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LA PHYSIQUE DES LEPTONS

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INTRODUCTION -

Quand j'ai été invité à cette conférence, ce n'était que pour parler de l'expérience "g-2" du CERN. Selon le programme, je dois traiter maintenant de tous les résultats expérimentaux concernant les leptons. On ne s'improvise pas facilement expert dans ce domaine, j'espère toutefois que cet exposé pourra être profitable à quelques uns. Je vais parler d'abord des muons, puis des électrons. Il est dommage que nous n'ayons pas encore les résultats des expériences entreprises avec les neutrinos de grandes énergies. Je ne dirai mot des désintégrations leptoniques des particules étranges et si quelqu'un ici peut dire quelque chose à ce sujet à ma place, il sera le bienvenu. Enfin, je me limiterai aux résultats apparus depuis la conférence de 1960 à Rochester.

LES MUONS -

Propriétés intrinsèques.

Rien ne suggère qu'il existe plus d'un type de meson mu et son antiparticule, liés par les conséquences de l'invariance par rapport à la transformation CP (changement de signe de la charge et du moment magnétique, inversion des directions relatives du spin et de l'émission de l'électron). C'est-à-dire que les faits expérimentaux peuvent être résumés par les lois d'invariance dont on se sert alors pour calculer des effets bien loin de l'expérience.

Le méson mu possède une charge électrique, une masse, un spin, un moment magnétique dipolaire M_1. Il pourrait avoir de plus un moment dipolaire électrique s'il était dans un état non-invariant dans l'inversion de l'axe du temps. Le spin $\frac{1}{2}$ lui interdit les moments d'ordres plus élevés. La valeur $\frac{1}{2}$ du spin est directement établie par les mesures des spectres des atomes mu-mésiques, par le facteur gyromagnétique, et surtout par la fréquence de précession dans un champ magnétique de l'atome mésique constitué par un mu négatif lié à un noyau de spin $\frac{1}{2}$ (Ag ou P^{31}). De toutes ces propriétés intrinsèques, il n'y en a qu'une seule qui peut se calculer à partir de la théorie. La charge et le spin sont d'une nature discrète. La masse ne peut être déduite d'aucune théorie permettant même d'affirmer que le méson mu est environ deux cents fois plus lourd que l'électron. Je pourrais sans doute inclure parmi les propriétés intrinsèques la "charge faible" qui est la mesure de la force de couplage du méson mu dans les interactions faibles. Mais comme il n'existe pas de champ statique d'interaction faible, je n'ai pas considéré la "charge faible" comme une propriété statique.

Finalement, seul le facteur gyromagnétique nous permet de contrôler nos idées sur le méson mu et sur son interaction avec le champ électromagnétique et c'est pourquoi depuis 1957 un grand intérêt était attaché à l'expérience appelée "g-2" parce qu'elle permettait d'atteindre le facteur gyromagnétique avec une précision assez considérable. Cet intérêt s'accrut dès qu'on eut trouvé que le moment magnétique était voisin de 2, l'unité prise étant le demi magnéton. La théorie de Dirac nous permet depuis 30 ans de calculer le moment magnétique d'une particule de spin $\frac{1}{2}$, de charge électrique unité et sans structure en négligeant l'interaction avec le champ électromagnétique. Ce dernier effet introduit une faible correction dont le calcul exact a été considéré comme un des grands triomphes de l'électromagnétique quantique. Le moment magnétique ainsi calculé peut s'écrire :

$$\mu = g\left(\frac{e}{2\,m\,c}\right)\left(\frac{h}{2}\right) \quad g \equiv 2\ (1 + a) \tag{1}$$

$$a = \frac{\alpha}{2\pi} - 0.33\ \frac{\alpha^2}{\pi^2} + 0\ (\alpha^3) + \ldots \tag{2}$$

Le premier terme du développement de a se calcule à partir du diagramme suivant :

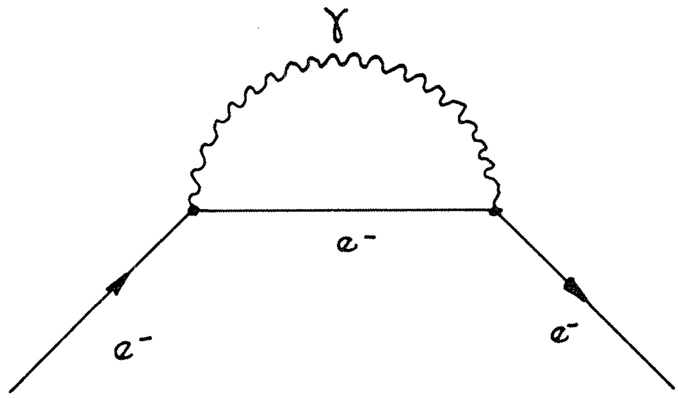


Figure 1

et provient uniquement de la présence du champ électromagnétique et du fait qu'il s'agit d'une particule de Dirac. Le deuxième terme vient des processus répétés et, en particulier, pour une petite partie, d'un diagramme :

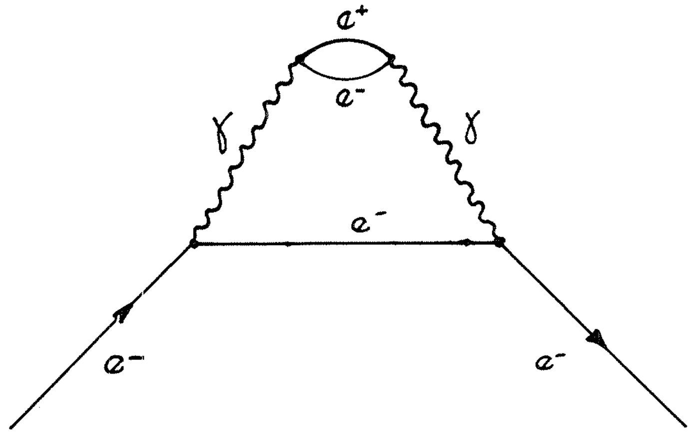


Figure 2

qui est responsable d'une certaine différence entre le méson mu et l'électron. Quand on veut connaître (g-2) pour une particule de Dirac de masse M au lieu de m, il faut ajouter au diagramme précédent qui comporte une paire de particules de masse M un autre diagramme similaire avec une paire d'électrons. Le résultat s'écrit :

$$a_{\text{théorique}} = \left(\frac{\alpha}{2\pi}\right) + 0,75\ \frac{\alpha^2}{\pi^2} + \ldots\ldots \simeq 0,001165 \tag{3}$$

Si on admet que l'électrodynamique quantique ne "fonctionne pas" pour des énergies supérieures à $\Lambda m c^2$ la valeur prédite est modifiée de la façon suivante :

$$\frac{\alpha}{2\pi} \rightsquigarrow \frac{\alpha}{2\pi}\left(1 - \frac{2}{3}\Lambda^{-2}\right) \qquad (4)$$

Notre connaissance de la valeur de "g" pour l'électron, conforme aux prédictions théoriques avec une précision de 2.10^{-6} (c'est-à-dire une précision sur (g-2) expérimental de 2.10^{-3}), nous apprend ainsi que l'électrodynamique quantique est valable jusqu'à 10 MeV ce qui, dans le cas du méson mu introduirait une incertitude sur g-2 d'un facteur 2 à 10. A vrai dire on sait, grâce aux résultats des expériences sur la diffusion e^-p, grâce aux mesures de sections efficaces des créations de paires e^+ et e^- à grand angle, etc. que rien ne change fondamentalement jusqu'à $q \sim 2\ f^{-1}$, c'est-à-dire 400 MeV ; mais on n'a aucune information sur la structure du muon à cette échelle. Un résultat sur (g-2) du muon en accord à 0.5 % près avec la théorie signifierait que l'électrodynamique quantique est valable jusqu'à 1,5 GeV et que le rayon quadratique moyen du muon est inférieur à 0,2 f.

Il faut se rappeler que l'on contrôle d'abord que le muon est une particule de Dirac, puis la validité de l'électrodynamique quantique et la structure du muon, et enfin le couplage éventuel du muon à des champs hypothétiques inconnus jusqu'à présent, qui contribueraient eux-mêmes au moment magnétique suivant le premier diagramme. Rappelons aussi que la série (2) ne converge point, quoique l'on a affirmé (je ne sais trop bien pourquoi) que les termes lointains ne contribuent pas trop. On peut mesurer (g-2) de plusieurs façons. Je ne vais décrire que la méthode et les résultats de l'expérience "g-2" du CERN, publiée au terme de sa première étape dans Phys. Rev. Letters du 1er février 1961 [1] avec une précision de 2 %. La précision définitive sera de 0,5 % environ. Il faut dire ici qu'une mesure de "g-2" pour l'électron a été obtenue par Schupp, Pidd et Crane [2] avec une précision de 0,2 %. C'est une belle expérience qui vérifie les mesures faites sur les électrons liés dans l'hydrogène.

Une expérience du type "g-2" est possible grâce un peu à un coup de chance : les équations relativistes purement classiques du mouvement du vecteur polarisation (lié au spin) sont telles que pour g = 2, les directions relatives de la polarisation et de la quantité de mouvement restent inchangées dans n'importe quel champ magnétique statique aussi compliqué soit-il dans l'espace. La valeur de g prédite dans la théorie étant très voisine de 2, on peut utiliser ce fait pour mesurer directement (g-2) jusqu'à une précision de 0,5 % qui mène immédiatement à une précision de 5.10^{-6} dans notre connaissance de g.

Notons que l'étude du mouvement du spin avec g = 2, même dans le cas d'un champ magnétique uniforme comporte nombre de chausse-trappes et qu'il est bon de se fier au papier de Bargmann Michel et Telegdi [3] qui s'appuie sur les propriétés classiques de la valeur moyenne temporelle du spin. Il suffit de dire que l'angle entre les vecteurs vitesse et spin d'une particule passant un temps t dans un champ toujours normal au plan de l'orbite s'écrit :

$$\vartheta_{lab.} = a \frac{e\,\overline{B}}{m_0 c} t$$

On essaie de travailler avec l'angle ϑ aussi grand que possible, qui implique le choix d'un champ magnétique aussi élevé que possible car le temps t est borné par la vie moyenne du muon.

La question de l'intensité est toujours importante ; elle est en partie conditionnée par les conséquences du théorème de Liouville appliqué ici à la conservation de l'espace de phase d'un faisceau de muons. Ce théorème nous apprend que l'intensité est plus élevée si on emploie une région plus étendue du champ magnétique. Tout cela nous conduit à utiliser un entrefer de 600 cm × 52 cm × 12 cm, avec 16 kilogauss. L'ordre de grandeur de l'angle spin-quantité de mouvement est :

$$\vartheta \sim 10^{-3} \times (8.10^{4}) \times (1.6\ .\ 10^{4}) \times (5.10^{-6}) = 360° \qquad (8)$$

La mesure de (g-2) à 1 % requiert une précision de 3° dans la mesure de ϑ et M. Liouville nous suggère d'accepter une grosse dispersion de la durée du stockage des muons dans l'aimant. On a fait cette expérience avec un champ purement statique de la manière illustrée dans la figure 3.

Les muons positifs injectés sont pour la plupart longitudinaux car ils proviennent de la désintégration en vol de pions près de la cible du cyclotron. Tout ce que je vous dis est assez approximatif, mais je vais décrire maintenant les éléments les plus importants de l'expérience. Les muons de 150 MeV/c entrent à gauche dans le gros aimant dont le cliché vous montre le plan médian. Leur quantité de mouvement est réduite par le ralentissement "M" en béryllium jusqu'à 90 MeV/c et les muons ne peuvent plus s'évader de l'aimant. Mais ils pourraient bien rencontrer "M" encore

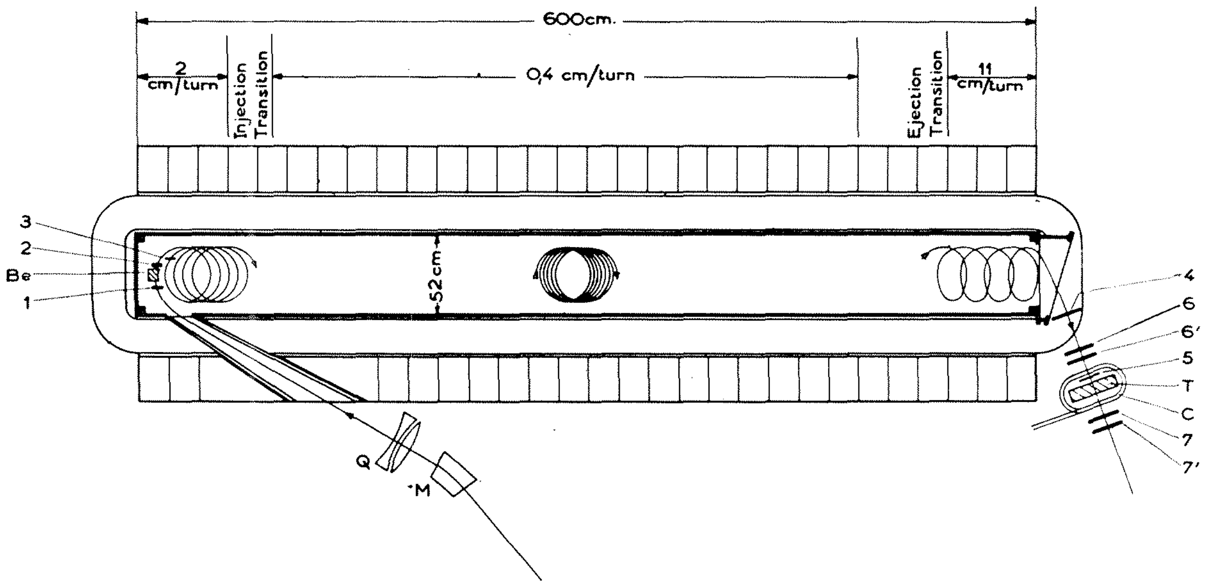


Figure 3

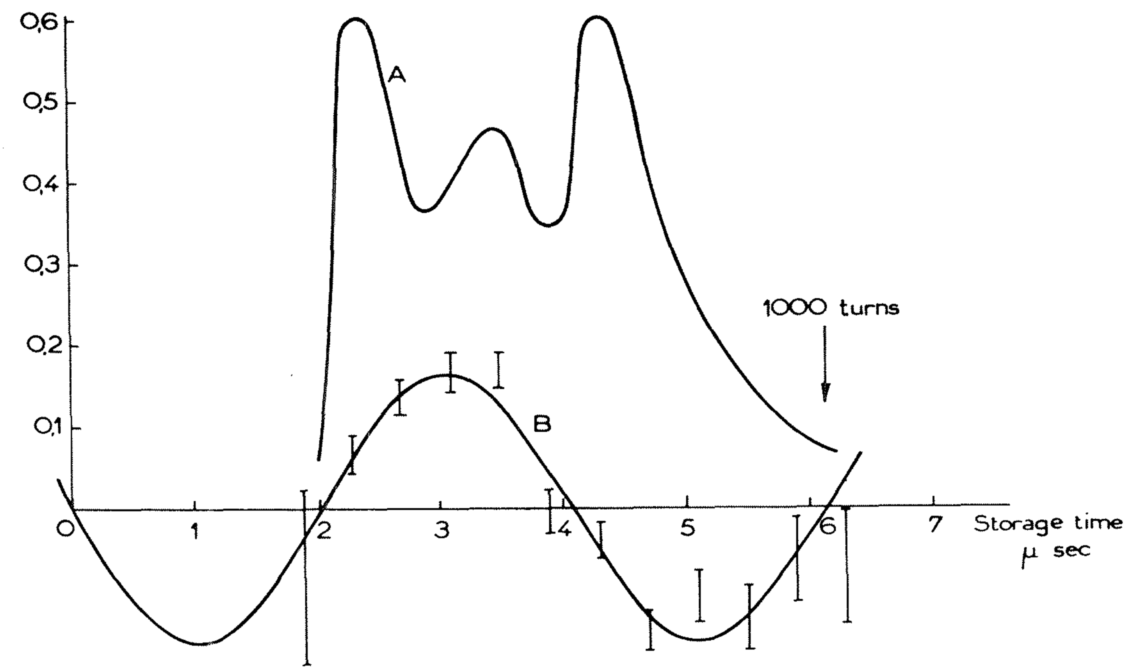


Figure 4

une fois et on se sert d'un gradient de champ pour persuader les orbites de longer l'aimant. La variation du champ magnétique a été soigneusement étudiée afin de donner au champ des propriétés focalisantes par rapport au plan médian. Mais la région d'opération est limitée par des bornes variées : "Stop-bands", etc., et les choses se compliquent un peu. Il faut, pour avoir une durée de stockage longue, maintenir une marche lente des orbites dans la majeure partie du trajet, sacrifiant ainsi de l'intensité à Monsieur Liouville, puis accélérer la marche afin que les muons puissent sortir de l'aimant et entrer dans l'analyseur de polarisation. Les scintillateurs 1, 2, 3, 4, 5, 6, 6', $\bar{7}$, nous donnent les signaux qui nous permettent de suivre la valeur de la polarisation en fonction du temps de stockage. Les circuits d'un "digitron" mesurent la durée du temps passé dans l'aimant et assurent la protection contre les erreurs fortuites (faux muon à l'entrée profitant de l'ouverture des circuits produite par un muon sorti de l'aimant, etc.). La composante transverse du spin est mesurée dans l'analyseur par une méthode de rotation alternée du spin de $\pm 90°$ à l'aide d'impulsions magnétiques. On obtient finalement une distribution de l'intensité en fonction du temps des muons qui s'arrêtent dans l'analyseur et dont on détecte l'électron.

Actuellement un muon par seconde entre dans l'analyseur dont un toutes les 4 secondes donne un électron détecté. A l'entrée du gros aimant on a 400 muons par seconde environ.

Sur la même figure on voit la composante transverse du spin en fonction du temps de stockage, obtenu par de nombreux "runs" à + 90° et - 90° alternés. L'évaluation du résultat apparaît maintenant simple : on détermine la meilleure valeur de la période de la sinusoïde, on évalue la polarisation initiale transverse des mésons mu acceptés dans le champ de stockage et qui entrent dans l'analyseur, on détermine la direction moyenne, en fonction du temps, des mésons injectés dans l'analyseur, on fait toutes ces intégrales et moyennes dont on doit réduire l'incertitude au-dessous de 2°, et on trouve comme ça une valeur pour "g-2" : On a fait ainsi pour arriver à une précision de 2 % et on fait de même actuellement pour essayer d'amener la précision à 0,5 % avec une intensité cinq fois plus grande et un solénoïde agissant sur le faisceau à l'entrée, afin de réduire d'un facteur 5 environ la composante horizontale transverse de la polarisation. Pour réduire la corrélation entre la polarisation transverse et la polarisation dans le faisceau et économiser le temps de mesure nécessaire à mesurer cette corrélation, on place une feuille de plomb mince à l'entrée de ce solénoïde. Le dernier chiffre publié a été :

$$a = a_{\text{théorique}}\,(0,983 \pm 0,019) = 0,001145 \pm 0,000022$$

Ce résultat confirme la validité de l'électrodynamique quantique (vertex - charge - photon) jusqu'à 300 MeV, il établit que le rayon moyen d'une structure éventuelle du muon doit être inférieure à 0,4 fermi, et enfin qu'il n'y a pas de couplage direct avec des champs autres que le champ électromagnétique avec une constante de couplage supérieure à 3.10^{-3} (pour un champ de même masse que le proton). Car le muon ne montre pas de structure jusqu'à telle énergie, on s'attend que le facteur de forme pour l'interaction faible sera constant jusqu'à une telle énergie, ce qui n'est pas tout à fait clair pour les nucléons. Il faut, bien entendu, tenir compte du fait que la croissance des sections efficaces peut toujours être limitée aux énergies élevées par l'existence d'un méson intermédiaire, le "boson vectoriel".

Maintenant, après avoir fait souffrir les anglais, aussi bien que les français sans doute, je vais continuer en anglais.

ELECTRIC DIPOLE MOMENT OF THE MUON -

The same apparatus has been used by the CERN group which at present consists of Charpak, Farley, Muller, Sens and Zichichi to improve upon previous upper limits to the magnitude of the electric dipole moment [4] of the mu$^+$. Aside from some intrinsic interest in indicating non-invariance under time reversal (an interest somewhat weakened because no EDM has been found), an EDM would change the anomalous precession frequency of the muon because of the interaction of the EDM $\left(\frac{f e \hbar}{m c}\right)$ with the motional electric field directed toward the orbit center ($E = \gamma\ \beta \times B \sim 3 \times 10^6$ V/cm in the rest system). The effect can be summarized in a vector diagram in the rotating frame (it helps to be also a nuclear resonance practitioner, to whom all properties are obvious in a suitable triply rotating frame).

For small f, the precession rate is affected quadratically but the spin moves in a plane inclined at angle $\vartheta = \frac{f\beta}{2a}$ to the horizontal. Thus to measure the EDM one needs only to look for a sinusoidal variation of vertical muon polarization with time.

The results obtained by rotating the pulsed flipping field in the polarization analyzer in order to detect vertical component of polarization are shown on figure 5. The amplitude of the asymmetry sets :

$$f = 3 \pm 6 \times 10^{-5} \tag{10}$$

small enough so as not to contribute to the error of a 0.5 % g-2 experiment.

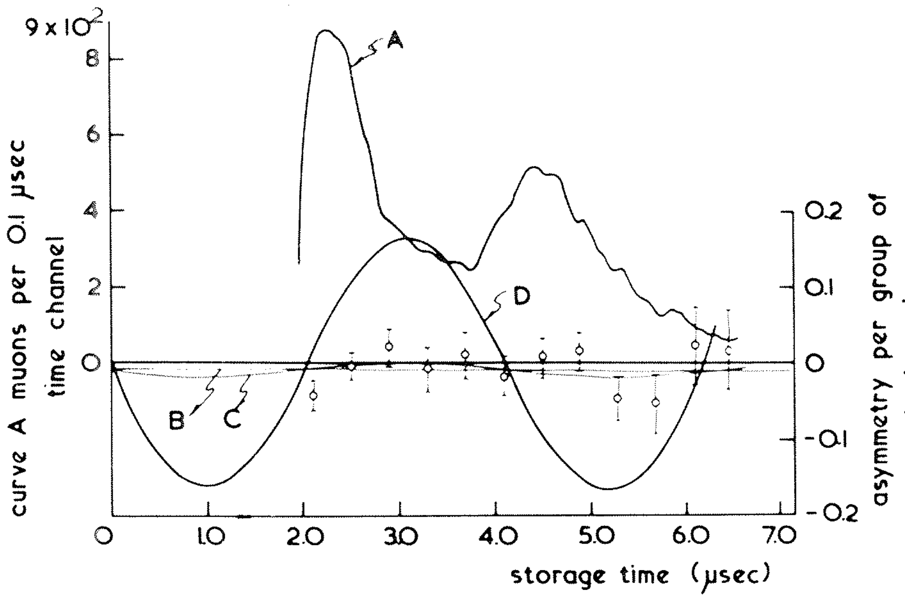


Figure 5

MAGNETIC DIPOLE MOMENT -

A direct comparison with a reference oscillator of the precession rate of stationary μ^+ in a 12 kGauss very homogeneous, perfectly stationary field has been performed by a group at Columbia [5]. The method used in this experiment is an extremely good one, but suffers from the disadvantage of being almost unexplainable. In fact what is measured is the "quasi-stationary" phase of the rotating decay-electron distribution, with respect to the reference frequency, for two groups of electrons-those from early decays and those from late. The phase difference thus obtained is plotted as a function of slight variation of magnetic field - the magnetic field for which this phase difference is zero being that field in which the muon precession rate is the same as the frequency of the reference oscillator. The proton spin precession rate in the same field (which is used to stabilize the field) is then a direct measure of the ratio of proton to muon moment.

