STRANGE ISOBARS

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Table 1

	Mass (Mev)	$\Gamma/2$ (Mev)	i-spin	J
Y_1^*	1380	25	1	?
Y_0^*	1405	10	0	?
	1525	20	0	$\geqslant 3/2$
	1815	60	0	$\geqslant 3/2$
K^*	885	8	$\frac{1}{2}$	?

The five strange resonances I am going to build my talk around are listed in Table 1. Four of these have half integer spins and one is a boson (K^*).

My major concern here is more with theoretical models and theoretical techniques which have been proposed for understanding these and similar structures (predictably coming in the near future) rather than with the five resonances themselves. This is principally because apart from position and isotopic spin there is little else known experimentally about the resonances. My talk therefore is bound to be qualitative.

Like almost everything else in elementary particle physics our thinking about these resonances is motivated either /1/ by symmetry properties and group theory or /2/ by dynamical methods and dispersion theory. Thus there will be two strands running through my talk.

On the one hand I shall discuss :

1/ global symmetry and

2/ unitary symmetry

and confront their predictions with Table 1.

On the other hand I shall discuss :

3/ the general dispersion frame-work in which dynamical resonances arise and also mention /4/ some specific dynamical mechanisms suggested by Frazer and Ball and Baz which give rise to resonances.

2 - GLOBAL SYMMETRY -

2.1 - The Global symmetry hypothesis assumes that :

1/ Λ, Σ relative parity is even ; (Gell-Mann and Schwinger)

2/ $g_{\pi\Lambda\Sigma} = g_{\Sigma\Sigma\pi}$ (restricted symmetry) ; (Gell-Mann and Schwinger)

3/ $g_{\pi NN} = g_{\pi\Lambda\Sigma} = g_{\Sigma\Sigma\pi} = g_{\Xi\Xi\pi}$ (global symmetry) ; (Gell-Mann and Schwinger)

4/ Hyperon-K-Couplings are appreciably weaker than π-couplings ; (Gell-Mann)

It is assumption /4/ which makes global symmetry a "useful symmetry" ; useful in the sense that one may read off the π-Y and N-Y potentials and scattering amplitudes from π-N and N-N amplitudes at corresponding momenta.

This correspondence is made as follows. With /1/, /2/ and /3/ and neglect of K-interactions, π-hyperon interaction can be written as :

$$g_{\pi NN} (Z_1^+ \underline{t}\, \underline{\pi}\; Z_1 + Z_2^+ \underline{t}\, \underline{\pi}\; Z_2) \qquad (1)$$

Here :

$$Z_1 = \begin{pmatrix} \Sigma^+ \\ \frac{\Lambda - \Sigma^0}{\sqrt{2}} \end{pmatrix} \qquad Z_2 = \begin{pmatrix} \frac{\Lambda^0 + \Sigma^0}{\sqrt{2}} \\ \Sigma^- \end{pmatrix}$$

are two isotopic doublets which replace a singlet (Λ) and a triplet Σ. Clearly Z_1 and Z_2 doublets possess the same interaction with pions as nucleons. In terms of physical matrix elements one may therefore expect relations like :

$$(\pi^+ p/\pi^+ p) = (\pi^+\Sigma^+/\pi^+\Sigma^+) = \frac{1}{2}\,[(\pi^+\Lambda^0/\pi^+\Lambda^0) + (\pi^+\Sigma^0/\pi^+\Lambda^0) + (\pi^+\Lambda^0/\pi^+\Sigma^0) + (\pi^+\Sigma^0/\pi^+\Sigma^0)] \qquad (2)$$

2.2 - This type of correspondence of hyperon data with that for the nucleons has been tested for the following :

1/ Λ-N potential

2/ $\Sigma^- + p \longrightarrow \frac{\Lambda^0 + n}{\Sigma^0 + n}$ ratio at threshold

3/ π - Y phases for $J = \frac{1}{2}$ state at low energies.

A critical evaluation of /1/ and /2/ has benn presented by Dalitz (*) recently. Summarizing his results :

1/ ^{1}S-wave amplitude for Λ-N system at low energies as computed from (a rather tricky extrapolation of) N-N interactions gives an equivalent central potential with volume integral 370 Mev f^3, in good agreement with the ^{1}S Λ-N potential strength deduced by Dalitz and Downs directly from data on Λ-hyper nuclei. Similar remarks apply to the triplet potential.

2/ The computed value of $\Sigma^- + p \longrightarrow \frac{\Lambda^0 + n}{\Sigma^0 + n}$ of 1.8 at threshold (computation by de-Swart and Dullemond, using low-energy N-N scattering amplitudes and the "global correspondence") agrees rather well with the experimental value of $2.0 \pm .5$. The phase-space ratio would have been 4.6.

3/ The π-Y $J = \frac{1}{2}$ phase shifts occur sensitively in the determination of $K^- + p \longrightarrow \frac{\pi^+ + \Sigma^-}{\pi^- + \Sigma^+}$ ratio between threshold and 200 $\frac{\text{Mev}}{c}$. Using Dalitz (complex scattering-length parameters in conjunction with π-Y $J = \frac{1}{2}$ phase-shifts (as computed by correspondence with known π-N phases), Σ^-/Σ^+ ratio was computed by Salam and Pati (Kiev Conference (1959). The results disagreed completely with experiment. With new Humphrey and Ross parameters (shown by Dalitz in the previous talk) the calculation has been done again by M. Islam. Global symmetry + Humphrey & Ross solution /1/ still disagree but solution /2/ is compatible with the known Σ^-/Σ^+ ratio up to 200 $\frac{\text{Mev}}{c}$. (It may be worth remarking however that solution /1/ gives the better fit to K^-p scattering-data, having $p(x^2)$ of 48 % as against $p(x^2)$ of only 8 % for solution /2/).

(*) R.H. Dalitz - Rev. of Mod. Phys. July 1961.

Thus all of the three tests still leave global symmetry as a "useful symmetry". G. Alexander & W. Laskar at Berkeley are currently studying Λ-p interaction between 70 and 380 Mev and determination of Λ^0 magnetic moment is in progress at Brookhaven. These will give additional tests for this symmetry, though perhaps the most crucial negative test is the determination of Λ, Σ relative parity. Of course it is very necessary to determine theoretically to what extent these "tests" do depend on the global hypotheses. This has not been done at all.

