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SU_3-Symmetry and the Existence of a Ninth Vectormeson*

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ABSTRACT

Baryon-baryon and baryon-antibaryon interactions are discussed within the framework of the SU_3-symmetric octet model without assuming R-invariance. Although this model, at least in the present version, contains a free parameter in the specification of the baryon-vectormeson coupling, it turns out that the predictions of the model are not consistent with the experimental determinations of the baryon-vectormeson coupling constants, if only exchanges of members of the vectormeson octet are considered. As the simplest way of removing such discrepancy the presence is proposed of a ninth vectormeson, a unitary singlet, for which some experimental support is available.

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Introduction

The aim of the present paper is to investigate the extent to which the values of the baryon-vectormeson coupling constants, as calculated from the octet model of the SU_3-symmetry scheme,[1] can be reconciled with the known empirical facts regarding baryon-baryon and baryon-antibaryon interactions. In the octet model the eight vector mesons, as well as the eight pseudoscalar mesons, may be thought of as bound states of a baryon-antibaryon pair or, strictly speaking, as superpositions of such states. The probability for a virtual decay of some meson into a particular baryon-antibaryon pair has the form of a product of a coupling constant, characteristic of the particular members of the baryon and meson octet in question, times a factor describing the general dependence on spin and energy. The coupling constant may again be factorized in a universal baryon-vector-meson (or baryon-pseudoscalar meson) coupling constant f and a coefficient G expressing the dependence on the generalized isobaric spin.

It is well known[1] that a certain ambiguity exists in the choice of coupling between three fields each of which belongs to the eight-dimensional representation of SU_3. In fact, the circumstance that the decomposition of the Kronecker product 8x8 contains the 8-dimensional representation itself twice allows of the construction of two orthogonal eigenstates of the Hamiltonian corresponding to the same spatial quantum numbers and the same values of the isospin T, T_3 and the hypercharge Y, the degeneracy being a consequence of the invariance of the Hamiltonian with respect to a permitation of the particles involved. Thus, from two single particle states $|a\rangle$ and $|b\rangle$, each of which is a carrier of the eight-dimensional representation, it is possible to form either an antisymmetrical or a symmetrical two-particle state $|c\rangle$:

$$|c\rangle_A = F_{ab}^{\ c}\ |a\rangle|b\rangle$$
$$|c\rangle_S = D_{ab}^{\ c}\ |a\rangle|b\rangle \qquad (1)$$

Each index a, b and c runs from 1 to 8 and the quantities F_{ab}^c are totally antisymmetrical in all the indices whereas the coefficients $D_{ab}^{\ c}$ are completely symmetrical.

Although in certain cases, where at least two of the particles involved are identical and thus Bose- or Fermi-statistics may be invoked, it is possible to make an unambiguous choice between the F and the D coupling, such uniqueness in the specification of the coupling requires in general the additional assumption of invariance of the interaction under the R-reflection operation.[1,2] More recent investigations seem, however, to raise rather large difficulties for the hypothesis of R-invariant interactions. In particular it has been pointed out by Cutkosky[3] that with the assumption of R-invariance the prediction of the existence of another baryon octet appears to be unavoidable. In the general case, when no assumption of R-invariance is made, an additional parameter, describing the "F-D-mixing ratio", enters for the characterization of the octet couplings. Indeed, the $B\bar{B}V$ vertex (figure 1)

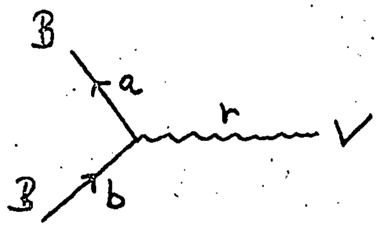


Figure 1. $B\bar{B}V$-vertex

has the following dependence on the unitary spin coupling matrices:

$$f \cdot G_{ab}^{r} = (a F_{ab}^{r} + b D_{ab}^{r}) f. \qquad (3)$$

Here B and V stands for baryon and vectormeson, respectively, whereas f denotes the universal BV-coupling constant. The F-D mixing ratio needs of course not be the same for the baryon-pseudoscalar coupling. The normalization of the matrices is fixed by the following introduction of the "mixing angle" Ψ :

$$a = \frac{1}{\sqrt{18}} \sin \Psi \qquad (4)$$
$$b = \frac{1}{\sqrt{10}} \cos \Psi .$$

1. The Existence of a Hard Core in the Baryon-Baryon Potential.

The low binding energy of the deuteron and the change of sign of the S-wave phase shifts in nucleon-nucleon scattering is generally traced to the presence of a repulsive core in the nucleon-nucleon potential. The experimental values available for the binding energies of the hyperfragments indicates that the idea of a repulsive core may apply to all baryon-baryon interactions.[4,5] Indeed, a rough comparison with the ground state of the hydrogen atom, taking appropriate values for the ratios between the coupling constants and masses involved, leads us (for the hyperfragments) to expect binding energies of hundreds of MeV. In contrast, the $_{\Lambda}H^2$ system is presumably not bound at all and the binding energy for the $_{\Lambda}H^3$ fragment is around 0.2 MeV. Similarly although the absence of Σ-fragments may be explained as a result of a conversion of the Σ^+ into a Λ in presence of neutrons, it is hard to understand why only a single Ξ^- fragment has been seen if these systems were reasonably strongly bounded (Wilkinson[6] report the finding of a ΞHe^8 with a binding energy about 6 MeV).

It was early suggested that the presence of a repulsive short range potential could be explained as a result of the exchange of a vector meson of appropriate mass between the two interacting baryons. In fact, the simple graph shown in figure 2

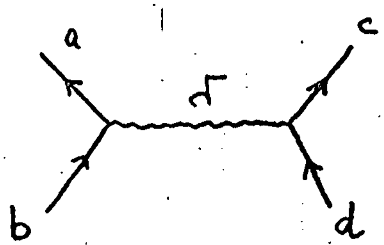


Figure 2

gives rise to a potential of the type[7)]

$$V = f^2 \cdot \frac{1}{r} e^{-Mr} \quad + \text{ spin dependent terms,} \qquad (5)$$

where M is the mass of the exchanged vector meson.

