold. Note that it is possible to cancel the $B\text{-}\Gamma_W$ correlation but that it is not possible to cancel the $B\text{-}M_W$ correlation.

Scan Strategies

It is clear that precise measurements of M_W and Γ_W require that LEP be operated in regions of small cross section. Since all other studies of the W -pair system require a large sample of data, there will be considerable pressure to operate the machine on the cross section plateau at the largest available energy. In order to estimate how precisely M_W and Γ_W could be measured in a 1-2 year run (500 pb⁻¹), we assume that 50% of the luminosity is dedicated to operating at the largest available energy (we assume that $\epsilon_b = 15$ GeV or $\sqrt{s} = 190$ GeV is achieved) and the remaining 50% is dedicated to operation in the threshold region.

It is instructive to first consider an extremely unrealistic scan scenario. We assume that we will measure only one parameter and that the other parameters are precisely known. In this case, we need only one scan point in the threshold region for a constrained fit. We choose to allocate the entire 250 pb⁻¹ luminosity to operation at the most mass-sensitive point ($\epsilon_b = 0.5$ GeV) or at the most width-sensitive point ($\epsilon_b = -1$ GeV). Using equation (4.17) we estimate the precision of these measurements to be

$$\delta M_W = 92 \text{ MeV} \quad \text{or} \quad \delta \Gamma_W = 286 \text{ MeV}.$$

The M_W measurement would be a very desirable result. The Γ_W measurement is not competitive with the recent indirect determinations that have been published by the CDF and UA2 collaborations,^[11,12]

$$\Gamma_W = (0.85 \pm 0.08) \cdot \Gamma_Z = 2.19 \pm 0.20 \text{ GeV (CDF)}$$

$$\Gamma_W = (0.89 \pm 0.08) \cdot \Gamma_Z = 2.30 \pm 0.20 \text{ GeV (UA2)}.$$

Since the width cannot be measured to an interesting level, it is clearly unwise to

design a scan to measure Γ_W . We therefore concentrate on the measurement of M_W .

A real measurement of M_W will require that the background constant B be varied as a fit parameter. Unfortunately, the $B\text{-}M_W$ correlation cannot be canceled by a clever choice of scan points. It is therefore necessary to measure both parameters well.

The number of scan points is somewhat arbitrary. A minimum of three points are required to constrain the two parameter problem. The presence of a high energy point implies that only two points are needed in the threshold region. Equation (4.15) implies that several closely spaced points in a region of large sensitivity are equivalent to a single point in the same region. We can therefore analyze the optimization of the M_W measurement by considering a two-point threshold measurement.

An optimal scan must include an energy point in a region of large background sensitivity $|S(E_b, B)|$ and a point near the maximum of the mass sensitivity function $|S(E_b, M_W)|$. We choose the scan point energies to be $\epsilon_b = -5$ GeV and $\epsilon_b = 0.5$ GeV, respectively.* The apportionment of the available luminosity between the two points is a straightforward problem in one-dimensional optimization. We find that the error δM_W has a very broad minimum about the ratio of luminosities, $L(0.5 \text{ GeV})/L(-5 \text{ GeV}) \simeq 2/1$. If the luminosities of the -5 GeV and 0.5 GeV points are 85 pb⁻¹ and 165 pb⁻¹, respectively, the minimum value of the error δM_W is approximately 155 MeV.

A two-point threshold scan is somewhat risky. It is safer to bracket the region of maximum M_W sensitivity with several scan points. We therefore construct an optimal four-point scan (a five-point measurement when the $\epsilon_b = 15$ GeV point is included) by assigning one third of the 165 pb⁻¹ (55 pb⁻¹) to each of three points: $\epsilon_b = 0$ GeV, 0.5 GeV, and 1.0 GeV. It is instructive to compare this scan

* Varying the energy of the second point about $\epsilon_b = 0.5$ GeV verifies that the $B\text{-}M_W$ correlation does not shift the point of maximum M_W sensitivity.

(Scan 1) with a slightly modified version. The modified version (Scan 2) is created by shifting the luminosity from the $\epsilon_b = 0$ GeV point to $\epsilon_b = -1$ GeV. We expect the second scan strategy to improve the width measurement at the expense of the mass measurement. Finally, we note that our modified scan strategy is similar to the scan strategy that was studied in Reference 10 (which we label Scan 3). The authors of Reference 10 assigned 100 pb^{-1} to each of the following five points: $\epsilon_b = -5 \text{ GeV}$, -1 GeV , 0 GeV , 1 GeV , and 15 GeV .

Using equation (4.15) and the sensitivity functions, the performance of each scan scenario can be estimated. The expected number of detected events and the expected precisions δM_W , $\delta \Gamma_W$, and δB are listed in Table VIII for each of the three scan strategies. The presence of a high energy point in each strategy reduces the M_W - Γ_W correlation sufficiently that the M_W precision obtained from the three parameter fit is essentially identical to that obtained from a two-parameter fit.

As one might expect, the third scan strategy which allocates 400 pb^{-1} to the threshold measurement provides the most precise M_W measurement, $\delta M_W = 150 \text{ MeV}$. The M_W precision obtained from the optimized mass scan (Scan 1) is worse by 7%. Note however, that Scan 1 produces nearly 60% more events than does Scan 3. Surprisingly, the second scan strategy provides a slightly better width measurement than does the third strategy. This occurs because the second scan produces a smaller B - Γ_W correlation than does the third scan strategy.

It is clear from equation (4.22) that the functions $S(E_b, a_j)$ are sensitive to the level of residual background and to the W -pair detection efficiency. We investigate these effects by reducing the background constant to $B = 0.5 \text{ pb} \cdot (M_W)^2$ and by increasing the detection efficiency to $\epsilon_{ww} = 0.70$. The results are listed in Table VIII. The error δM_W is improved by approximately 20 MeV in the case that the background is reduced by a factor of two. The mass error is improved by approximately 30 MeV when the efficiency is increased. Note that the optimal luminosity ratio $L(0.5 \text{ GeV})/L(-5 \text{ GeV})$ is nominally sensitive to both effects. However, the optimal region is so broad that the use of a 2/1 ratio degrades the

Table VIII

The predicted results of three different five-point measurements of the W -pair threshold. Scan 1 is optimized for the measurement of M_W . Scan 2 is an attempt to improve the measurement of Γ_W . Scan 3 is identical to the threshold scan used in Reference 10. The results are presented for several assumptions about the level of residual background B and the W -pair detection efficiency.

Quantity	Scan 1	Scan 2	Scan 3
$L[-5 \text{ GeV}] (\text{pb}^{-1})$	85	85	100
$L[-1 \text{ GeV}] (\text{pb}^{-1})$	0	55	100
$L[0 \text{ GeV}] (\text{pb}^{-1})$	55	0	100
$L[0.5 \text{ GeV}] (\text{pb}^{-1})$	55	55	0
$L[1 \text{ GeV}] (\text{pb}^{-1})$	55	55	100
$L[15 \text{ GeV}] (\text{pb}^{-1})$	250	250	100
$B = 1.0 \text{ pb} \cdot [2M_W]^2$ $\epsilon_{ww} = 0.53$			
Number of Events	2951	2912	1863
$\delta M_W (\text{MeV})$	160	176	150
$\delta \Gamma_W (\text{MeV})$	531	482	492
$\delta B (\text{pb} \cdot [2M_W]^2)$	9.12	0.12	0.12
$B = 0.5 \text{ pb} \cdot [2M_W]^2$ $\epsilon_{ww} = 0.53$			
Number of Events	2737	2698	1627
$\delta M_W (\text{MeV})$	137	154	130
$\delta \Gamma_W (\text{MeV})$	508	450	448
$\delta B (\text{pb} \cdot [2M_W]^2)$	0.096	0.098	0.098
$B = 1.0 \text{ pb} \cdot [2M_W]^2$ $\epsilon_{ww} = 0.70$			
Number of Events	3760	3709	2309
$\delta M_W (\text{MeV})$	130	144	123
$\delta \Gamma_W (\text{MeV})$	453	407	410
$\delta B (\text{pb} \cdot [2M_W]^2)$	0.12	0.13	0.13

result by less than 1%.

Systematic Errors

The measurement of the W -pair threshold is affected by systematic uncertainties on the energy scale and cross section normalization. The energy scale uncertainty affects the M_W measurement directly. Assuming that the fractional error on the beam energy scale is constant, the uncertainty on M_W should be comparable to the one that applies to the M_Z measurement. By 1994, this uncertainty is expected to be ~ 20 MeV.

The sensitivity of the results given in Table VIII to normalization errors can be estimated from equation (4.18). Taking the first scan strategy as an example, we estimate that the uncertainties on the parameters are related to an overall normalization uncertainty $\delta\sigma/\sigma$ as follows,

$$\begin{aligned}\delta M_W &= -2.26 \text{ GeV} \cdot \frac{\delta\sigma}{\sigma} \\ \delta \Gamma_W &= -19.3 \text{ GeV} \cdot \frac{\delta\sigma}{\sigma}.\end{aligned}$$

The normalization error must be controlled to the 3% level to avoid inflating the M_W error.

Sensitivity to Assumptions

Our analysis assumes that we have complete *a priori* knowledge of the W resonance parameters. Although the characteristic width in E_b space of the M_W -sensitive region is larger than the current uncertainty on M_W , our precision estimates are likely to be somewhat optimistic. It is possible to alter the results by $\lesssim 10\%$ by varying the resonance parameters over reasonable intervals.

Conclusions

Despite the uncertainties on the ultimate W -pair detection efficiency and residual background contamination, several conclusions can be drawn from this analysis:

1. The most sensitive scan region for the measurement of M_W is $e_b = 0.1$ GeV. *The mapping of the entire threshold shape would produce a less precise measurement.*
2. It is not possible to remove the correlation between the background parameter and M_W by a clever choice of scan point energies. This implies that a scan point of energy below the nominal threshold is quite important. If the energy is chosen to be $e_b = -5$ GeV ($E_b = 75$ GeV), an M_W -optimized scan strategy would allocate twice as much integrated luminosity to the M_W sensitive region as is allocated to the low energy point.
3. A measurement of M_W at the $\lesssim 160$ MeV level is possible with the dedication of a large integrated luminosity (250 pb^{-1}) and good control of the background contamination.
4. The measurement of Γ_W to an interesting level is difficult or impossible. It is probably unwise to attempt anything more than a cursory measurement.

5. The Search for New Particles

The Z^0 is the largest-mass neutral particle known to exist. Its couplings to fundamental particles are unambiguously determined from the quantum numbers of the particle in question. The strengths of these couplings are fairly uniform which implies that the Z^0 is remarkably democratic in its choice of final state. The branching ratio of the Z^0 into most hypothetical final states is typically larger than 1% unless it is suppressed by phase space or virtual intermediate states. The Z pole is therefore a good place to search for new particles.

The branching ratios of the Z^0 into hypothetical final states are large enough that the very modest Mark II data sample is adequate to perform new particle searches. A number of searches have been performed for new quarks and leptons,^[12] supersymmetric particles,^[13] and non-standard extensions to the Higgs sector.^[14]

Since most of these searches have also been performed by the LEP collaborations, I will discuss the only search which is unique to the Mark II, the search for doubly charged Higgs bosons.

5.1. DOUBLY CHARGED HIGGS BOSONS

There are currently two popular scenarios that give rise to doubly charged Higgs bosons. The first, known as the Gelmini-Roncadelli model,^[16-17] is a straightforward extension of the standard model to include a Majorana mass for the left-handed neutrino. The second scenario is the left-right symmetric extension of the standard model.^[18] Before describing these two models, it is worth reviewing the mass generation scheme of the Standard Model.

5.2. THE STANDARD MODEL

The Standard Model describes the masses of all leptons in terms of a trilinear Lagrangian of the form,

$$\mathcal{L}_{\text{Dirac}} = f_D^e l_L \Phi l_R + f_D^\nu \bar{l}_L (\tau_2 \Phi^*) \nu_R + \text{h.c.} \quad (5.1)$$

where: f_D^e, f_D^ν are dimensionless coupling constants; Φ is the ordinary (isodoublet) Higgs field; l_R is the right-handed charged-lepton field; τ_2 is the standard Pauli matrix; ν_R is the right-handed neutrino field; and where l_L is the left-handed doublet

$$l_L \equiv \begin{pmatrix} \nu \\ e^- \end{pmatrix}_L.$$

Note that the quantum numbers (I_w, Y) of the Higgs boson, $(\frac{1}{2}, -1)$, are equal to the quantum numbers of the bilinear product $\bar{e}_L e_R, (\frac{1}{2}, +1) \otimes (0, -2)$.

* There is a pair of constants f_D^e, f_D^ν for each lepton generation. For simplicity, generational labels are suppressed in this and the following expressions.

The actual mass terms are generated by the same spontaneous symmetry breaking that generates the gauge boson masses. The Higgs field is expanded about its non-zero minimum,

$$\Phi = e^{i\phi(x)} \begin{pmatrix} 0 \\ v + \eta \end{pmatrix} \quad (5.2)$$

where: $\phi(x)$ is a phase function; v is the vacuum expectation value of the Higgs field; and η is the physical Higgs field. Substituting equation (5.2) into equation (5.1) the usual Dirac mass terms for the charged and neutral leptons emerge

$$\mathcal{L}_{\text{Dirac}} = m_e (\bar{l}_L l_R + \bar{l}_R l_L) + m_\nu (\bar{\nu}_L \nu_R + \bar{\nu}_R \nu_L)$$

where the neutrino and lepton masses are $m_\nu = f_D^\nu v$ and $m_e = f_D^e v$, respectively.