2.3 - With this background one may inquire into the predictions of global symmetry regarding hyperon resonances. A simple correspondence argument with πN nucleon resonances has been developed by Kerth and Pais. Briefly Kerth and Pais show that corresponding to every pion-nucleon resonance, the global hypothesis predicts two pion hyperon resonances. When Λ, Σ mass difference is introduced the location of the resonances can be determined by the following procedure. From a group theoretic point of view the doublet representation of Λ and Σ corresponds to a representation of their isotopic spins I as sum of two half-integer isotopic spins t and k ; t is $\frac{1}{2}$ for all doublets (nucleons as well as Z_1 and Z_2) and k is $\frac{1}{2}$ for $\Lambda\Sigma$ and zero for nucleons :

$$\begin{aligned}(k_3 &= 0 \text{ for nucleons}\\ k_3 &= +\tfrac{1}{2} \text{ for } Z_1\\ k_3 &= -\tfrac{1}{2} \text{ for } Z_2).\end{aligned}$$

Generalising, if t^* is the isotopic value of any pion-nucleon resonance, there would correspond to it two pion-hyperon resonances with $I = \left(t^* + \frac{1}{2}\right)$ and $\left(t^* - \frac{1}{2}\right)$.

2.4 - Now three nucleon resonances (to be denoted as N^1, N^2, N^3) appear to be well established.

Table 2

	I	Q (Mev)	$\Gamma/2$	P^*_N $\frac{\text{Mev}}{c}$	State
N^1	$\frac{3}{2}$	160	45	230	$P_{3/2}$
N^2	$\frac{1}{2}$	430	30	450	$D_{3/2}$ (?)
N^3	$\frac{1}{2}$	600	50	570	$F_{5/2}$ (?)

Corresponding to N^1 there would be two hyperon resonances with I = 1 and 2, while N^2, N^3 would also give rise to two resonances each with I = 0 and 1. To obtain the masses, let us assume the following phenomenological mass formula for nucleons as well as Λ and Σ.

$$M = m(k^2) + \underline{k}.\underline{t} \ (*)$$

This leads to Table 3 :

(*) $M_N = m(0)$

$M_\Lambda = m\left(\frac{1}{2}\right) - \frac{3}{4}\Delta$ because $\underline{t}.\underline{k} = -3/4$ for Λ and $\frac{1}{4}$ for Σ.

$m_\Sigma = m\left(\frac{1}{2}\right) + \frac{1}{4}\Delta$

Table 3

		Mass (Mev)	Γ/2	p* (Mev/c)	Branching Ratio :
N^1	$I = 1,\ Y_1^1$	1380	23	$\frac{120(\Sigma)}{210(\Lambda)}$	10 . 1
	$I = 2,\ Y_2^1$	1530	70	270 (Σ)	0 : 1
N^2	$I = 0,\ Y_o^2$	1685	14	400 (Σ)	0 : 1
	$I = 1,\ Y_1^2$	1760	36	$\frac{460(\Sigma)}{510(\Lambda)}$	4 : 5
N^3	$I = 0,\ Y_o^3$	1855	33	530 (Σ)	0 : 1
	$I = 1,\ Y_1^3$	1930	82	$\frac{586(\Sigma)}{638(\Lambda)}$	1 : 1

In the pure doublet picture $\Delta = 0$:

$$\frac{\pi + \Lambda}{\pi + \Sigma} = 2 \qquad \text{for } Y_1^{(1)}$$

$$\frac{\pi + \Lambda}{\pi + \Sigma} = \frac{1}{2} \qquad \text{for } Y_1^{(2)},\ Y_1^{(3)}$$

The ratios $\frac{\Lambda}{\Sigma}$ as they appear in the last column of Table 3 are these intrinsic ratios multiplied by $\left(\frac{p^*}{p_N^*}\right)^{2l+1}$ to correct for phase-space. To obtain the width Γ for π-Y decay :

$$\frac{\Gamma}{\Gamma_N} = \frac{2}{3}\left(\frac{p_\Lambda^*}{p_N^*}\right)^{2l+1} + \frac{1}{3}\left(\frac{p_\Sigma^*}{p_N^*}\right)^{2l+1} \qquad \text{for } Y_1^1$$

$$\frac{\Gamma}{\Gamma_N} = \frac{1}{3}\left(\frac{p_\Lambda^*}{p_N^*}\right)^{2l+1} + \frac{2}{3}\left(\frac{p_\Sigma^*}{p_N^*}\right)^{2l+1} \qquad \text{for } Y_1^2,\ Y_1^3$$

Here Γ_N is the width for the nucleon resonances (*).

2.5 - Let us check with Table 1.

1/ The Y_1^* particle corresponds remarkably in mass, width as well as in Λ/Σ ratio to $Y_1^{(1)}$ of Table 3. Nothing is known about spin (J).

2/ The resonance at 1525 Mev (Table 1) has the right position as well as possibly the right spin ($D^{3/2}$) as $Y_2^{(1)}$ but the isotopic spin experimentally appears to be 0 rather than 2 and also the experimental width is too narrow. Even if the experimental i-spin assignment is wrong, it cannot be 2 since this resonance occurs in $K^- + p \longrightarrow \pi + Y$ reactions.

3/ There is some evidence from the work of Erwin, March and Walker and Stroot et al. (CERN) of existence of a "resonance" at 1580 Mev in the reactions $\pi^- + p \longrightarrow (\Sigma + \pi) + K$ with $I \geqslant 1$. The half-width of this bump is around 45 Mev. There is no evidence regarding its i-spin but it might conceivably correspond to $Y_2^{(1)}$ (I = 2).

4/ The resonance at 1815 Mev in Table 1 lies astride $Y_1^{(2)}$, $Y_o^{(3)}$.

5/ From global symmetry there is not even a suggestion of the existence of Y_o^* (mass 1405 Mev).