In case the exchanged vector meson is coupled to the isospin and hypercharge current the sign of the central potential will, of course, in general depend on the value of the isospin and hypercharge for the two-baryon state in question. In particular, within the framework of SU_3-symmetry the signs of the potential in the various unitarity spin states $|\lambda\mu\rangle$, corresponding to exchange of a member of the vector meson octet, can be evaluated. In the symbols of equation (3) the generalized isobaric spin dependence of the diagram in figure 1 is given by

$$\begin{aligned} \mathcal{M}_{ac,bd} &= G_{ab}^{r} \cdot G_{cd}^{r} \\ &= a^2 M_{ac,bd} + b^2 \bar{M}_{ac,bd} + 2ab \left[W_{ac,bd} \right]_{10} \end{aligned} \qquad (6)$$

where M, $\bar{M}$ and W are the combinations of matrices employed in reference 8. They are given by

$$M = F^r \cdot F^r, \quad \bar{M} = D^r \cdot D^r, \quad W = F^r \cdot D^r. \tag{7}$$

The square bracket with subscript 10 around W in (6) indicates that this term only contributes in the states (30) and (03). The eigenvalues of the operators M, $\bar{M}$ and $\bar{W}$ in the various states $(\lambda\mu)$ are given in reference 9. The corresponding eigenvalues of $\mathcal{M}$ are listed in Table 1.

d	State $(\lambda\mu)$	$\mathcal{M}$
1	$(00)_S$	$-12\, a^2 + \frac{20}{3}\, b^2 = \frac{2}{3} \cos^2 \psi$
8	$(11)_A$	$-6\, a^2 + \frac{10}{3}\, b^2 = \frac{1}{3} \cos^2 \psi$
8	$(11)_S$	$-6\, a^2 - 2\, b^2 = -(\frac{1}{3} \sin^2 \psi + \frac{1}{5} \cos^2 \psi)$
10	$(30)_A$	$-\frac{8}{3}\, b^2 + 8\, ab = -\frac{4}{15} \cos^2 \psi + \frac{4}{\sqrt{45}} \sin \psi \cdot \cos \psi$
10	$(03)_A$	$-\frac{8}{3}\, b^2 - 8\, ab = -\frac{4}{15} \cos^2 \psi - \frac{4}{\sqrt{45}} \sin \psi \cdot \cos \psi$
27	$(22)_S$	$4\, a^2 + \frac{4}{3}\, b^2 = \frac{2}{3} (\frac{1}{3} \sin^2 \psi + \frac{1}{5} \cos^2 \psi)$

Table 1. Eigenvalues of the operator $\mathcal{M}$.

$$\frac{a}{b} = \frac{\sqrt{5}}{3} \tan \psi$$

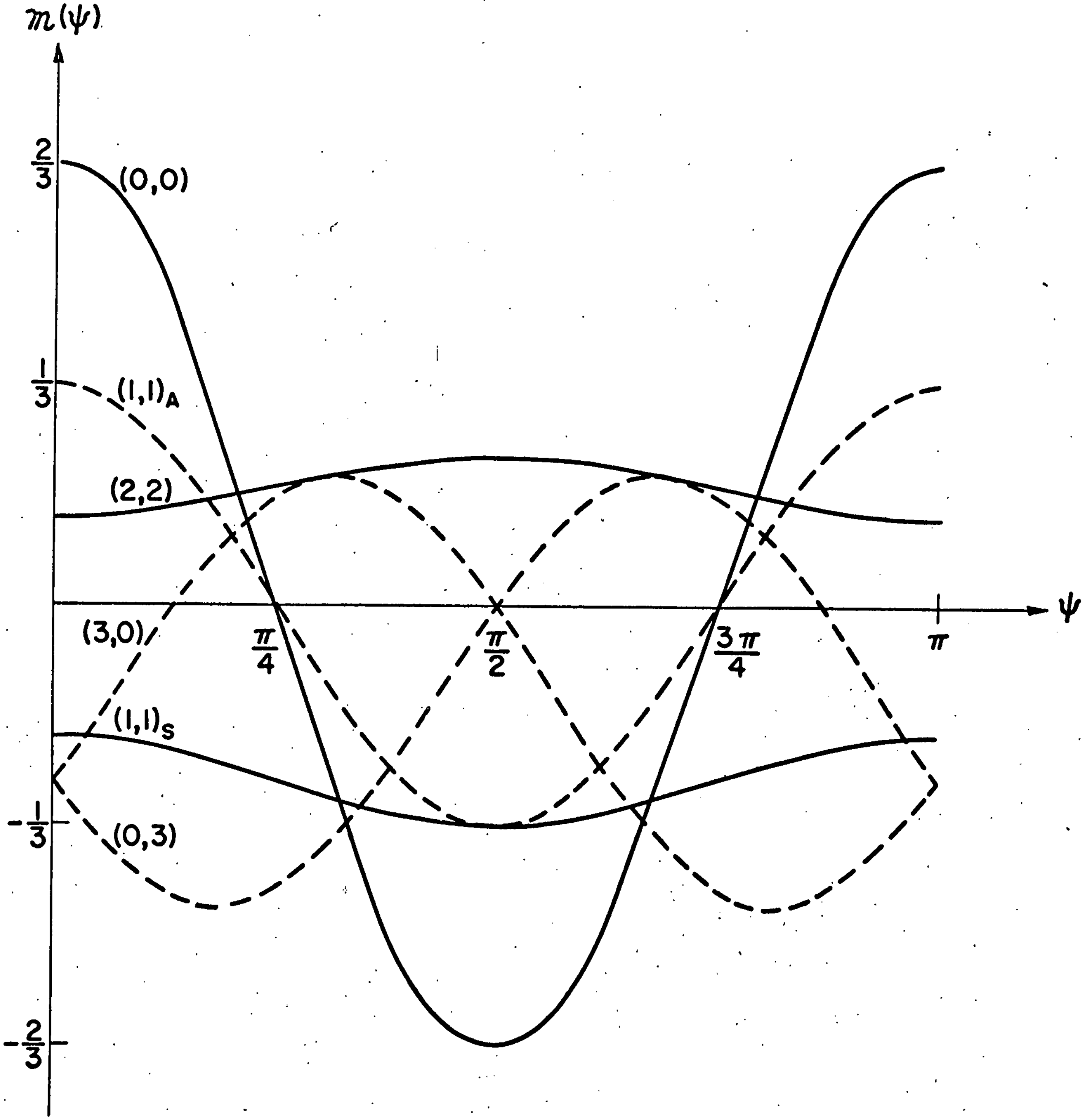


Fig. 3 BB-potentials in the different states $(\lambda\mu)$, plotted as functions of the mixing angle ψ.