5.3. THE GELMINI-RONCADELLI MODEL

The entire motivation of the Gelmini-Roncadelli model is to give the neutrino a mass without adding a right-handed (neutrino) field to the theory. The secret of doing this is to note that the charge conjugate of the left-handed lepton field doublet, l_L^c , is projected from the charge conjugate field $l^c = C \bar{l}^T$ (C is the charge conjugation matrix) by a *right-handed* projection operator,

$$l_L^c \equiv C \left[\frac{1}{2} (1 - \gamma_5) l \right]^T = \frac{1}{2} (1 + \gamma_5) C \bar{l}^T = \frac{1}{2} (1 + \gamma_5) l^c.$$

It is important to emphasize that l_L^c creates left-handed fermions and destroys right-handed anti-fermions exactly as $\bar{l}_L$ does. Some authors like to confuse everything by labelling l_L^c with an R subscript. It is clear that the bilinear $\bar{l}_L^c l_L$ does not vanish and represents a kind of mass term. In fact, it represents the mass term for a self-conjugate or Majorana field.

To generate a Majorana mass term from the vacuum expectation value of a Higgs field, we note that the Higgs field must have the same quantum numbers as

the bilinear $\ell_L^c \ell_L$,

$$\left(\frac{1}{2}, -1\right) \otimes \left(\frac{1}{2}, -1\right) = (1, -2) \oplus (0, -2).$$

The bilinear must be coupled to a weak isotriplet or a weak isosinglet. The charge of the isosinglet must be $Q = I_3 + Y/2 = -1$. It therefore cannot be coupled to the neutral neutrino pair. On the other hand, the charges of the isotriplet Higgs are $Q = I_3 + Y/2 = 0, -1, -2$. The presence of a neutral member allows one to construct a mass term of the form

$$\mathcal{L}_{\text{Majorana}} = f_M \ell_L^c \vec{H} \cdot (i\tau_2 \vec{\tau}) \ell_L + \text{h.c.} \quad (5.3)$$

where f_M is a dimensionless coupling constant; $\vec{H}$ are the three Higgs fields; and $\vec{\tau}$ are the three Pauli matrices.

The actual mass term is generated by giving the neutral member of the isotriplet a vacuum expectation value as follows

$$\langle \vec{\tau} \cdot \vec{H} \rangle = \left\langle \begin{pmatrix} H^- & \sqrt{2}H^{--} \\ \sqrt{2}H^0 & -H^- \end{pmatrix} \right\rangle = \begin{pmatrix} 0 & 0 \\ v_T & 0 \end{pmatrix}. \quad (5.4)$$

Substituting equation (5.4) into equation (5.3) produces the desired Majorana mass term

$$\mathcal{L}_{\text{Majorana}} = M_\nu (\nu_L^c \nu_L + \nu_L^c \nu_L^c) \quad (5.5)$$

where the Majorana mass, M_ν , is defined as $M_\nu = f_M v_T$. Note that the vacuum expectation value v_T must be small to avoid disturbing the ρ parameter,

$$\rho \equiv \frac{M_W^2}{M_Z^2 \cos^2 \theta_w} = \frac{v^2 + 2v_T^2}{v^2 + 4v_T^2}.$$

However, since we know that any left-handed neutrino is light as compared with the ordinary Higgs VEV (v is roughly 250 GeV), this condition is rather naturally satisfied.

Equation (5.3) also contains pieces that involve the coupling of the H^- to a lepton-neutrino pair (the H^- doesn't couple to quarks since $Y = -2$),

$$\mathcal{L}_{H^\pm} = -\frac{1}{\sqrt{2}} g_U \left[(\bar{\nu}_L \ell_L^c + \ell_L \nu_L^c) H^- + (\bar{\nu}_L^c \ell_L + \ell_L^c \nu_L) H^+ \right] \quad (5.6)$$

and the coupling of the H^{--} to a left-handed charged-lepton pair

$$\mathcal{L}_{H^{++}} = -g_U (\bar{\ell}_L \ell_L^c H^{--} + \bar{\ell}_L^c \ell_L H^{++}) \quad (5.7)$$

where $g_U \equiv \sqrt{2} f_M$ is a dimensionless coupling constant. Note that we would expect to g_U to increase with neutrino mass. The coupling of the H^{--} to $\tau^+ \tau^-$ could therefore be significantly larger than the couplings to $\mu^+ \mu^-$ or $e^+ e^-$. Although both the singly and the doubly charged Higgs bosons formally violate lepton flavor, only the doubly charged member visibly manifests the effect (because a light Majorana neutrino is virtually indistinguishable from a Dirac neutrino). Since the H^{--} couples only to charged lepton pairs, most existing searches for lepton flavor violation are fairly insensitive to doubly charged Higgs bosons. The existing limits are therefore relatively weak.^[19-21]

5.4. LEFT-RIGHT SYMMETRIC MODELS

Another model that incorporates doubly charged Higgs bosons is the so-called Left-Right Symmetric model of Pati and Salam.^[16] As its name implies, this model begins by treating both left- and right-handed fermions in a symmetric fashion. A right-handed weak isospin quantum number is added to the theory. All fermions are singlets of one isospin and doublets of the other. This permits one to discard the weak hypercharge quantum number and to replace it with a more physical one, $B - L$ ($B \equiv$ baryon number and $L \equiv$ lepton number). The electric charge of each state is then given by the relationship

$$Q = I_3^L + I_3^R + \frac{1}{2}(B - L)$$

where I^L and I^R are the left- and right-handed weak isospins, respectively. The

quantum numbers $(I^L, I^R, B - L)$ of the quark and lepton doublets are therefore:

$$\begin{pmatrix} u \\ d \end{pmatrix}_L \equiv \begin{pmatrix} \frac{1}{2}, 0, \frac{1}{3} \end{pmatrix} \quad \begin{pmatrix} \nu \\ e^- \end{pmatrix}_L \equiv \begin{pmatrix} \frac{1}{2}, 0, -1 \end{pmatrix}$$

$$\begin{pmatrix} u \\ d \end{pmatrix}_R \equiv \begin{pmatrix} \frac{1}{2}, \frac{1}{3} \end{pmatrix} \quad \begin{pmatrix} \nu \\ e^- \end{pmatrix}_R \equiv \begin{pmatrix} 0, \frac{1}{2}, -1 \end{pmatrix}$$

The gauge group of this model is expanded to $SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}$. Since we know that nature is left-handed in the low energy limit, the right-handed gauge symmetry must be broken at a significantly larger mass scale than its left-handed counterpart. The minimal Higgs sector that preserves the correct low energy phenomenology consists of a bidoublet Φ with quantum numbers $(\frac{1}{2}, \frac{1}{2}, 0)$ and two triplets, $\vec{H}_L$ and $\vec{H}_R$, with the quantum numbers $(1, 0, -2)$ and $(0, 1, -2)$, respectively. The vacuum expectation values of the Higgs sector are

$$\langle \Phi \rangle = \left\langle \begin{pmatrix} \phi_1^0 & \phi_1^+ \\ \phi_2^- & \phi_2^0 \end{pmatrix} \right\rangle = \begin{pmatrix} \kappa & 0 \\ 0 & \kappa' \end{pmatrix}$$

$$\langle \vec{\tau} \cdot \vec{H}_L \rangle = \left\langle \begin{pmatrix} H_L^- & \sqrt{2}H_L^{--} \\ \sqrt{2}H_L^0 & -H_L^- \end{pmatrix} \right\rangle = \begin{pmatrix} 0 & 0 \\ v_L & 0 \end{pmatrix}$$

$$\langle \vec{\tau} \cdot \vec{H}_R \rangle = \left\langle \begin{pmatrix} H_R^- & \sqrt{2}H_R^{--} \\ \sqrt{2}H_R^0 & -H_R^- \end{pmatrix} \right\rangle = \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix}$$

The Standard model phenomenology is preserved by choosing the vacuum expectation values κ, κ', v_R , and v_L to have the following hierarchy

$$v_R \gg [\kappa^2 + \kappa'^2]^{\frac{1}{2}} \gg v_L.$$

The Higgs sector now contains 20 degrees of freedom. Six degrees of freedom are used to provide mass for the $W_L^\pm, Z_L^0, W_R^\pm$, and Z_R^0 gauge bosons. The remaining 14 degrees of freedom contain four doubly-charged states, four singly-charged states,

and six neutral states. Since the $\vec{H}_L$ and $\vec{H}_R$ triplets already have the quantum numbers of the bilinear combinations $\bar{\ell}_L^c \ell_L$ and $\bar{\ell}_R^c \ell_R$, it is quite natural to include the following Majorana mass term for the neutrino sector (in addition to the normal Dirac mass term)

$$\mathcal{L}_{\text{Majorana}} = f_M^L \bar{\ell}_L^c \vec{H}_L \cdot (i\tau_2 \vec{\tau}) \ell_L + f_M^R \bar{\ell}_R^c \vec{H}_R \cdot (i\tau_2 \vec{\tau}) \ell_R + \text{h.c.} \quad (5.8)$$

There are now three masses for each neutrino generation: a left-handed Majorana mass, $M_L = f_M^L v_L$; a Dirac mass, $m_D = f_D^D \kappa$; and a right-handed Majorana mass, $M_R = f_M^R v_R$. Unless something very perverse is done with the coupling constants, the natural hierarchy of the masses is $M_L \ll m_D \ll M_R$. In fact, the most natural occurrence would be to have $m_D = m_\ell$. Making this assumption, the mass lagrangian for a neutrino generation can be written in the form^{*}

$$\mathcal{L}_{\text{Neutrino mass}} = \begin{pmatrix} \bar{\psi}_L & \bar{\psi}_R \end{pmatrix} \begin{pmatrix} M_L & \frac{1}{2}m_\ell \\ \frac{1}{2}m_\ell & M_R \end{pmatrix} \begin{pmatrix} \psi_L \\ \psi_R \end{pmatrix} \quad (5.9)$$

where the fields $\psi_L = \nu_L + \nu_L^c$ and $\psi_R = \nu_R + \nu_R^c$ are the left- and right-handed Majorana fields, respectively. The physical neutrino masses are given by the eigenvalues of equation (5.9),

$$\begin{aligned} m_1 &\simeq \frac{m_\ell^2}{4M_R} \\ m_2 &\simeq M_R. \end{aligned} \quad (5.10)$$

There is a light, left-handed neutrino (presumably the *standard* one) and a very heavy, right-handed neutrino (the magnitude of M_R is roughly that of the right-handed W boson). The beauty (sic) of this model is that light, left-handed neutrinos are generated without *diddling* coupling constants. Since the mass of each light neutrino generation increases with m_ℓ^2 , the coupling constants f_M^L and f_M^R do not necessarily increase with generation.

* The hierarchy $M_L \ll m_\ell \ll M_R$ is commonly known as the *Seesaw Mechanism*.

The same comment must therefore apply to the couplings of the doubly charged Higgs bosons to left- and right-handed charged-lepton pairs. Using equation (5.8), the $H^{\pm\pm} - \ell^{\pm} \ell^{\pm}$ couplings are

$$\mathcal{L}_{H^{\pm\pm}} = - \left[g_{\ell\ell}^L (\bar{\ell}_L \ell_L^c H_L^{\pm\pm}) + g_{\ell\ell}^R (\bar{\ell}_R \ell_R^c H_R^{\pm\pm}) + \tilde{f}_L^c \ell_L H_L^{\pm\pm} + \tilde{f}_R^c \ell_R H_R^{\pm\pm} \right] \quad (5.11)$$

where $g_{\ell\ell}^L = \sqrt{2} f_M^L$ and $g_{\ell\ell}^R = \sqrt{2} f_M^R$.

Unlike the Gelmini-Roncadelli model, there is no reason to expect that the $H^{\pm\pm}$ decays dominantly into $\tau^{\pm} \tau^{\pm}$ pairs.

5.5. EXISTING LIMITS

If we assume that the coupling of the doubly charged Higgs to lepton pairs is diagonal in lepton flavor, the best limits^[19-21] on the existence of doubly-charged Higgs bosons are derived from the Bhabha scattering data of several PEP and PETRA experiments and from searches for conversion of muonium ($\mu^+ e^-$) into antimuonium ($\mu^- e^+$). Both processes involve the t-channel exchange of a doubly charged Higgs boson. These searches have the unfortunate property that they limit the ratio g^*/M_H . Searches for real $H^{\pm\pm}$ production can extend the limits to very small values of the coupling constant for moderate mass values.

5.6. THE CROSS SECTION FOR $e^+ e^- \rightarrow H^{++} H^{--}$

The coupling of the Z^0 to a pair of doubly charged Higgs bosons of either the Gelmini-Roncadelli type^[17] or of the Left-Right Symmetric type^[22] is given by the following expression,

$$\mathcal{L}_{Z^0-H^{\pm\pm}} = -i\epsilon \left[\frac{I_3^L - Q_H \sin^2 \theta_w}{\sin \theta_w \cos \theta_w} \right] Z^\alpha \phi_H^\dagger (\overleftrightarrow{\partial}_\alpha - \overleftrightarrow{\partial}_\alpha) \phi_H \quad (5.12)$$

where: ϵ is the electric charge; Q_H is the charge of the Higgs boson; $\sin \theta_w$ is the sine of the electroweak mixing angle; Z^α is the Z^0 field; and ϕ_H is the Higgs field.

Note that the quantum numbers (I_3^L, Q_H) are (1,2) for the Gelmini-Roncadelli Higgs and for the left-handed Higgs of the Left-Right Symmetric model. The right-handed Higgs has the quantum numbers (0,2).