The recent improvements in this experiment consisted of a much more uniform and homogeneous field (10^{-5} static and 10^{-6} time-dependent variations), a higher frequency (12 kGauss $\simeq$ 170 Mc) for more place shift during the muon lifetime, and more stable circuits (all transistorized). With these innovations they have found :

$$\frac{\mu_{\mu^+}}{\mu_p} = \begin{matrix} 3.18334 \\ \pm \quad 5 \end{matrix} \tag{11}$$

a result which together with the g-2 experiment determines very accurately the muon mass :

$$m_\mu = \begin{matrix} 206.763 \\ \pm \quad 5 \end{matrix} \tag{12}$$

in agreement with the four times less precise direct mass measurements of Chicago and Columbia , reported at Rochester. We now leave the question of muon static properties and go on to the muon production in pion decay.

MUON HELICITY -

Two good measurements of the helicity of the muons from pion decay have been made during the last year. As it happens the sign found is consisted with the V-A interaction - the opposite handedness would have required an interaction S, T, P [6]. At CERN, Backenstoss, Hyams, Knop, Marin and Stierlin [7] made use of a beam of ~ 8 Gev muons to produce knock-on electrons (Möller scattering) in iron magnetized at 30° to the beam direction. Reversal of the magnetization then gives a change in the intensity of scattering on the two magnetic electrons of the 26 in iron.

Their result is :

$$\mathcal{H}(\mu^-) = 1.17 \pm 0.32 \qquad (13)$$

A beem of transversely polarized μ^- of low energy was used by Bardon, Franzini and Lee [8] at Columbia to determine by Coulomb scattering (spin-orbit interaction, Mott scattering) the helicity of the μ^- in the pion rest frame. Their apparatus is shown in figure 6. Intensity is a real problem in this experiment, and the experimental apparatus was replicated ten times about the pion-beam. A long pulse of mesons from the cyclotron (vibrating target in a coasting beem) minimized accidental background and enabled them to obtain a result :

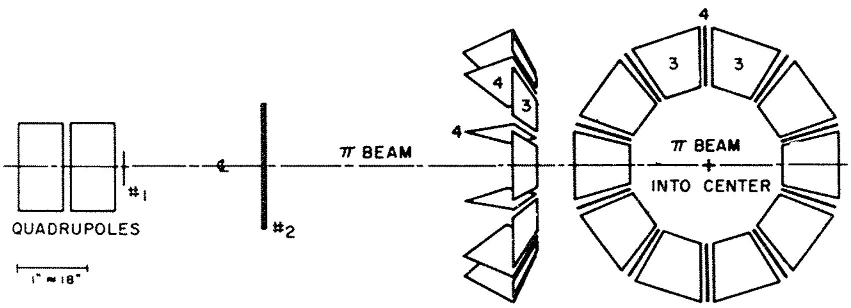
$$A = -0.090 \pm 0.031 \text{ as compared with} \qquad (14)$$

$$A_{th} = -0.080 \text{ predicted for positive muon}^- \text{ helicity.}$$

These two experiments prove that in :

$$\pi^- \longrightarrow \mu^- + \nu_\mu \qquad (15)$$

the mu$^-$ and therefore the $\bar{\nu}_\mu$ is right-handed just as is the $\bar{\nu}_e$, giving no indication that $\bar{\nu}_\mu \equiv \bar{\nu}_e$.



Experimental setup showing disposition of beam counter No. 1, collimating counter No. 2, and wheel arrangement of ten sets of counters and lead scatterers.

Figure 6

MUON DECAY -

A new measurement of the total decay rate of the positive muon :

$$(\mu^+ \longrightarrow e^+ + \nu + \bar{\nu}) \qquad (16)$$

has been made by Lundy [9]. The result is :

$$\tau = (2.204 \pm 0.004) \times 10^{-6} \text{ sec} \qquad (17)$$

and is an exercise in the use of the digitron and in the study of backgrounds. This result is in agreement with previous work, and does nothing to explain the 2-4 % difference in muon-decay and beta-decay coupling constants.

RARE MODES OF μ DECAY -

$\mu^+ \longrightarrow e^+ + \gamma$. The weakness of this decay mode (previously known to be $< 10^{-6}$) is one of the principal arguments against the identity of μ and e (or against that of ν_μ and ν_e). The decay through an intermediate scalar boson is expected to give a branching ratio $\frac{\alpha}{24\pi} \sim 10^{-4}$ wile that through a charged vector boson should give $R \sim 8 \times 10^{-4}$ for reasonable values of the mass and magnetic moment of the boson. On the other hand, it is claimed by Ebel and Ernst that an anomalous moment of 0.7 for the intermediate beson would reduce $R < 10^{-6}$. It is difficult to see why the intermediate boson, if it exists, should have as its primary aim in life to bring its anomalous moment within ~ 2 % of this particular critical value, and it is more reasonable to believe that we are seeing the effect of a new selection rule. I am told that a spark-chamber result is now available from Berkeley (Pensylvania group and Chamberlain group) with $R < 3 \times 10^{-8}$. The increased sensitivity is directly attributable to the improved discrimination as to collinearity, etc., available with spark chambers.

CATALYTIC DECAY - $A + \mu^- \longrightarrow A + e^-$ with A = Cu.

This coherent process has been sought again by Conversi, di Lella, Penso, Toller and Rubbia at C.E.R.N. [10]. Improved sensitivity was obtained by the use of a multiplate spark chamber, allowing a thick target for the stopping of a large fraction of muons, while retaining the electron energy resolution of a thin target (since the number of plates traversed by the electron is observed). Their result in searching for these 103 Mev e^- is $R < 1.6 \times 10^{-7}$ relative to capture.

This number is based on four events, all of which may be due to accidentals. Had the result of Sard et al. been correct, 100 real events should have been seen. This limit restricts the possible parameters of the intermediate boson in a different way than does the absence of $\mu^+ \longrightarrow e^+ + \gamma$, but in fact the extreme rarity of this letter indicates either two types of neutrinos or a selection rule of different type.

MUON CAPTURE -

In investigating the universality of the weak interactions one wants to observe also the process :

$$\mu^- + p \longrightarrow N + \nu \tag{17'}$$

and to measure its rate. This is a very complicated business, because the accuracy desired is ~ 10 %. What is worse is that the liquid hydrogen which for all other processes of high-energy physics is a source of free protons, is for this experiment a source in addition of great complication. For one has the rapid (μ^-p) formation by initial Auger transition and by radiation :

$$\mu^- + (e^-p) \longrightarrow (\mu^-p) + e^- \tag{18}$$

Of course the mesic atom can be singlet or triplet, with a few tenths ev difference in binding energy. ($kT = 2 \times 10^{-3}$ ev) so that at equilibrium one has only the singlet state. The cross section for proton exchange by the neutral (μ-p) is large ($\sim 10^{-8}$ sec.). The singlet (μ^-p) then form ($p\mu^-p$) molecules almost entirely in the $S_p = 1$ ortho state E = -2623 ev (since the Auger process is electric dipole for this P state, while it is monopole for the para-formation E = -2771 ev. These things have been calculated by Weinberg [11] and by Cohen, Judd and Riddell. We have only time here to remark that even the ortho ($p\mu p$) are in various states $|J,S>$ separated by a few tenths of a volt.

(1/2, 1/2), (1/2, 3/2), (3/2, 1/2), (3/2, 3/2), (5/2, 3/2)

the rate from the S = 1/2 state is :

$$W_{1/2} = 2\gamma_o \,(3/4\, W_o + 1/4\, W_1) \tag{19}$$

$$W_{3/2} = 2\gamma_o \,(W_1)$$

where W_o and W_1 are the (μ^-p) atom absorption rates in the S = 0 and S = 1 states. The fraction of S = 1/2 molecules formed is in principle calculable, but is denoted as $\frac{1}{2} < \xi < 1$.

What is new is that the elementary capture (17') has now been observed by Hildebrand in a bubble chamber. He finds :

$$W = (490 \pm 170)\ \mathrm{sec}^{-1}$$

as compared with $300 < W_{th} < 565\ \mathrm{sec}^{-1}$ predicted by V-A with $\xi = 1/2$ and 1 respectively. The experimental result already shows that the interaction is not V + A, for instance, and can probably be carried to the precision required to determine accurately the coefficient of A. The capture is seen of course by the recoil protons from those product neutrons which scatter in the hydrogen of the bubble chamber. The range of the proton is given by the angle with the line joining the μ-stop and the recoil proton. The spectrum of the 12 recoil protons is clearly different from the calibration taken with π^-, and the apparatus is obviously a very good neutron spectrometer.

At Columbia, Lederman et al. have counted capture neutrons from μ^- stopped in a pure H_2 target, but no capture rate is as yet available to me.

μ^- CAPTURE IN COMPLEX NUCLEI -

Both the moderate hyperfine conversion rates and their confirmation in the negative curvatures in the decay curve versus time, which were discussed at Rochester 1960 have now disappeared [12]. The rate of Auger radiation of the hyperfine energy is larger than had been calculated, because of the previous neglect of the lowest electrons state which can be ionized, which has a large density at the nucleus. The experimental evidence for conversion was removed by a better experimental technique, including the elimination of all carbon (as in scotch tape) and with improvements in the "digitron" used for timing the events and for the resolution of background. The higher expected conversion rates produce very small effects in the electron decay time-spectra, but a time-spectrum of the capture neutrons is sensitive to the capture rate in upper (F = 1) and lower (F = 0 for P^{31}) hyperfine states and should show the hyperfine transition time as well as the capture rates in the two states. Unfortunately no data is yet available for this experiment.

An interesting logical problem is posed by the following set of experiments :

In stopping polarized μ^- in P^{31} (red allotropic form) Ignatenko et al. [13] have observed precession of the electron distribution at half the free muon frequency and with about half the asymmetry observed in spin 0 nuclei. This is proof that the muon has spin 1/2 and shows that the population of the F = 1 state seen in precession is ~ 50 %. It shows also that the conversion rate is much less than predicted by Winston and Telegdi [14], in direct contradiction to a precession measurement at Chicago, which shows zero assymmetry. The group of Ignatenko et al. has recently shown that the precession is absent in black phosphorous (a conductor) and have jumped to the conclusion that the hyperfine transitions are catalyzed by the presence of conduction electrons. Since ~ 120 ev of hyperfine energy must be transferred to the Auger electron it seems completely unreasonable to expect any large difference in hyperfine conversion rates in red and black phosphorous. I believe that black phosphorous <u>depolarizes</u> the F = 1 states not by transition to the F = 0 hyperfine state, but just by re-orientation among the various m_F values, because of the presence of free electron spins and paramagnetic centers, and therefore it seem probable to me that all the conclusions of reference [13] are in error, as <u>no</u> conclusions can be drawn. Granting this, there remains the direct experimental contradiction of the two experiments on red phosphorous, about which I don't have any good idea.

"Fe - ANOMALY" -

The ~ 20 % peak in <u>decay</u> rate for a mu⁻ bound to Fe, which had been exhibited by Yovanovitch, has been contradicted by a group at CERN [15], by whom the decay rate is shown by an absolute experiment and by comparison with positive muons to be within ~ 2 % of the free μ rate. No convincing explanation has been given for the source of an error of the magnitude corresponding to the difference between these two rates.

MESIC X-RAYS -

The Stearns and Stearns anomaly has been removed by a careful repetition of the experiment at Chicago [16]. Rather than a linear yield of mesic x-rays, rising from 0 at 0 energy to ~ 1 at

100 kev and remaining ~ 1 above that energy, the yield remains ~ 1 down to 30 kev. Again, no convincing explanation is given as to the source of this error in the original experiment, and this unknown source of error may still affect the Chicago measurements in some measure.

WEAK MAGNETISM -

The conserved vector current (C.V.C.) hypothesis of Feynmann and Gell-Mann was used by Gell-Mann to propose a sensitive experiment as to the existence of the "weak magnetism", a striking consequence of the C.V.C. In the comparison of the $B^{12} \longrightarrow C^{12}$ and the $N^{12} \longrightarrow C^{12}$ β-decay there should be corrections to the shape of the transition to the ground state, of the form :

$$\frac{S(E, B^{12})}{S(E, N^{12})} = \text{const.} \ [1 + (A + \delta A) E] \ f(E)$$

f(E) is the form of the correction for internal bremsstrahlung. The value of A was given by Gell-Mann from the rate of a γ-decay in C^{12} as $(1.33 \pm 0.15) \times 10^{-2}\,\text{Mev}^{-1}$; while $\delta.A = -(0.25 \pm 0.15) \times 10^{-2}$ is the magnitude of the shape - dependent correction in the absence of weak magnetism. Note that the difference in the shape of two beta-spectra is used in this experiment, which eliminates many systematic errors. A very careful experiment has now been done by Mayer-Kuckuck and Michel [17] with the result that $(A + \delta A) = (1.13 \pm 0.25) \times 10^{-2}/\text{Mev}$, in agreement with the C.V.C. prediction of (1.08 ± 0.22).

This is certainly a very significant experiment, and I don't want to imply by my brief treatment of it, that it lacks importance. The authors note that a branching ratio of 3.5 % was assumed for the decay of N^{12} to the 7.6 Mev state in C^{12}. Should this be 3.2 %, $(A + \delta A) \longrightarrow (1.25 \times 10^{-2})$, etc. Obviously this branching ratio must be determined more accurately.

ACKNOWLEDGMENTS -

I should like to thank my colleagues for their help in the preparation of this talk, in particular Drs. V.L. Telegdi and G. Charpak.

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LEPTON PHYSICS

T. D. LEE

I - At present, the term "leptons" covers only three members : $\mu^{\pm}$; $e^{\pm}$ and ν ($\bar{\nu}$). Our known knowledge about their intrinsic properties and the interactions between these particles have just been summarized in the excellent talk by Dr. Garwin. This leaves me with a very easy task, since besides these experimental results there is very little and, in fact, almost no progress on the theoretical side of lepton physics during the last year. In this talk I will only summarize very briefly what are the known theoretical descriptions and emphasize several of the outstanding questions that have been asked up to the present time.

The Lagrangian describing the leptons can be written as :

$$\mathcal{L} = \mathcal{L}\,(\text{free}) + \mathcal{L}\,(\text{electromagnetic}) + \mathcal{L}\,(\text{weak}) \tag{1}$$

All known results are <u>compatible</u> with the following theoretical descriptions :

i) 2 component theory of neutrino :

$$(1 - \gamma_5)\ \psi_\nu = 0 \tag{2}$$

ii) Conservation of leptonic number, l.

Assign to each elementary particle a leptonic number l where :

$$\begin{array}{rcl} l = +1 & \text{for} & \mu^-,\ e^-,\ \nu \\ -1 & '' & \mu^+,\ e^+,\ \bar{\nu} \\ 0 & '' & \text{others.} \end{array} \tag{3}$$

The algebraic sum of l is conserved in all reactions.

iii) Time reversal invariance (or C P invariance)

and

iv) Except for the mass term, $\mathcal{L}$ is symmetric with respect to μ and e.

At present, i) and ii) are reasonably well established. The best evidence on iii) for weak interactions is still the measurement on polarized neutron decay done by the Chicago and Argonne group. The most impressive evidence of iv) for electromagnetic interaction comes from the recent (g-2) measurement and that for weak interactions is the π_{e2} and $\pi_{\mu 2}$ branching ratio. Both were done at CERN.

II - Both $\mathcal{L}$(free) and $\mathcal{L}$(electromagnetic) are known explicitly. The $\mathcal{L}$(weak), however, exists only in a phenomenological form. The weak interactions of the leptons can be separated into three groups :

i) Those involve only leptons ; e.g., μ-decay.

ii) Those involve also non-leptons but conserve strangeness ; c.g., β-decay.

iii) Those involve non-leptons but do not conserve strangeness ; e.g. K-decays and hyperon decays.

The effective Lagrangian for these three groups can be written as :

$$[\mathcal{L}\,(\mathrm{weak})]_{\mathrm{eff}} = \frac{G}{\sqrt{2}}\left\{(j_\lambda)_e\,(j_\lambda)_\mu^* + \sum_{l=e,\mu}(j_\lambda)_l\,[\mathcal{J}_\lambda + \mathcal{S}_\lambda]^*\right\} + \mathrm{h.c.} \qquad (4)$$

where :

$$(j_\lambda)_l = i\,\psi_l^+\,\gamma_4\,\gamma_\lambda\,(1+\gamma_5)\,\psi_\nu \qquad (l = e, \mu)$$

G is the Fermi constant for β and μ-decay $\left(G \cong \frac{10^{-5}}{m_p^2}\right)$ and $\mathcal{J}_\lambda$, $\mathcal{S}_\lambda$ are, respectively, the strangeness conserving and the strangeness non-conserving currents for the non-leptons. Because of the presence of strong interactions, the detailed structure of $\mathcal{J}_\lambda$ and $\mathcal{S}_\lambda$ are known only partially. Each observed decay process gives a direct measurement on certain matrix elements of these operators. While these elements are important for the structures of non-leptons, it appears that the dominant question concerning leptons and weak interactions is to seek for those features of $\mathcal{J}_\lambda$ and $\mathcal{S}_\lambda$ that do not depend on the detailed properties of the strong interactions.

i) The current $\mathcal{J}_\lambda$ has both a vector part V_λ and an axial vector part A_λ.

$$\mathcal{J}_\lambda = V_\lambda + A_\lambda \qquad (5)$$

One of the most remarkable achievements of last year in lepton physics is the experimental confirmation of the theoretical suggestion made by Feynman and Gell-Mann that :

$$\frac{\partial V_\lambda}{\partial x_\lambda} = 0 \qquad (6)$$

and V_λ, V_λ^* and the iso-spin vector part of the electromagnetic current $J^\lambda_{\mathrm{electr.}}$ (divided by e) together form an isotopic spin triplet.

ii) From the definition of $\mathcal{J}_\lambda$, it is obvious that the following commutation relation holds :

$$[\mathcal{J}_\lambda\,,\,S] = 0$$
$$[\mathcal{J}_\lambda\,,\,Q] = \mathcal{J}_\lambda \qquad (7)$$

$$\text{i.e.}\quad \Delta S = 0,\ \ |\Delta Q| = 1, \qquad (8)$$

where S = strangeness operator and Q = charge operator.

$\mathcal{J}_\lambda$ does not commute with the isotopic spin operator $\underset{\sim}{I}$. It has been suggested that under an isotopic spin rotation :

$$\mathcal{J}_\lambda \ \text{and}\ \mathcal{J}_\lambda^*$$

behave like the two $I_z = \pm 1$ members of a single isotopic triplet ($I = 1$). Therefore, the charge of isotopic spin of the non-leptons for processes satisfying (8) is restricted to :

$$|\Delta \underset{\sim}{I}| = 1 \qquad (9)$$

A consequence of this rule is that the rates :

$$\Sigma^+ \longrightarrow \Lambda^\circ + e^+ + \nu$$

and :

$$\Sigma^- \longrightarrow \Lambda^\circ + e^- + \bar{\nu}$$

are related to each other. Another consequence is that the high energy ν and $\bar{\nu}$ cross-sections; e.g.

$$\nu + p \longrightarrow n + l^- + 2\pi^+$$

and :

$$\bar{\nu} + n \longrightarrow p + l^+ + 2\pi^-$$

are related to each other.

That this rule $|\Delta \underset{\sim}{I}| = 1$ may be correct is based on the simple assumption that V_λ and A_λ, while differ under space reflection, may have similar properties under I-spin rotation.

iii) From its definition, $\mathcal{S}_\lambda$ satisfies :

$$[\mathcal{S}_\lambda , Q] = \mathcal{S}_\lambda \qquad \text{and} \qquad [\mathcal{S}_\lambda, S] \neq 0 \tag{10}$$

α) It has been suggested by Feynman than Gell-Mann that :

$$[\mathcal{S}_\lambda, (Q-S)] = 0 \tag{11}$$

i.e. $\Delta Q = \Delta S$ rule.

β) It has been further suggested by Marshak and collaborators that under a general isotopic spin rotation, $\mathcal{S}_\lambda$ behaves like a single $I = 1/2$ spinor ; i.e.

$$|\Delta \underset{\sim}{I}| = 1/2 \ \text{rule} \tag{12}$$

The validity of (α) is necessary for that of (β), but not vice versa.

Consequences of (α) are, e.g., $K^0 \not\rightarrow e^- + \pi^+ + \bar{\nu}$ and $\Sigma^+ \not\rightarrow n + e^+ + \nu$.

Consequences of (β) are, e.g., the decays :

$$K^+ \longrightarrow \pi^0 + l^+ + \nu \qquad (l = e, \mu)$$

and :

$$K^0 \longrightarrow \pi^- + l^+ + \nu$$

are related to each other.

iv) The $\Delta Q = \Delta S$ rule suggestion was first obtained by Feynman and Gell-Mann by assuming the $\mathcal{L}$(weak) for the weak leptonic process to be of a very simple form :

$$\mathcal{L}\,(\text{weak}) = \frac{G}{\sqrt{2}} J_\lambda \, J_\lambda^* \tag{13}$$

where :

$$J_\lambda = \mathcal{J}_\lambda + \mathcal{S}_\lambda + \sum_{l = e, \mu} (j_\lambda)_l$$

The further condition that :

$$|\Delta S| \neq 2,$$

as is required by K_1^0, K_2^0 mass difference, demands then :

$$\Delta S = \Delta Q \ \text{rule holds.}$$

Another type of interesting consequences of (13) is, e.g.,

$$e + \nu \longrightarrow e + \nu$$

These suggestions such as $\Delta Q = \Delta S$ rule, $|\Delta \underset{\sim}{I}| = 1$ rule, etc., deal with the general characteristics of $\mathcal{L}$(weak). Confirmation or negation of these proposals is, therefore, of great interest. At present, the experimental status of these rules is still unclear. There are however some new results presented yesterday on $K^0 \longrightarrow \pi^\pm + e^\mp + \bar{\nu}\ (\nu)$ reactions. Perhaps, Dr. Fry may discuss these results after this talk.