The potentials are plotted as functions of ψ in figure 2 in the $(\lambda\mu)$-representation. This representation corresponds to scattering states which do not contain baryons of a well defined type. The situation is in fact somewhat similar to that met with in the Paschen-Back effect in atomic spectra representing the transition from the "weak coupling" (or j j coupling) to the strong coupling (or LS-coupling) limit. Thus, as long as the baryons are so far apart from each other that they may be considered as free particles, we may characterize the state by the individual isospin and hypercharge quantum numbers t_1, t_{1z}, t_2 t_{2z}, Y_1 and Y_2 of the baryons. In contrast, when the particles enter into the region of strong interaction, which is non-diagonal in the individual quantum numbers of the particles, the state is characterized exclusively by the conserved quantities $T = |\vec{t}_1 + \vec{t}_2|$, $T_z = t_{1z} + t_{2z}$, $Y = Y_1 + Y_2$ and $(\lambda\mu) \subset (11) \otimes (11)$, and it is built up by superposition of states containing two baryons of specified type.

If we consider the scattering of two such particles of a definite type, N and Λ, say, originally far apart from each other, we may, therefore conclude from the experimental observations that no strong attraction can exist in any of the states $|\lambda\mu\rangle$ which can be formed from the two particle system in question. However, the N Λ -state is characterized by the quantum numbers $T = \frac{1}{2}$, $Y = 1$ and, thus, contains non-vanishing components of the states $|22\rangle$, $|(11)_A\rangle$, $|(11)_S\rangle$ and $|03\rangle$, of which the state $|(11)_S\rangle$ exhibits an attractive core for all values of ψ. For NN-scattering ($Y = 2$) the T=0 state is a pure (03)-state whereas the T=1 state is identified as the (22)-state. For pp-scattering we have, thus, in particular

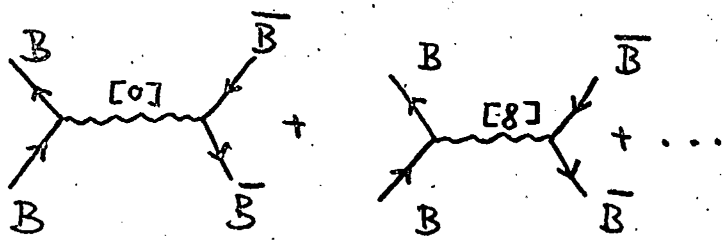
$$\mathcal{N}_{pp}(\psi) = \frac{2}{3}\left(\frac{1}{3}\sin^2\psi + \frac{1}{5}\cos^2\psi\right). \qquad (8)$$

Since this quantity is positive definite, we see that SU_3-symmetry, for any value of Ψ, reproduces the repulsive core in pp-scattering. In general the distribution of attraction and repulsion among the various states $|\lambda\mu\rangle$ will of course depend on the value of the mixing angle Ψ. However, as shown in the appendix of reference 3 and also directly checked from table 1, the weighted mean value over all representations of the potential $\mathcal{M}$ vanishes, i.e.

$$\sum_R d_R \mathcal{M}(R) = 0 \qquad (9)$$

where d_R denotes the dimension of the representation R. Hence it follows that the ratio between the total amount of attraction and the total amount of repulsion is fixed. All we achieve by varying Ψ is to shift the attractive core from some states to others.

The simplest way of attaining the necessary repulsion in all states is by postulating the existence of a ninth vectormeson, a unitary singlet. Being a singlet this vectormeson would mediate the same amount of repulsion in all $B\bar{B}$-states. Since the decomposition of the Kronecker product 8x8 includes the 1-dimensional representation, it is possible within the scheme of SU_3-symmetry to describe the singlet as a baryon-antibaryon state in analogy with the octet mesons. Such possibility rest on the assumption that a strong attraction exists in the (o,o)-state of the $B\bar{B}$-system. This attraction is partly provided by the singlet itself and partly (for suitable chosen mixing angle) by the octet, i.e. we may symbolically write

[0] =  (10)

The coupling constant $g_B[o]$ is, of course, quite independent of the coupling constant $g_B[8]$ as far as SU_3-symmetry is concerned. We shall later return to the possibilities of estimating the ratio $g_B[o]/g_B[8]$. In the present context it suffices to say that in order to compensate the strong attraction we may presumably have that

$$g_B[o] \gg g_B[8] \tag{11}$$

In that case the second graph in (10) may cause merely a small perturbation of the mass spectrum as derived from the singlet exchange graph alone. This implies that we, along with the singlet itself, would encounter 63 other degenerate vector meson states. Clearly, this means that the exchange of other multiplets must contribute in important manner to the structure of the singlet and the whole question becomes a matter of examining the self-consistency equations for the vector meson multiplets. We hope to be able to return to this problem in a future publication and restrict ourselves to remembering that it might well happen that the Regge trajectories associated with the ~~27~~ higher-dimensional multiplets, under influence of the symmetry breaking forces (which are rather strong after all) bend over without ever reaching the value $j = 1$. In that case this multiplets would not turn up as actual resonances although a particularly strong baryon-baryon interaction would be expected at energies where a member of this family of trajectories approached the value $j = 1$. Recently Sakurai[10] has directed the attention to a report by the Brookhaven-Syracuse group who has found some evidence for the existence of a narrow resonance ($\Gamma \leq 20$ MeV) in the $K\bar{K}$-system with a mass about 1020 MeV. If the spin-isospin assignment of this resonance turns out to be 1, 0, we have here a ninth vector meson, which we shall refer to as φ . The ω-meson is usually considered as the T=0 member of the unitary octet. In that case the φ is to be identified with the unitary singlet. The $K\bar{K}$ decay of the φ will then be forbidden by SU_3-invariance. Alternatively, the φ may be regarded as the T=0 member of the octet, thus identifying the ω with the unitary singlet. It should, however, be noticed that since the ω and the φ has in common the values of all the usual quantum numbers like spin, parity, isospin, G-parity and hypercharge, a particularly strong mixing of the SU_3-invariant states is to be expected. In the very crude approximation where the interference with all the other vector states is neglected one may write