Using equation (5.12) it is straightforward to calculate the cross section for the process $e^+ e^- \rightarrow H^{++} H^{--}$,

$$\frac{d\sigma}{d\cos\theta} = \frac{9\pi}{M_Z^2} \frac{\Gamma_{HH}^0 \Gamma_{ee} s}{(s - M_Z^2)^2 + s\Gamma_Z^2} \left[1 - \frac{4M_H^2}{s} \right] (1 - \cos^2\theta) \quad (5.13)$$

where: M_Z, Γ_Z are the mass and width of the Z^0 , respectively; Γ_{ee} is the leptonic width of the Z^0 ; s is the square of the center-of-mass energy; and Γ_{HH}^0 is the partial width for the decay $Z^0 \rightarrow H^{++} H^{--}$ (unsuppressed by phase space),

$$\Gamma_{HH}^0 = \frac{G_F M_Z^2}{6\pi\sqrt{2}} (I_3^L - 2\sin^2\theta_w)^2. \quad (5.14)$$

The ratio of the decay rates of the Z^0 into Higgs pairs and into a single neutrino species can therefore be written as

$$\frac{\Gamma(Z^0 \rightarrow H^{++} H^{--})}{\Gamma(Z^0 \rightarrow \nu\bar{\nu})} = 2(I_3^L - 2\sin^2\theta_w)^2 \beta^3 \simeq \beta^3 \cdot \begin{cases} 0.57, & H_L^{\pm\pm}; \\ 0.43, & H_R^{\pm\pm} \end{cases}$$

where the Higgs velocity β is defined as

$$\beta \equiv \left[1 - \frac{4M_H^2}{M_Z^2} \right]^{\frac{1}{2}}.$$

The branching ratio of the Z^0 into Higgs pairs is roughly one half of the branching ratio for a neutrino species (multiplied by a β^3 factor).

5.7. THE DECAY OF THE DOUBLY CHARGED HIGGS BOSON

The conservation of quantum numbers greatly restricts the available decay modes of triplet higgs bosons. A member of the triplet can decay into a lepton pair or into a (perhaps virtual) W boson and another member of the Higgs triplet. To fully understand the branching ratios of a doubly charged Higgs boson, we must know the hierarchy of the triplet masses. In order to simplify the search, it is assumed that the doubly charged member of the triplet can decay only into lepton pairs. The decay rate of a doubly charged Higgs boson into a same-sign pair of leptons, $\Gamma_{\ell\ell}$, is given by the following expression:^[17,23]

$$\Gamma_{\ell\ell} = \frac{g_H^2}{8\pi} M_H \left[1 - \frac{2m_\ell^2}{M_H^2} \right] \left[1 - \frac{4m_\ell^2}{M_H^2} \right]^{1/2} \quad (5.15)$$

where m_ℓ is the lepton mass. The Higgs bosons are therefore short-lived (in an experimental sense) unless the coupling constants g_H are *very small* (less than 10^{-7}).

5.8. THE MARK II SEARCH

In order to place limits on the decay $Z^0 \rightarrow H^{++}H^{--}$, we must search for final states of the form $\ell^+\ell^+\ell^-\ell^-$ where ℓ and ℓ may or may not be the same flavor of lepton. Of the six possible four-lepton final states, the most difficult to detect is the one consisting of four τ leptons. The strategy of this analysis is to define a set of topological selection criteria that can identify the four- τ final state with high efficiency. It is clear that such criteria select four-lepton final states that contain two or more stable leptons with comparable or larger efficiency.

Since 90% of all four- τ events decay into six or fewer charged particles, we consider only those event candidates that contain six or fewer charged tracks that project into a cylindrical volume of 2 cm radius and 6 cm length that is centered on the interaction point of the SLC. In order to suppress two-photon events and

badly accepted hadronic final states, we require that the scalar sum of the track momenta be at least 10 GeV/c.

Isolated τ leptons appear as isolated single tracks or as low-mass clusters of tracks. The 4-vectors of the charged tracks are thus subjected to a mass-based clustering algorithm. Initially, each track is defined to be a cluster. The pair of clusters with the smallest invariant mass is merged if its mass is less than 2.0 GeV/c². The procedure is repeated until all pairs of clusters have invariant masses larger than 2.0 GeV/c².

We expect most $H^{++}H^{--}$ events to appear as four-cluster events. There is a reasonable probability, however, that a cluster occurs in one of the forward regions, $|\cos\theta| > 0.80$, and is not detected (approximately 30% of all events fall into this category). We therefore require that each event candidate contain either three or four clusters of energy larger than 1.0 GeV. The net charge of each cluster must be unity. The event must not contain any clusters with charges larger than unity. The net event charge must be zero for four-cluster candidates or unity for three-cluster candidates.

None of the events in the Mark II data sample pass the selection criteria. Using a Monte Carlo simulation, we can predict the number of events that would have been observed if doubly charged Higgs bosons were present in the data. The detection probabilities for electron and muon final states are essentially identical. The limit is therefore a function of the Higgs boson to $\tau\tau$ branching ratio, B_τ , only. In the regions of parameter space that the expected number of observed events is larger than 2.3 and 3.0, we can exclude the presence of the $H^{\pm\pm}$ with 90% and 95% confidence. In the short lifetime region, the intervals of M_H that are excluded with 90% confidence and with 95% confidence are listed in Table IX for several values of B_τ and I_3^L . The upper limits on M_H are due to the 3^3 suppression of the cross section. The lower limits are due to the loss of efficiency as M_H becomes small. The efficiency function for stable leptons falls sharply at the cluster mass of 2.0 GeV/c². The 90% and 95% confidence limits occur quite close to this point. For $B_\tau = 0.5$,

the number of events with four stable leptons is sufficient to exclude values of M_H down to the τ -pair threshold. The short lifetime constraint requires that there be at least one coupling constant in the region $g_{\tau\tau} \gtrsim 7.4 \times 10^{-7} / \sqrt{M_H}$ (M_H in GeV/c^2). This implies that the dominant coupling(s) be larger than $\sim 5 \times 10^{-7}$ in the small mass region and $\sim 1 \times 10^{-7}$ in the high mass region.

Table IX

The intervals of M_H that are excluded at 90% confidence and at 95% confidence for left-handed ($I_3^L = 1$) and right-handed ($I_3^L = 0$) doubly charged Higgs bosons. The excluded intervals are tabulated as a function of the τ branching ratio B_τ . They are valid in the region of coupling constant $g_{\tau\tau} \gtrsim 5 \times 10^{-7}$.

I_3^L	B_τ	90% Limit (GeV/c^2)	95% Limit (GeV/c^2)
0	1.0	$6.5 < M_H < 36.4$	$7.3 < M_H < 34.3$
0	0.5	$3.6 < M_H < 37.7$	$3.6 < M_H < 36.0$
0	0.0	$2.0 < M_H < 38.5$	$2.0 < M_H < 36.7$
1	1.0	$5.9 < M_H < 38.2$	$6.5 < M_H < 36.6$
1	0.5	$3.6 < M_H < 39.2$	$3.6 < M_H < 37.9$
1	0.0	$2.0 < M_H < 39.7$	$2.0 < M_H < 38.4$

The excluded regions overlap and significantly extend the existing small-coupling limits on M_H (which are independent of I_3^L).^[21] The 90% confidence limit for $B_\tau = 0$ is extended from approximately $21.5 \text{ GeV}/c^2$ to $39.7 \text{ GeV}/c^2$ ($38.5 \text{ GeV}/c^2$) for left-handed (right-handed) Higgs bosons. For $B_\tau = 1$, the 90% limit is extended from $14 \text{ GeV}/c^2$ to $38.2 \text{ GeV}/c^2$ ($36.4 \text{ GeV}/c^2$) for left-handed (right-handed) Higgs bosons.

In order to illustrate the dependence of the limits upon the mass and coupling constants, the least restrictive 90% limit ($B_\tau = 1$, $I_3^L = 0$) is plotted in $g_{\tau\tau} - M_H$ space in Figure 13 (the solid curve). Note that it extends to $g_{\tau\tau} = 7.2 \times 10^{-8}$. The limit is compared with two rather specific limits from Reference 21.

6. Polarization Physics at the SLC

6.1. THE POLARIZED SLC

One of the advantages of linear colliders is that they are relatively straightforward to polarize. A layout of the polarized SLC is shown in Figure 1. The orientation of an electron spin vector is shown as the electron is transported from the electron gun to the interaction point.

A gallium arsenide based photon emission source produces pulses of up to 10^{11} longitudinally polarized electrons at repetition rates of up to 120 Hz. The electrons are then accelerated in the first sector of the linac. The beam pulse achieves an energy of 1.21 GeV as it arrives at the entrance of the LTR (Linac To Ring) transfer line.

The electrons must be stored in the North Damping Ring for one machine cycle (the cycle time is $\simeq 8$ ms). A system consisting of the LTR bend magnets and a superconducting solenoid is used to rotate the spins into the vertical direction that is necessary for storage in the damping ring. After one machine cycle, the bunch is extracted and passed through another spin rotation system consisting of the bend magnets of the RTL (Ring-To-Linac) transfer line and two superconducting solenoids. The system is sufficiently flexible to provide essentially any spin orientation as the bunch reenters the linac at the beginning of sector 2.

The beam pulse is then accelerated to 46.5 GeV in the linac (due to synchrotron radiation losses, the energy at the end of the LINAC is ~ 1 GeV larger than it is at the interaction point). To insure that the spin gymnastics in the damping ring have worked properly and to study many of the potential sources of depolarization, a Møller polarimeter is located at the end of the linac near the PEP injection line. This polarimeter is used primarily for diagnostic purposes.

The beam pulse is then transported through the north machine arc and the final focus section to the interaction point. At full energy, the spin vectors precess roughly 26 times. Vertical precession also occurs in the nonplanar arcs. Since

longitudinal polarization is required at the interaction point, the total precession angle must be calculated for the exact machine energy and the polarization at the arc entrance must be adjusted appropriately.

After colliding with the unpolarized positron bunch, the electron beam is transported through the south final focus system where a Compton polarimeter is located. The beam continues to the south extraction line where a second Møller polarimeter is located. The bending magnets of the final focus and extraction line cause an additional spin precession of roughly 540° between the interaction point and the Møller target. Both polarimeters continuously monitor the beam polarization.

6.2. THE POLARIZED SOURCE

The SLC polarized electron source is based upon polarized photoemission from Gallium Arsenide (GaAs). The band structure of GaAs at the energy maximum of the valence band and energy minimum of the conduction band is shown in Figure 14. The band energy versus momentum is shown on the left-hand side and the energy level structure is shown on the right-hand side of the figure. The band gap of the material is $E_g = 1.52 \text{ eV}$. At the minimum of the conduction band and the maximum of the valence band, the electron wavefunctions have S and P symmetry, respectively. Spin-orbit splitting causes the $P_{3/2}$ states to reside in energy above the $P_{1/2}$ states by an amount $\Delta = 0.34 \text{ eV}$. The absorption of single photons proceeds via an electric dipole transition. The selection rules for the absorption of right- and left-handed circularly polarized photons are $\Delta m_j = +1$ and $\Delta m_j = -1$, respectively. They are indicated by the solid and dashed arrows in Figure 14. Since the electric dipole operator changes the orbital angular momentum of the initial state by one unit, the spin of the electron remains unchanged in the process.

Let's consider what happens when a right-circularly polarized photon is incident upon a GaAs crystal. The photon direction is the only vector in the system. All angular momentum projections refer to the incident photon direction. If

the photon energy E_γ is in the range $E_g \leq E_\gamma \leq E_g + \Delta$, then transitions can only occur from the $P_{3/2}$ states to the $S_{1/2}$ states. Specifically, the P state with $m_j = -3/2$ can make a transition to the S state with $m_j = -1/2$ and the P state with $m_j = -1/2$ can make a transition to the S state with $m_j = +1/2$. In the former case, the emitted electron has spin antiparallel to the incident photon direction (or parallel to its ejected direction). In the latter case, the spin of the emitted electron is parallel to the incident photon direction (antiparallel to its ejected direction). Due to Clebsch-Gordon coefficients (the P state with $m_j = -3/2$ is a pure spin state whereas state with $m_j = -1/2$ is not), the former transition is three times more likely than the latter. The relative transition rates are indicated by circled numbers in Figure 14. This implies that the absorption of a right circularly polarized photon produces a right-handed electron with a polarization

$$P = \frac{3-1}{3+1} = 50\%.$$

Actually, all that's been shown so far is that polarized electrons can be pumped into the conduction band with a beam of circularly polarized photons. In order to make a polarized source, the electrons must leave the material. In normal GaAs, the energy gap from the bottom of the conduction band to the free electron state is approximately 2.5 electron volts. Even with a large applied electric field, pure GaAs is a poor photoemitter. The magic that is necessary to make it an efficient photoemitter is shown in Figure 15. The energy of the various bands is shown as a function of depth near the surface for several materials: pure GaAs, GaAs with a cesiated surface, and GaAs with a surface layer of Cs_2O . The energy of the free electron state is shown as E_∞ . The addition of cesium to the surface causes the energy gap between the conduction band and the free electron state to decrease to zero. The addition of Cs_2O to the surface causes the gap to become negative! Quantum efficiencies as large as 5% have been observed for GaAs photocathodes that have been treated with Cs_2O (actually CsF is currently used instead).

Figure 16 shows the electron polarization that was measured^[23] for several photocathodes as a function of the photon wavelength. The photocathodes are

composed of pure gallium arsenide and several forms of gallium aluminum arsenide. The gallium aluminum arsenide is made by substituting aluminum atoms for fraction of the arsenic atoms. At the wavelength of the SLC polarized light source, 715 nm, the pure GaAs sample produces a polarization of only 35%. The electron polarization will be improved to 42% by the use of GaAl_{0.1}As_{0.9}.

6.3. THE SPIN ROTATION SYSTEM

The second major element of the polarized SLC is the spin rotation system. As was mentioned at the beginning of this chapter, the spin rotation system has two functions:

1. To rotate the (initially longitudinal) polarization vector of the electron bunch into the vertical direction for storage in the North Damping Ring.
2. To allow the orientation of the electron polarization vector to be controlled as the bunch reenters sector 2 of the linac. This is necessary to compensate for precession in the machine arcs.