III - From a theoretical point of view the Lagrangian $\mathcal{L}$(weak) is unsatisfactory for still another reason. As is well known if we regard the $\mathcal{L}$(weak) as a bona-fide Lagrangian then the higher order terms in G are highly divergent. Since infinity makes very little sense, the rule is, therefore, changed. The $\mathcal{L}$(weak) is regarded only as a phenomenological (or effective) Lagrangian,

$$[\mathcal{L}\,(\text{weak})]_{eff}.$$

The lowest order perturbation calculation using $[\mathcal{L}\,(\mathrm{weak})]_{\mathrm{eff.}}$ gives directly the physical result ; and the higher order divergent terms are simply taken to be un-related to nature. It is easy to show that this naive rule, while being adequate at low energy region, must be modified at higher energy region. Otherwise, it would lead to results violating unitarity. For example, according to this simple rule, the S-wave cross-section for :

$$e^- + \nu \longrightarrow \mu^- + \nu \quad \text{is given by}$$

$$\sigma(e^- + \nu \longrightarrow \mu^- + \nu) = \frac{4G^2}{\pi}\, p_\nu^2 \tag{14}$$

where p_ν is the initial energy of neutrino in the C.M. system. For $p_\nu \gtrsim 300$ Gev, becomes bigger than the limit

$$\frac{1}{2}\,\pi\,\lambda\!\!\!^{-}{}^2$$

as demanded by unitarity. Therefore we expect the <u>effective</u> Lagrangian must not be a strict point interaction, say, between $(j_\lambda)_e$ and $(j_\lambda)_\mu$.

In a phenomelogical description, certain amount of non-locality must be present in this effective Lagrangian. For example, instead of $(j_\lambda)_e\,(j_\lambda)_\mu^*$ we may have in (4) :

$$[j_\lambda(x)]_e\; F_{\lambda\lambda'}(x-x')\;[j_{\lambda'}^*(x')]_\mu \tag{15}$$

where F has an extension

$$d > \sqrt{G} \qquad (\text{perhaps}, \gg \sqrt{G})$$

An immediate consequence of (15) is that in μ-decay :

$$\mu^- \longrightarrow e^- + \nu_2 + \bar{\nu}_1 \tag{16}$$

if $\nu_1 = \nu_2$ then through the charged current distribution connected with $F(x)$, the following process becomes possible :

$$\mu^- \longrightarrow e^- + \nu + \bar{\nu} + \gamma \longrightarrow e^- + \gamma \tag{17}$$

The computation of this radiative decay rate depends, of course, on the detailed assumption of this current distribution. However, the general form and order of magnitude of the branching ratio can be estimated to be :

$$\frac{\mu^- \longrightarrow e^- + \gamma}{\mu^- \longrightarrow e^- + \nu + \bar{\nu}} \approx \alpha\, f_{(d\,p_{max})} \tag{18}$$

where α = fine structure constant, d the extension of F, p_{max} the maximum momentum of the virtual ν and $\bar{\nu}$ pair in (17) and f depends on the variable $p_{max}\,d$ and the form of F. The reasonable choice of :

$$(p_{max}\, d) \sim 1$$

makes it difficult to understand why $\mu^- \longrightarrow e^- + \gamma$ appears to be highly forbidden for both real γ and virtual γ.

A more appealing suggestion is that in (16) :

$$\nu_1 \neq \nu_2 \tag{19}$$

Consequently there exists a selection rule which forbids $\mu^- \longrightarrow e^- + \gamma$. The crucial test at present would come from the high energy neutrino experiment by taking ν from, say, $\pi_{\mu 2}$ decay and trying to see if it can produce electrons upon collision with nucleon. If $\nu_1 \neq \nu_2$, then this allows a second conservation law which may be called conservation of muonic number m :

$$m = +1 \quad \text{for} \quad \mu^-, \nu_2$$
$$m = -1 \quad \text{"} \quad \mu^+, \bar{\nu}_2 \qquad (20)$$
$$m = 0 \quad \text{"} \quad \text{all others}$$
$$\text{(including, e. } \nu\text{, etc.)}$$

The algebraic sum of m is conserved in all reactions.

Clearly, any linear combination of the leptonic number l and the muonic number m is also conserved.

IV - Another crucial question concerning weak interactions is whether all observed weak reactions are actually second order processes through emissions and absorptions of a set of bosons called W. For exemple μ-decay is generated through the coupling :

$$\mu^- + \bar{\nu}_2 \rightleftarrows W^- \rightleftarrows e^- + \bar{\nu}_1$$

and β-decay through that of :

$$n + \bar{p} \rightleftarrows W^- \rightleftarrows e^- + \bar{\nu}_1$$

These bosons, if exist, must consist of at least one W^+ and one W^-. If $\Delta Q = \Delta S$ rule does not hold, then more than one kind of charged bosons might be present. It must also have the following properties :

i) Spin = 1 in order to transmit the observed vector and axial vector form of weak interactions

ii) $(\text{mass})_W \gtrsim (\text{mass})_K$ in order $K \not\rightarrow W + \gamma$

iii) W can decay in to

$$W^- \longrightarrow e^- + \bar{\nu}_1$$
$$\longrightarrow \mu^- + \bar{\nu}_2$$

and :

$$\longrightarrow \pi^- + \pi^\circ$$
$$\longrightarrow K + \pi \quad (\text{if } m_W > m_K + m_\pi) \text{ etc. with a life time} < 10^{-17} \text{ sec.}$$

$W^\pm$ can be produced by

strongly interacting particles, e.g.

$$\pi^+ + p \longrightarrow W^+ + p$$

with a cross-section $\sigma \approx 10^{-32}\,\text{cm}^2$. Recently Bernstein and Feinberg have investigated the detailed cross-section of this reaction. They found that if $m_W \approx 750$ Mev then the cross-section can be greatly enhanced by the existing 2π resonance. In such a case, the decay of $W^\pm \rightarrow \pi^\pm + \pi^\circ$ will also be enhanced.

$W^\pm$ can also be produced by leptons, e.g.

$$\nu + Z \longrightarrow W^+ + l^- + \begin{cases} Z \\ Z^* \end{cases}$$

The theoretical cross-section have now been completely calculated. It is found that if $m_W \sim m_p$ or less then the neutrino flux of the existing high energy machine could perhaps be used to produce such a particle.

With the rapid advance of experimental technique and the possibility of using high energy leptons as incident beams it is to be anticipated that in the near future answers to these questions will be known. The subject of lepton physics, and with it the entire field of weak interactions may then acquire a new dimension.

ELECTRIC STRUCTURE OF NUCLEONS

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In an age of giant accelerators, of complex experiments and of mystifying theories it is a pleasure to report on some simple experiments, made with simple equipment and having a simple interpretation - simple, that is, if one doesn't look too closely. The electron energies now available are about 1 Gev and the de Broglie wavelength of such electrons is about 0.2 fermis, hence we can expect to make out some of the details of the proton and neutron.

Last year at the Rochester Conference, as a result of scattering experiments made at Stanford University [1] with their Linac and at Cornell University [2] with our electron synchrotron, we were both able to report the beginning of a detailed structure in that for large momentum transfers the electric and magnetic form factors were no longer nearly equal - as had previously been believed to be true for energies less than about 500 Mev.

During the past year, the measurements at Stanford [3], [4] and at Cornell [5], [6] have been considerably refined and extended. As a result a remarkably clear and simple picture of the electric structure of the proton and neutron is developing. The most dominant feature of the nucleon according to these experiments is a meson cloud that is the same for the neutron and the proton except that it has a charge of $+e/2$ for the proton and $-e/2$ for the neutron. This isovector cloud seems to correlate well with the two-pion resonant state, $T = 1$, $J = 1$ that is also evident in meson experiments. Less clear is another mesonic cloud of larger radius but of smaller charge which is positive for both neutron and proton. This isoscalar cloud may be related to the three-pion resonant state, the one with $T = 0$, $J = 1$ although its size does not seem to correlate too well with the mass recently discovered for a similar state that is revealed in nucleon-antinucleon annihilation. Finally we may just be able to distinguish the charge, the size, and the magnetic moment of the central core of the nucleon strictly speaking, however, these properties of the core are just emerging from the shadows of the experimental and theoretic background. I will now discuss the scattering measurements, how they yield the form factors, and then make some remarks about the interpretation of these form factors in terms of the structure of the nucleon.

Figure 1 shows the experimental arrangement for the experiments at Cornell. My collaborators in the early measurements were Cassels, Berkelman, and Olson and in the present measurements they are Schopper, Littauer and Rouse. The electron source is the 1.3 Gev electron synchrotron. A thin target made of CH_2 or CD_2 is put directly in the beam in a straight section, where many traversals of the target can be made. The electrons being scattered at a particular angle are momentum analyzed by the single quadrupole magnet which forms a horizontal line image of the nearly point target. Along this line image is placed a long narrow scintillation counter that detects the electron : the narrow dimension of the counter together with the dimensions of the obstacle at the center of the magnet determine the momentum resolution of the magnet. A total absorption glass Cerenkov counter behind the thin scintillation counter separates electrons from protons and mesons : the electron makes a large pulse proportional to its energy but neither the proton or meson can give a very large pulse. In fact there are two such independent magnet and counter systems so that data can be taken simultaneously at two different angles. The solid angle of one of the quadrupoles, a so-called current-sheath quadrupole is 17 milli-steradians and its momentum resolution is 6 %. The solid angle of the magnets have been computed from their dimensions and in addition the magnets have been cross-calibrated by placing each of them at 90° but on opposite sides of the electron beam. The electron beam is monitored by measuring the absolute number of bremsstrahlung produced in the target. This gives just the proper product of the number of electrons times the effective target thickness that is obviously needed for calculating a cross section, however, the method is hazardous in that electrons striking anything except the target can contribute to a spurious reading.

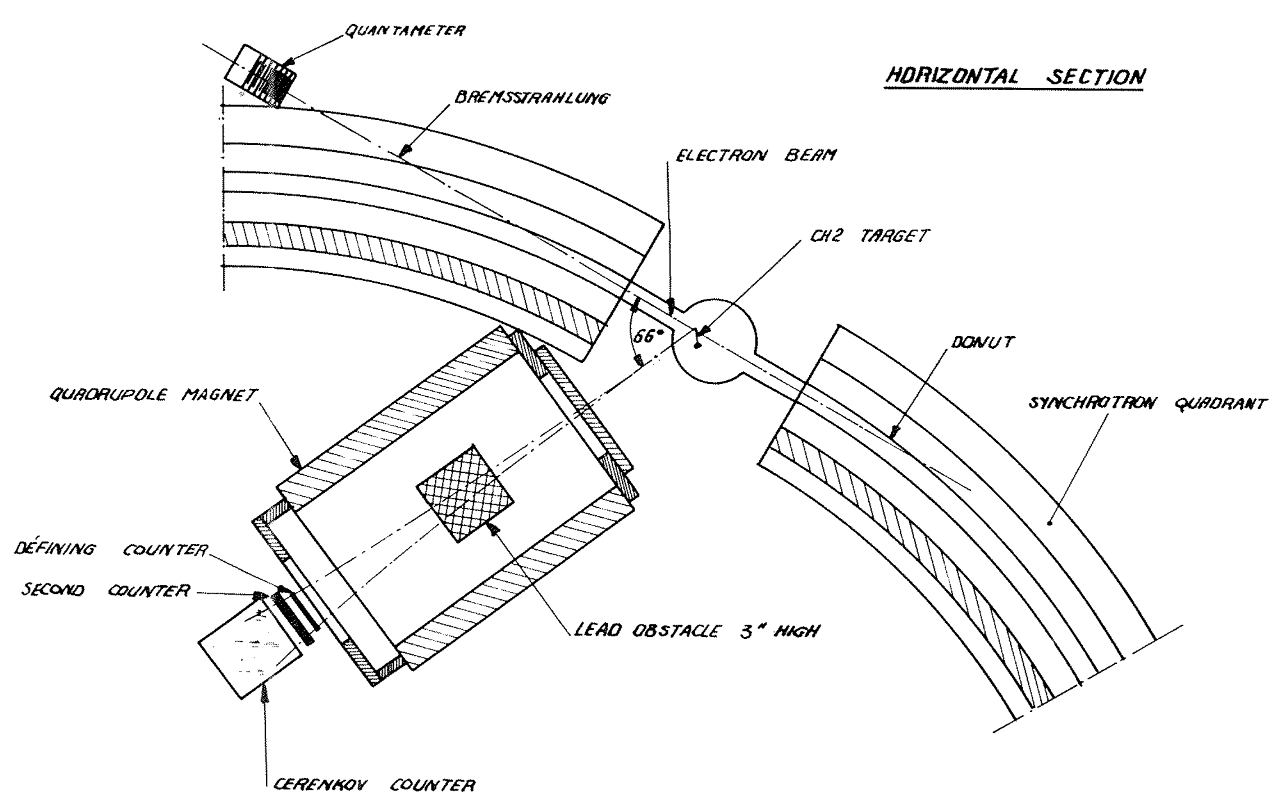

Figure 1 - Plan view of the scattering experiment.

The experiments carried out at Stanford [1], [3] at energies up to about one Gev have more of a quality of elegance. There they luxuriate in the direct beam of a linac and there they use larger magnets having better momentum resolution.

At a particular angle the procedure is to trace out a curve of counting rate v.s. magnet current or momentum of the electrons. For a target of $C\,H_2$ one typically recognizes a peak corresponding to a momentum of an elastically scattered electron. A carbon target does not show the narrow peak but does enable one to subtract the background in $C\,H_2$ due to C. Reversing the magnet allows us to examine the discrimination against mesons, something which becomes increasingly difficult at large angles and at high energies.

Figure 2a shows the Cornell measurements of R.M. Littauer, H.F. Schopper and R.R. Wilson [5] for the scattering by hydrogen of electrons at 45, 90, 112, and 135° where the cross-section is plotted against the angle of scattering. Our data do not differ appreciably from the Stanford data [3], [4] even though we do not use quite the same radiation correction. We have applied the Schwinger correction [7]- not too different from the correction of Tsai.

The determination of the cross section of scattering from deuterium is complicated by the internal motion of the nucleons in the deuteron. This causes the peak that is observed in the counting rate v.s. magnet current to be spread out, of course. Although one might think it best to measure a complete counting rate curve and then to integrate it in order to get a cross section, it turns out instead to be best to measure the value of the counting rate at the energy where the electrons scattered from hydrogen give an elastic peak. Then the total deuteron cross section is calculated from this peak counting rate using the impulse approximation as given by Goldberg [8]. Recently, Durand [9] has derived Goldberg's formula more rigorously and has shown it to be accurate to within a few percent. I have heard that at very low energy, the final state interactions may become important (See the Hofstadter paper of this meeting). The deuteron cross sections obtained using the simple impulse approximation are plotted in figure 2b - again our data are in good agreement

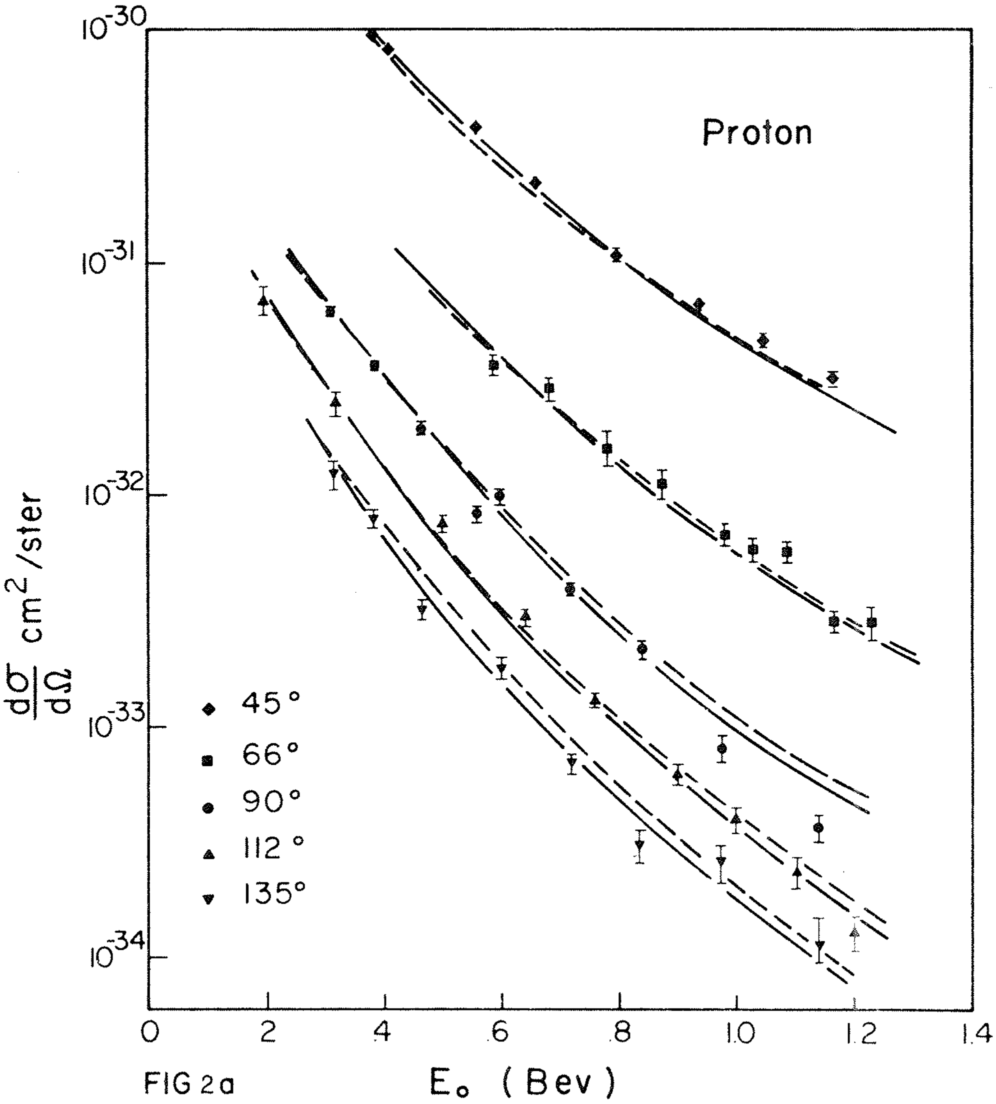

FIG 2a

Figure 2 - Differential elastic scattering cross section for the proton (a) and the deuteron (b), as a function of incident electron energy and laboratory scattering angle. Full curves computed with core model, dashed curves according to BSFV (cf. the following Letter).

with Stanford where they overlap - at low energy our data are not as accurate as the Stanford data and have not been corrected for final state interactions.

The procedure used to reduce the cross sections to form factors is straightforward. The Rosenbluth formula for the differential scattering cross section is almost as simple as Rutherford's formula ; it can be written :

$$\sigma = \sigma_M \left\{ G_1^2 + \frac{q^2}{4M^2}\left[2(G_1 + G_2)^2 \tan^2\frac{\vartheta}{2} + G_2^2 \right] \right\} \tag{1}$$

where σ_M is Mott scattering per unit solid angle from a point charge, q is the momentum-energy transfer (approximately equal to E $\sin\frac{\vartheta}{2}$), G_1 is the Dirac electric form factor normalized to a unit charge at $q^2 = 0$ and G_2 is the Pauli magnetic form factor which is normalized at $q^2 = 0$ to

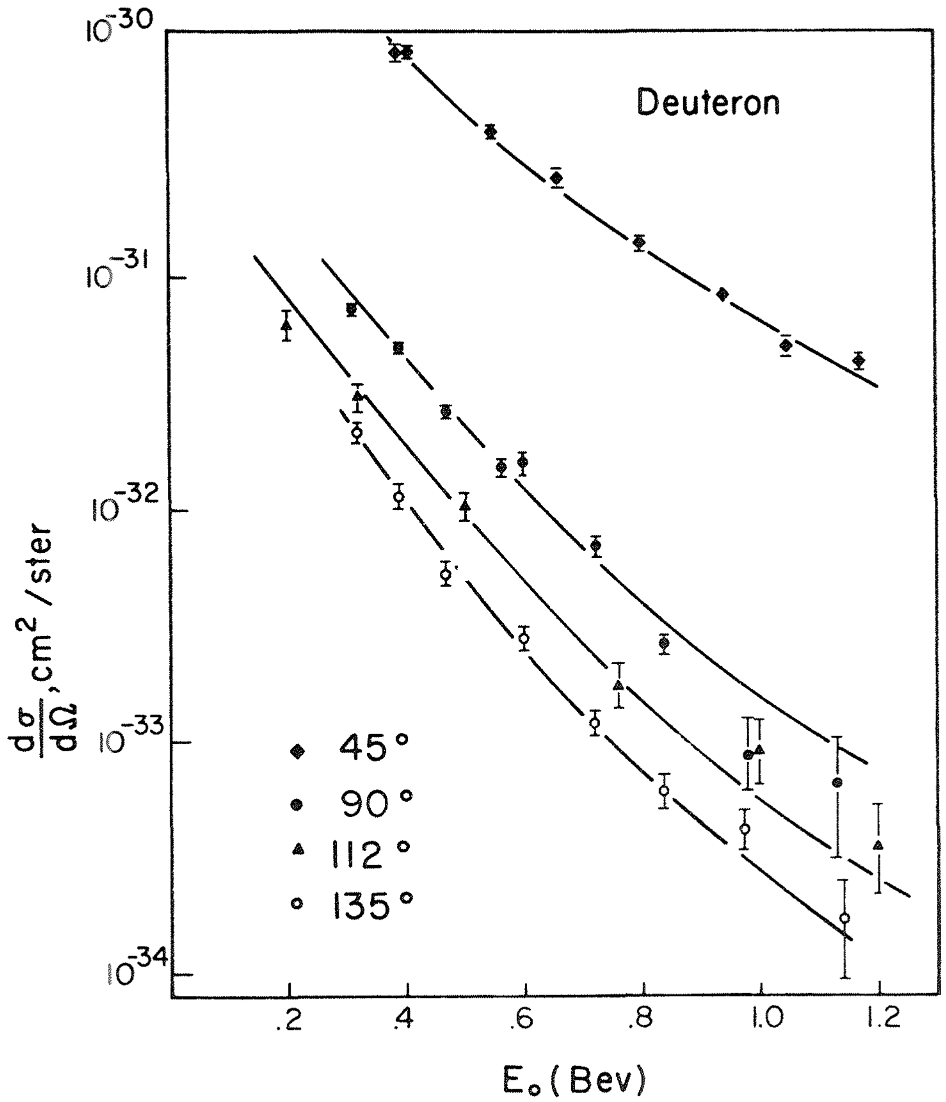

Figure 2 b

the anomalous magnetic moment in nuclear magnetons. In the static model it is just the Fourier transform of the form factors that gives the radial distribution of charge or magnetic moment in the nucleon. Theory instructs us furthermore that G_1 and G_2 should be functions only of q^2, if they are to be relativistically invariant a tremendous simplification if true.