$$
\begin{aligned}
|\varphi\rangle &= \lambda_1 \,|[0]\rangle - \lambda_2 \,|[8], T=0\rangle \\
|\omega\rangle &= \lambda_2 \,|[0]\rangle + \lambda_1 \,|[8], T=0\rangle \\
&\lambda_1^2 + \lambda_2^2 = 1.
\end{aligned}
\tag{12}
$$

A similar relation holds for the strength of the φB- and the ωB-couplings and we get in virtue of eq. (11):

$$\begin{aligned} f &= \lambda_1 g_0 - \lambda_2 g_1 \approx \lambda_1 g_0 \\ g &= \lambda_2 g_0 + \lambda_1 g_1 \approx \lambda_2 g_0 \end{aligned} \qquad (13)$$

where

$$f = g_{B\varphi} ,\ g = g_{B\omega} ,\ g_0 = g_{B[0]} \text{ and } g_1 = g_{B[8]_{T=0}} .$$

It is known that the Okubo mass formula[11])

$$m = m_0 \left[1 + \alpha Y + \beta \{ T(T+1) - \tfrac{1}{4} Y^2 \} \right] \qquad (14)$$

which works so remarkably well for the baryon and pseudoscalar meson octets predicts the mass of the T=0 member of the vectormeson octet to be 930 MeV when the parameters α and β are adapted to the values 750 MeV and 835 MeV for the ρ and M mass, respectively.*) One could now take the attitude that in absence of the $\omega - \varphi$ mixing the mass of the T=0 member of the octet should be given exactly by the formula (14).

We have already seen that the octet and the singlet is expected to be nearly degenerate in absence of symmetry violating interactions. With this assumption it is possible from the actual values 1020 MeV and 787 MeV for the masses of the φ and the ω, respectively, to estimate the coefficients λ_1 and λ_2 of equation (12) as the eigenvectors of the 2x2 potential matrix. One obtains in this way

$$\begin{aligned} \lambda_1 &= 0.63 \\ \lambda_2 &= 0.78 \end{aligned} \qquad (15)$$

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*) We use throughout this paper the letter M rather than K^* for the vector meson doublet.

showing that the φ contains a larger amount of the octet state whereas the ω is mainly the singlet state. For this reason we adapt the notation ω_0 for the singlet and φ_0 for the T=0 member of the octet. We see indeed that the admixture is rather strong. From (13) and (15) we get the ratio between the effective coupling constants:

$$g^2/f^2 = g^2_{B\omega}/g^2_{B\varphi} = 1.56 . \qquad (16)$$

It need hardly be emphasized that this estimate is not to be taken too literally. Attempts of fitting the isoscalar part of the nucleon form factors with an expression of the type (compare figure 4)

$$F_i^{(s)} = \frac{\lambda_2 g_0 \cdot \gamma}{M_\omega^2 - t} + \frac{\lambda_1 \cdot g_0 \cdot \gamma}{M_\varphi^2 - t} + C_i^{(s)} \qquad (17)$$

have turned out to be inconclusive owing to the large experimental uncertainties in the determination of the form factors. In fact, it seems at the present that two poles are needed for the fitting of the data for the isovector as well as the isoscalar part of the nucleon form factor. The circumstance that the mass of the second pole in both cases comes out more or less equal to that of the φ-meson can therefore hardly be attached any deeper significance. Presumably the form factors can be about as well fitted for any position between 1000 MeV and 1500 MeV of the second pole.

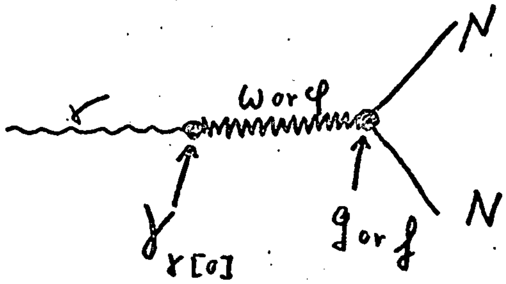

Figure 4. Vector meson contribution to nucleon form factor.

2. High Energy Behaviour of the Nucleon Scattering Cross Sections

The high energy behaviour of the nucleon-nucleon as well of the nucleon-antinucleon scattering cross sections is believed to be governed by Regge trajectories representing bound $N\bar{N}$-pairs in the t-channel (fig. 5). In the physical region of the S-channel ($S \geq 4N^2$, $t < o$) the invariant amplitude $F(s,t)$ describes the scattering process

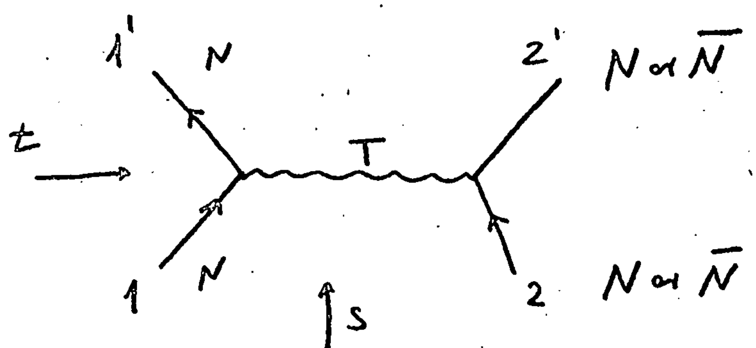

Figure 5.