A detailed representation of the north damping ring, the north LTR transfer line, and the north RTL transfer line is shown in Figure 2. The orientation of the polarization vector at various places is shown by the double arrow. The electron bunch arrives at the entrance to the LTR transfer line with an energy of 1.21 GeV. At this energy, the spins precess by 90° for each 32.8° that the electron trajectories are bent by a transverse magnetic field. The initial bend angle of the LTR has been chosen to be 5 × 32.8°. The longitudinal polarization of the beam emerging from sector 1 of the linac is therefore rotated into the horizontal direction. A superconducting solenoid of strength 6.34 T-m is introduced into the LTR optics after the first bend. The solenoid has only a small effect on the optics of the transport system but causes a rotation of the spin vector about the beam axis by 90°. The spins are therefore rotated into the vertical (either upward or downward) direction. After one machine cycle (≈ 8 ms), the electron bunch is extracted from the damping ring with a horizontal kicker magnet and passed

through a second superconducting solenoid magnet. The horizontal bend magnets of the RTL transfer line then deflect the beam by an angle of 3 × 32.8° before it reenters the linac at the beginning of sector 2. A third superconducting solenoid is introduced into the linac lattice just downstream of the reentry point. If the second (RTL) solenoid is adjusted to have the same strength as the first (LTR) solenoid has, the system will restore the longitudinal beam polarization. If it is not energized, the beam polarization will be vertical upon reentry into the linac. The third (linac) solenoid can then rotate the polarization vector to any transverse orientation. The combination of the two solenoids and the RTL bending magnets permits the selection of any orientation of the polarization vector.

6.4. POLARIMETRY AT SLC

The polarization of the SLC electron beam will be monitored by three polarimeters. Two of the polarimeters are based upon polarized electron-electron scattering (Møller scattering) and one is based upon polarized electron-photon scattering (Compton scattering).

The Møller Polarimeters

The Møller polarimeters measure the elastic scattering of the polarized electron beam from the polarized atomic electrons in a magnetized foil target. The cross section for this process has the form (in the limit of zero electron mass),

$$\frac{d\sigma}{d\Omega} = \frac{\alpha^2}{s} \frac{(3 + \cos^2\theta)^2}{\sin^4\theta} \left\{ 1 - \mathcal{P}_z^1 \mathcal{P}_z^2 A_z(\theta) - \mathcal{P}_t^1 \mathcal{P}_t^2 A_t(\theta) \cos(2\phi - \phi_1 - \phi_2) \right\} \quad (6.1)$$

where: s is the square of the total energy in the cm frame; θ is the cm frame scattering angle; ϕ is the azimuth of the scattered electron (the definition of $\phi = 0$ is arbitrary); $\mathcal{P}_z^1, \mathcal{P}_z^2$ are the longitudinal polarizations of the beam and target, respectively; $\mathcal{P}_t^1, \mathcal{P}_t^2$ are the transverse polarizations of the beam and target, respectively; ϕ_1, ϕ_2 are the azimuths of the transverse polarization vectors; and

$A_z(\theta)$ and $A_t(\theta)$ are the longitudinal and transverse asymmetry functions which are defined as

$$A_z(\theta) = \frac{(7 + \cos^2\theta)\sin^2\theta}{(3 + \cos^2\theta)^2}$$

$$A_t(\theta) = \frac{\sin^4\theta}{(3 + \cos^2\theta)^2}.$$

The differential cross section is the product of the unpolarized cross section and the sum of one and two polarization dependent terms. The first is the product of the longitudinal polarizations of the beam and target particles and the function $A_z(\theta)$. The second is the product of the transverse polarizations of the two electrons, an azimuthal factor, and the function $A_t(\theta)$. Both asymmetry functions are maximal for 90° scattering. The longitudinal asymmetry function becomes quite large ($A_z(90^\circ) = 7/9$) whereas the transverse asymmetry function never exceeds $1/9$. The analyzing power of any polarimeter scales as the product of the unpolarized cross section and the square of the asymmetry. This combination is also largest at $\theta = 90^\circ$ but has a rather broad maximum. Because the longitudinal and transverse asymmetries have maxima at the same scattering angle, it is quite straightforward to build three-axis polarimeters.

The polarimeters function by measuring the asymmetry in the scattering rate when either the beam or target polarizations are reversed.

$$A_{ee} \equiv \frac{\sigma(\mathcal{P}_1^1 \mathcal{P}_2^2) - \sigma(-\mathcal{P}_1^1 \mathcal{P}_2^2)}{\sigma(\mathcal{P}_1^1 \mathcal{P}_2^2) + \sigma(-\mathcal{P}_1^1 \mathcal{P}_2^2)} = -\mathcal{P}_2^1 \mathcal{P}_2^2 A_z(\theta) - \mathcal{P}_1^1 \mathcal{P}_1^2 A_t(\theta) \cos(2\phi - \phi_1 - \phi_2). \quad (6.2)$$

The beam polarization $\mathcal{P}_1^1$ is inferred from the measured asymmetry A_{ee} , the measured target polarization $\mathcal{P}_2^2$, and the theoretical asymmetry $A_x(\theta)$ (where x is z or t).

The precision of the Moller polarimeters is limited by the uncertainties on the target polarization and on the subtraction of background from the radiative process $e^- N \rightarrow e^- N' \gamma$ where N is a nucleus in the target material. We expect that a precision $\delta\mathcal{P}/\mathcal{P}$ of 5% is possible.

The Compton Polarimeter

The Compton polarimeter is the primary monitor of the polarization of the SLC electron beam. It measures the large asymmetry in the scattering of polarized electrons from circularly polarized laser photons. The light source is a frequency doubled Nd:YAG laser which produces 2.23 eV photons. The backscattered electrons are separated from the SLC beam by the first large dipole magnet of the SLC final focus system after it has passed through the interaction point.

The kinematical properties of the scattering of a high energy electron with an optical photon seem quite strange to those accustomed to working in reference frames that are nearer the center-of-mass frame. The energy of the electron is typically 10 orders of magnitude larger than that of the photon. It is clear that all final state particles are swept into the forward direction (along the incident electron direction). It is therefore convenient to define all angles with respect to the incident electron direction. The direction of the outgoing photon, θ_K , differs from the normal definition of the scattering angle by 180° (if the colliding $e^- \gamma$ are collinear). Let E , E' , K , and K' be the incident electron energy, scattered electron energy, incident photon energy, and scattered photon energy, respectively. The maximum energy of the scattered photon K'_{max} and the minimum energy of the scattered electron E'_{min} can then be written as

$$K'_{max} = E(1 - y)$$

$$E'_{min} = E y$$

where the parameter y is defined as

$$y \equiv \left(1 + \frac{4EK}{m^2}\right)^{-1}.$$

The emission angle of the scattered photon θ_K is related to the scattered photon energy by the following expression,

$$K' = K'_{max} \left[1 + y \left(\frac{E\theta_K}{m}\right)^2\right]^{-1}$$

$$= K'_{max} \cdot x$$

where the definition of x is obvious. The parameter x varies from unity at zero emission angle to zero at larger angles. The scale of the angular range is set by the angle for which the energy has been reduced by a factor of two. This occurs when $E\theta_K\sqrt{y}/m = 1$ or at the angle $\theta_K = m/E\sqrt{y}$. For the SLC Compton polarimeter operating with a 46 GeV beam, the value of the parameter y is 0.389. Therefore, the maximum photon energy is 28.1 GeV and the minimum electron energy is 17.9 GeV. The angle at which the photon energy has been decreased by a factor of two is 1.8×10^{-5} radians. The scattered electron and photon both remain along the beam direction.

The polarized cross section can be expressed in terms of the laboratory variables x , y , and the azimuth of the photon with respect to the electron transverse polarization ϕ as follows:^[23]

$$\left(\frac{d^2\sigma}{dx d\phi}\right)_{\text{Compton}} = \left(\frac{d^2\sigma}{dx d\phi}\right)_{\text{unpol}} \left\{ 1 - \mathcal{P}_z^e [\mathcal{P}_z^e A_z^{\text{e}\gamma}(x) + \mathcal{P}_t^e \cos \phi A_t^{\text{e}\gamma}(x)] \right\} \quad (6.3)$$

where: the unpolarized cross section is defined as

$$\left(\frac{d^2\sigma}{dx d\phi}\right)_{\text{unpol}} = r_0^2 y \left\{ \frac{x^2(1-y)^2}{1-x(1-y)} + 1 + \left[\frac{1-x(1+y)}{1-x(1-y)} \right]^2 \right\};$$

$\mathcal{P}_z$, $\mathcal{P}_t$ are the longitudinal and transverse polarizations of the electron; $\mathcal{P}_\gamma$ is the circular polarization of the photon; and where the longitudinal and transverse asymmetries are defined as

$$A_z^{\text{e}\gamma}(x) = r_0^2 y [1-x(1+y)] \left\{ 1 - \frac{1}{[1-x(1-y)]^2} \right\} \cdot \left(\frac{d^2\sigma}{dx d\phi}\right)_{\text{unpol}}^{-1}$$

$$A_t^{\text{e}\gamma}(x) = r_0^2 y x(1-y)^{1/2} \cdot \left(\frac{d^2\sigma}{dx d\phi}\right)_{\text{unpol}}^{-1}$$

For $y = 0.389$, the unpolarized cross section is very large (several hundred millibarns) and peaked at $x = 1$. The longitudinal asymmetry has a maximum of 75% also at $x = 1$. Note, however, that as x is decreased, $A_z^{\text{e}\gamma}$ decreases rapidly

and becomes negative near $x = 0.72$. It reaches a minimum of -25% near $x = 0.47$ and returns to zero at $x = 0$. The transverse asymmetry is zero at both endpoints and reaches a maximum of 33% near $x = 0.75$.

The acceptance of the SLC polarimeter integrates over the entire azimuth. The polarimeter therefore measures only the longitudinal asymmetry. Using equation (6.2), the longitudinal beam polarization is measured. It appears that the polarimeter is capable of measuring the polarization with a precision of 1% to 2%.

6.5. DEPOLARIZATION EFFECTS

There are numerous possible sources of electron beam depolarization. None of them are expected to be serious. The following is a summary of the most important.

Depolarization in the Linac

The depolarization of a longitudinally polarized electron beam by the SLAC linac has been calculated to be very small.^[24] This has been verified by several experiments.^[24,25] The polarized SLC does differ in two respects from the old SLAC linac:

1. The electron bunches are much smaller than they were for the fixed target experiments. It was pointed out by W.I.H. Panofsky^[26] that the intra-bunch fields could depolarize the bunch via an effect that is analogous to Thomas precession. More detailed calculations indicate that this effect causes less than a one percent depolarization of the beam.
2. The SLC must accelerate beams with transverse components of the polarization vector. This is not expected to be a problem, however, detailed calculations and experimental verification are still needed.

Depolarization in the Damping Ring

Since the helicity of each electron pulse is determined by a random number generator at the source, half of the electron pulses stored in the north damping

ring will have their polarization vectors aligned with the guide field and half will have anti-aligned polarization vectors. The natural polarization of a storage ring by the Sokolov-Ternov effect^[30] causes the spins to anti-align themselves with the guide field. This would cause the depolarization of the aligned bunches if they were stored in the ring for an appreciable fraction of a polarizing time. The polarizing time for the damping ring is approximately 15 minutes. Since the storage time of an electron bunch is only 8 milliseconds (at a 120 Hz repetition rate), this effect is negligible. Indeed, the short storage time (which is several damping times) implies that the only process that could cause a serious problem is resonant depolarization. The resonance condition is

$$\nu = N + I\nu_x + J\nu_y + K\nu_s$$

where: ν is the spin tune of the damping ring (the number of spin precessions per orbit); N , I , J , K are integers; ν_x and ν_y are the horizontal and vertical betatron tunes, respectively; and ν_s is the synchrotron tune of the damping ring. The SLC damping ring is designed to operate at an energy $E = 1.21$ GeV. The spin tune at this energy is given by the expression

$$\nu = \frac{g-2}{2} \cdot \frac{E}{m_e} = \frac{E}{440.65 \text{ MeV}} = 2.746$$

(where $(g-2)/2$ is anomalous magnetic moment of the electron). The horizontal and vertical betatron tunes are $\nu_x = 7.20$ and $\nu_y = 3.20$, respectively (the $\nu_x - \nu_y = 4$ coupling resonance is used to produce round beams). The synchrotron tune is very small ($\nu_s \simeq 0.04$). Therefore, the nearest spin depolarizing resonance occurs when $N, I, J = 6, 0, -1$ (the synchrotron tune is ignored since only relatively weak resonances are associated with it). The right hand side of the resonance equation is equal to 2.80 in this case. Since the natural width of this sideband resonance is expected to be less than 0.001, no serious resonant depolarization is expected.

Depolarization in the Arcs

The SLC arcs are fairly achromatic transport systems (they can transport a momentum interval $\Delta P/P = 5\%$). Since the total precession angle is a sensitive function of the beam energy, the finite energy spread of the beam ($\Delta P/P = 0.3\%$) causes a spread in the final spin directions of the electrons. The average longitudinal polarization at the interaction point is reduced by a factor 0.93.

Depolarization from Beam-Beam Interactions

Because the SLC beams are very small at the interaction point, each beam is subjected to very strong electromagnetic fields during the collision. These fields cause some depolarization of the electron bunch. The size of the effect is given by the expression

$$\Delta\theta_s = \frac{g-2}{2} \cdot \frac{E}{m_e} \cdot \theta_d$$

where: $\Delta\theta_s$ is the average precession angle of beam particles; E is the beam energy; and θ_d is the disruption angle of the beam. Since the disruption angle at SLC is roughly one milliradian, the average depolarization is less than one percent.

Systematic Effects

It is possible that the average beam polarization as measured by the two downstream polarimeters be different from the luminosity weighted average polarization. There are two possible causes for this effect.

1. The beam-beam interaction obviously changes the polarization before it arrives at the Compton and extraction line Møller polarimeters. The size of this effect is estimated to be less than 0.5%.
2. If the electron beam at the interaction point has a non-zero dispersion function, it is possible that a beam-beam targeting error could cause the luminosity weighted beam energy and polarization to differ from the average beam energy and polarization. The beam-beam deflection process allows the beam to be targeted to within a small fraction of the beam sizes. Therefore, even if

the dispersion function at the interaction point were as large as 3 mm (which is quite large), the fractional deviation of the monitored polarization from the average one is less than two percent. If the dispersion function is the more normal 1 mm, this effect is a few tenths of one percent.