(*) Isobars $N^{(2)}$ and $N^{(3)}$ decay substantially in $N + 2\pi$. The estimate of $Y_1^{(2)}$, $Y_1^{(3)}$ widths has implicit the assumption that baryons + 2 pions width is about the same fraction for the nucleon as for the hyperon case. This and the complete neglect of $\bar{K}N$ channel makes the estimates for $\Gamma^{(2)}$ and $\Gamma^{(3)}$ highly tentative.

2.6 - To summarise, from earlier tests global symmetry appears to be a useful symmetry. What we are testing now is if it can provide a useful correspondence between nucleon and hyperon resonances. The most crucial tests are the spin of Y_1^* and the search for I = 2 resonance $Y_2^{(1)}$ around 1540 Mev. The circumstance that there exist resonances not predicted by global symmetry is not an argument against the existence of the symmetry because even relatively weak K-force may produce these through mechanisms we shall consider later.

3 - DYNAMICAL MECHANISMS -

3.1 - The work of Kerth and Pais is phenomenological. However the same predictions as these authors were made so far as $Y_1^{(1)}$ and $Y_2^{(1)}$ particles are concerned (earlier) by Amati, Stanghellini and Vitale (*) using a static model of pion-hyperon interaction and in fact the Pais & Kerth formula with its linear dependence on $\underline{t}.\underline{k}$ was tailored fo fit the Amati, Vitale and Stanghellini conclusions.

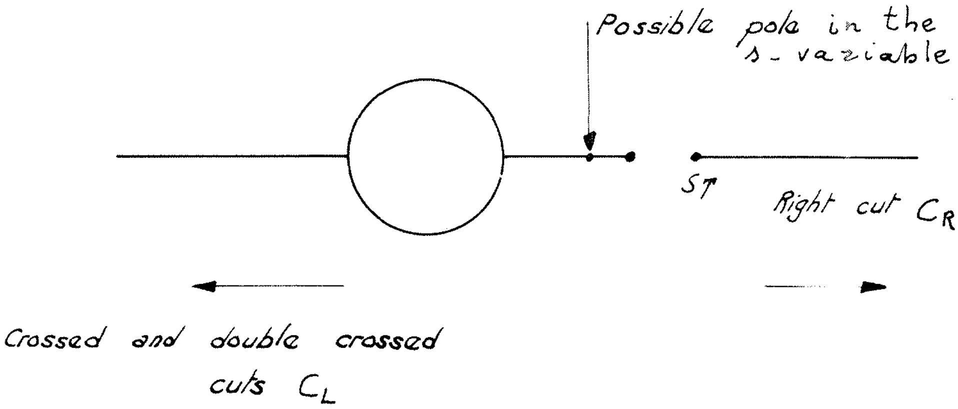
The object of this section is not to present these calculations but simply to review qualitatively the essential resonance formulation covering both πY and $\bar{K}N$ interactions. The pattern in all these cases is that of the first Chew-Low calculation of the 3,3 resonance rewritten in terms of the inverse T-matrix (T^{-1}). To bring out the essentials the same simplifying approximations are made as those by P.T. Matthews in an earlier talk.

3.2 - Let s, u and t stand for the three Mandelstam variables (s corresponds to energy and u and t to momentum transfer). The partial wave amplitude $T_l(s)$ is an analytic function of s except for three cuts :

1/ The right cut extending from the threshold S_T to $+\infty$.

2/ The "crossed cut" arising from poles and branch cuts of the amplitude T(s,u,t) associated with the variable u. The poles in u give rise to logarithmic singularities.

3/ The "double crossed cut" (which arises from poles and branch cuts of T(s,u,t) associated with the variable t). The crossed and double crossed cuts lies to the left of S_T.



The Mandelstam conjecture tells us that a dispersion relation of the following form can be written :

$$T(s) = B_l(s) + \frac{1}{\pi}\int_{C_R,C_L} \frac{\text{Im}\, T_l(S')}{S' - S + i\varepsilon}\, dS'$$

where B(s) contains a possible s-pole contribution as well as the logarithmic singularities arising from any poles of T(s,u,t) in the variables u and t. Consider Im T(s) on the left cuts. Chew and Mandelstam have suggested that one may approximate to Im T(s) on C_L by $g^2\, \delta_2(s - s_0)$. Thus $\int_{C_L}$ may be replaced by a term $\frac{1}{\pi}\frac{g^2}{s - s_0}$ which (as discussed by Matthews) we incorporate in B. The dispersion now reads :

(*) Amati, Stanghellini and Vitale, Nuovo Cimento 13, 1143, (1959) ; Phys. Rev. Letters 5, 524 (1960).

$$T_l(s) = B_l(s) + \frac{1}{\pi}\int_{C_R} \frac{\text{Im } T_l(s')\, dS'}{S' - S + i\varepsilon}$$

when finally $B_l(s)$ contains :

1/ The s-pole contribution of $T(s,u,t)$.

2/ Logarithmic singularities arising from possible t- and u-poles of $T(s,u,t)$. As a rule these logarithms can themselves be well approximated by pole of varying orders.

3/ Pseudo-Poles which simulate the contributions from the left cuts.

3.3 - From the above dispersion relation, one can write down an even simpler one for $T_l^{-1}(s)\, B_l(s)$. This is the relation :

$$T_l^{-1} B_l = 1 + \frac{1}{\pi}\int \frac{\text{Im}(T^{-1}B)}{S' - S + i\varepsilon}\, dS'$$

where we have assumed $T^{-1}(\infty) = B^{-1}(\infty)$. Now from unitarity (see Matthews) :

$$\text{Im } T_l^{-1}(s) = -k \qquad s > s_T$$

so that with all the approximations made,

$$T_l^{-1}(s) = B_l^{-1}(s) - F_l - ik$$

Here k is the channel momentum which can easily be expressed in terms of s, and :

$$F_l = \frac{p.v}{\pi} B^{-1} \int \frac{k' B(s')\, dS'}{S' - S}$$

For a single channel :

$$\text{Re } T_l^{-1} = B_l^{-1} - F_l = k \cot \delta_l$$

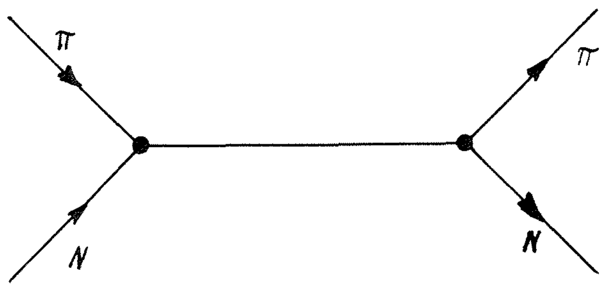
The condition for a resonance is that δ_l increases through $\left(n + \frac{1}{2}\right)\pi$ at $s = s_r$. Thus at the resonance energy :

$$B_l^{-1} - F_l = 0 \qquad s = s_r \qquad s_r > s_T$$

3.4 - Let us see how Chew-Low theory of 3,3 pion-nucleon resonance works out in this formulation. For $I = 3/2$, $J = 3/2$,

$$B(\omega) = f^2 \frac{4k^2}{3\omega}$$

(Instead of s we are using the variable w, the meson energy).