$$1 + 2 \longrightarrow 1' + 2'$$

whereas the same function $F(s,t)$ when analytically continued to the physical region of the t-channel ($t \geq 4N^2$, $S < o$), accounts for the scattering process

$$1 + \bar{1}' \longrightarrow \bar{2} + 2' \quad .$$

Since the object exchanged between 1 and 2 is assumed to possess a definite value for the isobaric spin, we will be interested in such scattering states in the S-channel which correspond to a definite isospin state of the pair $1\bar{1}'$ in the t-channel. The forward scattering amplitude diagonal in these states may then be written as a sum of Regge poles with a particular value for the isospin[12,13]. The relative signs of the different terms are partly determined by the appropriate Clebs-Gordon coefficients, partly by the signature of the trajectories in question. In fact, defining the coupling constant for general values of t as

$$g^2(t) = \frac{\beta(t)}{\pi \alpha'_R(M^2)}$$

and assuming that this expression does not change appreciably and, at least, remains positive definite when extrapolated from $t = M^2$ to $t = 0$, we may write

$$\mathrm{Im}\, F(s,0) = -\sum_r \delta_r \left(\frac{s}{2N^2}\right)^{\alpha_r(0)} \cdot \pi g_r^2(0)\, \varepsilon_r \qquad (18)$$

where δ_r is the signature and $\varepsilon_r \equiv \left[\frac{d}{dt}\mathrm{Re}\,\alpha_r(t)\right]_{t=M_r^2}$.
With these conventions which seem to differ from those of reference 12) we get, remembering that the pseudoscalar trajectories do not contribute at $t = 0$, the relations given by Drell.[13]

In particular:

$$\frac{1}{2}\left(\sigma_{\bar{p}p} - \sigma_{pp}\right) = \pi \varepsilon_\omega g^2_{pp\omega}\left(\frac{s}{2N^2}\right)^{\alpha_\omega - 1} + \pi \varepsilon_\varphi g^2_{pp\varphi}\left(\frac{s}{2N^2}\right)^{\alpha_\varphi - 1} + \pi \varepsilon_\rho g^2_{pp\rho}\left(\frac{s}{2N^2}\right)^{\alpha_\rho - 1}$$

$$\frac{1}{2}\left(\sigma_{np} - \sigma_{pp}\right) = \pi \varepsilon_\rho g^2_{pp\rho}\left(\frac{s}{2N^2}\right)^{\alpha_\rho(0) - 1} \qquad (19)$$

Following Drell's analysis we note that in the case that no unitary singlet existed, thus identifying the ω_0 with the T=0 member of the octet, the experimental figures for the two differences in equation (19) ($\frac{1}{2}(\sigma_{\bar{p}p} - \sigma_{pp}) \sim 20\,\mathrm{mb}$, $\frac{1}{2}(\sigma_{np} - \sigma_{pp}) < 2\,\mathrm{mb}$ at an energy around 10 BeV) would require