6.6. THE LEFT-RIGHT POLARIZATION ASYMMETRY

In order to understand the utility of a polarized electron beam, we must consider the cross section for the (longitudinally) polarized process $e^+e^- \rightarrow f\bar{f}$. The beam polarizations, $\mathcal{P}^-$ and $\mathcal{P}^+$, are described in terms of a helicity basis ($\mathcal{P} = +1$ describes a right-handed beam, $\mathcal{P} = -1$ describes a left-handed beam). We can then write the tree-level cross section in the cm frame as follows,

$$\frac{d\sigma_f}{d\Omega} = \frac{\alpha^2 N_f^2}{64s \sin^4 2\theta_w} \cdot \left\{ (1 - \mathcal{P}^+ \mathcal{P}^-) [\sigma_u^Z + \sigma_u^Z] + (\mathcal{P}^+ - \mathcal{P}^-) [\sigma_p^Z + \sigma_p^Z] \right\} \quad (6.4)$$

where: the unpolarized partial cross sections due to γZ interference and pure Z exchange are defined as,

$$\begin{aligned} \sigma_u^{\gamma Z} &= -8Q_f \sin^2 2\theta_w \text{Re}[\Gamma(s)] \left[(1 + \cos^2 \theta^*) v v_f + 2 \cos \theta^* a a_f \right] \\ \sigma_u^Z &= |\Gamma(s)|^2 \left[(1 + \cos^2 \theta^*) (v^2 + a^2) (v_f^2 + a_f^2) + 8 \cos \theta^* v a_f a_f \right]; \end{aligned}$$

θ^* is the angle of the outgoing fermion relative to the incident electron; the polarized partial cross sections due to γZ interference and pure Z exchange are defined as.

$$\begin{aligned} \sigma_p^{\gamma Z} &= 8Q_f \sin^2 2\theta_w \text{Re}[\Gamma(s)] \left[(1 + \cos^2 \theta^*) a v_f + 2 \cos \theta^* v a_f \right] \\ \sigma_p^Z &= -|\Gamma(s)|^2 \left[(1 + \cos^2 \theta^*) 2v a (v_f^2 + a_f^2) + 2 \cos \theta^* (v^2 + a^2) 2v a_f a_f \right]; \end{aligned}$$

the constant N_f^2 is the color factor (3) for quark final states; $\Gamma(\hat{s}) = \hat{s}/(\hat{s} - M_Z^2 + i\Gamma_Z \hat{s}/M_Z)$ is the normalized Z propagator; and where the coupling constants without subscript, v and a , refer to the electron vector and axial-vector coupling constants. Note that we've assumed that the masses of all final state fermions are

small as compared with $\sqrt{s}$ and that the unpolarized cross section for pure photon exchange is small as compared with the pure Z and interference terms.

The forward-backward asymmetries are defined to select the part of the e^+e^- cross section that is odd under spatial reflection. The left-right polarization asymmetry is designed to select the part of the cross section that is odd in difference of the beam polarizations $\mathcal{P}^+ - \mathcal{P}^-$. It is therefore useful to define a generalized beam polarization $\mathcal{P}_g$ that is proportional to $\mathcal{P}^+ - \mathcal{P}^-$ and has a convenient normalization,

$$\mathcal{P}_g \equiv \frac{\mathcal{P}^+ - \mathcal{P}^-}{1 - \mathcal{P}^+ \mathcal{P}^-}. \quad (6.5)$$

Note that $\mathcal{P}_g$ is positive whenever the electron beam is left-handed and/or the positron beam is right-handed. It is negative whenever the reverse is true. The generalized polarization becomes unity when either beam is completely polarized. The positron beam of the SLC is unpolarized. The generalized polarization therefore has the simple form, $\mathcal{P}_g = -\mathcal{P}^-$.

The left-right polarization asymmetry is defined as the ratio of the difference of the total Z^0 production rates with left-handed and right-handed beams to the total rate. This can be expressed more precisely as,

$$A_{LR} \equiv \frac{\sum_f \left\{ \int_{-x_f}^{x_f} d\sigma_f(c, \mathcal{P}_g = +1) - \int_{-x_f}^{x_f} d\sigma_f(c, \mathcal{P}_g = -1) \right\}}{\sum_f \left\{ \int_{-x_f}^{x_f} d\sigma_f(c, \mathcal{P}_g = +1) + \int_{-x_f}^{x_f} d\sigma_f(c, \mathcal{P}_g = -1) \right\}}, \quad (6.6)$$

where: $c \equiv \cos \theta^*$; $\sigma_f(c, \mathcal{P}_g)$ is shorthand for the differential cross section $d\sigma_f/d\Omega$; $\pm x_f$ are integration limits that depend upon fermion type; and where the sum is taken over all visible final state fermions except electrons (to exclude the t-channel scattering process). Note that the integrals must be taken over symmetric limits (which is a natural property of most e^+e^- detectors).

Substituting equation (6.4) (actually, the version of equation (6.4) with finite final state masses) into equation (6.6) it is straightforward to show that the left-

right asymmetry takes the following form on the Z^0 pole,

$$A_{LR} = \frac{-2va \sum_f \int_{-z_f}^{z_f} d\epsilon[(v_f^2 + a_f^2)(1 + \beta_f^2 c^2) + (v_f^2 - a_f^2)(1 - \beta_f^2)]}{(v^2 + a^2) \sum_f \int_{-z_f}^{z_f} d\epsilon[(v_f^2 + a_f^2)(1 + \beta_f^2 c^2) + (v_f^2 - a_f^2)(1 - \beta_f^2)]},$$

where β_f is the velocity of the final state fermion in the $f\bar{f}$ center-of-mass frame. Cancelling the common factor, we recover a familiar expression,

$$A_{LR} = \frac{-2va}{v^2 + a^2} = \frac{2(1 - 4\sin^2\theta_w)}{1 + (1 - 4\sin^2\theta_w)^2}. \quad (6.7)$$

A number of conclusions can be drawn from this derivation:

1. A_{LR} depends upon the Z^0 -electron couplings alone. The dependence on the final state couplings cancels in the ratio.
2. A_{LR} is independent of the detector acceptance. This remains true even if each final state fermion is accepted differently.

3. A_{LR} is independent of final state mass effects (which would cause β_f to differ from unity).

4. All of the visible final states except the electron pairs can be used to measure A_{LR} . The measurement therefore utilizes about 96% of the visible decays. The various other Standard Model tests that are performed on the Z^0 pole make use of much smaller fractions of the event total ($\sim 4\%$ for the muonic forward-backward asymmetry, $\sim 0.9\%$ for the τ polarization measurement, and $\sim 4\%$ for the b-quark forward-backward asymmetry).

5. A_{LR} is very sensitive to the electroweak mixing parameter $\sin^2\theta_w$. This is shown graphically in Figure 17. Small changes in A_{LR} are related to changes in $\sin^2\theta_w$ by the following expression,

$$\delta A_{LR} \simeq -8\delta\sin^2\theta_w. \quad (6.8)$$

For $M_Z = 91.17$ GeV, the asymmetry is expected to be in the range 13%-15%.

Radiative Corrections

The left-right asymmetry has the property that it is insensitive to a large class of relatively uninteresting real and virtual radiative corrections and is very sensitive to an interesting set of virtual electroweak corrections. This behavior can be summarized as follows:

1. The left-right asymmetry is very insensitive to initial state radiative corrections. The emission of real photons by the incident electron and positron causes a smearing of the center-of-mass energy of the $f\bar{f}$ system ($\sqrt{s}$). The left-right asymmetry is quite insensitive to small changes in $\sqrt{s}$. The energy dependence of A_{LR} is compared with those of several forward backward asymmetries in Figure 18. The size of the initial state radiative correction to A_{LR} is calculated to be^[32] $\delta A_{LR} \simeq 0.002$ (this is a 2% correction to the asymmetry). The uncertainty on the correction to A_{LR} is smaller by an order of magnitude.
2. The QCD corrections to the left-right asymmetry vanish entirely to all orders in the strong coupling constant α_s at the leading order in the electromagnetic coupling constant α . The leading QCD corrections to A_{LR} are the (extremely small) corrections to the weak vector boson box diagrams.
3. The theoretical uncertainty on A_{LR} is completely dominated by the uncertainty on the renormalization of the electromagnetic coupling constant to the Z^0 mass scale. The current value of this uncertainty is^[33] $\delta A_{LR} \simeq 0.002$.
4. The left-right asymmetry is quite sensitive to virtual electroweak corrections and to the presence of new particles. The sensitivity of the asymmetry to the top quark mass (m_{top}) and the Higgs boson mass (m_{Higgs}) will be discussed in the last section of this document.

Experimental Errors

At the SLC, the measurement of A_{LR} is performed by randomly flipping the sign of the beam polarization on a pulse-to-pulse basis and by counting the number

of Z^0 events that are produced from each state. The measured asymmetry, A_{LR}^{exp} , is related to the theoretical asymmetry, A_{LR} , by the following expression,

$$A_{LR}^{exp} \equiv \frac{N_Z(\mathcal{P}_g = +\mathcal{P}_0) - N_Z(\mathcal{P}_g = -\mathcal{P}_0)}{N_Z(\mathcal{P}_g = +\mathcal{P}_0) + N_Z(\mathcal{P}_g = -\mathcal{P}_0)} = \mathcal{P}_0 A_{LR}. \quad (6.9)$$

where $\mathcal{P}_0$ is the magnitude of the beam polarization ($\mathcal{P}_0 \sim 0.40$), and $N_Z(\mathcal{P})$ is the number of Z^0 events logged with beam polarization $\mathcal{P}$. Since the left-handed and right-handed Z^0 cross sections are measured simultaneously, any systematic effects due to variations in detector livetime, luminosity, beam energy, beam position, etc., are cancelled in the ratio of the cross sections. This technique was used successfully to measure a very small polarized asymmetry ($\sim 10^{-5}$) in electron-deuteron scattering in 1978.^[24] The dominant systematic error is expected to be the uncertainty on the beam polarization measurement. We expect that the SLC Compton polarimeter is capable of measuring the beam polarization with a precision of 1-2% ($\delta\mathcal{P}_0/\mathcal{P}_0 = 1-2\%$).

There are a number of consistency checks that can be made with the SLC polarization hardware. It is possible to reverse the circular polarization optics of the electron source laser to search for systematic problems in that system. The polarity of the spin rotation system can be reversed to check for systematic problems in the damping rings. The polarization direction of each polarimeter target is reversible. The beam polarization can be measured separately with each target polarization direction (and must be consistent). Finally, the left-right asymmetry for small-angle Bhabha scattering is *very small* ($\sim 10^{-4}$). The luminosity monitors therefore provide an important check that the left-handed and right-handed luminosities are equal (the left-right asymmetry of the Bhabha signal must be consistent with zero).

Assuming that the dominant systematic error is the beam polarization uncertainty, the combined statistical and systematic uncertainty on A_{LR} is given by the following expression.

$$\delta A_{LR} = \left[A_{LR}^2 \left(\frac{\delta\mathcal{P}_0}{\mathcal{P}_0} \right)^2 + \frac{1 - (\mathcal{P}_0 A_{LR})^2}{\mathcal{P}_0^2 N_{tot}} \right]^{1/2}, \quad (6.10)$$

where N_{tot} is the total number of Z^0 events. The expected precision of the A_{LR} measurement and the corresponding precision on $\sin^2\theta_w$ are listed in Table X for several values of N_{tot} . Note that the statistical uncertainty dominates the total error in the region $N_{tot} \leq 10^6$. At $N_{tot} = 3 \times 10^6$, the statistical and systematic components are comparable.

Table X

The expected error on A_{LR} and $\sin^2\theta_w$ as a function of the number of Z^0 events. The left-right asymmetry is assumed to be $A_{LR} = 0.135$ (which is in the middle of the range that is expected for $M_Z = 91.17$ GeV). The beam polarization is assumed to be $\mathcal{P}_0 = 0.40$ and the precision of the polarization monitoring is assumed to be $\delta\mathcal{P}_0/\mathcal{P}_0 = 0.01$.

N_{tot}	δA_{LR}	$\delta \sin^2\theta_w$
100K	0.008	0.0010
300K	0.005	0.0006
1M	0.003	0.00035
3M	0.002	0.00025

The experimental confidence intervals that are presented in Table X are compared with the theoretical expectation for A_{LR} in Figures 19 and 20. The solid curves in Figure 19 enclose the 68.3% confidence region that is expected for $m_{Higgs} = 500$ GeV and m_{top} varying between 60 GeV and 240 GeV. The finite width of the region is due to a ± 20 MeV uncertainty on the Z^0 mass (we assume $M_Z = 91.17 \pm 0.02$ GeV). The solid curves in Figure 20 enclose the 68.3% confidence region that is expected for $m_{top} = 150$ GeV and m_{Higgs} varying from 100 GeV to 900 GeV. The size of the theoretical error on A_{LR} (± 0.002) is shown as the dotted vertical error bar in each figure. The sizes of the experimental 68.3% confidence intervals that correspond to the various values of N_{tot} are indicated by the solid vertical error bars. Since the τ polarization asymmetry is formally equivalent to A_{LR} , we plot the confidence region that is expected from a measurement with a

6M Z^0 sample. It is clear that A_{LR} is quite sensitive to m_{top} . A measurement with 300K Z^0 events constrains the top quark mass to a region of roughly $\delta m_{top} = \pm 17$ GeV which is comparable to a 100 MeV determination of M_W . The sensitivity to m_{Higgs} is clearly much smaller. A very high statistics measurement of A_{LR} could provide, at best, an indication of m_{Higgs} .