The first objective of the experiment should be to ascertain the validity of the Rosenbluth formula. This is especially important for electron energies above one Gev : the formula in addition to assuming the validity of quantum electrodynamics also assumes the absence of fourth-order processes in which, for example, two photons are exchanged between the electron and nucleon ; this latter assumption becomes particularly suspect at high energy. Does it not seem strange that photon-proton scattering is so completely dominated by nucleonic processes such as the (3-3) resonant nucleon state and yet that these effects should be absent from electron-proton scattering ? Drell [10] assures us that such is the case, to within a few percent at least. Nevertheless, we should be anxious to test the validity of his calculation.

Let us rewrite (1) in the form :

$$\sigma/\sigma_M = \left(G_1^2 + \frac{q^2}{4M^2} G_2^2\right) + \frac{q^2}{2M} (G_1 + G_2)^2 \tan^2 \frac{\vartheta}{2} \tag{2}$$

then we see that if we measure cross sections at different angles while keeping q^2 constant by varying the electron energy, and if we plot the resulting values of σ/σ_M as a function of $\tan^2 \vartheta/2$ as has been done for some typical values in figure 3, then the points should fall on a straight line. If indeed the plot is linear, then we have at least some evidence for the validity of the formula so that we can then go on and determine values of G_1 and G_2. This kind of plot is most useful for such a determination inasmuch as the intercept at $\tan^2 \vartheta = 0$ which is equal to $\left(G_1^2 + \frac{q^2}{4 M^2} G_2^2\right)$ turns out to be approximately equal to G_1^2 because the second term in the parenthesis is always quite small. The measured slope of the line then gives the combination $(G_1 + G_2)^2$ and hence G_2.

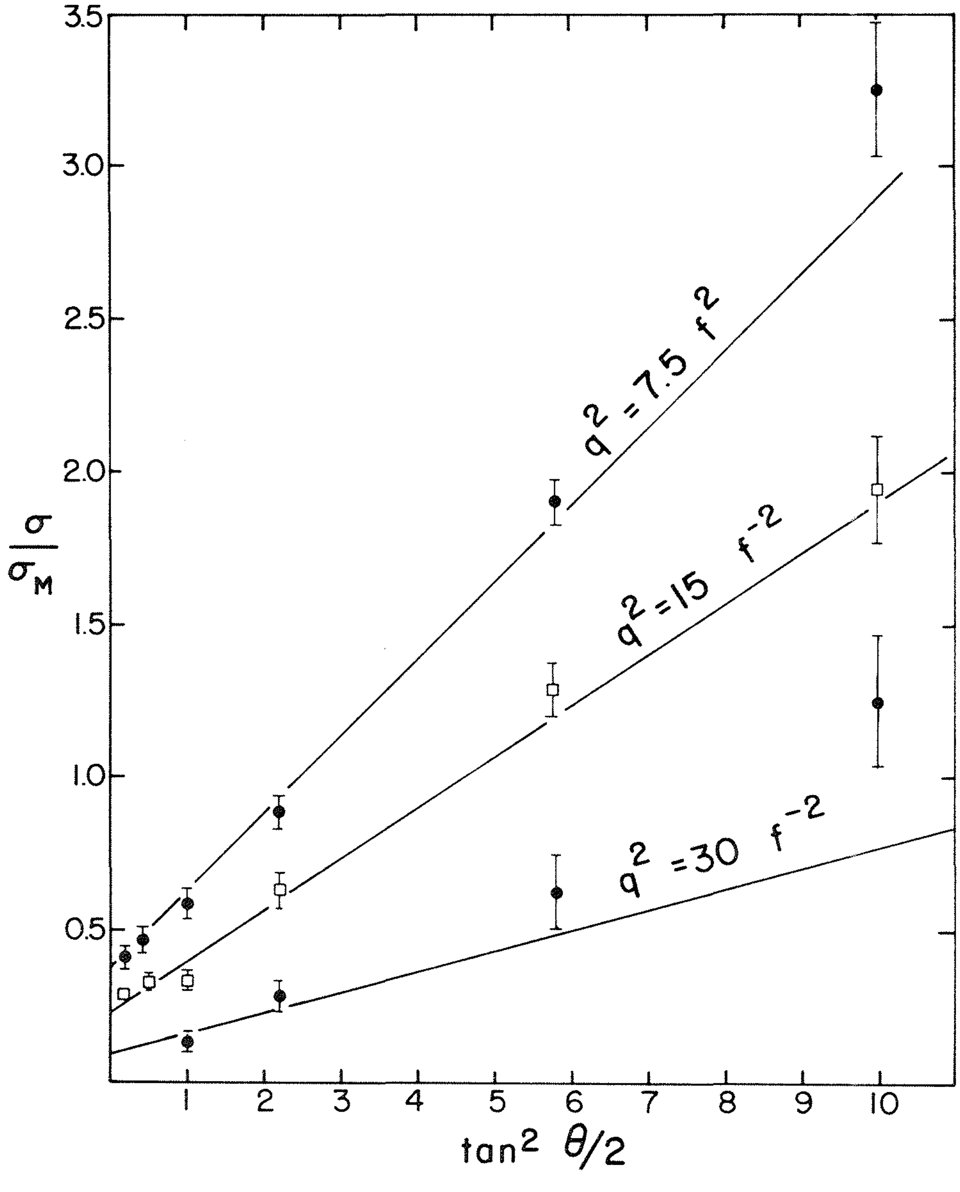

Figure 3 - The ratio σ/σ_M plotted as a function of $\tan^2 \frac{\theta}{2}$

Professor Hofstadter has introduced an equivalent method [11] in which he plots ellipses in G_1 and G_2 space that correspond to cross sections given by the Rosenbluth formula. Each experimental cross section gives one ellipse and the intersection of two ellipses give specific values of G_1 and G_2 ; a test of the formula is to notice whether more than two ellipses intersect at a point. We have also used this method in reducing our data. Generally speaking, the ellipses do intersect in a point for $q^2 < 20$ f^{-2} ; furthermore, as figure 3 shows, the experimental values of σ/σ_M do fall along straight lines. However, for $q^2 > 25$ the ellipses do not all intersect and there are inconsistencies in the straight lines. Thus for $q^2 = 30$ of figure 3, the intercept $\left(G_1^2 + \frac{q^2}{4M^2} G_2^2\right)$ is nearly zero, only possible if both G_1 and G_2 are both zero. Nevertheless, the slope of the line, proportional to $(G_1 + G_2)$, is still quite large - obviously not consistent.

This possible deviation from the Rosenbluth formula was first pointed out by the Stanford workers on the basis of their work combined with ours at somewhat higher energies. They have also found an interesting trend in their data taken at 145° which indicate a dramatic flattening-out of the cross sections above an energy of about 850 Mev. This summer, we decided to test this trend by extending their measurements, which stop a little below one Gev, to our top energy of 1.3 Gev. Our measurements are still in progress but are in accord with the Stanford work where they overlap. At our highest preliminary point at 1 120 Mev, the cross section is nearly a factor two above what one would expect from the form factors as determined at smaller angles. This is an indication, very tentative, that the Rosenbluth formula is no longer entirely valid. If this trend develops, life can be especially exciting for the physicists building high energy electron machines.

Finally, the values of G_1 and G_2 corresponding to the straight lines drawn in figure 3 are given in figure 4 as a function of q^2. The Stanford values for G_1 and G_2 are also indicated on the curve for the most recent Stanford results for the neutron see their contribution of this meeting.

Basic to our interpretation of the form factors has been the idea that the neutron and proton have structural components that are essentially the same except for a change of the sign of the charge of some particular component. We express this by analyzing the form factors into an isoscalar component which is the same for neutron and proton and an isovector component for which the sign changes. Thus we write :

$$\begin{aligned} G_{1P} &= G_{1S} + G_{1V} \quad ; \quad G_{1N} = G_{1S} - G_{1V} \\ G_{2P} &= G_{2S} + G_{2V} \quad ; \quad G_{2N} = G_{2S} - G_{2S} \end{aligned} \tag{3}$$

clearly these components can be obtained directly from the experimental values of G_1 and G_2 for example,

$$2\,G_{1S} = G_{1P} + G_{1N} \qquad 2\,G_{1V} = G_{1P} - G_{1N} \tag{4}$$

and similarly for G_{2S} and G_{2V}. Having done this we then try to interpret these components with structural details of the nucleons.

As illustrative of this approach let us assume a very simple model of the nucleon and compare it to the data. Consider one, for example, in which both the neutron and the proton have a point core of positive charge $+\frac{e}{2}$. This will give rise to an isoscalar from factor G_{1S} of value 1/2 which remains a constant as q^2 is changed. Let us also postulate that surrounding the point core is an extended meson cloud of charge + 1/2 for the proton and - 1/2 for the neutron. This gives rise to a G_{1V} term whose value at $q^2 = 0$ is 1/2 and we expect G_{1V} to decrease as q^2 increases in a manner characteristic of the particular charge distribution that we have assumed. Because the form factor can be expanded for small q^2 in the static approximation :

$$G = 1 - \frac{q^2 a^2}{6} + \dots \tag{5}$$

where a is the rms radius of the distribution, the derivative $\partial G/\partial(q^2)$ at $q^2 = 0$ should be equal to $-a^2/6$.

Comparing this model to the experimental values of G_{1S} and G_{1V} given in figure 4, we see that it fails on a number of counts. G_{1S} is not a constant of value 1/2, rather it decreases to about 0.25 by the time q^2 has reached 20 or 30 f^{-2}. The variation of G_{1V} does not seem to be inconsistent with our model and its form corresponds to an exponential charge distribution of r.m.s. radius

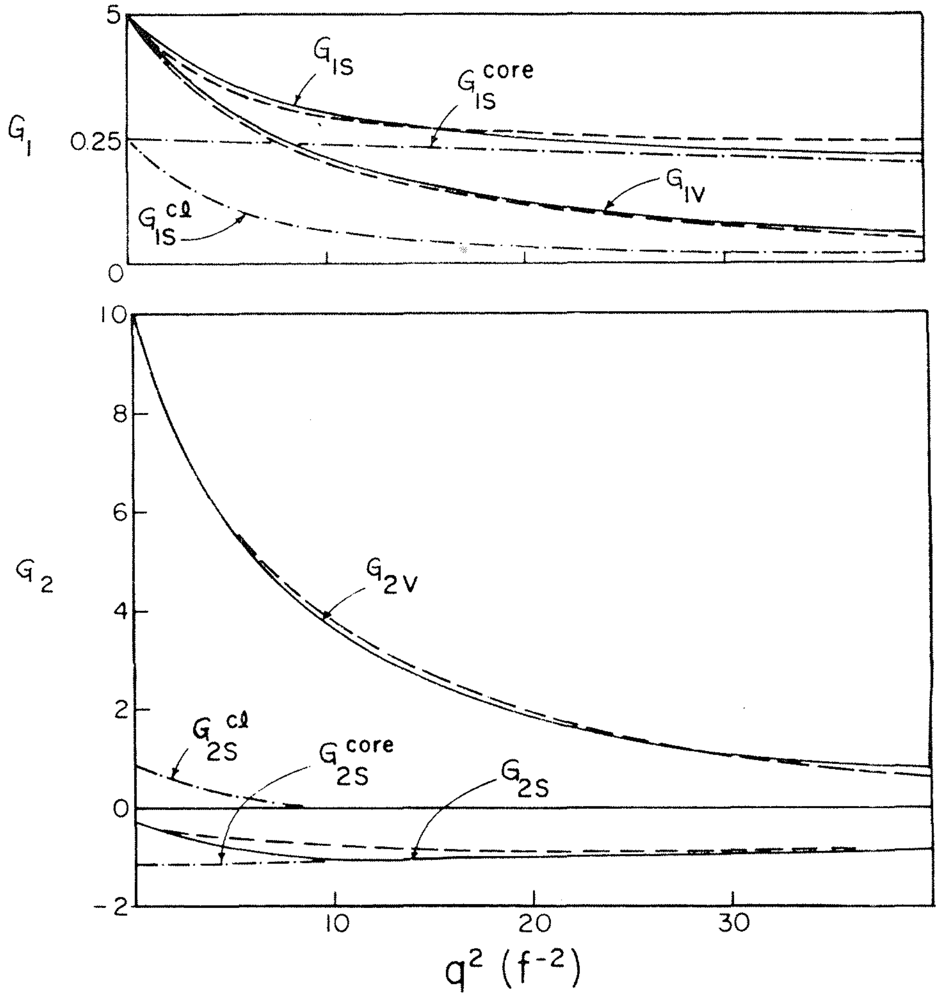

Figure 4 - Partial form factors for G_1 and G_2/μ, each form factor is resolved into isoscalar and isovector parts, whose sum and difference give, respectively, the neutron and proton form factors. The solid curves indicate the fit according to the core model, where the scalar partial form factors have been further split into terms corresponding to a core and an extended cloud, each of exponential distribution. The dashed curves indicate the best fit obtained by the Clementel-Villi form.

equal to 0.80 f. The simple model also conflicts with the result of experiments on the scattering of low energy neutrons by electrons. Foldy [11] has shown that these experiments can be interpreted to mean that the mean squared radius of the neutron is zero or, more accurately, 0 ± 0.006 f^2, whereas our model would give 0.84 f^2 for the neutron. Then, you might ask, what changes can be made in this model to bring it into agreement with these considerations. In the first place, we can assign the core a radius ; this causes the form factor to decrease at large values of q^2 just as observed. In addition to this, we can say that the isoscalar part of the form factor also has some kind of meson cloud associated with it.

With this more complicated core model we have six parameters, i.e. the partial charges and the radii of the core and of the two meson-like clouds that we have postulated. However, we have four conditions : two that are set by the charge of the neutron and proton, and two by the radii of the particles, thus the radius of the neutron is known to be zero from the neutron-electron scattering while the r.m.s. radius of the proton is given by the slope of the G_1^p curve at $q^2 = 0$. This leaves two parameters to be determined from our data. Our procedure has been to assume that the behavior of G_{1S}^p at very large q^2 is dominated by the properties of the core and so we fit this part of the curve by assigning a partial charge and a radius to the core. Then all the other parameters are determined by means of simple algebraic relations [5]. Table I gives the values of the various parameters that best fit our data, and the resulting partial form factors are plotted in figure 5. I must hasten to add, however, that we have also assumed rather arbitrarily that all the radial charge distributions are simple exponentials. Of the forms without a singularity at the

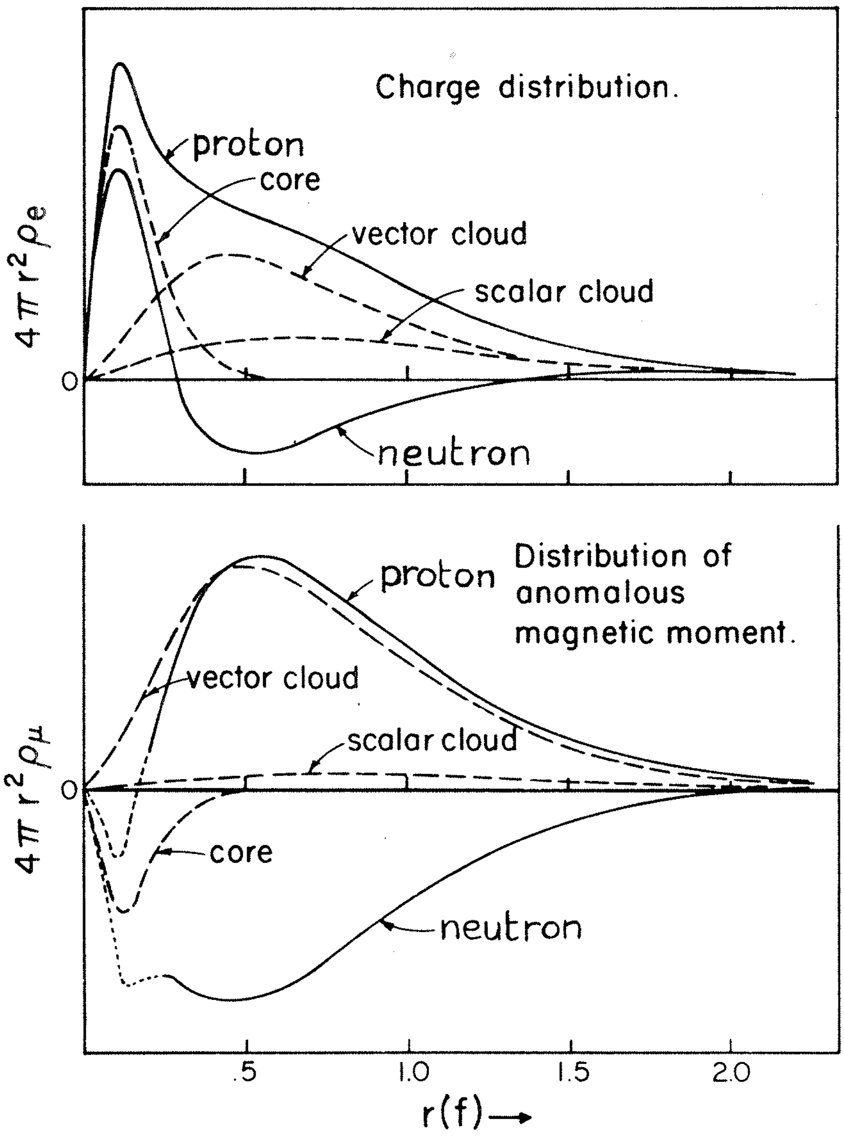

Figure 5 - Spatial distribution of charge and anomalous magnetic moment for proton and neutron, according to the core model, are shown for the sentimentalists only.

origin, we have tried gaussian distributions, but then the overall fit of the data is worse. Of course nothing at all can be said about the shape of the core, the main question here being whethe it is spread out with an r.m.s. radius of about 0.2 f or not.

Table I

Best-Fit Parameters for Core Model with Exponential Density Distributions. See Text and Reference [6] for Definition of Symbols.

e_s^{core}	= 0.25 e	$a_{s,core}$	= 0.2 f
e_s^{cl}	= 0.25 e	$a_{s,cl}$	= 1.13 f
e_v	= 0.5 e	a_v	= 0.80 f
μ_s^{core}	= -0.22 n.m.	$b_{s,core}$	undetermined
μ_s^{cl}	= 0.16 n.m.	$b_{s,cl}$	= 1.30 f
μ_v	= 1.853 n.m.	b_v	= 0.89 f
a_p	= 0.80 f	b_p	= 0.98 f
a_n	= 0	b_n	= 0.79 f

In exactly the same manner one can apply independently the same model to fit the Pauli magnetic form factors. Again there are six parameters consisting of the three partial magnetic moments and their three ranges that must be assigned to the core and the two meson clouds. But again algebraic conditions relating the anomolous magnetic moments of the neutron and proton, the values of the radii (determined by $\partial G/\partial(q^2)$ at $q^2 = 0$) leave only two parameters to fit to the experimental curve at high q^2. Again we have chosen to fit the partial magnetic moment of the core and its radius to the G_{2S} curve at large values of q^2. All the values pertaining to G_2 are also shown in Table I and figure 5.

The qualitative result of applying the core model to the magnetic form factor is that it is necessary to attribute a small magnetic moment of negative sign to the core. A point moment would give a satisfactory fit to the data as would an extended core with a radius of 0.2 f. The negative sign of this magnetic moment does not seen unreasonable : for if the angular momentum of the meson cloud is unity ; then necessarily the core will be left with a spin of - 1/2, which then might rather naturally give rise to a negative magnetic moment, for example, by dissolution of the core into a $K^+\Lambda$ system in direct analogy to the pionic origin of the Pauli moment. It is interesting that the isovector magnetic cloud has a comparable range 0.89 f to that of the isovector charge cloud, and that it gives rise to most of the magnetic moment of the nucleon, as one would expect from the near equality of the Pauli moments of the neutron and proton. The isoscalar magnetic and charge clouds, although far less well determined, also show a similar range.

The most exciting note that has been sounded on the subject of nucleonic structure has been the hypothesis of Frazer and Fulco [13] that the radial extent of the isovector cloud is due to the T = 1, J = 1 resonant state of the two pion system. Fubini will discuss his own and his collaborators more refined calculations [14] based on dispersion theory and a similar two-pion interaction ; however, for a comparison with the experimental data I will simply say that his theory yields expressions for the four partial form factors of the Clementel - Villi form :

$$G_1^{S \text{ or } V} = \frac{e}{2}\left[(1 - \alpha) + \frac{\alpha}{1 + q^2/q^2_{S \text{ or } V}}\right] \tag{6}$$

The constant α specifies the fraction of the partial charge that is in the cloud or the core and has the value α_S or α_V for G_1^S or G_1^V. The two constants q_S or q_V also refer to G_1^S or to G_1^V : q_V is the resonant energy of the two-pion state with T = 1, J = 1 ; and q_S is the resonant energy of a postulated three-pion state with T = 0, J = 1. The same form of equation applies for F_2 except that the charge e/2 is now replaced by the partial magnetic moment 1.85 n.m. for G_2^V and by -0.06 for G_2^S. The constant α in (6) becomes β_V for G_2^V and β_S for G_2^S. These six parameters are reduced to five because of the extra condition that the neutron charge radius is zero. Table II gives values of the five independent parameters that best fit our data plus the r.m.s. radii of the neutron and proton. The same parameters have also been determined at Stanford [6] and their values in our nomenclature are shown for comparison ; the general agreement is good. Slight discrepancies arise mainly because our experimental values for G_{2p} at high q^2 - values are close to zero whereas an extrapolation of the Stanford form factors gives negative values. Our experiments indicate also a somewhat smaller difference between G_{2p} and G_{2n}.

Table II

Parameters for Best Fit According to BSFV [14] Form

	Cornell	Stanford (old values)
α_v	1.10	1.20
β_v	1.14	1.20
α_s	0.58	0.56
β_s	-1.5	-3.0
a_v	0.85 f	0.77 f
a_s	1.16 f	1.13 f
a_p	0.88 f	0.85 f
b_p	0.95 f	0.94 f
b_n	0.87 f	0.76 f
q_v^2	$8.3\ f^{-2} = 16\ (m_\pi/c)^2$	$10\ f^{-2} = 19.6\ (m_\pi/c)^2$
q_s^2	$4.4\ f^{-2} = 8.5\ (m_\pi/c)^2$	$4.7\ f^{-2} = 9\ (m_\pi/c)^2$

The experimental data are fit equally well by the BSFV [14] form factors or by those following from our simple core model.