$$g^2_{pp\omega} \gtrsim 10\, g^2_{pp\rho} \qquad (20)$$

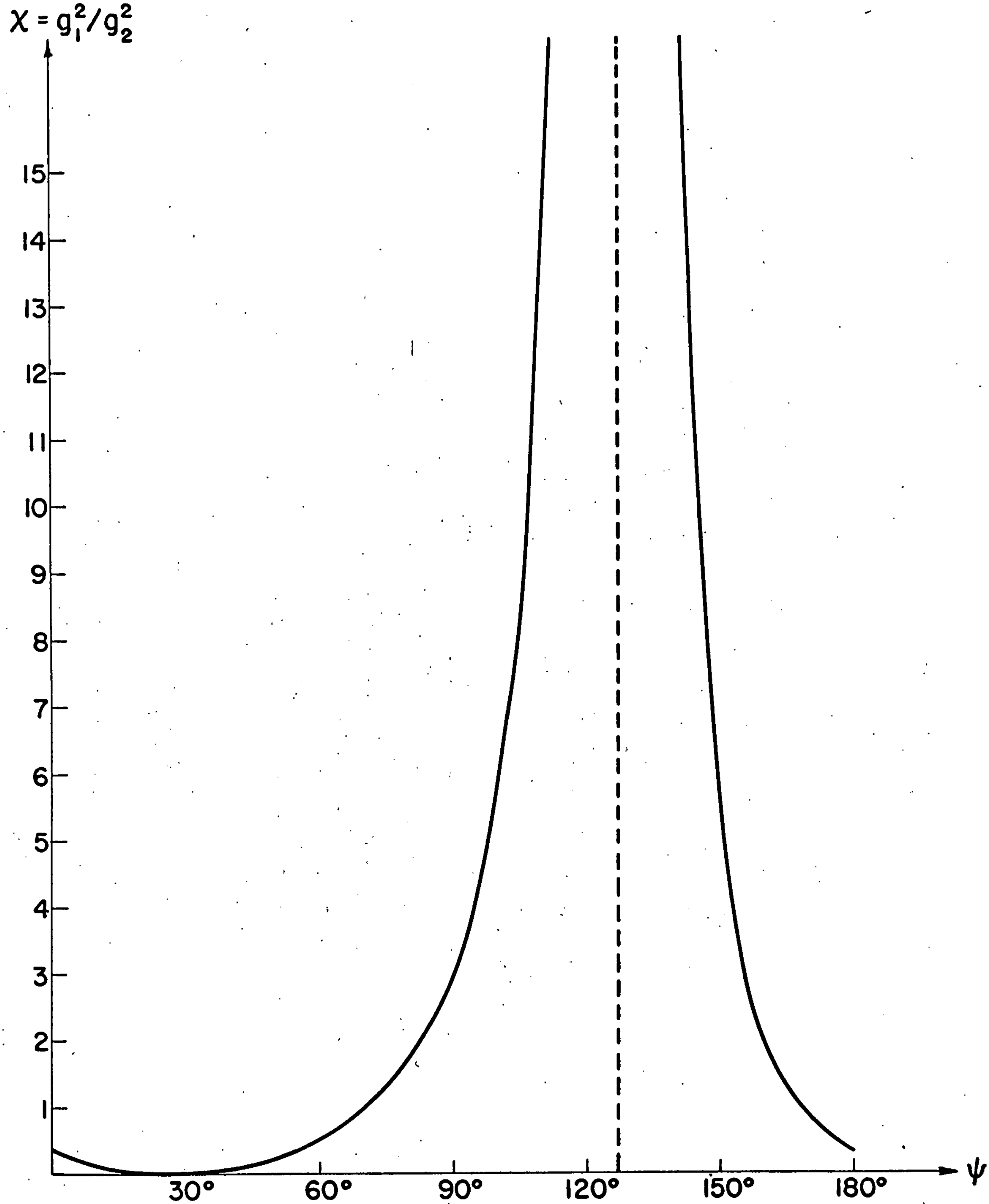

Fig. 6 The ratio $\chi = g_1^2/g_2^2$ as calculated from SU_3-symmetry as function of the mixing angle ψ.

since the ω and ρ trajectories are not supposed to deviate considerably from each other.

SU_3-symmetry gives the following relation between the two coupling constants (putting $g_{pp\rho} = g_2$):

$$X = \frac{g_1^2}{g_2^2} = \frac{1}{3}\left(\frac{3a-b}{a+b}\right)^2 = \frac{15\sqrt{5}\sin^2\psi + 3\sqrt{5}\cos^2\psi - 30\sin\psi\cos\psi}{5\sqrt{5}\sin^2\psi + 9\sqrt{5}\cos^2\psi + 30\sin\psi\cos\psi} \qquad (21)$$

Figure 6 shows this ratio as a function of ψ. The relation (20) restricts the possible values of the mixing angle ψ to the range

$$110^\circ \lesssim \psi \leq 145^\circ \qquad (22)$$

Although a comparison with figure 6 discloses the upper end of the interval (22) to be consistent with a repulsive core in most of the two baryon states, the $(11)_S$ and (30) states are still attractive.

The presence of the singlet invalidates the relation (20). Indeed, if we make the assumption that the ω-trajectory to a reasonably good approximation coincides with those of the φ and ρ, although it does not belong to the family of octet trajectories, we get instead of (20)

$$h^2 \equiv g^2 + f^2 = g_0^2 + g_1^2 \gtrsim 10 g_2^2 \qquad (23)$$

where we have used (13). This inequality allows the coupling between the ω and the baryons to be compatible with or weaker than the strength of the corresponding ρ-coupling. From low energy S-wave πN - scattering and the decay rate of the ρ, Sakurai [14] estimates $g^2_{\rho pp}/4\pi$ to be of the order of 2 or 3. ⟶ The $g^2_{\omega pp}$ may be determined for instance from the leptonic ω-decay $\omega \to e^+ + e^-$ or from the two pion decay branching ratio. It would be very interesting if the ratio x) turned out to be smaller than ten.

3. Inelastic Nucleon-Antinucleon Scattering.

So far we have only considered processes in which no exchange of hypercharge took place. However, it is possible to gain some additional support for the hypothesis of the existence of the φ by comparing the values of the ratio $g^2_{PYM}/g^2_{PP\omega}$, indirectly obtained from experiments, with the values predicted by SU_3-invariance. For the two diagrams in figure 7 we write, respectively,

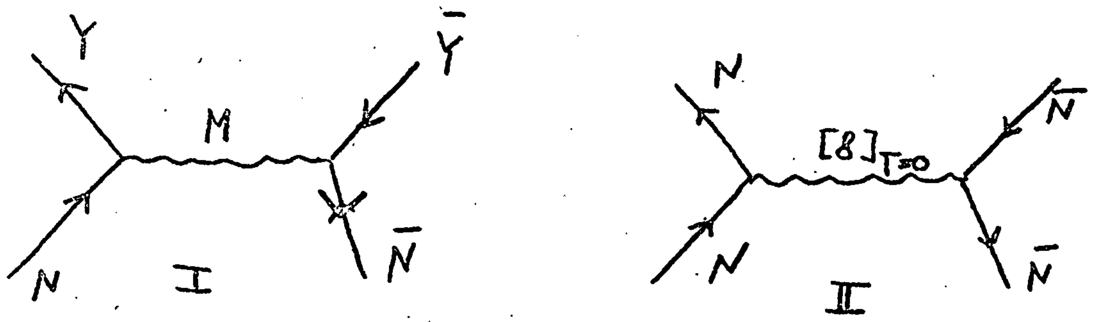

Figure 7

$$F_I(\nu) = g^2_{NYM} \cdot \alpha \cdot A(\nu)$$

$$F_{II}(\nu) = g_1^2 \cdot \beta \cdot A(\nu) \qquad (24)$$

where α and β are appropriate isospin factors. At sufficiently high energies we have, if no singlet exists, that

$$\Delta\sigma^{II} = \frac{1}{2}\left[\sigma_{\bar{p}p} - \sigma_{pp}\right] \approx g_1^2 \cdot \beta \cdot \mathrm{Im}\, A(0)\, \frac{4\pi}{p_c} \qquad (25)$$

where P_c is the center of mass momentum.