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FIGURE CAPTIONS

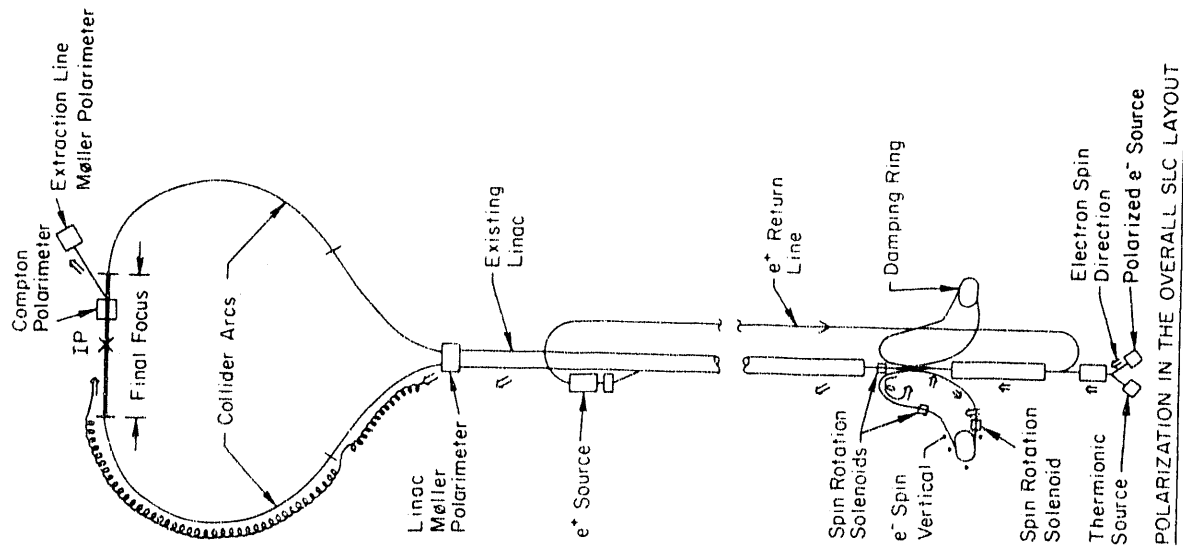
- 1) A layout of the SLAC Linear Collider. The orientation of an electron spin vector is shown as the electron is transported from the electron gun to the interaction point.
- 2) The spin rotation system as incorporated into the north damping ring complex. The orientation of the polarization vector at several points is shown by the double arrow.
- 3) The Mark II detector.
- 4) The sensitivity function for M_Z as a function of center-of-mass energy about the Z pole, $E - M_Z$.
- 5) The sensitivity function for Γ_Z as a function of center-of-mass energy about the Z pole, $E - M_Z$.
- 6) The sensitivity function for $\sigma_{had}^0(M_Z^2)$ as a function of center-of-mass energy about the Z pole, $E - M_Z$.
- 7) The Z^0 lineshape as measured by the Mark II Collaboration.^[4] The dashed curve is the result of a single parameter fit (for M_Z). The results of two and three parameter fits are indistinguishable and are shown as the solid curve.
- 8) The north (electron) energy spectrometer of the SLC. The beam is focused by a quadrupole doublet to a point at the detector plane. The beam passes through a small horizontal bend dipole magnet, a large vertical bend dipole magnet, and a second small horizontal bend magnet. The passage of the beam through the horizontal bend magnets produces flat distributions of synchrotron radiation which are detected by a phosphor screen detector. The separation of the flat distributions is proportional to the beam energy.
- 9) The cross section for the process $e^+e^- \rightarrow W^+W^-$ as a function of $E_b - M_W$. The mass and width of the W are assumed to be 80 GeV and 2.1 GeV, respectively. Note that three curves are plotted: the dashed curve is the

- basic tree-level cross section; the dashed-dotted curve is the cross section including the effect of initial state radiation; and the solid curve is the cross section including initial state radiation and the effect of a finite W width.
- 10) The sensitivity function for M_W as a function of the single beam energy about the W pair threshold $E_b - M_W$.
 - 11) The sensitivity function for Γ_W as a function of the single beam energy about the W pair threshold $E_b - M_W$.
 - 12) The sensitivity function for the background parameter B as a function of the single beam energy about the W pair threshold $E_b - M_W$.
 - 13) The 90% confidence contours of M_H versus the leptonic coupling strength $g_{\ell\ell}$ that are obtained from several processes. The excluded regions are indicated by the shaded side of each contour. The result of this search is shown as the solid contour (the limit is independent of lepton flavor). The limit^[21] that is obtained from the limit on muonium to antimuonium conversion is shown as a dotted line ($\sqrt{g_{ee}g_{\mu\mu}}$ is plotted along the horizontal axis). The limit^[22] that is obtained from the Bhabha scattering data of several PEP and PETRA experiments is shown as a dashed curve (g_{ee} is plotted along the horizontal axis). For reference, the sizes of the coupling constants g , g' , and e are indicated. The strong coupling limit occurs at the value $\sqrt{4\pi}$.
 - 14) The band structure of GaAs near the bandgap minimum.^[23] The energy levels of the states are shown on the right. Allowed transitions for the absorption of right (left) circularly polarized photons are shown as solid (dashed) arrows. The circled numbers indicate the relative transition rates.
 - 15) The band structure of Gallium Arsenide near its surface^[23] for: (a) pure GaAs, (b) GaAs with a cesiated surface, and (c) GaAs with a layer of Cs_2O on its surface.
 - 16) The polarization of electrons emitted from several GaAs photocathodes as functions of photon wavelength.^[23] The cathodes consisted of pure gallium

arsenide and several compositions of gallium aluminum arsenide.

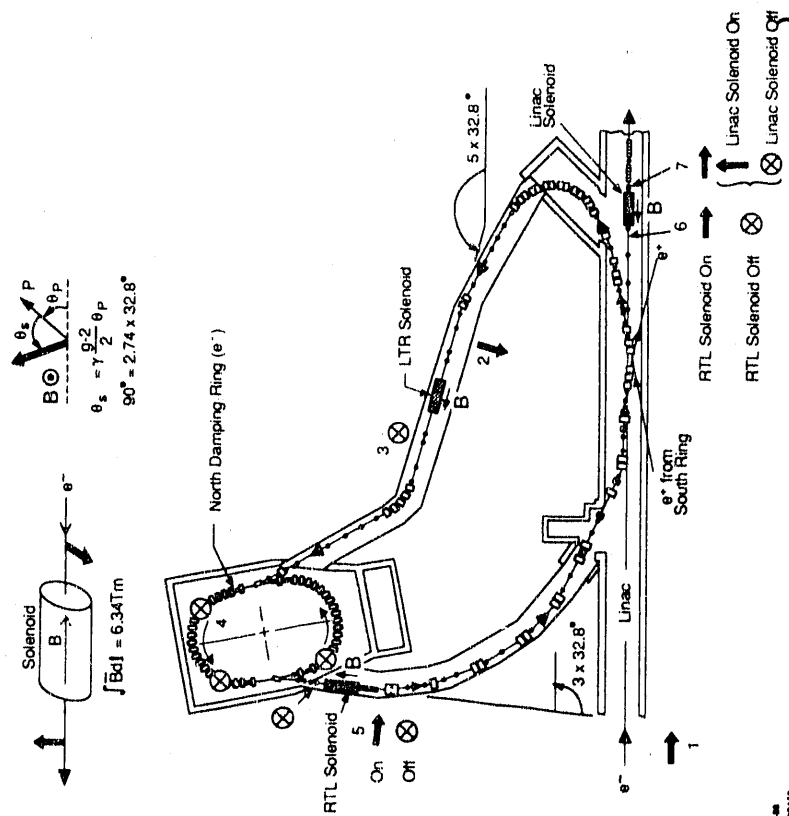
- 17) The left-right asymmetry A_{LR} is plotted as a function of $\sin^2\theta_w$ and M_Z (for some choice of m_t and m_h). The leptonic forward-backward asymmetry A_{FB}^{ℓ} is shown for comparison.
- 18) The forward-backward asymmetries for leptons, u -quarks, and d -quarks are plotted as a functions of the center-of-mass energy about the Z^0 pole. The asymmetries are also shown for the case that the incident beams are polarized. The energy dependence of the left-right asymmetry and an improved polarized forward-backward asymmetry $\bar{A}_{FB}^{\ell}$ (from Reference 31) are also shown. The Z^0 mass is assumed to be 94 GeV. Note that the unimproved forward-backward asymmetries are much more sensitive to the center-of-mass energy than are the improved ones and the left-right asymmetries.
- 19) The left-right asymmetry as a function of the top quark mass (m_{top}). The Higgs boson mass (m_{Higgs}) is assumed to be 500 GeV. The solid curves enclose the 68.3% confidence region that is expected for a ± 20 MeV uncertainty on M_Z (we assume $M_Z = 91.17 \pm 0.02$ GeV) as m_{top} is varied from 60 GeV to 240 GeV. The dotted vertical error bar shows the size of the theoretical error (± 0.002) on A_{LR} . The sizes of the experimental 68.3% confidence intervals that are expected for the various values of N_{tot} are indicated by the solid vertical error bars. The confidence interval that is expected from a measurement of the τ polarization asymmetry with 6M Z^0 events is also shown.
- 20) The left-right asymmetry as a function of the Higgs boson mass. The top quark mass is assumed to be 150 GeV. The dashed curves enclose the 68.3% confidence region that is expected for a ± 20 MeV uncertainty on M_Z (we assume $M_Z = 91.17 \pm 0.02$ GeV) as m_{Higgs} is varied from 100 GeV to 900 GeV. The dotted vertical error bar shows the size of the theoretical error (± 0.002) on A_{LR} . The sizes of the experimental 68.3% confidence intervals that are expected for the various values of N_{tot} are indicated by

the solid vertical error bars. The confidence interval that is expected from a measurement of the τ polarization asymmetry with 6M Z^0 events is also shown.

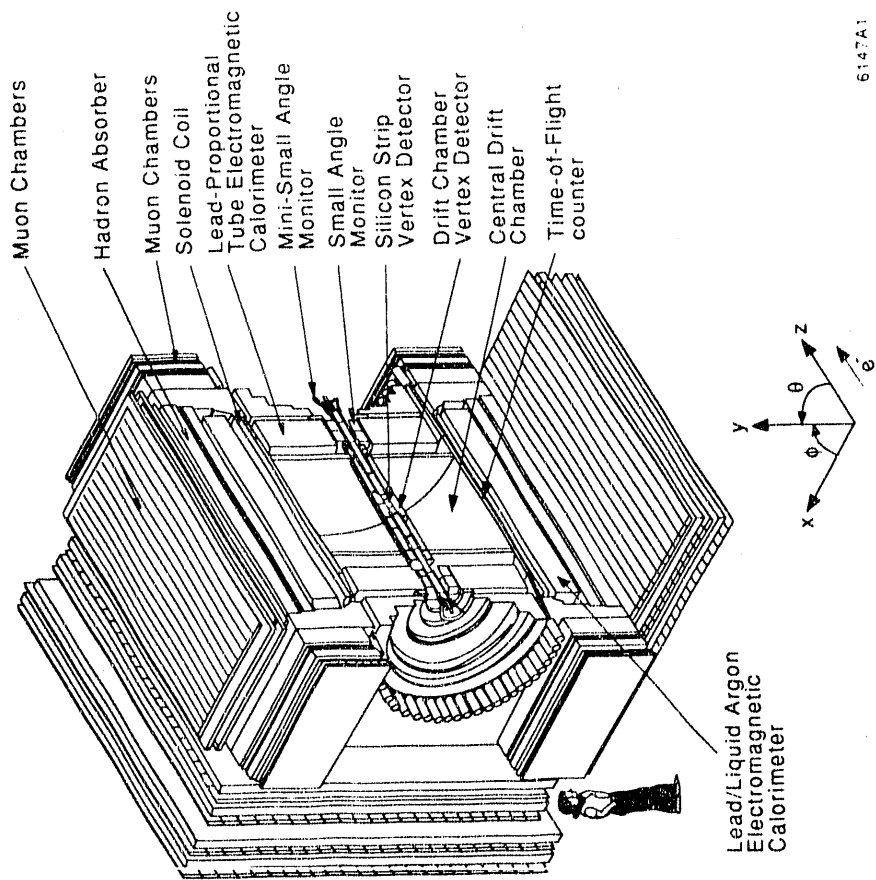


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MARK II AT SLC



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Fig. 3

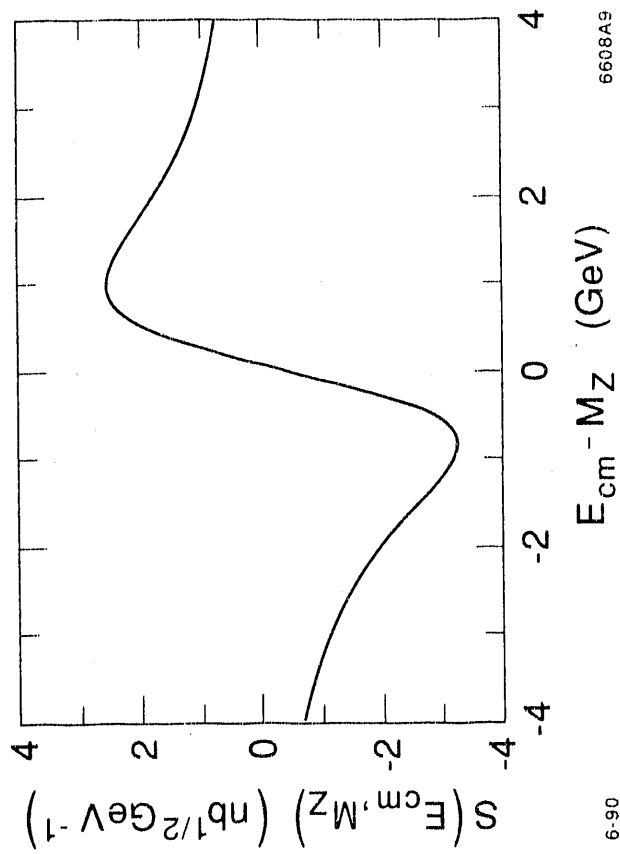


Fig. 4

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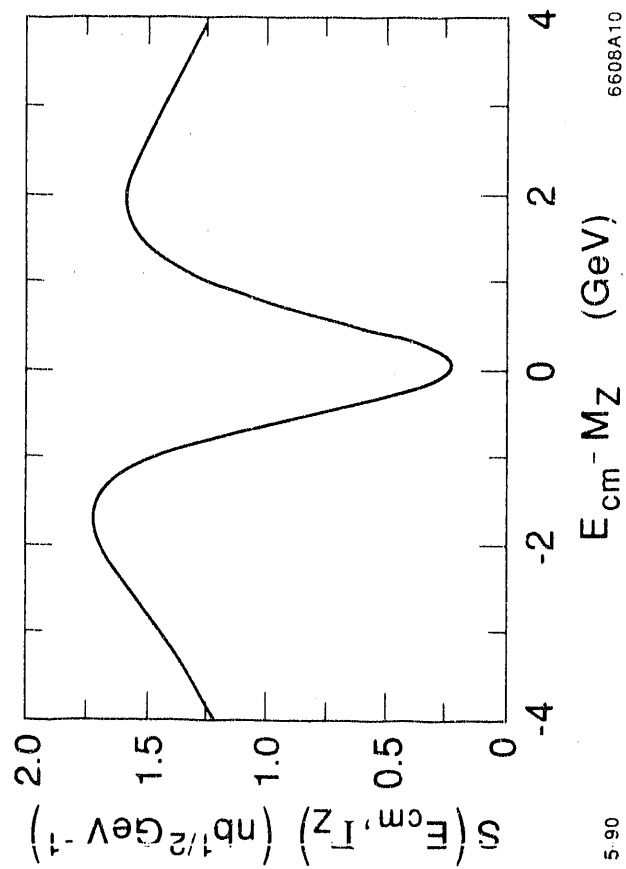