The existence of the resonance and its position ω_r is therefore given by :

$$0 = 1 - \left(\frac{4f^2}{3\pi}\right)\omega_r \int \frac{k'^3 \, d\omega'}{\omega'^2(\omega' - \omega_r)}$$

If this equation has a root for $\omega > m_\pi$ a resonance exists, otherwise not. For the pion-hyperon system we must generalise the above procedure to take account of the multi-channel nature of the problem. If T, B, k etc. are considered as matrices (**).

the resonance condition takes the form :

$$\det \operatorname{Re} T^{-1} = 0^{(*)}$$

3.5 - Using this formalism with a static approximation and neglecting $\bar{K}N$ channel, Amati Vitale & Stanghellini proved the following results :

$g_{\pi\Sigma\Sigma} = g_{\pi\Lambda\Sigma} = g_{\pi NN}$ (global symmetry)	$g_{\pi\Lambda\Sigma} = g_{\pi NN}$ $g_{\pi\Lambda\Sigma} \neq g_{\pi\Sigma\Sigma} \neq 0$ $\delta = \frac{g^2_{\Sigma\Sigma} - g_{\Lambda\Sigma}}{g^2_{\Sigma\Sigma} + g_{\Lambda\Sigma}}$	$g_{\pi\Sigma\Sigma} = 0$
J = 3/2 I = 1 $E_1 = m + \frac{3}{2}\Delta + \Omega$ I = 2 $E_2 = m + \frac{1}{2}\Delta + \Omega$ $\left(\frac{\Sigma}{\Lambda}\right)_1 = \frac{1}{2}\left(\frac{p^*_\Sigma}{p^*_\Lambda}\right)^3$ $\Omega = \frac{1}{g_1^2}\frac{1}{12\pi} \times \int \frac{d\omega' \, k'^3}{\omega'^2(\omega' - \omega_r)}$	J = 3/2 $E_2 - E_1 = 2\Delta + \frac{4}{3}\delta\Delta$ $\left(\frac{\Sigma}{\Lambda}\right)_1 = \frac{1}{2}\left(\frac{p^*_\Sigma}{p^*_\Lambda}\right)^3 \frac{1}{(1+\delta)^2}$	J = 3/2 Three resonances I = 1 and I = 0,2 degenerate

Wentzel has done a strong coupling calculation for the case $g_{\pi\Sigma\Sigma} = 0$ and finds the same result as Amati, Vitale & Stanghellini.

3.6 - Franklin (pre-print) has suggested that if a weak $\bar{K}N$ channel were included it may be possible to remove the degeneracy of I = 0 and I = 2 states. This way the 1385, 1520 and the possible 1580 Mev resonance may get identified with J = 3/2 <u>I = 1, 0, 2 resonances</u> respectively.

3.7 - Let us briefly look at the influence of the $\bar{K}N$ channel and also the inclusion of other singularities in B. The figure below shows the position of the physical region as well as of the Baryonic pole Y and the π-π cut. Also marked in the figure (taken from a paper of Feldman, Fulton & Wali, preprint) is the location of the singularity which may arise from an exchange of ρ and ω particles.

(*) R.H. Dalitz (Reviews of Modern Physics) and Feldman, Fulton & Wali (to be published) show that if a resonance lies between the thresholds for two channels, one must make the replacement $k \longrightarrow i|k|$ before taking the real part of T^{-1} for the closed channels.

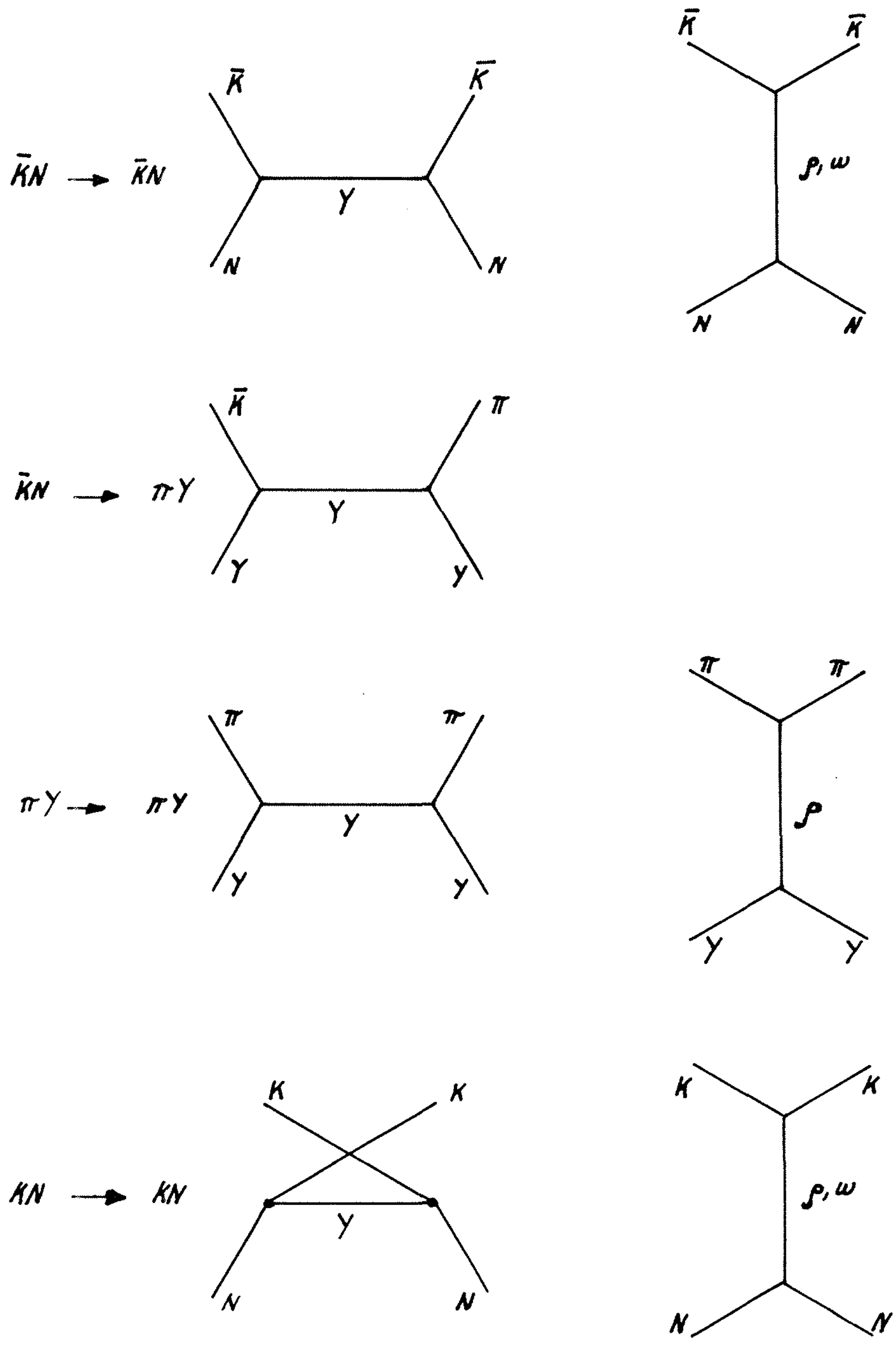
(**) e.g. for I = 0, T is a 2 × 2 matrix