According to the interpretation of BSFV, our values of q_v and q_s would imply that the resonant energy of the two-pion state is 4 m_π and that of the three-pion state is about 3 m_π. The first value is not too different from that obtained from the meson experiments of Walker [15] et al. which give 5.5 m_π. Professor Hofstadter has already remarked in his paper that a choice of 5.3 for the isovector mesonic state and 4.5 for the isoscalar state fits the data almost as well. It is relevant to remark that as the isovector meson mass is raised, then the isoscalar mass must also be raised to fit the data. Our best values, though, are those shown in Table II.

It is reasonable at this point to become skeptical of how well the various parameters in the above analyses habe been determined or indeed if they are even unique. It is gratifying that the Stanford and Cornell measurements do give such similar results. Nevertheless, the measurements are not easy, and with regard to the Cornell measurements Johnson's famous comment about the walking dog applies : "The wonder is not that she does it well, but that she does it at all'". We must refine our work considerably before we can honestly stand behind any detailed conclusions that are to be drawn from it. Even worse, I fear, is that the basic concepts of the interpretation are subject to grave doubts and that the way abounds with dangerous pit falls. Particularly is the core model suspect. In discussing spatial distributions, we are flying exactly in the face of dire warnings from many of our theoretical friends who will consider nothing except the raw form factors. Recoil of the nucleon, they point out, has become perilously relativistic at these energies. Sachs [16] has warned of this and points out that the only quantities that have meaning if we are to make spatial models are what he calls the electric form factor G_e defined as $G_1^2 + \frac{q^2}{4\,M^2} G_2^2$ and magnetic form factor defined as $(G_1 + G_2)$. In fact he asserts that a Fourier transform of G_e and G_M will give the spatial distributions of the charge and magnetic moments. In this regard perhaps it is relevant that the Rosenbluth formula can be rewritten in terms of Sach's form factors [17] :

$$\frac{\sigma_{exp}}{\sigma_M} = \frac{\left(G_e^2 + \frac{q^2}{4\,M^2} G_M^2\right)}{\left(1 + \frac{q^2}{4\,M^2}\right)} - \frac{q^2}{2\,M^2} G_M^2 \tan^2 \frac{\vartheta}{2} \tag{6}$$

This has the same pleasing simplicity as the form involving G_1 and G_2. But Yennie, Levy, and Ravenhall tell us that it is ambiguous as to whether to use G_1 and G_2, G_e and G_M, or, worse yet, even something else as characteristic of static distributions. The theorists would be kind if they were to explain this mystery in gentle terms.

The procedure of Fubini in using a dispersion relation to calculate G_1 and G_2 directly is surely on firmer ground. I hope that he will reassure us that a different and unique prescription would be used were he to calculate instead say, G_e and G_M.

Despite this caveat, it seems evident that shining out from the form factors, no matter how they are considered, are three gross characteristics : (1) that the neutron and proton are states of a fundamental nucleon (charge independence) ; (2) the existence of meson-like clouds ; (3) the mysterious core itself. If we have only "through a glass seen darkly" these characteristics of the nucleon, still our esthetic appetite has been whetted for the revelations that future and more precise measurements will bring.

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ELECTROMAGNETIC STRUCTURE OF PIONS AND NUCLEONS

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INTRODUCTION -

Giving a report on the electromagnetic structure of pions and nucleons is at the same time very pleasant and rather dangerous. It is very pleasant since in the last year very important progress both experimental and theoretical has been achieved, on the other side many questions are still open and I have to take the risk that some of the things I am going to say might soon become out of date.

In the first part of my talk I shall recall the definitions of the form factors, their connection with different kinds of experiments and those general properties which can be rigorously deducted from theory and that will constitute the basis of the subsequent phenomenological treatment.

Part two deals with the different theoretical approximations which are necessary to extract the form factors from experiment especially in the case of unstable particles like the neutron and the pion.

Finally, in Part three I shall discuss the recent phenomenological attempts to understand the nucleon form factors and to correlate those data with the ones coming from different experiments.

In the preparation of this talk I had much help from Dr. A. Stanghellini to whom I wish to express my most sincere thanks.

I - GENERAL PROPERTIES

1 - ELECTRON SCATTERING -

Let us consider elastic electron scattering by a nucleon or a pion. Since the electron interacts only electromagnetically, the scattering matrix will depend on the exchange of one or many photons between the electron and the nucleon (pion).

The approximation of keeping only the one photon exchange term is well justified at the energies and momentum transfers reached in the present experiments, the evaluation of the two photon exchange contribution will be discussed in part II.

The one photon exchange graph is represented in figure 1 :

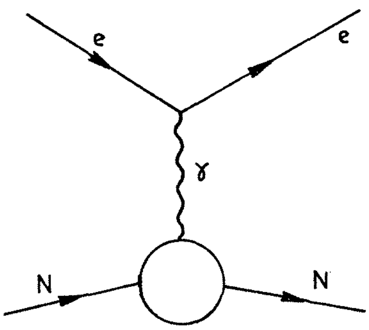

Figure 1

It gives the (eN) scattering amplitude as the product of the electron-photon vertex function, of the photon propagator and of the nucleon-photon vertex function :

$$< e_2 N_2 | T | e_1 N_1 > = \frac{(\bar{u}_2 \gamma_\mu u_1) < N_2 | J_\mu | N_1 >}{k^2} \tag{I.1}$$

where u_2 and u_1 are the initial and final electron spinors and :

$$k^2 = (e_1 - e_2)^2 = (N_2 - N_1)^2 = t \tag{I.2}$$

is the four momentum transfer between electron and nucleon.

The only unknown of the problem depending in an essential manner on the strong interactions responsible for the complicated structure of the nucleon is the photon-nucleon vertex $< N_2 | J_\mu | N_1 >$ which represents the expectation value of the electromagnetic current in the physical nucleon state.

Invariance under Lorentz transformations and under time and space inversions allows to express the photon-nucleon vertex in terms of two functions of the momentum transfer t as follows :

$$< N_2 | J_\mu | N_1 > = i \, \bar{v}_2 \, [G_1(t) \gamma_\mu + G_2(t) \sigma_{\mu\nu} k_\nu] v_1 \tag{I.3}$$

where v_1 and v_2 are the initial and final nucleon spinors.

The form factors G_i are still operators in the isotopic spin space. One defines the isovector and isoscalar form factors as follows :

$$G_i = G_i^S + G_i^V \tau_3 \tag{I.4}$$

and so the proton and neutron form factors are given by :

$$\begin{cases} G_i^P = G_i^S + G_i^V \\ G_i^N = G_i^S - G_i^V \end{cases} \tag{I.5}$$

From Eqs. (I.1) and (I.3) one obtains immediately the well-known expression for the electron-nucleon scattering cross-sections :

$$\sigma_n = \frac{\sigma_N^o}{1^2} \left[G_1^2 + \frac{t}{2 M^2} (G_1 + 2 M G_2)^2 \, tg^2 \, \frac{\vartheta}{2} + 4 M^2 G_2^2 \right] \tag{I.6}$$

where σ_N^o is the cross-section due to a point nucleon.

This is the Rosenbluth formula giving the (eN) cross-section in terms of the nucleon form factors. We wish to point out that the fact that the cross-section, which depends both on the electron energy and angle, is given in terms of two functions of one single variable, gives a very strong restriction implying, e.g. that the results of three different experiments with the same value of the momentum transfer must depend only on two numbers.

This restriction depends essentially on the fact that only one particle is exchanged between electron and nucleon. The two photon exchange corrections will depend both on the electron energy and angle and so by making several measurements at different energies but at the same t one has a way of testing experimentally the validity of the one photon exchange approximation. For the electron-pion scattering matrix element one has a formula analogous to (I.1) which will now contain the pion-photon vertex function. Using the same invariance arguments of the nucleon case one can write :

$$< \pi_1 | J_\mu | \pi_2 > = \frac{(\pi_1 + \pi_2)_\mu}{(4 \, \pi_1^o \pi_2^o)^{1/2}} G_\pi(t) \; T_3 \tag{I.7}$$

where π_1 and π_2 are the initial and final pion momenta (and π_1^o, π_2^o their time components), $G_\pi(t)$ is the electromagnetic form factor of the pion and T_3 is the third component of the isospin of the pion. Eq. (I.7) shows that the π^o has no electromagnetic structure unlike the neutron. This important difference depends on the fact that the π^o is a self-charge conjugate particle (on the other hand the K_o like the neutron would have electromagnetic structure).

The electron-pion scattering cross-section is given by :

$$\sigma_{l\pi} = \frac{\sigma_\pi^\circ}{l^2} G_\pi^2(t) \qquad (I.8)$$

where σ_π° is the cross-section due to a point pion charge distribution. At zero momentum transfer the functions $G_i(t)$ tend to the well-known limits given by the nucleon and pion total charge and of the nucleon anomalous magnetic moment :

$$\begin{aligned}
&G_1^P(O) = e, \qquad G_1^N(O) = 0, \qquad G_1^\pi(O) = e \\
&G_2^P(O) = e\,\frac{g'_P}{2M} \qquad\qquad G_2^N(O) = \frac{e g_N}{2M} \\
&G_1^S(O) = G_1^V(O) = \frac{e}{2} \\
&G_2^S(O) = \frac{e g_S}{2M} \qquad\qquad G_2^V(O) = \frac{e g_V}{2M}
\end{aligned} \qquad (I.9)$$

One often finds in the literature the form factors $F_v(t)$ normalized as follows :

$$\begin{aligned}
G_2(t) &= G_2(O)\ F_2(t) \\
G_1^P(t) &= G_1(O)\ F_1^P(t) \\
G_1^N(t) &= e\ F_1^N(t)
\end{aligned} \qquad (I.10)$$

2 - THE FORM FACTORS FOR TIME-LIKE MOMENTUM TRANSFERS -

Until now we have considered the determination of the form factors from electron scattering. In this case the photon four momentum k, being the difference of the initial and final electron momenta, will always be space-like and thus (•) $t < 0$. The possibility of determining experimentally the form factors for positive values of t is given by the reactions :

$$e^- + e^+ \begin{cases} \nearrow \pi^- + \pi^+ \\ \searrow N + \bar{N} \end{cases}$$

The many interesting possibilities given by (e^-e^+) reactions have been studied in detail by R. Gatto and N. Cabibbo.

In the case of the e^-e^+ reactions the virtual photon momentum is the <u>sum</u> of the electron and positron momenta and will therefore be time-like, the momentum transfer t will now be positive and larger than $4m_\pi^2$, $4M_N^2$ for the $(\pi^-\pi^+)$ and $(N\bar{N})$ cases.

In the one photon exchange approximation these reactions will be represented by the same graphs as in figure 1 but viewed from the reverse side. Now our matrix elements will depend on the $\langle O|J_\mu|\pi^-\pi^+\rangle$, $\langle O|J_\mu|N\bar{N}\rangle$ vertex functions. Expressions for them can be obtained by applying trivial changes to Eqs. (I.3), (I.7).

$$\langle O|J_\mu|\pi^-\pi^+\rangle = \frac{(\pi_- - \pi_+)_\mu}{(4\,\pi_-^\circ\,\pi_+^\circ)^{\frac{1}{2}}} G_\pi(t) \qquad (I.11)$$

$$\left.\begin{array}{l} \langle O|J_\mu|P\bar{P}\rangle \\ \langle O|J_\mu|n\bar{n}\rangle \end{array}\right\} = i\,(\bar{v}_{\bar{N}}[G_1^{(P,n)}(t)\gamma_\mu + G_2^{(P,n)}(t)\ \sigma_{\mu\nu} k_\nu]v_N) \qquad (I.12)$$

where now $k = e_+ + e_-$. The cross-sections for the $(e^- + e^+ \longrightarrow \pi^- + \pi^+)$ reaction is given by (see Cabibbo and Gatto) :

(•) We use a metric for which $t = k_o^2 - k^2$.

$$\frac{d\sigma_{\pi^-\pi^+}}{d\cos\vartheta} = \frac{\pi}{16}\,\frac{\beta^3}{E^2}\left|G_\pi(t)\right|^2 \sin^2\vartheta \qquad (I.13)$$

where $E = \sqrt{\frac{t}{4}}$ $\beta = \sqrt{\frac{t-4}{t}}$.

There is an important difference between the form factors as defined in the space-like or time-like region : the form factors appearing in electron scattering are always real (both positive and negative) whereas in the time-like region the form factors will in general be complex, the imaginary part depending on the interaction between the two particles in the final state. More precisely, using invariance under time reversal, one can obtain for the imaginary part of the pion vertex the following expression :

$$\mathrm{Im} < 0 \mid J_\mu \mid \pi^-\pi^+ > = \Sigma_\alpha < 0 \mid J_\mu \mid \alpha >< \alpha \mid T^+ \mid \pi^-\pi^+ > \qquad (I.14)$$

where α are all states with nucleon number zero and total $J = 1$ (since the current is a vector operator) and T^+ is the scattering matrix. An analogous expression can be obtained for the $N\bar{N}$ final state.

Eq. (I.14) means that for a given energy the imaginary part of the electromagnetic form factor of the pion is large only when there is a large pion-pion interaction for $J = 1$ at that energy. For values of $t < 16\ m_\pi^2$; below the threshold for production of 4 pions Eq. (I.14) takes a particularly simple form :

$$\mathrm{Im}\ F_\pi(t) = F_\pi(t)\ e^{-i\delta_1} \sin\delta_1 \qquad (I.15)$$

This means that $F_\pi(t)$ will be given by a real quantity multiplied by $e^{i\delta_1}$ or equivalently :

$$\mathrm{Im}\ F_\pi(t) = \mathrm{Re}\ F_\pi(t)\ \mathrm{tg}\,\delta_1 \qquad (I.15')$$

where δ_1 is the $T = J = 1$ pion-pion phase shift.

An equation analogous to Eq. (I.14) can be written in the nucleon case. Its direct use is not so simple and illuminating since in $N\bar{N}$ annihilation kinematics allows the production of states from 2 to 13 pions ! This equation however will be very useful when analytically continued for values of $t \ll 4\ M^2$. This will be discussed in detail in connection with the dispersion relations. The use of the final state theorem has shown the connection of the form factors in the time-like region with pion-pion interaction. I wish here to recall a very important symmetry property of the pion-pion system which will be useful in the study of form factors.

It is well-known that the π° is a self charge conjugate particle so that :

$$C\,\pi^\circ = \pi^\circ$$

If we introduce the real pion fields π_1 and π_2 related to π^+ and π^- by :

$$\pi^+ = \pi_1 + i\pi_2$$
$$\pi^- = \pi_1 - i\pi_2$$

we have :

$$C\pi^\circ = \pi^\circ, \qquad C\,\pi_1 = \pi_1, \qquad C\,\pi_2 = -\pi_2 \qquad (I.16)$$

We are looking for a transformation under which all three kinds of pions transform in the same manner. Let us consider the charge symmetry operator :

$$S = e^{i\pi\,\tau_2} \qquad (I.17)$$

representing a rotation of 180° around the y axis in the isotopic spin space. We have :

$$S\pi^\circ = -\pi^\circ, \qquad S\pi_1 = -\pi_1, \qquad S\,\pi_2 = \pi_2 \qquad (I.18)$$

So that if one considers the product of both transformations ;

$$G = CS \qquad (I.19)$$

one has for all three kinds of pions :

$$G\pi_i = -\pi_i \tag{I.20}$$

The pion state is an eigenstate of G corresponding to eigenvalue -1 (*). Invariance under G tells us that an initial state with an even (odd) number of pions must always lead to an even (odd) number of pions in the final state.

Let us now apply G invariance to electromagnetic transitions. On the basis of the behaviour in isospin space, the electromagnetic current operator can be separated into two parts :

$$J_\mu = J_\mu^S + J_\mu^V \tag{I.21}$$

where J^S behaves like an isoscalar and J^V like the third component of an isovector. Since J_μ is odd under charge conjugation we have the following properties under G transformation.

$$\begin{aligned} G\, J_\mu^S G^{-1} &= -J_\mu^S \\ G\, J_\mu^V G^{-1} &= -J_\mu^V \end{aligned} \tag{I.22}$$

Eq. (I.21) shows that the isoscalar current must always be coupled to an odd number of pions and the isovector current to an even number. This explains for example why the current appearing in the pion form factor is only <u>isovector</u> (recall the operator T_3 in the expression (I.7) for the $(\pi\pi\gamma)$ vertex function).

3 - DISPERSION RELATIONS FOR THE FORM FACTORS -

I wish now to discuss the connection between the form factors defined in the space-like and in the time-like regions. The general principles of quantum field theory tell us that since those form factors depend on the same $(NN\gamma)$, $(\pi\pi\gamma)$ vertex functions, the two kinds of reactions $e\pi \longrightarrow e\pi$ and $e + \bar{e} \longrightarrow \pi + \pi$ lead to the values of the <u>same</u> functions in the intervals $-\infty < t < 0$, $4\,m^2 < t < +\infty$ respectively. Since the regions defined by the two kinds of experiments are separated by the gap $(0 - 4\,m^2)$ the statement that we have to deal with the same function is still rather academic unless we do not give a well-defined procedure to connect directly the data we can obtain in both regions.

Such a procedure is given by the use of dispersion relations which allow to continue analytically from one region to the other.

Let us first consider the simpler case : the pion form factor. It is possible to prove, starting from the general principles of quantum field theory, that G(t) is an analytic function of t whose only singularity is a cut on the real axis from $4\,m^2$ to $+\infty$. Moreover perturbation theory suggests that for $t \longrightarrow \infty$ G(t) does not behave worst than a constant. Thus applying the Cauchy theorem to the function $\frac{G(t)}{t}$ (the denominator t is to ensure convergence at infinity) to the contour drawn in figure 2 we obtain for G(t) the following representation valid for any complex value of t.

$$\frac{G_\pi(t)}{t} = \frac{G_\pi(O)}{t} + \frac{1}{\pi}\int_{4\,m_\pi^2}^{\infty} \frac{g_\pi(t')}{t'\,(t' - t)}\, dt'$$

where the real function g(t') represents the discontinuity of G(t) along the cut.

Using the low energy limit given by Eq. (I.9) we obtain :

$$G_\pi(t) = e + \frac{t}{\pi}\int_{4\,m_\pi^2}^{\infty} \frac{g_\pi(t')}{t'\,(t' - t)}\, dt' \tag{I.23}$$

If t is on the positive real axis on the integration path, the physical interpretation of Eq. (I.23) is well-known : by taking $t = t_0 \pm i\eta$ where η is a small positive number, one obtains $G_\pi(t)$ and $G_\pi^*(t)$ respectively.

(*) For the nucleon the situation is of course much more complicated : under C $P \longrightarrow \bar{P}$ and under S $P \longrightarrow N$ so that G transforms $P \longrightarrow \bar{N}$.

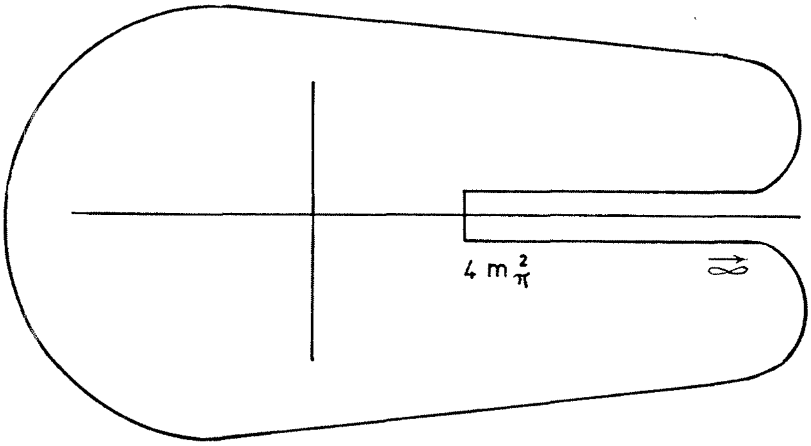

Figure 2

This leads immediately to :

$$g_\pi(t) = \frac{G_\pi(t) - G_\pi^*(t)}{2\,i} = \mathrm{Im}\; G_\pi(t) \qquad (I.24)$$

Eqs. (I.23) and (I.24) give the relation between form factors for space-like and time-like momentum transfers.

The spectral function $g_\pi(t) = \mathrm{Im}\; G_\pi(t)$ can in principle be obtained starting from (e^+e^-) experiments, its knowledge allows to obtain an unambiguous prediction on the form factors for space-like t. Recalling the expressions (I.14) and (I.15) for g (t) given by the final state theorem one can see the close connection between the electromagnetic form factors of the pion and pion-pion interaction.

If low energy pion-pion interactions were unimportant one would expect $g_\pi(t')$ to be small for low t'. Because of the double denominator in Eq. (I.23) this would cause the dispersion integral to be small. In other words, with a small pion-pion interaction we would ewpect $G_\pi(t) \approx e$ for $|t| < M^2$. On the other hand, it will be shown in the next section that the recently discovered $T = J = 1$ $\pi - \pi$ resonance causes a rapid variation of the form factors for low values of t. In the case of the nucleon form factors the situation is more complicated. Let us first of all write down explicitly the expression for the imaginary analogous to Eq. (I.14) fiven by the final state theorem :

$$\mathrm{Im} < 0\,|\,J_\mu\,|\,N\bar{N} > \; = \; \Sigma_\alpha < 0\,|\,J_\mu\,|\,\alpha > < \alpha\,|\,T^+|N\bar{N} > \qquad (I.25)$$

Eq. (I.25) has only a direct physical meaning for $t \geqslant 4\,M^2$ which is the minimum mass for a physical $N\bar{N}$ state. However, the lightest state which has the right quantum number to contribute to the sum in Eq. (I.25) is the two pion state with a minimum mass $4\,m^2$.

This fact strongly suggests that the imaginary parts of the form factors are actually different from zero for $t > 4\,m^2$. This is well illustrated in figure 3 which shows the graphs in which two pions are exchanged between the $N\bar{N}$ system and the photon.

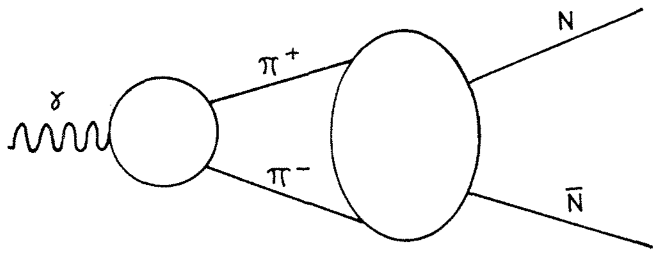

Figure 3

The two blubs in figure 3 represent the $<0|J_\mu|\pi^+\pi^->$ and $<\pi^+\pi^-|T^+|N\bar{N}>$ matrix elements in Eq. (I.25).