In the expression for the amplitude F_I we approximate the angular distribution by an exponential:

$$|A(\vartheta)|^2 = g(\vartheta)\cdot|A(0)|^2 \approx e^{\lambda(1-\cos\vartheta)}\cdot|A(0)|^2 \qquad (26)$$

where the quantities $|A(0)|$ and λ can be determined from a logarithmic plot of the measured angular distribution. Next we observe that

$$\frac{\mathrm{Im}\,A(0)}{|A(0)|} = \frac{-\delta \sin \pi\alpha(0)}{|1+\delta e^{-i\pi\alpha(0)}|} = \frac{-\delta}{\sqrt{2}}\,\frac{\sin\pi\alpha(0)}{(1+\cos\pi\alpha(0))^{\frac{1}{2}}} \approx \frac{1}{\sqrt{2}} \qquad (27)$$

assuming $\alpha(0) \approx \frac{1}{2}$ and noticing that the signature δ is negative. If the diagram I in figure 6 is the dominant one for the production of hyperon-antihyperon pairs in nucleon-antinucleon collisions we may identify

$$\sigma_{tot}^{I} = \int_{-1}^{+1} |A(\cos\vartheta)|^2\, d(\cos\vartheta). \qquad (28)$$

From (24)-(28) we then get

$$\frac{g^2_{NYM}}{g_1^2} = \frac{4\pi\sqrt{\lambda}}{\sqrt{2}\,P_c}\cdot\frac{\beta}{\alpha}\cdot\frac{\sqrt{\sigma^{I}_{tot}}}{\Delta\sigma_{\pi}} \qquad (29)$$

Experimental values for the total cross sections as well as for the angular distributions are available for the processes

$$N+\bar{N} \longrightarrow \begin{cases} \Lambda+\bar{\Lambda} \\ \Sigma^{+}+\overline{\Sigma^{+}} \\ \Lambda+\overline{\Sigma^{0}} \\ \Sigma^{0}+\bar{\Lambda} \end{cases}$$

at laboratory momenta of 3 BeV/c[15] corresponding to a total energy of 2.7 BeV in the center of mass system. Unfortunately, this energy is hardly high enough to justify the application of the asymptotic formulae essential for the derivation of equation (29). The observed total cross section were $\sigma_{\Lambda\bar{\Lambda}} = 73,5 \pm 23\mu B$, $\sigma_{\Sigma^+\bar{\Sigma}^+} = 38 \pm 7\mu B$, $\sigma_{\Lambda\bar{\Sigma}^0} = \sigma_{\Sigma^0\bar{\Lambda}} = 45.5 \pm$ [illegible] μb. The values for the coupling constants are not too sensitive to the precise value of λ. A choice of $\lambda \sim 7$ seems to fit the angular distributions reasonably well. Inserting the isospin factors ($\beta=1$, $\alpha=1, 2, 1$ for $\Lambda\bar{\Lambda}$, $\Sigma^+\bar{\Sigma}^+$ and $\Sigma^0\bar{\Lambda}$, respectively) and using the value 15 mb for $\Delta\sigma^{II}$ [16] we obtain from (29)

$$\begin{aligned} r_1 &= \frac{g^2_{\rho\Lambda M}}{g_1^2} = 0.14 \\ r_2 &= \frac{g^2_{\rho\Sigma^+ M}}{g_1^2} = 0.05 \\ r_3 &= r_1/r_2 = 2.8 \end{aligned} \qquad (30)$$

From SU_3-symmetry the corresponding ratios are given as functions of the mixing parameter:

$$r_1 = \frac{(3a+b)^2}{(3a-b)^2} = \frac{5\sin^2\psi + \cos^2\psi + 2\sqrt{5}\sin\psi\cos\psi}{5\sin^2\psi + \cos^2\psi - 2\sqrt{5}\sin\psi\cos\psi} \qquad (31)$$

$$r_2 = \frac{6(b-a)^2}{(3a-b)^2} = \frac{2}{3}\,\frac{15\sin^2\psi + 9\cos^2\psi - 6\sqrt{5}\sin\psi\cos\psi}{5\sin^2\psi + \cos^2\psi - 2\sqrt{5}\sin\psi\cos\psi}. \qquad (32)$$

These functions are plotted on figure 8, together with the ratio $r_3 = \frac{g^2_{\rho\Lambda M}}{g^2_{\rho\Sigma^+ M}}$.