$$\begin{pmatrix} T_{\bar{K}N \to \bar{K}N} & T_{\pi\Sigma \to \bar{K}N} \\ T_{\bar{K}N \to \pi\Sigma} & T_{\pi\Sigma \to \pi\Sigma} \end{pmatrix}$$

$$k = \begin{pmatrix} k_{\bar{k}\to\bar{k}} \\ 0 \end{pmatrix} \begin{pmatrix} 0 \\ k_{\pi\to\pi} \end{pmatrix}$$

In a static approximation note that the direct Y pole for $\pi Y \longrightarrow \pi Y$, $\bar{K}N \longrightarrow \bar{K}N$ and $\bar{K}N \longrightarrow \pi Y$ appears in $J = \frac{3}{2}$ state if $p(\Lambda\Sigma) = +1$ (assuming with Dalitz that $p(KN\Lambda) = -1$) and in $J = \frac{1}{2}$ state if $p(\Lambda\Sigma) = -1$. Notice that $\pi\pi$ cut approaches the closest to the physical region for $\bar{K}N$ - $\bar{K}N$ and $KN \longrightarrow KN$ processes. However due to the heavy mass of ρ and ω particles, the position of the singularity arising from the exchange of these particles is not so close. There are two points of view one may adopt here :

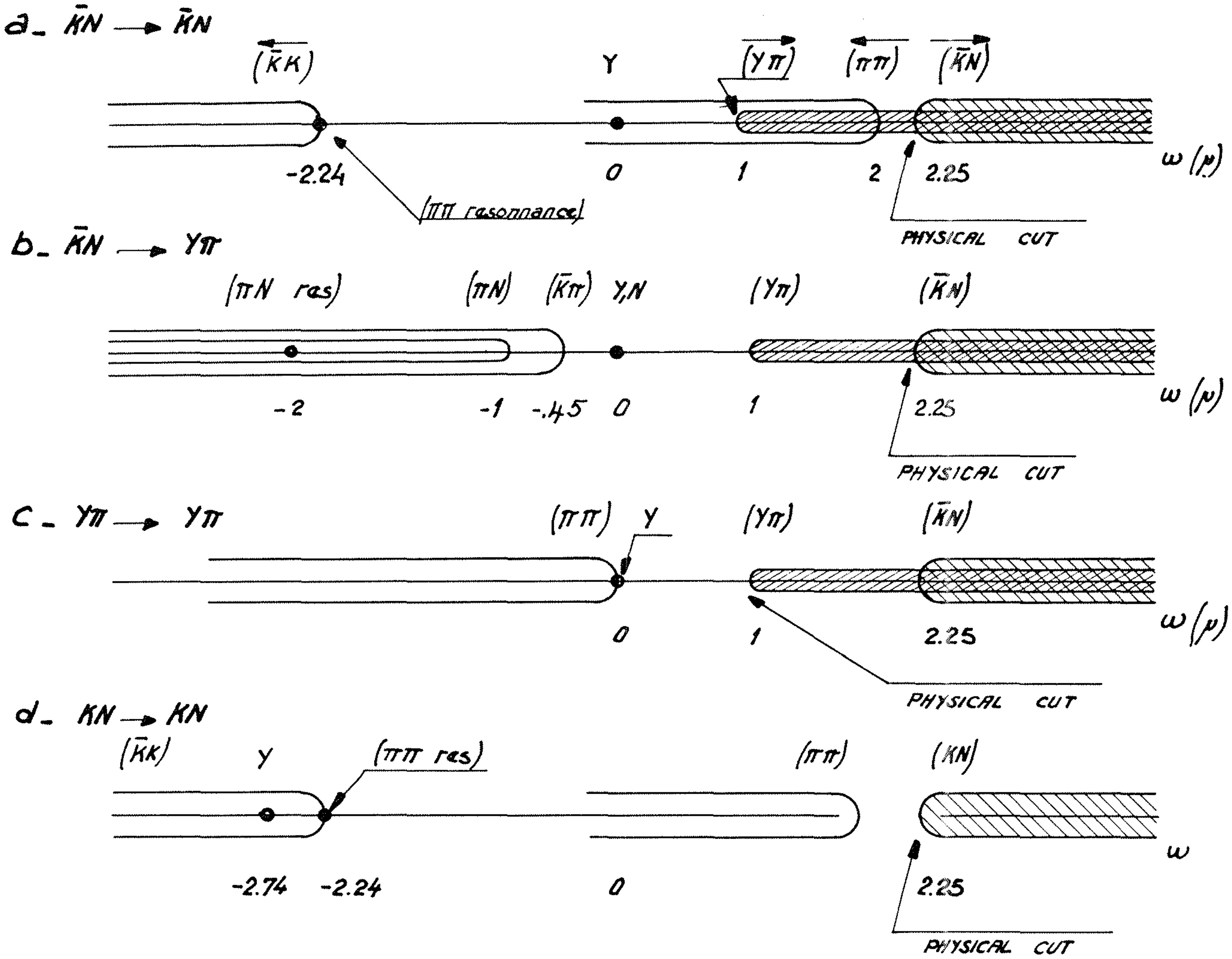
1/ Since ρ, ω singularities are fairly far, the expression $\frac{R}{S - S_{\rho,\omega}}$ may be approximated by the constants $\frac{R}{S_T - S_{\rho,\omega}}$.