Fig. 5

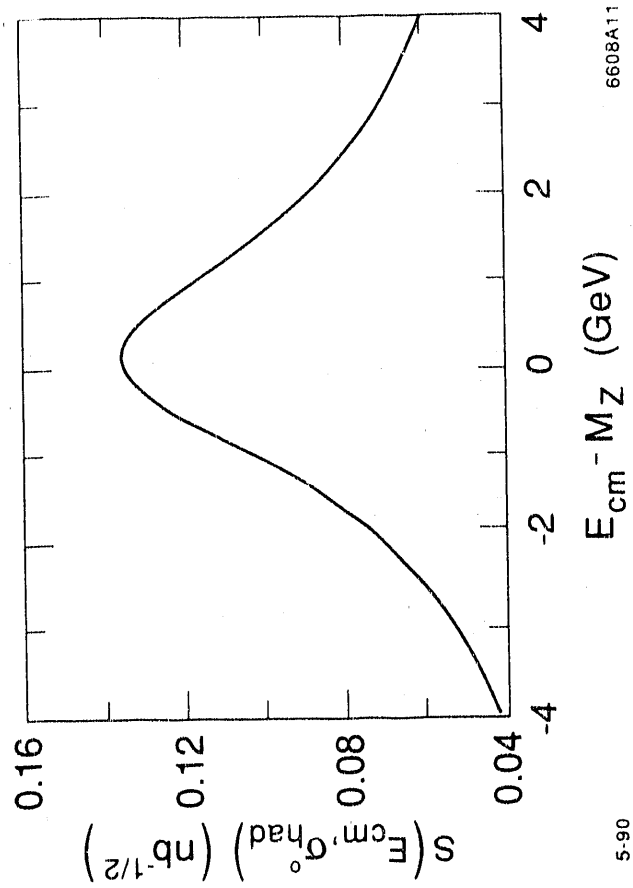


Fig. 6

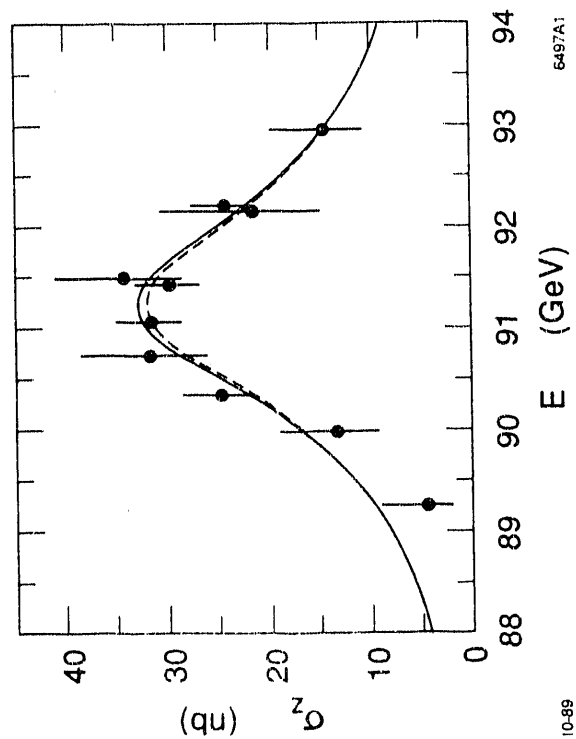


Fig. 7

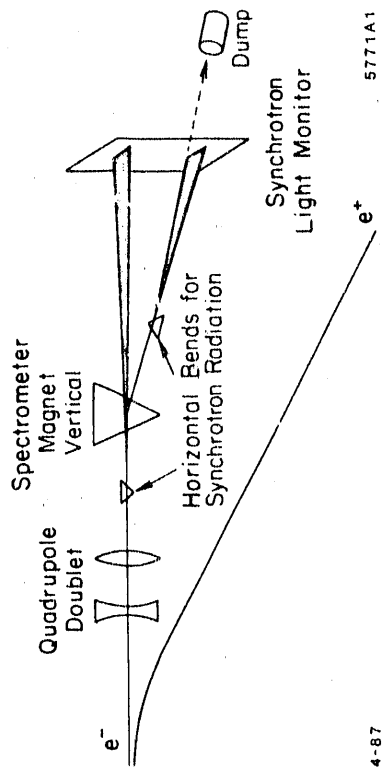


Fig. 8

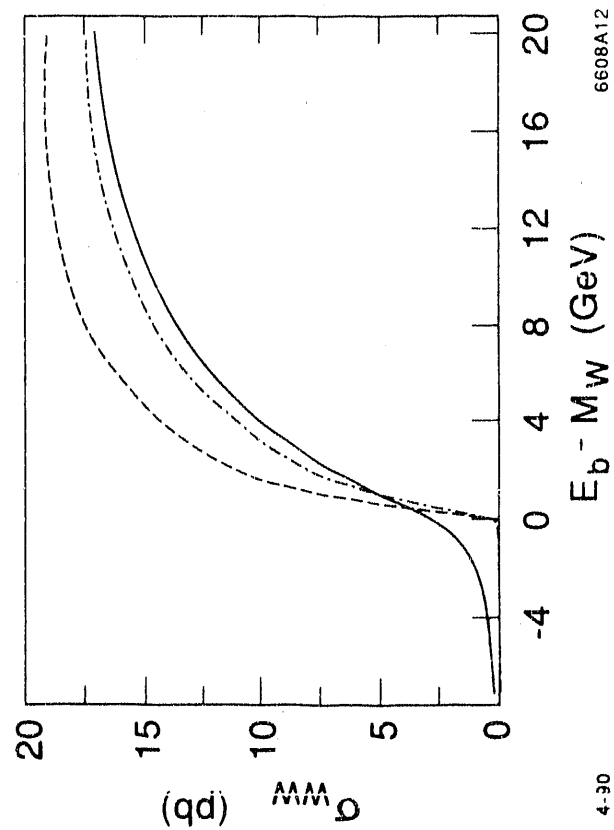


Fig. 9

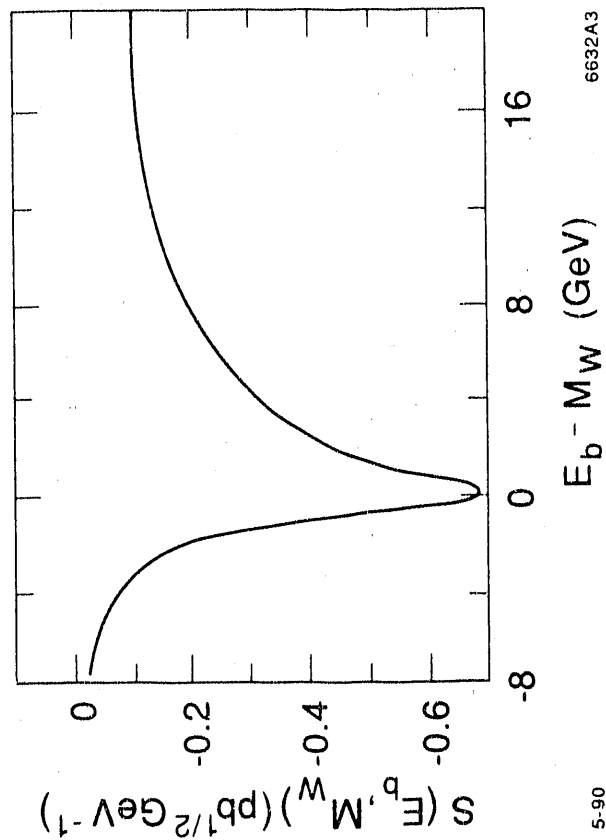


Fig. 10

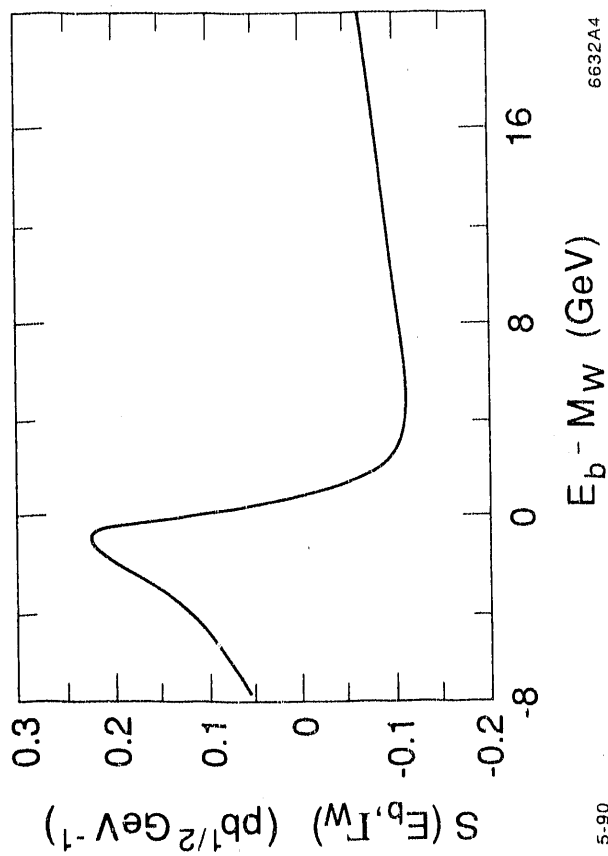


Fig. 11

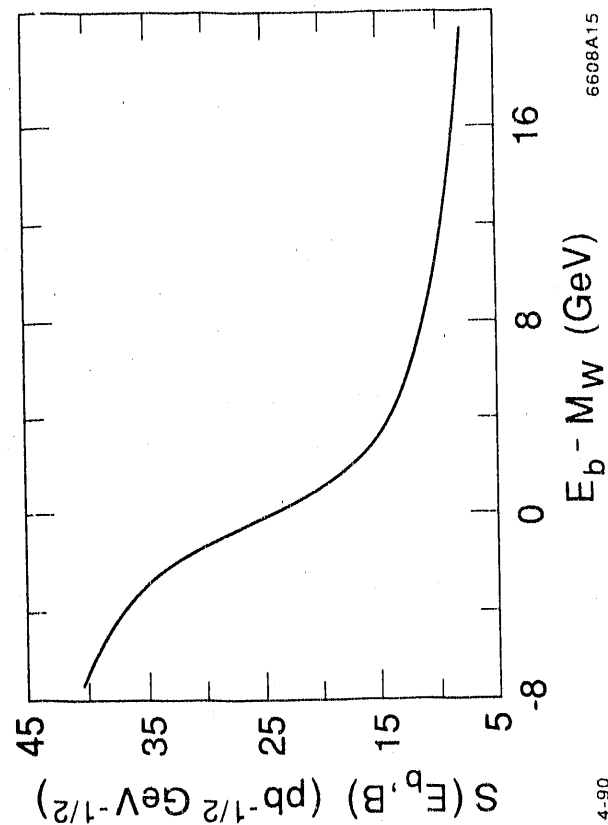
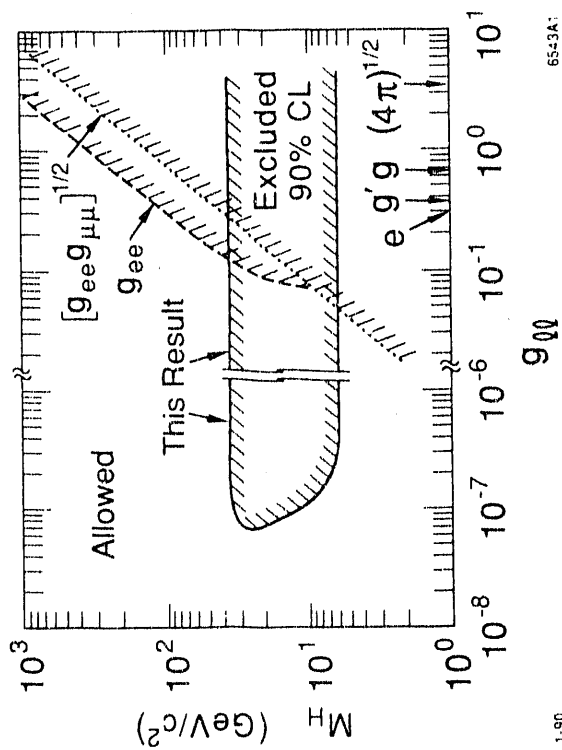