If we now try to write down dispersion relations analogous to Eq. (I.23) one expects the integration on t' to start at $4\,m^2$ not at $4\,M^2$. One indeed obtains :

$$G_v^{1,2}(t) = G_v^{1,2}(O) + \frac{t}{\pi}\int_{4m_\pi^2}^{\infty} \frac{g_v^{1,2}(t')}{t'(t'-t)}\,dt' \tag{I.26}$$

$$G_s^{1,2}(t) = G_s^{1,2}(O) + \frac{t}{\pi}\int_{9m_\pi^2}^{\infty} \frac{g_s^{1,2}(t')}{t'(t'-t)}\,dt' \tag{I.27}$$

The lowest limits in Eqs. (26) and (27) are $4\,m^2$ and $9\,m^2$ respectively. This is due to the fact that, as discussed in Section 2., the isovector current is coupled to an even number of pions (and so can lead to the 2π state), and the isoscalar current is coupled to an odd number (and so the lowest mass state is the three pion state).

The dispersion relations for the nucleon form factors have therefore the big difficulty that the spectral functions $g_v(t)$ $g_s(t)$ are not directly connected with experiment in the interval $t < 4\,m^2$. In this same interval (called "unphysical region") the unitarity relation (I.25) cannot be used unless one does not obtain a theoretical prescription which gives it a meaning and which allows to compute matrix elements of the type $<\pi^+\pi^-|N\bar{N}>$ for $t < 4\,M^2$. This is indeed possible by means of the Mandelstam representation and will be discussed in the next section.

The appearance of such large unphysical regions has not yet allowed a derivation of Eqs. (I.26) and (I.27) from quantum field theory. However, one has verified their validity to all orders of perturbation theory.

II - THE DERIVATION OF THE FORM FACTORS FROM EXPERIMENT

Before explaining the specific models which have been recently proposed for the form factors, I wish to discuss the methods and theoretical approximations which are necessary in order to extract the different form factors from experiment.

First of all one has to consider the limits of validity of the simple one photon exchange approximation. Secondly, since the neutron and pion are unstable particles the determination of their form factors cannot be made directly by electron scattering but one has to use more complicated phenomena like (eD) interactions and electron production of pions so that we need some kind of a theory for these phenomena.

1 - TWO PHOTON EXCHANGE CONTRIBUTIONS -

The graph giving the two photon exchange contribution to electron scattering is shown in figure 4 where the lower blub represents Compton effect of virtual photons on the nucleon.

Of course, the graph representing two photon exchange is of order e^4 in the electromagnetic coupling constant whereas the leading one photon term is of order e^2. So the first correction to the Rosenbluth formula for e scattering comes from the interference between the two photon and the one photon contributions to the scattering matrix.

Let us now consider in some detail the graph in figure 4 : unlike the one photon term which contains only real form factors the two photon term will have both a real and an imaginary part.

The imaginary part will be related through unitarity to inelastic electron nucleon scattering :

$$\mathrm{Im} <e'N'|T|eN> = \Sigma <e'N'|T^+|e''N''\ \text{pions}> \times <e''N''\ \text{pions}|T|eN> \tag{II.1}$$

This means that the imaginary part will have a large enhancement factor coming from resonant pion-nucleon intermediate states in Eq. (II.1). On the other hand, the real part of the amplitude will be related to the imaginary part by a dispersion relation : it will therefore change of sign passing through zero at the energies for which the different resonances are produced.

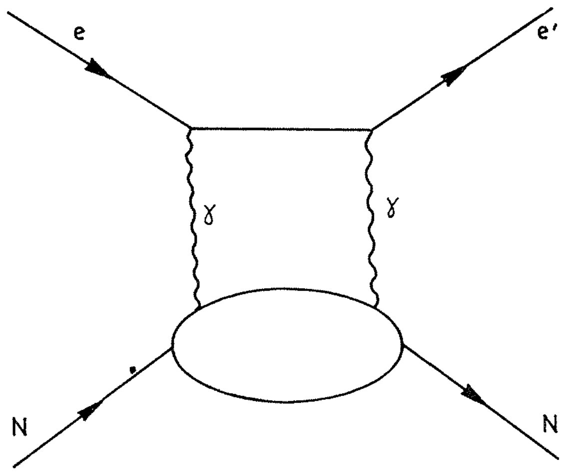

Figure 4

Now we have already seen that the first correction to the Rosenbluth cross-section comes from interference between e^2 and e^4 terms. The e^2 term being real, it does only interfere with the real part of the e^4 term which is not enhanced by any resonance.

This effect together with the $\frac{1}{137}$ factor make the corrections to the Rosenbluth formula less than 1 % in the region of experimental interest.

2 - THE DEUTERON AS A SOURCE OF INFORMATION ABOUT THE NEUTRON -

Since neutron targets are not easily available the main source of information about the neutron comes from the use of deuteron targets which due to the small binding energy B and to the corresponding large spread of the wave function is a very good source of quasi-free neutrons.

The first reactions one looks at is :

$$e + D \longrightarrow P + N + e \qquad \text{(II.2)}$$

or equivalently :

$$(\gamma) + D \longrightarrow P + N \qquad \text{(II.2')}$$

where (γ) indicates the virtual photon exchanged between the electron and the deuteron. The experiments are usually carried out by looking only at the spectrum of the recoil electron for a fixed value of the scattering angle. This means, according to (II.2'), that one determines experimentally the four momentum k_o, $\vec{k}$ (in the lab. system) of the virtual photon but one does not observe the P and N in the "final" state.

The process in which we are particularly interested is the one in which one of the nucleons (for example the neutron) interacts directly with the virtual photon whereas the other particle plays the role of a spectator. Therefore, one has to select the appropriate kinematical circumstance to reveal such a quasi-free particle scattering. This is obtained by considering that the virtual nucleon in the deuteron (N) can be considered as a particle having total energy M-B and narrow momentum distribution of width $\sqrt{MB}$. One has in the continuum a quasi-elastic peak corresponding to the reaction :

$$(\gamma) + (N) \longrightarrow N$$

The position of the peak will be given by :

$$(M - B + k_o)^2 - \vec{k}^2 = M^2$$

or equivalently :

$$k_o = B - \frac{t}{2M} \tag{II.3}$$

with a width of the order of $\sqrt{MB}$. The electron deuteron cross-section on the peak can be simply determined by means of impulse approximation. Indeed, for large enough values of t, the wavelength of the virtual photon is small as compared with the deuteron size and thus one can write the total cross-section as the incoherent sum of a proton and a neutron contribution. One obtains :

$$\frac{d^2\sigma}{d\Omega_e dE'_e} = C(\sigma_p + \sigma_n) \tag{II.4}$$

where C is a coefficient depending on the kinematical variables and on the deuteron wave function and σ_p and σ_N are the free proton and neutron cross-sections given by the corresponding Rosenbluth formulae. Eq. (II.4) has been first obtained by Goldberg and is now the basis of the derivations of the neutron cross-section from experiment.

However, the Goldberg derivation makes use of non-relativistic quantum mechanics, in this framework the distinction between three dimensional and four dimensional momentum transfers and between electric and magnetic effects is rather ambiguous. Thus one had some doubts on the reliability of Eq. (II.4) which only recently has become the basis of the determination of the neutron form factor.

Therefore, it seems instructive to discuss the interpretation, due to Durand, of Eq. (II.5) using relativistic Feynman graphs.

The graphs shown in figure 5 represent the contribution from the quasi-free proton and neutron to the deuteron cross-section :

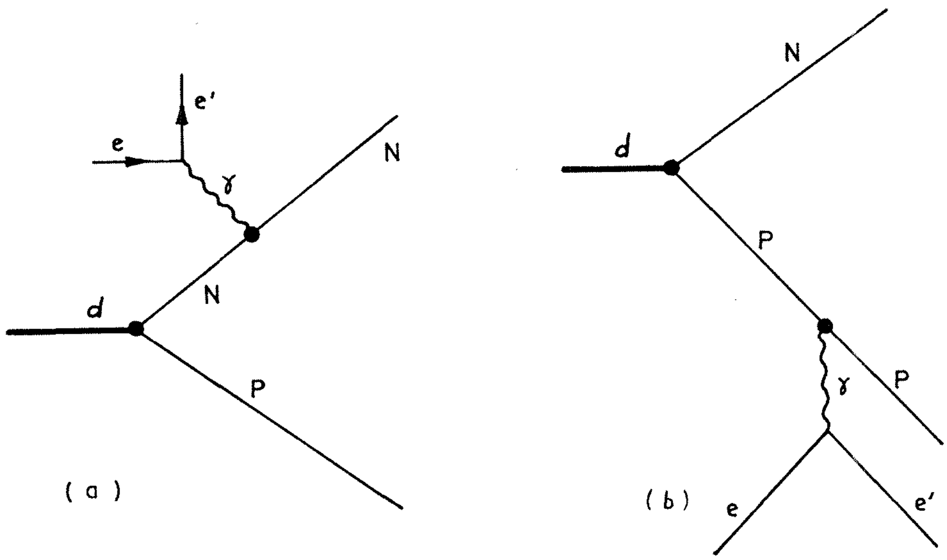

Figure 5

The complete expressions of the matrix elements corresponding to graphs (a) and (b) are given in the paper by Durand. They contain the product of three terms :

1/ the (DPN) vertex function whose size depends on the asymptotic normalization constant of the deuteron wave function N,

2/ the nucleon propagator containing denominators :

$$D_n = (d - p)^2 - M^2, \qquad D_p = (d - n)^2 - M^2 \tag{II.5}$$

for graphs (a) and (b) respectively (d p n are the deuteron and the final proton and neutron four momenta).

3/ the (γPP) and (γNN) vertex functions containing the usual combination of the G_1 G_2 form factors.

The appearance of the two denominators D_n and D_p in the expression for the matrix elements is of the utmost importance and determines the main features of quasi-elastic scattering. Indeed, in the laboratory system we can simply write :

$$\frac{1}{2\,M_d}\,D_{p,n} = -\left(\frac{B}{2} + T_{n,p}\right) \tag{II.6}$$

where Tp,n are the kinetic energy of the final nucleons related by energy conservation.

$$T_p + T_n = k_o - B \tag{II.7}$$

This equation shows that the two matrix elements will vary very rapidly with T_p, T_n and become very large in narrow regions of the phase space around $T_p \sim 0$. In other words : scattering by a quasi-free nucleon will be most important when the other "spectator" nucleon comes out with a small recoil.

The above discussion shows that for k_o sufficiently large the interference effect between proton and neutron is quite small since the regions of importance of the two graphs are very different and well separated. Therefore, the deuteron cross-section is simply given in terms of proton and neutron cross-sections :

$$\frac{\partial^3\sigma}{\partial k\,\partial k_o\,\partial T_p} = \left[\frac{\sigma_n}{\left(T_p + \frac{B}{2}\right)^2} + \frac{\sigma_p}{\left(T_n + \frac{B}{2}\right)^2}\right] \times \text{const} \tag{II.8}$$

Equation (II.8) shows that the best procedure to measure the neutron form factors is to select these events corresponding to small proton recoils. This has not yet been done and one has, for the moment, to rely on the comparison with experiment of the cross-section integrated on all final nucleon states.

The integration will, of course, be made between T_{min} and T_{max} (functions of t and k_o) which are the maximum and minimum values of the nucleon kinetic energy allowed by the kinematics of the problem.

Now Eq.(II.8) shows that quasi-elastic scattering is important only in a limited range of values of T_p i.e., for $T_p < B$ and $T_N < B$. This means that we shall have a narrow maximum in the cross-section when k_o and t are related in such a manner that the range of integration extends between 0 and $k_o - \frac{B}{2}$ and therefore covers completely the two important regions $T_p < B$ $T_N < B$. This maximum corresponds to the position of the quasi-elastic peak in the spectrum of the recoil electron. The peak relation between k_o and t is immediately given by energy momentum conservation :

$$k_o - B = \sqrt{M^2 + \vec{k}^2} - M$$

which coincides with Eq. (II.3).

Now, integrating Eq. (II.8) between 0 and k_o - B one obtains the expression of the peak cross-section which coincides with the Goldberg formula (*).

The foregoing has made it clear that the range of applicability of the simple Goldberg formula is limited to large (> 15 μ^2) values of t. For smaller values of t the approximation of summing the free proton and neutron effects becomes less accurate and may introduce considerable errors in the determination of nucleon form factors. In particular one important correction comes from the effect of final state P N interaction. A very elegant method of dealing with those corrections has been recently developped by B. Bosco. Consider the transition matrix element M(E) to a final P N state

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(*) Actually the relativistic derivation of the constant C appearing in Goldberg formula is made by multiplying the nucleon propagator by the damping factor : $\frac{T_o}{T + T_o}$ which simulates the effect of the inner part of the deuteron wave function given by Hulthen model.

with a well-defined angular momentum J and energy E and let us call $M_0(E)$ the same matrix element computed without final state corrections.

Bosco's analysis is based on the following properties :

1/ M(E) satisfies the final state theorem $\mathrm{Im}\, M(E) = \mathrm{Re}\, M(E)\, \mathrm{tg}\,\delta$ (δ being the nucleon phase shift),

2/ M(E) has simple analytic properties deduced for a general class of potential models,

3/ for large values of E $M(E) \longrightarrow M_0(E)$.

These properties allow to obtain a linear integral equation whose solution is an expression of the matrix element M(E) depending directly on the experimental NN phase shifts. The final state corrections turn out to be rather large (around 15 %) and give very important changes in the determination of the form factors for small t.

The result of Bosco shows very clearly that at small values of t simple models are not adequate to describe the whole situation. It is very important in the future to consider very carefully other possible corrections like the ones due to interference effects.

Finally, I wish to mention elastic eD scattering as a source of information on the coherent combination of proton and neutron amplitude.

The interest in elastic deuteron scattering lies in the fact that only the isoscalar form factors appear. The difficulty of theoretical interpretation of these data is due to the fact that quasi-free nucleon effects are summed coherently with other effects (for example of meson currents) which are quite difficult to evaluate.

3 - ELECTRON PRODUCTION OF PIONS -

A source of information on the neutron and electron form factors is given by inelastic electron nucleon scattering, in particular with the production of one pion.

The study of electron-production of pions as a method to obtain the electromagnetic form factors has been started by Fubini, Nambu and Wataghin using the dispersion method.

However, the techniques used in this first investigation are rather crude (since they were based on a $\frac{1}{M}$ expansion) and allow only a first determination of the isovector form factor of the nucleon.

An improved calculation of electron production has been recently carried out by P. Dennery. This calculation is completely covariant and makes use of the Cini-Fubini approximation to the Mandelstam representation. The result of Dennery can be simply illustrated by means of the following relativistic graphs.

The entire electroproduction amplitude is given by the sum of the graphs drawn in figure 6. Let us discuss separately each contribution :

a) the graphs represent the usual nucleon pole terms : they contain the $(NN\gamma)$ vertex function which is given in terms of the four nucleon form factors,

b) the graph represents the direct interaction of the photon with the virtual pion. It contains the $(\pi\pi\gamma)$ vertex function i.e., the electromagnetic form factor of the pion,

c) the graphs represent the effect of the excitation of the (33) resonant state N^*. They contain the $(N^*N\gamma)$ vertex function. This vertex function can be analyzed (using the general invariance arguments) in terms of three new form factors G_{N^*}.

Of course, the excitation of the $\frac{3}{2}$ N^* state requires that the photon behaves as an isovector.

So the first step in Dennery's work is to express the whole amplitude in terms of the five pion and nucleon form factors and of the three G_{N^*} related to the excitation of the isobaric state. The second step is to note that, since the (33) resonant state is composed of pion and nucleon, the G_{N^*} will be linear combinations of G_π and G_N in the same manner as the deuteron form factor is a linear combination of the proton and neutron ones. This determination can be carried out by means of the dispersion method by imposing the consistency of the dispersion relations with unitarity.

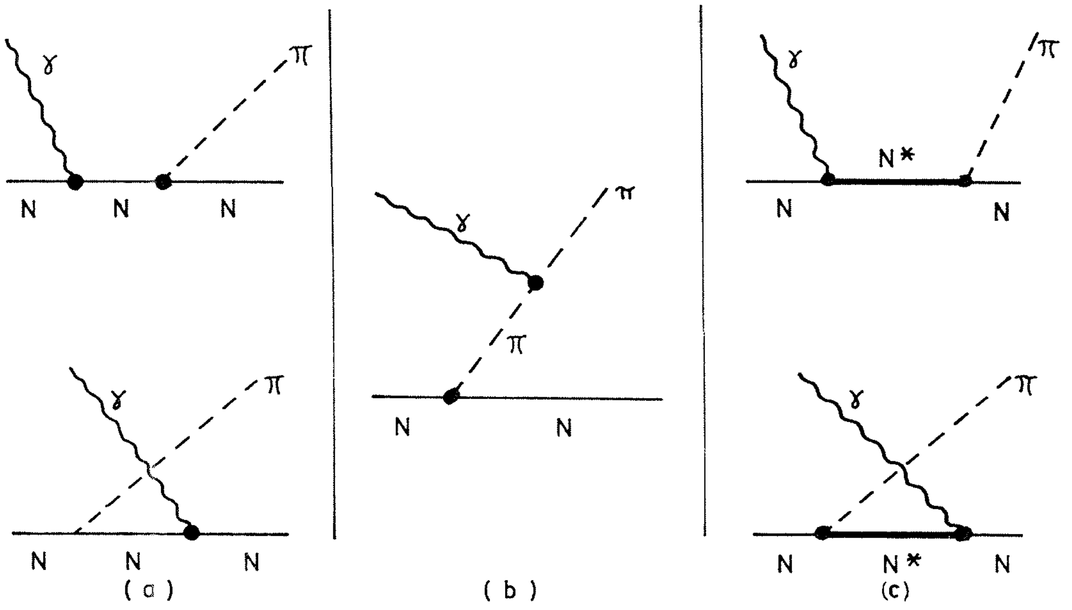


Figure 6

This combination of analyticity and unitarity gives rise to an integral equation whose solution allows to evaluate the three G_{N^*}.

Therefore, it is hoped that Dennery's work will offer a powerful tool for the determination of the form factor. His expression for the matrix elements although simple in principle, is algebraically rather involved because of the unpleasant fact that the N has spin 1/2 the photon spin 1 and the N* spin 3/2 . Numerical calculation of the different cross-sections are now in progress.

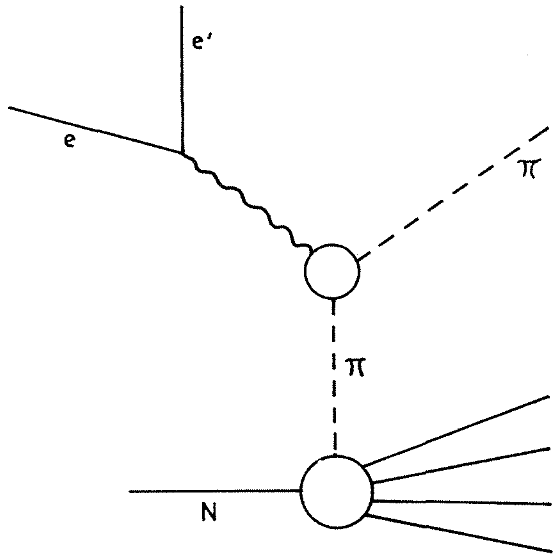


Figure 7

We have seen that electron-production of mesons offers in principle a way of determining $G_\pi(t)$. The practical determination of $G_\pi(t)$ is rather difficult because (due to the appearance of γ_5 in the π N vertex) the graph (b) gives a small contribution to the total production cross-section.

As pointed out by Drell the effect of direct electron pion interaction will become much more important in reactions in which several pions are produced. The situation is illustrated in figure 7. In the lowest vertex many pions are produced in the interaction between the virtual pion and the target nucleon. This effect enhances the contribution of the peripheral graph of figure 1. Fokion Hadjioannou at Stanford has investigated this process in detail and found for an initial energy of 10 GeV, a reasonably large cross-section for the process. The competing diagrams turn out to be sufficiently small not to interfere with the determination of G_π.

III - THEORETICAL MODELS FOR THE NUCLEON FORM FACTORS

The theoretical attempts to understand the nucleon form factors have followed two different points of view. The first one is to try to compute directly the expectation value of the electromagnetic current in the physical nucleon state starting from a field theoretical Hamiltonian.

The main difficulty of these approaches is that reliable methods of computation are not available when perturbation theory fails. The attempt in this direction which has had a certain amount of success is the one based on the static Hamiltonian of Chew and Low [1]. Using only two parameters the coupling constant and the cut-off (whose values were already determined from π N scattering) one has obtained a rough agreement with the magnetic structure of the nucleon, whereas the model failed completely in explaining the charge structure of the neutron. In later investigations, the validity of the static approximation as applied to the nucleon structure has been questioned because the recoil corrections have been found to be large and going in the wrong direction [2], [3].

The second, more recent, point of view has been based on the use of dispersion relations. Its aim is more modest in the sense that what one tries to do is to obtain sufficiently simple expressions for the form factors in terms of a few parameters and to connect those parameters with the ones appearing in the theory of different phenomena like π-π, π-N and N-N scattering. The general philosophy of this approach will be discussed by Matthews and I Shall discuss here only the specific application to the nucleon form factors.

The first dispersion calculations were performed in a scheme in which one assumed no pion-pion correlation. The values of the spectral functions g(t) obtained in such a manner were much too small in the low t region and so they excluded any possibility of agreement with experiment. Even by performing two subtractions the disagreement between theory and experiment was still very large. Therefore one becomes convinced that some of the starting hypothesis had to be changed, in particular Frazer and Fulco have shown that the introduction of a strong pion-pion interaction could enhance the spectral function in such a manner that theory can become consistent with experimental data.

The results I am going to discuss are based on the idea of a strong correlation between pions. This idea has received in the last few months several experimental confirmations.

1 - THE FORM FACTOR OF THE PION -

Let us first discuss briefly the application of dispersion methods to the pion form factor [2], [4].