2/ One may however include an expression like $\frac{R'}{S - S_{ABC}}$ to take account of the closeness of the 2π-cut (possibly the existence of a low mass Abashian Booth & Crowe (ABC) resonance at $S_{ABC} = 5\,m_\pi^2$ or the T = 0, J = 1 structure at $s = 10\,m_\pi^2$ which Matthews & Fubini have spoken about in connection with the electromagnetic form factors).

To date to my knowledge no calculations with all these singularities included are reported for $J = \frac{3}{2}$ state. For $J = \frac{1}{2}$ state two groups have computed matrix elements for $\bar{K}N$ and KN scattering. These are calculations of :

1/ Feldman, Fulton & Wali for $J = \frac{1}{2}$, $p(\Lambda\Sigma) = -1$. These authors emphasise the Y-pole while the ρ, ω pole terms are replaced by constants. They claim to have fitted all existing data on low energy scattering and absorption with reasonable coupling constants and to have found on extrapolation $k_K \longrightarrow i|k_K|$ below the physical $\bar{K}N$ threshold, the correct position and width of the Y_1^* resonance.



2/ Costa, Frye, Ferrari and Pusterla consider the case $p(\Lambda\Sigma) = +1$ so that for $J = \frac{1}{2}$ case, the Y poles do not make a strong contribution. However the ρ, ω singularities as well as a possible ABC singularity around $s \sim 10\,m_\pi^2$ is included. The scattering amplitudes computed by these authors are highly energy dependent. Again agreement with experiment and prediction of the Y_1^* state parameters are claimed.

3.8 - To relate the above with what Dalitz said in his talk let me repeat some of his remarks in respect of the $J = \frac{1}{2}$ scattering matrices. Ross and Humphrey have been able to fit all known scattering and absorption data by making the zero-range approximation. The amplitude for $\bar{K}N \longrightarrow \bar{K}N$ scattering is then given by :

$$T_{\bar{K}-\bar{K}} = \frac{1}{X - iY - ik_K}$$

where $\frac{1}{X - iY} = a + ib$ and $(a + ib)$ is the (constant) complex scattering length. The extrapolation $k_k \longrightarrow i|k_k|$ would give :

$$T_{\bar{K}\to\bar{K}} = \frac{1}{(X + |k|) - iY}$$

This is a resonance-like expression if $X < 0$ and approximates to :

$$T_{\bar{K}\to\bar{K}} \simeq \frac{X/\mu}{\left(\frac{X^2}{2\mu} - E\right) - i\,\frac{XY}{\mu}}$$

where $E \sim \frac{|k|^2}{2\mu}\ \left(\mu = \frac{m_K m_N}{m_K + m_N}\right)$.

Thus $\frac{X^2}{2\mu}$ gives the position of a possible pole and the half width $\frac{\Gamma}{2}$ is given by $-\frac{XY}{\mu}$. Using Humphrey & Ross data, one finds that there is no pole of the above type in $T_{\bar{K}\to\bar{K}}$ below threshold for I=1 state for either solution /1/ or /2/. However for I = 0 case solution /2/ may give a resonance with a mass of 1415 Mev. The width however is far too large $\left(\frac{\Gamma}{2} \approx 44 \text{ Mev}\right)$.

There are two conclusions one may draw if Y_1^* and Y_0^* do have $J = \frac{1}{2}$.

1/ The simple extrapolation procedure $k \longrightarrow i|k|$ is unjustified,

or :

2/ X and Y are strongly energy dependent. This is the Costa et al. point of view,

or alternatively Y_1^* and Y_0^* have $J = 3/2$.

Thus at the present time the only attitude one can take to the dispersion calculations I have spoken about is that it is in some ways too premature to try to "predict" theoretically the locations and widths of Y_1^* and Y_0^* or even to try to settle their existence. Experiment must decide the questions and then one may use the parameters to obtain reliable information about the relevant coupling constants.

4 - K^*-MESON -

Before considering the fourth (1815 Mev) resonance on Table 1, let us consider the K^* particle. Apparently this particle has i-spin $\frac{1}{2}$. The most eagerly awaited parameter about it is its spin. The clinching evidence which would distinguish its spin value would be provided by the decay mode :

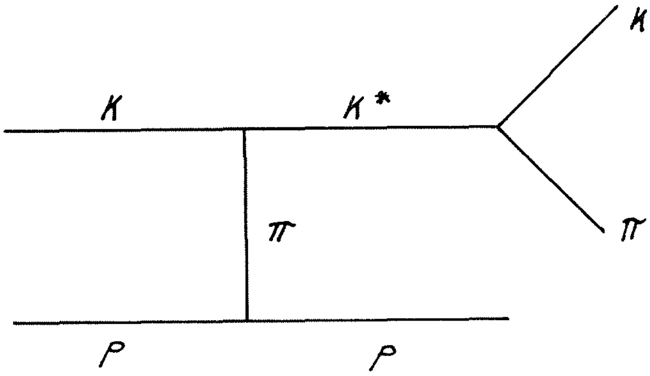
$$K^* \longrightarrow K + \gamma$$

Its absence would indicate J = 0. If J = 1 the expected decay probability is 1 % of the $K^* \longrightarrow K + \pi$ mode.