Fig. 12



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Fig. 13

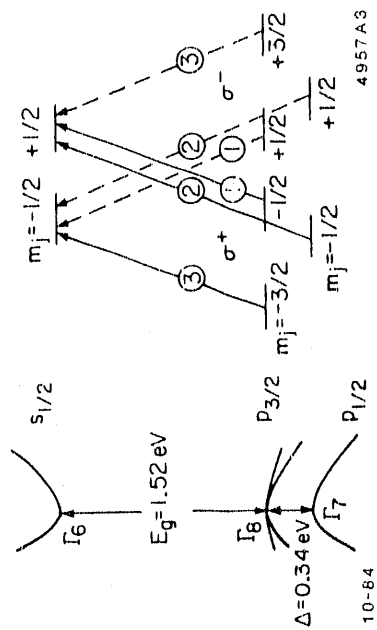


Fig. 14

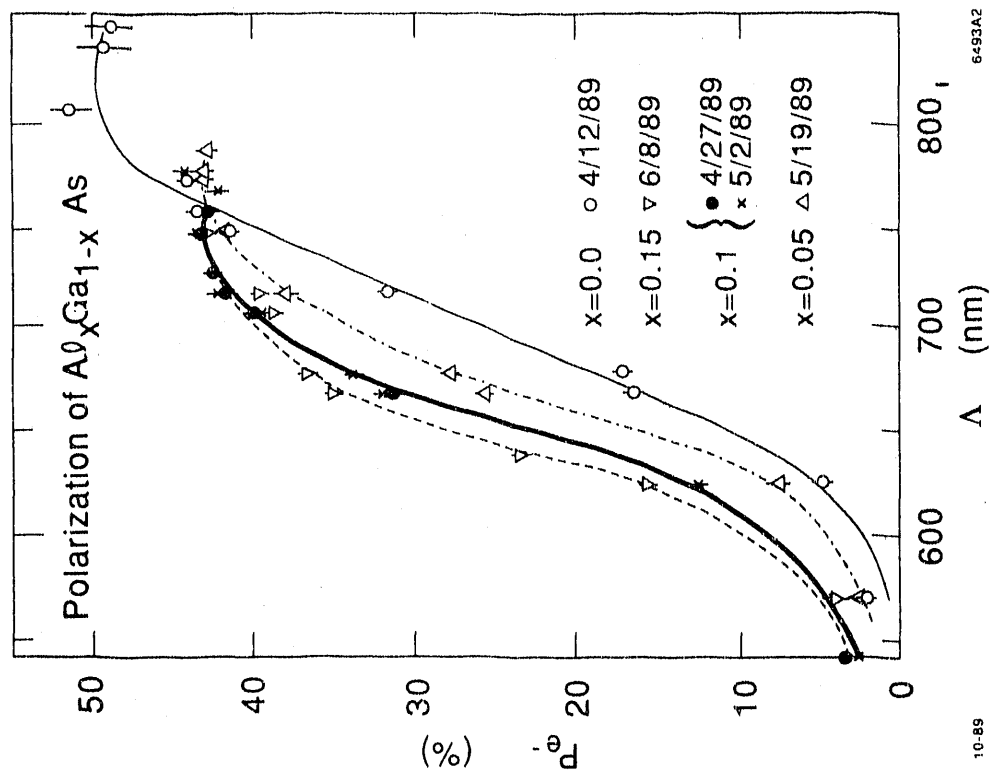


Fig 16

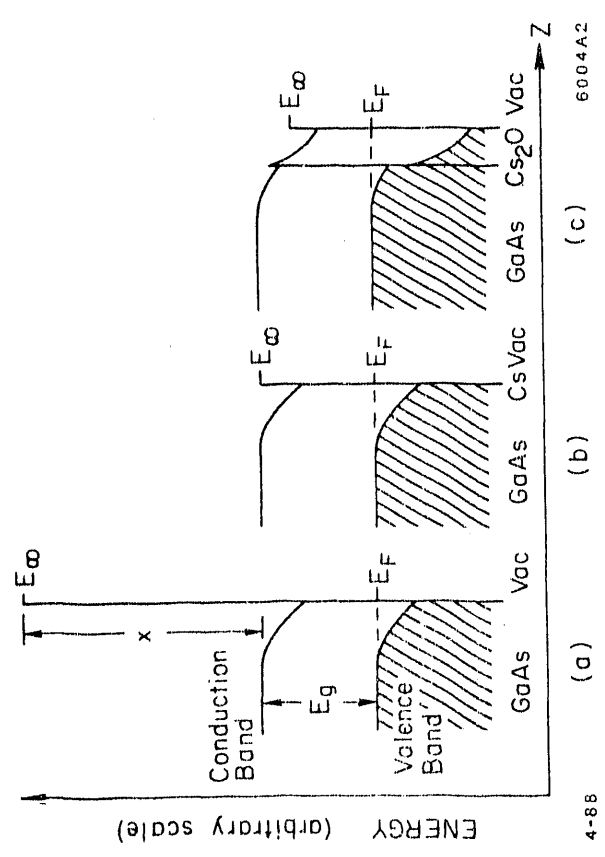


Fig. 15

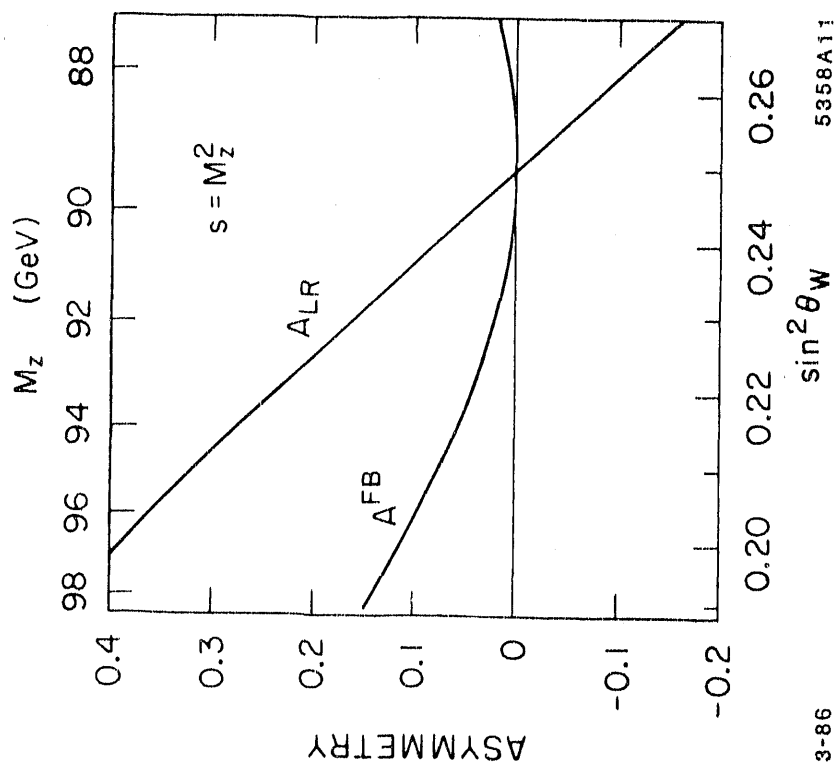


Fig. 17

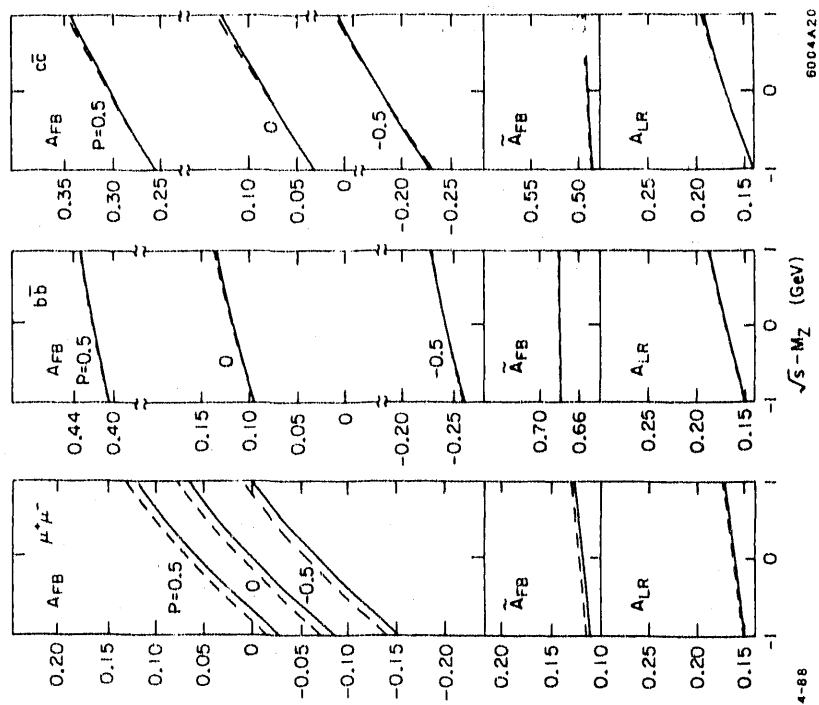


Fig. 18

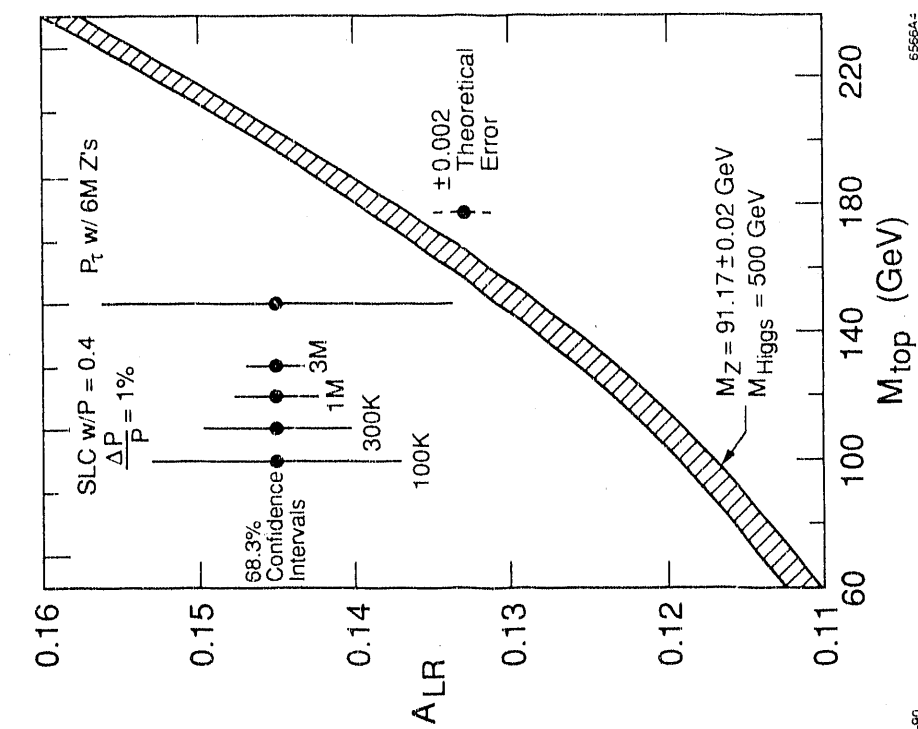


Fig. 19

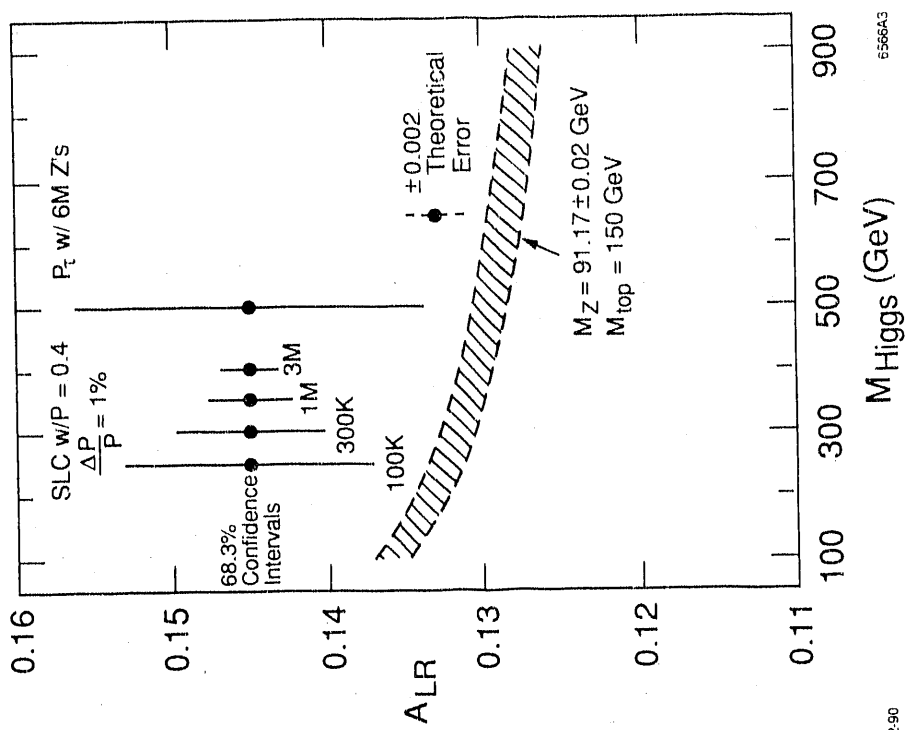


Fig. 20

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