We recall from part I that $G_\pi(t)$ could be written in the spectral form :

$$\operatorname{Re} G_\pi(t) = e + \frac{t}{\pi} \int_{4m_\pi^2}^{\infty} \frac{\operatorname{Im} G_\pi(t')}{t'\,(t' - t)}\, dt' \qquad \text{(III. 1)}$$

for those values of t for which only the 2π contribution is important (i.e., when $\pi\pi$ scattering is mainly elastic).

$$\operatorname{Im} G_\pi(t) = \operatorname{Re} G_\pi(t)\ \mathrm{tg}\,\delta_1 \qquad \text{(III. 2)}$$

In the "two pion approximation" one can therefore write :

$$\mathrm{Re}\, G_\pi(t) = e + \frac{t}{\pi} \int_{4m_\pi^2}^{\infty} \frac{\mathrm{Re}\, G_\pi(t')\mathrm{tg}\,\delta\,(t')}{t'(t'-t)}\, dt' \qquad \text{(III. 3)}$$

The integral equation (III. 3) is of a well-known type and its general solution has been given by Omnès [5]. The application of Omnès' method to Eq. (III. 3) gives, in the no-subtraction philosophy :

$$G_\pi(t) = e \exp\left[\frac{t}{\pi} \int_{4m_\pi^2}^{\infty} \frac{\delta\,(t')}{t'\,(t' - t - i\,\eta)}\, dt'\right] \qquad \text{(III. 4)}$$

Eq. (III. 4) solves in principle the problem of the pion form factors in the two pion approximation. This solution requires a rather good knowledge of the T = J = 1 pion-pion phase shift which is at present not yet available. If one approximates the pion-pion amplitude by means of a simple Breit-Wigner resonance formula :

$$e^{i\delta}\sin\delta = \frac{\Gamma\, k^3}{(t_v - t) - i\Gamma\, k^3} \qquad \text{(III. 5)}$$

then a simple approximation to Eq. (III. 4) is :

$$G_\pi(t) = C\,\frac{e^{i\delta}\sin\delta}{k^3} \qquad \text{(III. 6)}$$

where the constant C is approximately given by the condition G (O) = e

$$C \cong \frac{e\Gamma}{t_v} \qquad \text{(III. 6')}$$

The general trend of the real and imaginary parts of $G_\pi(t)$ in the case of a π-π resonance is sketched in figure 8.

Figure 8 shows clearly the physical situation. The most interesting features of the pion form factors are exhibited in the time-like region whereas in the space-like region one merely sees the decrease of a Breit-Wigner resonance.

We wish to stress the very qualitative character of Eqs. (III. 4) and (III. 5) : we now know that the position of the π-π resonance is for $t \cong 3\, m^2$ in a region in which the 4 pion channel is already open so that the validity of the two pion approximation may be questioned, especially for Eq. (III. 6') . There is, however, one important feature which is model independent, i. e. the spectral function is mainly concentrated in the π-π resonance region. This feature already allows to obtain a general form for $G_\pi(t)$ valid in the space-like region. Since for t space-like we are very far from the resonance position, a simple approximation is obtained by applying the mean value theorem.

$$G_\pi(t) = e + \frac{Bt}{t - t_v} = e\left(1 - a + \frac{a}{1 - \frac{t}{t_v}}\right) \qquad \text{(III. 7)}$$

where :

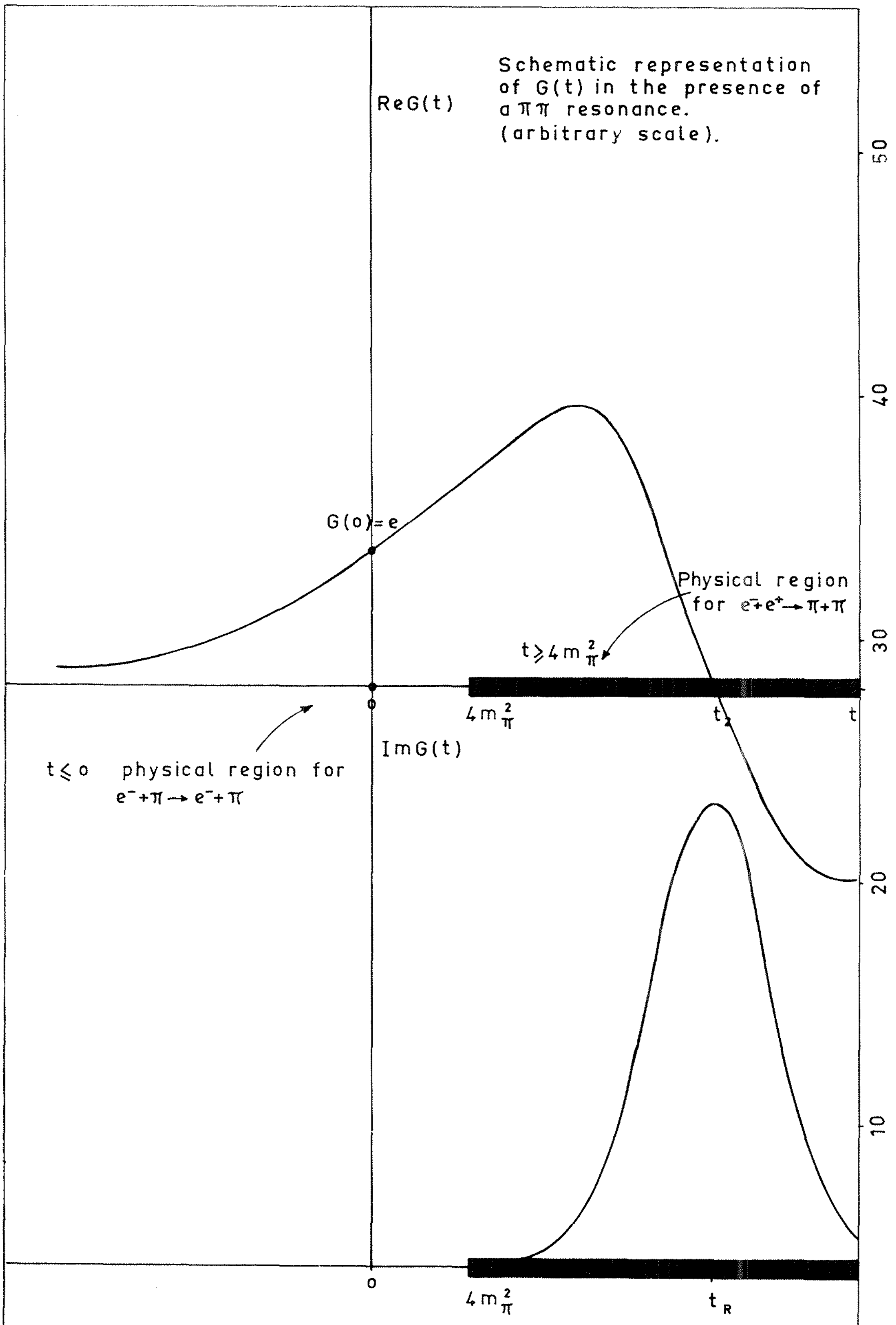
$$B \cong \int \frac{g_v(t')\, dt'}{t'}.$$

The corrections due to the widths of the resonance are of order $(\Gamma/t_v)^2$ and therefore amount to a few percents.

Unfortunately the simple formulae obtained in this section refer to a process for which experimental data are not yet available. Many of our difficulties come from the fact that many of the processes which are simple from the theoretical standpoint are difficult from the experimental one and vice-versa.

2 - THE FORM FACTORS OF THE NUCLEON -

In section I. 3 we have seen that the isovector form factors of the nucleon can be written in the spectral form :



Schematic representation of G(t) in the presence of a ππ resonance. (arbitrary scale).
ReG(t)
G(o)=e
Physical region for e⁻+e⁺→π+π
t≥4m²π
0
4m²π
t₁
t
ImG(t)
t≤o physical region for e⁻+π→e⁻+π
0
4m²π
tR
0
10
20
30
40
50

Figure 8

$$G_v^{1,2}(t) = G_v^{1,2}(O) + \frac{t}{\pi} \int_{4m_\pi^2}^{\infty} \frac{g_v^{1,2}(t')}{t'(t'-t)} dt' \qquad \text{(III. 8)}$$

and that the spectral functions were related through unitarity to the sum :

$$\sum_\alpha < O | J_\mu | \alpha > < \alpha | T^+ | N\bar{N} > .$$

In the region where the two pion contribution is the dominant one, the spectral function $g_v(t)$ depends from the product of two terms :

a) the $< O | J_\mu | \pi\pi >$ vertex function which has been studied in section 1.

b) the $< \pi\pi | T^+ | N\bar{N} >$ amplitude.

We have already seen that this amplitude appears for a ($N\bar{N}$) energy t much lower than the physical threshold for the process. Therefore we actually need a theoretical procedure which allows to give a meaning to such an amplitude.

This procedure is given in principle by the Mandelstam representation [6] which allows to continue analytically the scattering amplitude from the physical region to the unphysical values of energy and angle.

I am not going to discuss here the mathematical details of the derivation of this amplitude and its use in connection with the calculation of $g_v(t)$ which can be found in the paper by Frazer and Fulco [4] but only report on the main physical results which are intuitively very simple. (See figure 9).

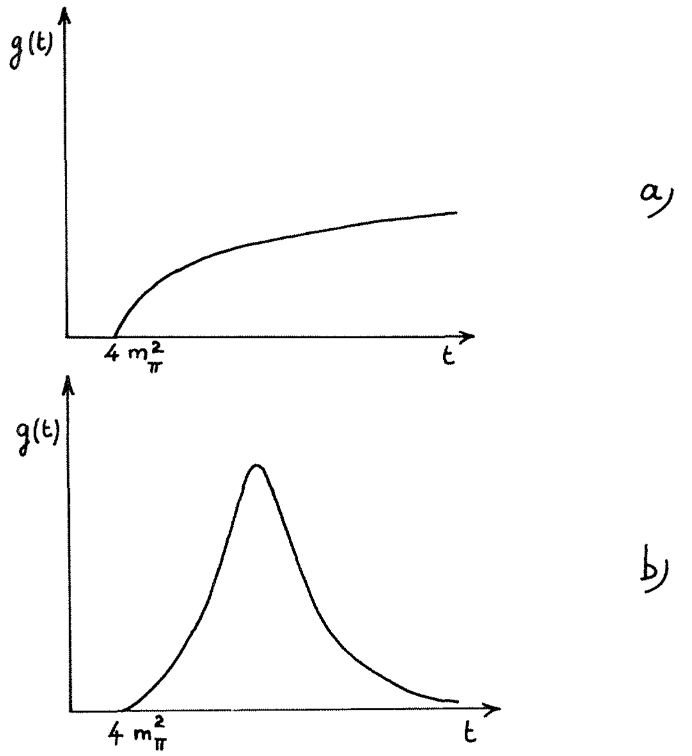


Figure 9 - Schematic representations of g(t) in arbitrary scale.
a) uncorrelated pions ; b) strong pion-pion resonance.

Without π-π interaction the spectral function is a smoothly varying function of t starting from zero at $t = 4\,m^2$. On the other hand, in the presence of a strong pion-pion resonance at $t = t_v$ the spectral function will exhibit a maximum at $t = t_R$ (as in the case of the pion form factor).

This is intuitively clear since $g_v(t)$ represents the weight with which the different two pion states contribute to the nucleon structure. Therefore a strong π-π correlation will concentrate the weight function in the resonance region.

Thus, in the case of a strong π-π interaction, we can write as for Eq. (III.7) :

$$G_v^{1,2}(t) = G_v(O)\left[(1 - a_v^{1,2}) + \frac{a_v^{1,2}\, t_v}{t_v - t} \right] \qquad \text{(III. 9)}$$

Eq. (III.9) has first been proposed by Clementel and Villi [7] and used in comparison with the early experiment on electromagnetic structure of the proton with a value $t_v = 23\ m_\pi^2$. This was interpreted by Bowcock, Cottingham and Lurié [8] in terms of a dipion mass of $\sim 4.7\ m_\pi$ not far from the recent experimental value of $5.4\ m_\pi$ found in pion production.

Let us now discuss the isoscalar part of the form factors. In this case the spectral representation is :

$$G_s(t) = G_s(O) + \frac{t}{\pi} \int_{9m_\pi^2}^{\infty} \frac{g_s(t')}{t'(t' - t)}\, dt' \qquad \text{(III. 10)}$$

As already discussed in part I, the first contributor to $g_s(t)$ will be the three pion state. This contribution is the product of two terms :

a) the $(\gamma\, 3\pi)$ vertex $\langle O | J_\mu | \pi\pi\pi \rangle$

b) the $\langle \pi\, \pi\, \pi | T^+ | N\bar{N} \rangle$ matrix element.

The treatment of the isoscalar form factor is thus much more complicated than in the isovector case. However, some of the qualitative features are the same as in the isovector case.

If there exists a strong correlation between three pions the spectral function will exhibit a maximum at the resonance position t_s, otherwise it will be a smooth function of t.

The present experimental evidence suggests a strong concentration of the spectral function in the low t region ; one was therefore led to assume also the existence of a 3π resonance and to write :

$$G_s^{1,2}(t) = G_s(O)\left[(1 - a_s^{1,2}) + \frac{a_s^{1,2}\, t_s}{t_s - t} \right] \qquad \text{(III. 11)}$$

The three pion resonance t_s is the translation in the dispersion language of the neutral vector meson suggested in 1957 by Nambu [9], [10] in order to explain the apparent lack of neutron charge structure at that time.

Eqs. (9) and (10) constitute a well-defined model for the four nucleon form factors [11] . It contains six parameters : the positions of the two resonances and the four constants representing the effect of these resonances in the electromagnetic structure of the nucleon. These constants are related through theory to similar constants giving the effect of pion-pion interaction in different phenomena like πN and NN scattering.

So we have arrived to a stage that, although we are still far from a fundamental understanding of the nucleon structure, we have a very simple description of both proton and neutron structure and we can correlate the (eN) scattering experiments with other experiments in pion physics.

3 - THE PION-PION RESONANCES -

The model described in the last section requires the existence of two metastable particles :

a) a T = 1 J = 1 two π resonance,

b) a T = 0 J = 1 three π resonance.

The experimental fits of Eqs. (III.9) and (III.10) made both at Stanford [12] an d at Cornell [13] were giving for the masses of these resonances the following values :

$$t_s \sim 10\ m_\pi^2 \qquad\qquad t_v \sim 20\ m_\pi^2$$

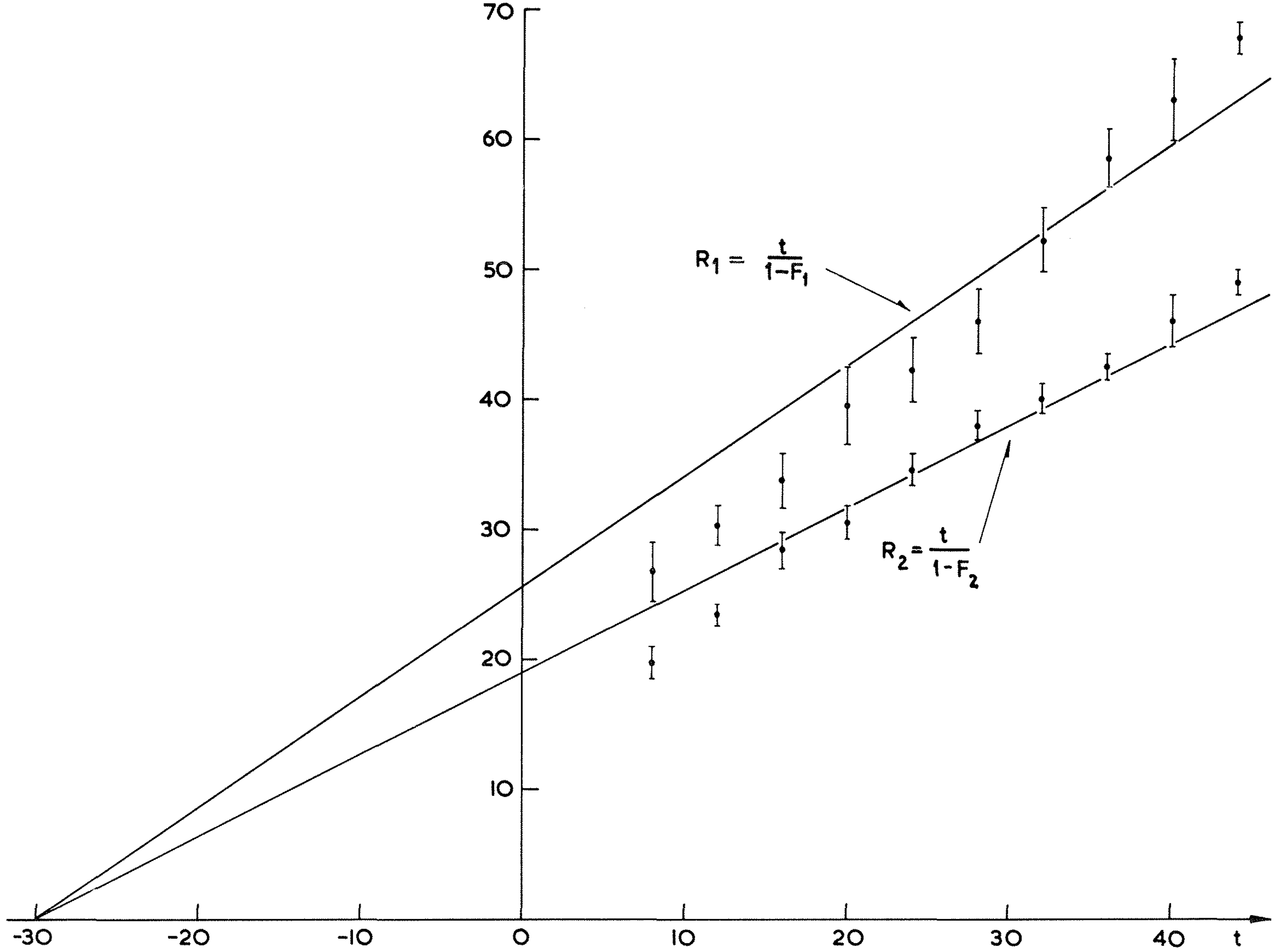
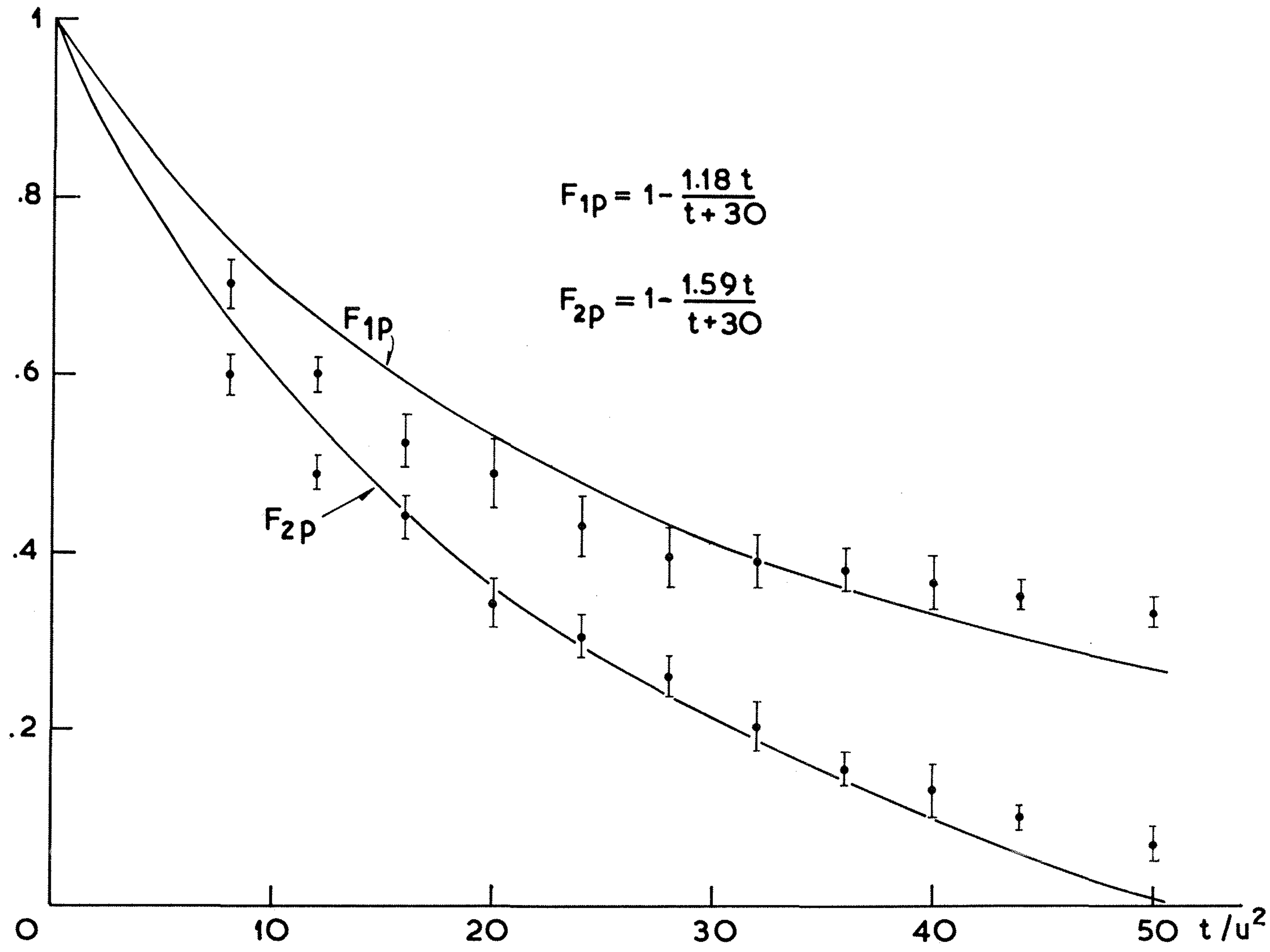

70
60
50
40
30
20
10
$R_1 = \frac{t}{1-F_1}$
$R_2 = \frac{t}{1-F_2}$
-30
-20
-10
0
10
20
30
40
t

Figure 10



$F_{1p} = 1 - \frac{1.18\,t}{t+30}$
$F_{2p} = 1 - \frac{1.59\,t}{t+30}$
F_{1p}
F_{2p}
1
.8
.6
.4
.2
0
10
20
30
40
50
t/u^2

Figure 11

A statistical analysis performed by Bergia and Stanghellini [14] shows that the acceptable values of t_s and t_v can very between rather wide limits.

Very recently direct evidence for the existence of pion pion resonances was found.

First of all a T = 1 π - π resonance has been found in $\pi + N \longrightarrow \pi + \pi + N$ experiments both at Wisconsin [15] and at Bologna, Saclay and Orsay [16]. An angular distribution analysis [16] gives strong evidence in favour of J = 1. The mass of the particle is 5.4 m_π, i.e., $t_v = 29\ m_\pi^2$ which is certainly compatible with the results of the analysis of the isovector form factor.