There is some slight evidence in favour of J = 1. Beg and De Celles and independently Chan have argued that the width of K^* together with cross section for the production process :

$$K^- + p \longrightarrow K^* + p \longrightarrow \begin{array}{l} K^- + \pi^\circ + p \\ \bar{K}^\circ + \pi^- + p \end{array}$$

determine whether J = 0 or 1 provided it is assumed that the lowest order diagram dominates :



In fact at 1.15 $\frac{\text{Bev}}{c}$ K^- incident momentum these authors compute :

$$\sigma = 1.32 \text{ mb for } J = 1$$
$$= .105 \text{ mb for } J = 0$$

assuming $\Gamma/2$ = 8 Mev.

The experimental value is ~ 2 mb favouring (within the context of such a calculation) J = 1.

A rather ingenious but tough proposal for spin determination has been made by Schwartz who showed that for the S-state annihilation of :

$$p + \bar{p} \longrightarrow \begin{array}{l} K^\circ + \bar{K}^{\circ *} \\ \bar{K} + K^{\circ *} \end{array} \longrightarrow K^\circ + \bar{K}^\circ + \pi$$

J = 0 K^* would imply the outgoing mesons must both be in K_1° or K_2° mode but not a mixture while for J = 1, a mixture is permissible.

The importance of the spin determination lies in the fact that the existence of vector K^*-meson together with the recently discovered dipion ρ ($\longrightarrow 2\pi$) and tripion ω ($\longrightarrow 3\pi$) (possibly vector particles) all nearly of the same mass (880 Mev, 750 and 780 Mev respectively) would complete the set of gauge particles associated with the unitary symmetry.

I do not wish to go here into a long discussion of gauge theories. Briefly the idea is this. Given a set of elementary particles which form a multiplet under some symmetry property, the gauge transformations (of the second kind) which may be associated with these symmetry properties can give rise to a set of vector mesons interacting with the source set of particles we started from. The electromagnetic field is the best known example of a gauge field. As is well known it arises from gauge transformations associated with conservation of electric charge.

The first serious attempt to build a gauge theory of strong interactions was made by J.J. Sakurai, who considered the conservation laws of baryons, hypercharge and i-spin and postulated the existence of three types of vector mesons, B_o, ω and ρ associated with these conservation laws. The experimental appearance of ω and ρ mesons seem to provide a definite encouragement to the gauge ideas.

In Sakurai's theory there is no direct connection between ρ and ω mesons. There is no reason why their masses should be nearly equal. Furthermore, although one of the important aspects of his theory is the near universality of all ρ couplings and all ω couplings, there is no prediction regarding the relative magnitudes of these two types of couplings.

In the first part of my talk I spoke of global symmetry as one of the possible higher symmetries which have been thought up ; higher in the sense that it goes beyond conservation of i-spin and hypercharge. Now in so far as the group of rotations in the isotopic space is the unitary group in a 2-dimensional space, a direct extension of the isotopic group in a search for higher symmetries is provided by the unitary group in a 3-dimensional space. Putting it another way if one wanted one type of symmetry which includes both conservation of i-spin as well as hypercharge, we would arrive naturally at the unitary symmetry in three dimensions. Gauge transformations associated with the unitary symmetry lead to just three types of vector mesons and these indeed are the ρ particle, the ω particle and a particle carrying $S = \pm 1$, $I = \frac{1}{2}$. K^* meson if it had unit spin would ideally fill the role.

If the unitary symmetry were an exact symmetry the masses of ρ, ω and K^* would be identical as also the coupling parameters of these mesons to all other particles. As will be apparent below the extent to which the symmetry may be expected to be violated is the extent to which Λ mass differs from nucleon mass and this appears also to be the extent to which ρ, ω and K^* appear to differ.

The connection of nucleons and Λ particles with the unitary group arises in the following manner. If one assumes with Sakata that the basic elementary set of particles consists of the triplet :

$$\chi = \begin{pmatrix} P \\ n \\ \Lambda \end{pmatrix}$$

the natural group of transformations under which the kinetic energy part of the free Lagrangian :

$$\chi^+ \gamma_4 \gamma_\mu \frac{\partial}{\partial x_\mu} \chi$$

remains invariant is just the set of unitary transformations in a (3)-space. Thus the unitary symmetry may also be considered as the natural symmetry arising from the Sakata model (•):

Clearly the unitary group will lead to other multiplets besides the basic multiplet $\begin{pmatrix} P \\ n \\ \Lambda \end{pmatrix}$, both for bosons as well as fermions. π and K mesons together with the elusive π^{oo} form a spin zero "tensorial" multiplet. For fermions there can exist a $J = 3/2$ multiplet incorporating :

$$N^*(I = 3/2), \quad N^*\left(I = \frac{1}{2}\right), \quad \Xi^*\left(I = \frac{1}{2}\right),$$

$$Y_1^*(I = 1), \quad Y_o^*(I = 0) \quad \text{and a triplet} \quad X^*(I = 1, \quad S = +1)$$

As an alternative to the Sakata triplet as providing the basic set of particles from which (together with the vector mesons) all other particles are formed, one may equally well take the octet of baryons consisting of Λ, Σ, N, Ξ particles as the elementary set. In this case Λ, Σ relative parity must be even. Also the allowed set of higher multiplets will be different and will contain 27 particles :

$$N^*(I = 3/2); \quad N^*\left(I = \frac{1}{2}\right); \quad \Xi^*\left(I = \frac{1}{2}\right); \quad \Xi^*(I = 3/2),$$

$$Y_1^*(I = 1) \quad ; \quad Y_o^*(I = 0) \; ; \; Y_2^*(I = 2);$$

$$X^*(I = 1) \; (S = +1) \quad ; \quad Z^*(S = 3, \; I = 1)$$

This is the 8-fold way of Gell-Mann and Neeman.

(•) A Salam & J.C. Ward, Il Nuovo Cim. (1961).