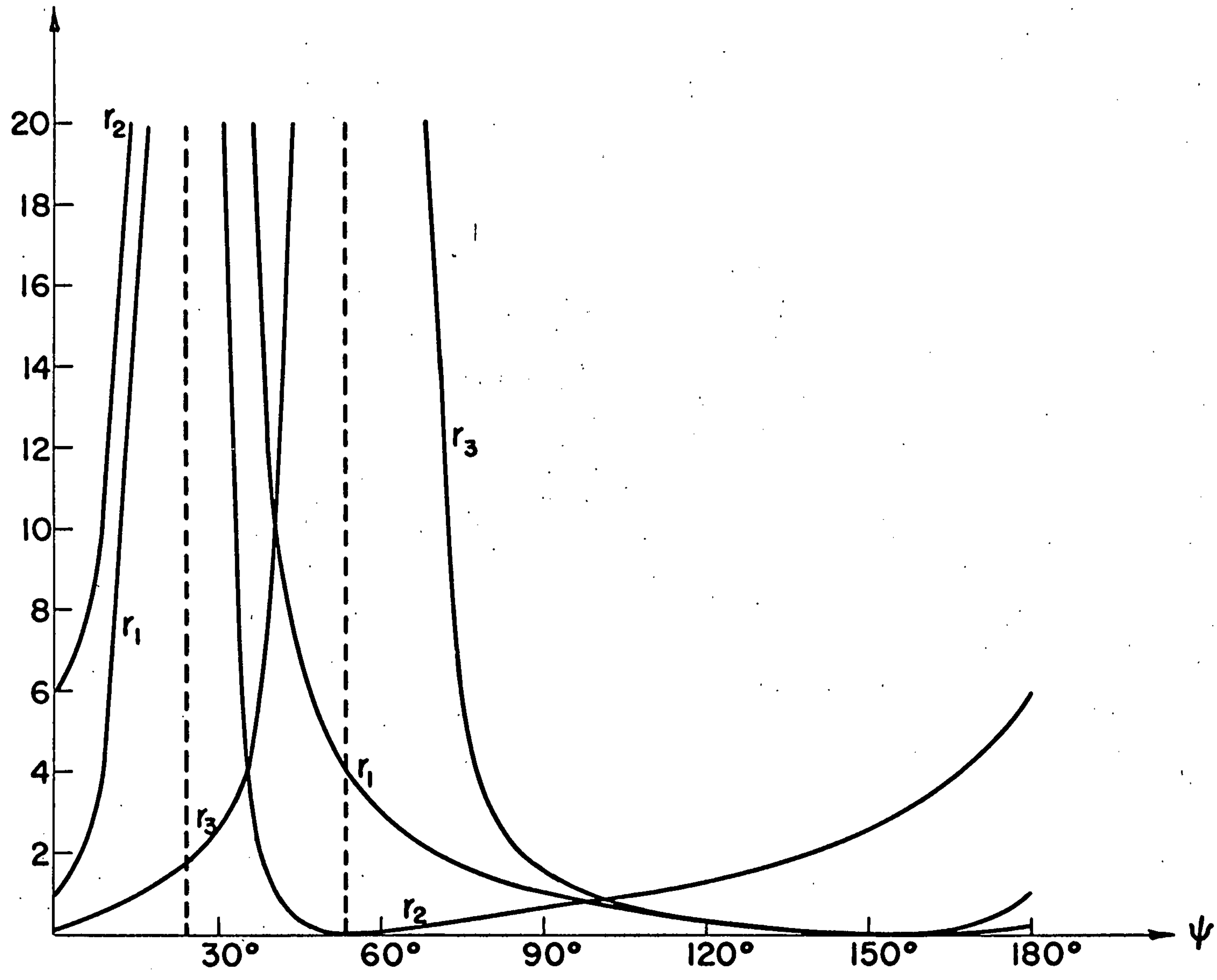

Fig. 8 The ratios r_1, r_2 and r_3 calculated from SU_3-symmetry.

It is apparent that no choice of the angle Ψ can make the predicted values of the ratios (31) and (32) consistent with the values (30). In presence of the singlet, however, the denominators in (30) are to be replaced by $h^2 = g_0^2 + g_1^2$ and, consequently, the numbers quoted there multiplied by h^2/g_1^2.

It would be tempting to assume that the ratio $r_3 = \frac{g^2_{\rho\Lambda N}}{g^2_{\rho\Sigma N}}$, as predicted by (31) and (32), is fairly well comparable to the value given by equation (30) provided no other graphs involving exchange of other multiplets turn out to be important. Indeed, we might determine the mixing angle ψ by adapting r_3 to the value 2.8 given by (32). The value $\psi = 31^\circ$ for the mixing angle thus obtained exhibits a suggestive coincidence with the value $\vartheta = 33^\circ$ which is obtained for the mixing angle in the case of baryon-pseudoscalar coupling.[3, 17] For $\psi = 31^\circ$ (30) is consistent with (31) and (32) if h^2/g^2 is around 400 or

$$\frac{g_0^2}{g_1^2} \sim 400 \qquad (35)$$

which at least agree with (24). Such large difference in the strength of the two interactions could of course hardly be justified if it did concern the actual particles and may be taken as a confirmation of the hypothesis of $\omega - \varphi$ - mixing. It may be noticed that the "renormalized" ratio g^2/g^2 is independent of the "bare" ratio g_0^2/g_0^2 if this last number is $\gg 1$ (cp. eq. (18)). Taking $\psi = 31^\circ$ we have

$$\frac{g_1^2}{g_2^2} \simeq 0.02 \qquad (36)$$

The fact that the two mixing angles ψ and ϑ tend to be of the same order of magnitude makes one suspect that they may be proved to be exactly equal, at least in a certain limit. Indeed, in a self-consistent model as that discussed by Cutkosky[3] the ratio between the mixing angles is in principle fixed by the self-consistency requirement.

In conclusion, the arguments presented above seem to prove that the predictions of SU_3-symmetry are not even qualitatively consistent with the experimental observations if no other vector mesons than the members of the octet are assumed to exist. The inconsistencies can be removed by the introduction of a unitary singlet vectormeson for which some empirical support is available. However, it might well be that other more complicated interactions, such as the exchange of members of other multiplets, could play an important role.

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REFERENCES

1) M. Gell Mann, Cal. Tech. Report CTSL-20 (Unpubl.), Phys. Rev. 125, 1067 (1962).

Y. Ne'eman, Nuclear Phys. 26, 222 (1961).

(Further references to basic papers on SU_3-symmetry are listed in ref. 3 and 8.)

2) M. Gell Mann and S. L. Glashow, Ann. of Phys. 15, 437 (1961).

3) R. E. Cutkosky, preprint, Carnegie Institute of Technology, (to be published).

4) H. P. Noyes, Proc. 1960 Ann. Conf. High Energy Phys.. Rochester.

H. Stapp, M. Moravcsik and H. P. Noyes, loc. cit.

G. Breit, loc. cit.

5) R. D. Dalitz, Proc. 1962 Intern. Conf. High Energy Phys. Cern, Geneva.

6) D. H. Wilkinson, Phys. Rev. Lett. 3, 397 (1959).

7) R. S. McKean, Phys. Rev. 125, 1391 (1962).

P. T. Matthews, Aix-en-Provence Conf. 1962.

8) R. E. Cutkosky, J. Kalckar and P. Tarjanne, Phys. Lett. 1, 93 (1962).

9) P. Tarjanne, Ann. Acad. Scient. Fennicae VI, 105, Helsinki 1962.

10) J. J. Sakurai, Phys. Rev. Lett. 9, 472 (1962).

11) S. Okubo, Progr. Theoret. Phys. (Japan) 27, 949 (1962).

12) B. M. Udgaonkar, Phys. Rev. Lett. 8, 142 (1962).

13) S. D. Drell, Rapporteur's talk, Proc. 1962 Intern. Conf. High Energy Phys. Cern, Geneva.

14) J. J. Sakurai, Proc. International School of Physics "Enrico Fermi", Varenna 1962.

15) Armentos et al. Proc. 1962 Intern. Conf. High Energy Phys. Cern, Geneva.

16) G. V. Dardel et al. Phys. Rev. Lett 5, 333 (1960).
S. Lindenbaum etal. Phys. Rev. Lett 7, 182 (1961).

17) P. Tarjanne (to be published).