5 - BALL-FRAZER MECHANISM -

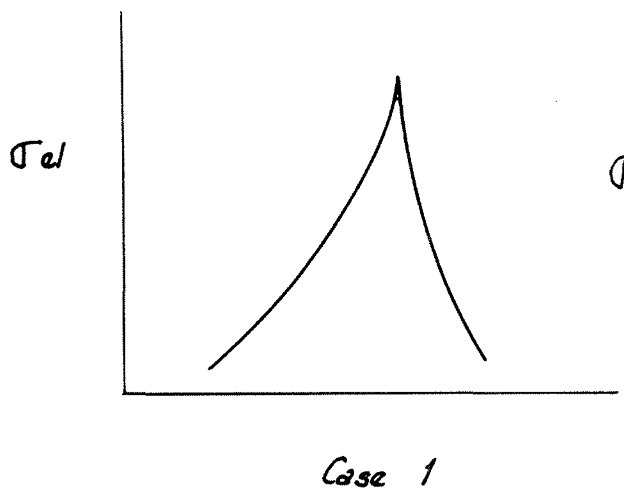
The discussion of K^* leads us naturally to the consideration of the highest known hyperon resonance at 1815 Mev. The most striking feature of this resonance is that it occurs at $K^* + N$ threshold in the $K^- + p$ channel.

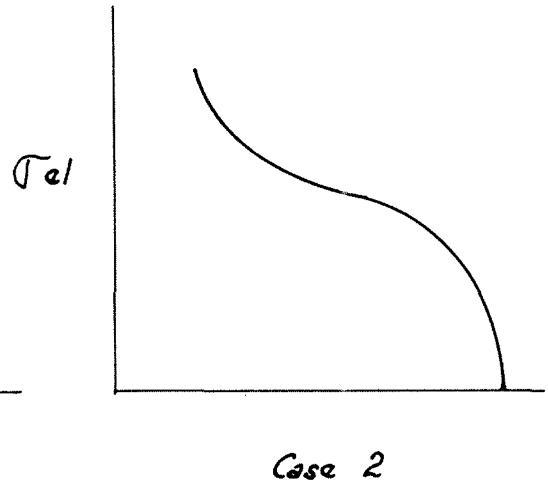
The existence of a resonance near an inelastic threshold has been noted before. The third pion-nucleon resonance $N^{(3)}$ occurs close to $N + \rho$, $N + \omega$ as well as $\Sigma + K$ thresholds (Baz, Kiev Conference, 1959), while $N^{(4)}$ the I = 3/2 resonance at 1900 Mev falls at $Y_1^* + K$ threshold.

Baz stated a general theorem to the effect that if at any given energy a new inelastic channel opens and the particles in the final state themselves possess a long range interaction, the cross-section for the reaction will show resonance-like bumps. I shall not discuss Baz's work because arguments given by him in support of his result were qualitative. I wish to describe in somewhat more detail some recent related remarks of Ball and Frazer.

Ball and Frazer prove the following result. A rapidly rising inelastic contribution to a single partial wave which attains a value near total absorption and remains large over a considerable energy range will produce a sharp peak in the elastic amplitude in the energy region where the inelastic cross section is rising. Ball and Frazer stress that (1) the inelastic cross-section itself need not be sharply peaked to produce a sharp sizable peak in the elastic. They also stress that (2) a large elastic peak does not occur at all inelastic thresholds ; the condition of a rapid rise to near-total absorption must be satisfied.

The result itself may be made plausible in the following manner. Whenever a new channel opens the well known cusp phenomena occur :





in the elastic channel of the shape shown in the figure. Ball and Frazer's conditions ensure the occurence of Case 1 so that the resonance they speak about essentially is an enlarged cusp.

The fact that the fourth K^-p resonance occurs at $\bar{K}^* + n$ threshold makes it plausible that Ball and Frazer mechanism may indeed be operative. To apply their theorem they must show that $\sigma_{in}(K^- + p \longrightarrow K^* + N)$ does rise to saturation. They (•) claim to show this by essentially a perturbation type of argument provided K^* has J = 1 and the final state is $D^{3/2}$.

(•) Frazer and Ball have invoked the same mechanism to give an explanation for the occurence of the second and third pion-nucleon resonances. According to this picture the resonances in the elastic peak come about on account of the rapidly rising $N + \rho$ production in $D^{3/2}$ state in the energy region around the second resonance (final state S-wave production of the ρ-particle) and also higher up in $F^{5/2}$ state (with final state p-wave production of ρ).

$I = \frac{1}{2}$ state is favoured over I = 3/2 by a factor of 4 to 1. The weakness of this argument as stressed by Sakurai is that the threshold for $\rho + N$ production seems to lie around 1680 Mev while thesecond pion nucleon resonance has a mass of 1510 Mev. Ball and Frazer in their paper used a mass value for ρ considerably smaller than the Wisconsin-Berkeley value of ($\approx$ 750 Mev). However, the third pion nucleon (mass value $\sim$ 1680) may indeed be a Ball-Frazer cusp-like resonance.

The assignment of an I-value to this resonance appears to be a tricky problem. Frazer and Ball argued on the basis of Clebsch-Gordon coefficients that I = 0 should be favoured over I = 1 state by a factor of 9 : 1. They would thus expect the cusp-like phenomena to occur in $K^+ + n$ cross-section but not in $K^- + n$. However a search by the spark chamber group at Berkeley seems to show a discontinuity of slope in $K^+ + p$ (I = 1, S = +1), but no sizable bump in $K^+ + n$ scattering unless the resonance is exceedingly narrow.

To correlate K^-p and K^+p bumps, baryonic intermediate state (I = 1, S = ± 1) are certainly possible as an alternative explanation but no theory to-date accommodates these with nearly the same mass. It seems to me the Frazer-Ball mechanisms with its emphasis on K^* production a phenomenon symmetrical for S ± 1 provides the simplest basis for comprehending the problem and it should be looked into more to state the complete set of conditions under which cusps may rise to the eminence of resonances.

Concluding then, we have a number of promising suggestions for understanding the Hyperon resonances. The most favoured one today seem to be the Franklin-Wentzel model ($p(\Lambda, \Sigma) = +1$, $g_{\pi\Sigma\Sigma} = 0$) and the Frazer-Ball mechanism. But quite honestly, a theorist can only express his humility and no more. If the data of Table 1 change, I would be surprised if any of these models will survive.