Secondly a 3π resonance of T = 0 has been recently discovered in two different experiments [17], [18] with a mass slightly larger than the one of the two-pion resonance.

In this case the experimental result is very far from the prediction coming from the analysis of the form factor. In order to give an idea of the discrepancy, Figures 10 and 11 show the comparison with experiment of an attempt of fit taking $t_s \sim 32\ m_\pi^2$ as suggested by the recently found 3π resonance. This comparison shows that even if the 3π resonance recently found will turn out to have J = 1, the electromagnetic structure of the nucleon still needs a significantly lower mass contribution to the isoscalar spectral functions.

It is very difficult to make any prediction on the issue of future experiments, I wish, however, to recall that there exists another low energy T = 0 effect which in my opinion is not yet understood. I am referring to the He_3 recoil experiment in P + D reaction [19]. The bump found in the recoil spectrum of He_3 was first interpreted as a T = 0 J = 1 state. This interpretation was then dismissed on the basis of the experimental width and on arguments based on the K meson decay [20]. I wish to point out that the interpretation now favoured on the basis of a T = 0 J = 0 low energy π - π interaction (with a scattering length of $a = 2.5 \frac{1}{m_\pi}$) gives also rise to very serious difficulties.

First of all it is very hard to understand why such a strong $\pi\pi$ correlation at threshold does not strongly perturb the decay spectrum of the τ meson [21] which shows a rather small variation from phase space.

Another difficulty of such a large scattering length has been pointed out by C. Ceolin and R. Stroffolini [22].

If one tries to compute the effect of π-π interaction with $a = 2.5 \frac{1}{m_\pi}$ on $\pi + N \longrightarrow \pi + \pi + N$ one finds a production cross-section which is about twenty times larger than the experimental ones. The situation is illustrated in figure 12, which shows that $a_s > 1$ is inconsistent with production data.

So both interpretations of the He_3 recoil experiment are, in my opinion, equally unsatisfactory and maybe the ABC particle might still be considered as a possible candidate for the T = 0 J = 1 role.

4 - CONCLUSIONS -

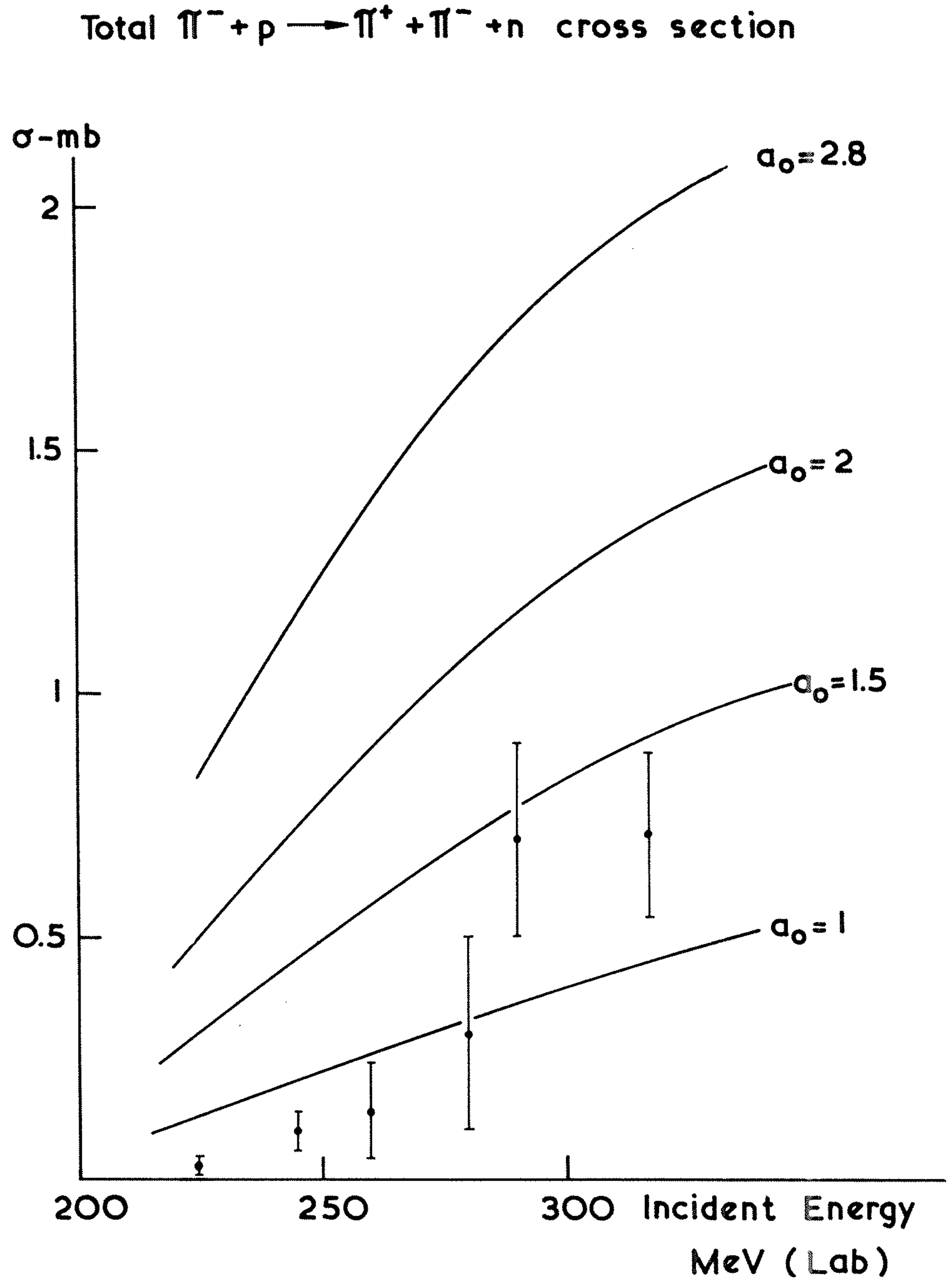
We have seen that the situation concerning the electromagnetic structure of the nucleon, although still unsettled, is certainly very interesting.

We have now a general model which tries to interpret both proton and neutron structures in terms of a unique physical phenomenon : strong interaction between bosons.

Such a strong interaction has now been revealed experimentally with features which are not very far from the ones that were guessed on the basis of the interpretation of form factors.

Now the main theoretical problem for the future is to study and correlate between themselves the different features of the vector particles and to predict other possible higher resonances which might have an effect on the inner structure of the nucleon.

Very interesting work on this programme is now in progress, but since it still is in a rather preliminary stage, I shall only refer to the main features. There are two independent and complementary lines of thought. The first one interprets the vector unstable particles as resonant state due to some kind of pion-pion force. Dispersion methods have been applied both to study the 3π resonance in terms of the 2π (R. Blankenbecler and J. Tarski) or the 2π resonance in terms of the 3π (R. Blankenbecler). These calculations also give predictions about the existence of higher resonances. A potential calculation has been performed by L. Schiff, the parameters of the π - π force are extracted from the experimental information about the dipion. It is then shown that this attraction leads to a three pion quasi-bound state with the right quantum numbers.



Total π⁻ + p ⟶ π⁺ + π⁻ + n cross section
σ-mb
2
1.5
1
0.5
a₀=2.8
a₀= 2
a₀=1.5
a₀= 1
200
250
300
Incident Energy
MeV (Lab)

Figure 12

The second approach, followed by Sakurai [23] and Gell-Mann and Zachariasen gives to the vector particles a more fundamental role in elementary particle physics.

The vector particles can be considered as "heavy photons" and interact with current which (as the electromagnetic current) are conserved, namely the currents related to conservation of baryon charge and hypercharge. One difficulty is that the possible experimental tests of this theory refer to vector mesons having zero momentum and zero total energy. The vector particles experimentally found until now have such a large mass so that we are very far from the limit in which the test of the theory can be done.

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LES ISOBARES DES NUCLEONS ET LES ETATS RESONNANTS DES SYSTEMES DE PLUSIEURS MESONS π

J. F. DETŒUF

INTRODUCTION

Depuis quelques années les évidences expérimentales de l'existence d'états résonnants dans le système π-nucléons et dans les systèmes de deux ou plusieurs pions se sont accumulées.

Quoiqu'il soit encore trop tôt pour affirmer que tous les états "résonnants" ou "isobares" dont je vais parler sont en réalité de vraies résonances correspondant à des déphasages passant par 90° (ou à des pôles de la matrice S), il sera commode pour l'exposé d'utiliser un de ces deux vocables, même s'il n'est pas entièrement justifié.

Ces états résonnants jouent un rôle important dans toute la physique des hautes énergies : aussi ai-je dû restreindre mon exposé aux expériences relatives au scattering π-nucléon (élastique ou inélastique), à la photoproduction et au scattering nucléon-nucléon en me limitant aux énergies comprises entre quelques centaines de Mev et quelques Gev. Je ne parlerai en particulier ni de la photoproduction aux basses énergies ni du scattering π-nucléon aux basses énergies.

Même en le restreignant ainsi, le sujet est vaste et je n'ai pas la prétention d'être complet.

Nous considèrerons successivement : 1) les expériences qui permettent de caractériser les isobares π-nucléons c'est-à-dire qui permettent de fixer leur énergie et dans la mesure du possible de leur assigner des nombres quantiques (chap. II), 2) le rôle de ces isobares π-nucléon dans les réactions de production (chap. III), 3) les évidences expérimentales d'états résonnants des systèmes de plusieurs π (multipion) (chap. IV).

II - CARACTERISTIQUES DES ISOBARES π-NUCLEONS

1 - LISTE DES ISOBARES

La figure 1 montre les sections efficaces totales π^+-p et π^--p avec les maxima maintenant bien connus A, B, C et D pour des énergies cinétiques T_π du π incident respectivement de 200, 605, 890 et 1 300 Mev. Les maxima A, B et C se retrouvent dans la section efficace totale de photoproduction des π^0 pour la même énergie dans le centre de masse, c'est-à-dire pour une énergie dans le laboratoire du γ incident E_γ = T_π + 150 Mev. La section efficace totale de photoproduction des π^+ présente un maximum à 700 Mev qui peut être attribué à B, le décalage de 50 Mev en énergie étant dû à l'interférence du terme photoélectrique.

De ces maxima seul A peut être attribué avec certitude à une résonance. C'est la résonance 3/2 3/2 bien connue, c'est-à-dire qu'elle a un spin isotopique I = 3/2, un spin total J = 3/2, un moment angulaire 1 (onde P 3/2) correspondant à une parité totale +, elle correspond à une énergie des π incidents T_π = 180 Mev et sa largeur est de 100 Mev environ.

En ce qui concerne les maxima B, C et D, le premier problème si l'on veut parler d'isobares ou de résonances est d'essayer de leur assigner une énergie et des nombres quantiques de spin, spin isotopique et parité. A la conférence sur les interactions fortes tenue à Berkeley en décembre, deux articles traitaient de la question, l'un de Moyer [1] l'autre de Falk-Vairant et Valladas [2]. Depuis lors, il y a eu quelques données expérimentales nouvelles, mais qui ne changent pas sensiblement leurs conclusions de telle sorte que sans entrer dans les détails de la discussion, je

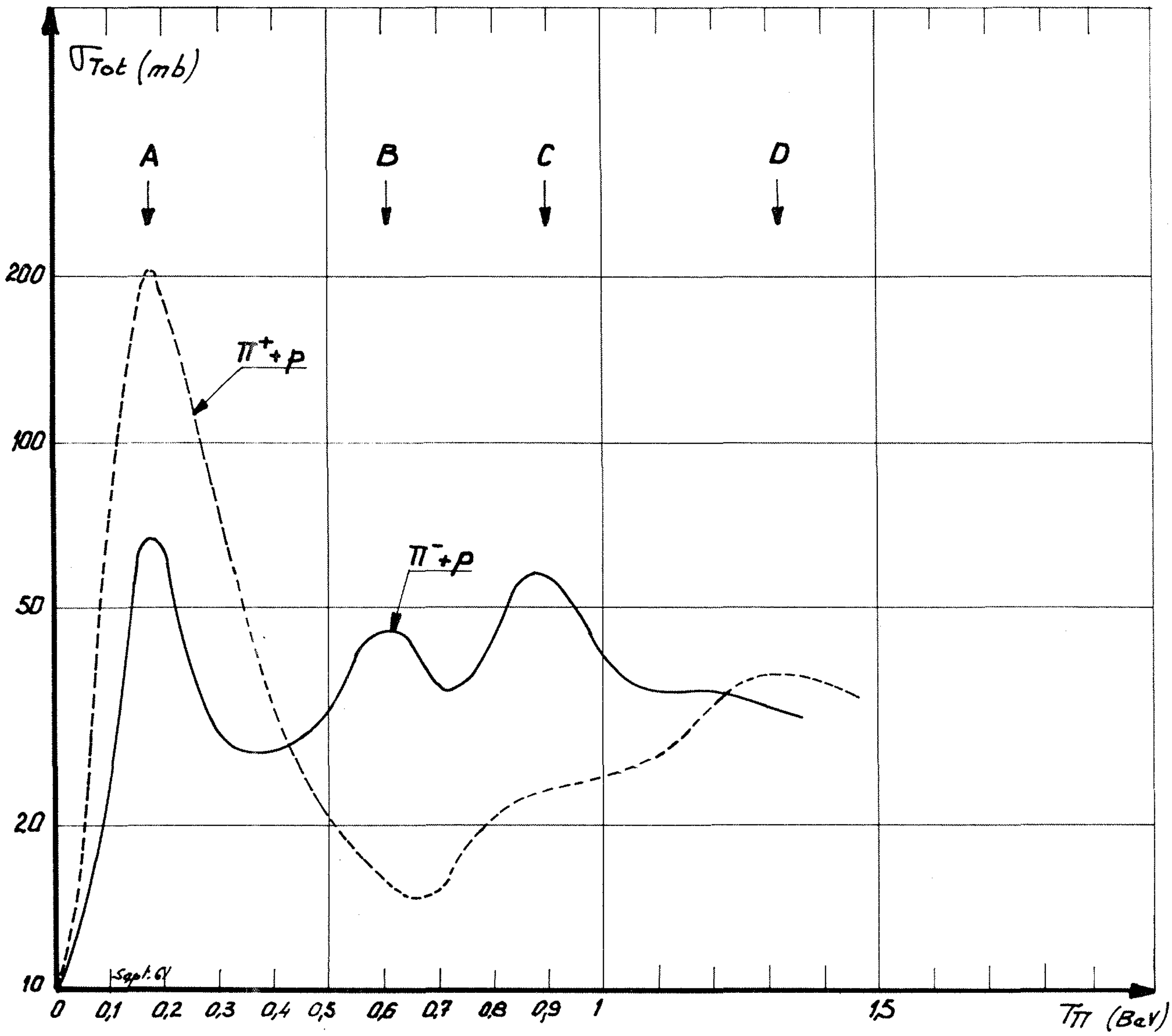


Figure 1 - Sections efficaces totales π^+ - p et π^- - p. En abscisse l'énergie cinétique du π dans le laboratoire.

mentionnerai ces conclusions en rappelant les nouveaux résultats expérimentaux. On trouvera une discussion plus détaillée dans le rapport d'Omnès [3].

2 - CARACTERISTIQUES DE L'ISOBARE B

Les principales informations concernant les isobares B et C viennent de la diffusion π-p et de la photoproduction. L'état B qui apparaît dans σ_{tot} (π^--p) pour T_π = 605 Mev avec une largeur de 100 à 150 Mev n'apparaît pas dans σ_{tot} (π^++p), il est donc certainement dans l'état de spin isotopique I = 1/2. La distribution angulaire de photoproduction des π^0 à E_γ = 750 Mev est en :

$$5 - 3 \cos^2\vartheta$$

ce qui permet de lui attribuer le spin total J = 3/2 [4]. La section efficace élastique π-p dans l'état I = 1/2 pour T_π = 600 Mev est compatible avec cette valeur du spin [2].

Comme Sakurai [5] l'a suggéré, une forte polarisation du proton émis à 90° dans le centre de masse CM au cours de la réaction de photoproduction des π° : $\gamma + p \longrightarrow p + \pi^\circ$ ne peut résulter que de l'interférence entre 2 états de parité opposée. Pour des énergies comprises entre A et B on peut espérer que les seules amplitudes importantes sont celles de A et B et la mesure de la polarisation doit permettre de choisir pour B entre les états $P_{3/2}$ et $D_{3/2}$. La forte polarisation trouvée de l'ordre de 0,5 à 0,6 [6] [7] (figure 2) est donc en faveur d'un état de parité négative c'est-à-dire de $D_{3/2}$. Sur la figure 2 est reporté le point à 910 Mev obtenu à Frascati [7] et présenté à cette conférence par Querzoli.

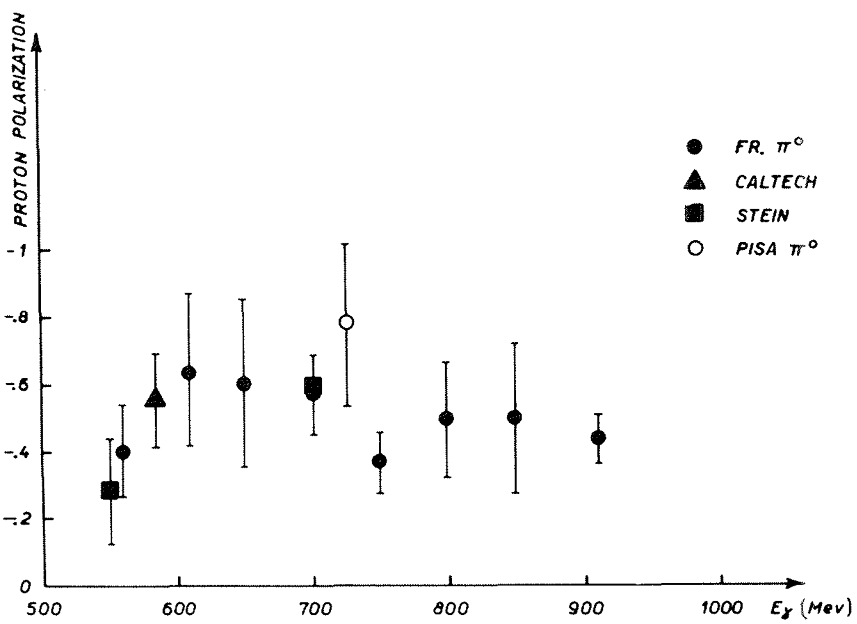


Figure 2 - Polarisation du proton émis à 90° CM dans la photoproduction des π°.

Maloy [6] a regardé si cette polarisation pouvait résulter d'une interférence de B dans l'état $P_{3/2}$ soit avec un terme E_{11} correspondant à une onde $S_{1/2}$ non résonnante, soit avec un terme M_{25} correspondant à un état $F_{3/2}$ pour la résonance C. Utilisant pour les termes résonnants ABC de simples formules de résonance, ses calculs ont montré que seule l'hypothèse $D_{3/2}$ pour B était en accord avec les résultats expérimentaux. Cependant, les hypothèses de Maloy ne sont peut-être pas suffisamment générales pour exclure complètement tout autre possibilité.

Avant de quitter la photoproduction, je voudrais mentionner 2 résultats expérimentaux concernant la photoproduction des π^+ à 180° et à 0°.

La figure 3 montre la section efficace différentielle à 0° et 180° telle qu'elle était connue en 1960 à Rochester. Comme R. L. Walker l'a montré, ces sections efficaces sont intéressantes parce que le terme photoélectrique s'y annule, les seuls termes pouvant encore contribuer étant les amplitudes qui peuvent être associées aux résonances. On voit que comparée à l'influence de la résonance A a E_γ = 350 Mev, celle de la résonance B a E_γ = 750 Mev n'affecte pas beaucoup la section efficace. Walker a remarqué que ceci était explicable pour une résonance $D_{3/2}$ parce que le terme E_{11} $(S_{1/2})$ domine et le comportement résonnant de E_{13} n'affecte pas beaucoup la variation de la section efficace. Par contre, une résonance $P_{3/2}$ nécessiterait un fort terme M_{13} (S) pour lequel il n'y a pas d'autre évidence.

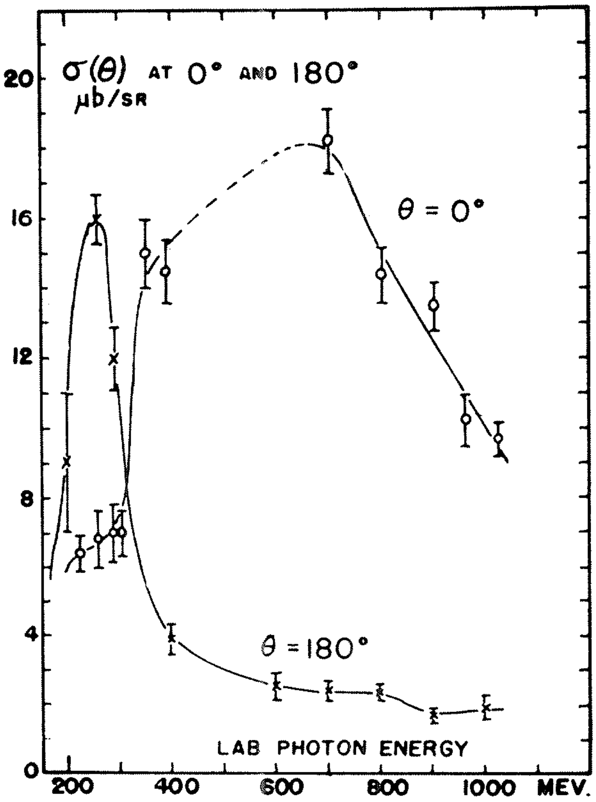


Figure 3

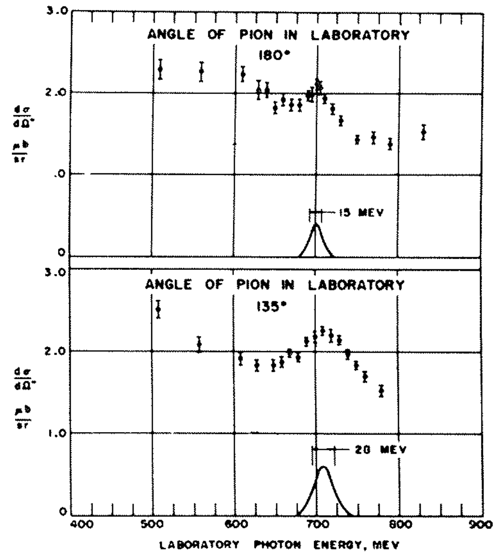


Figure 4

Figure 3 - Section efficace différentielle de photoproduction des π^+ à 0° et 180°.

Figure 4 - Section efficace différentielle de photoproduction des π^+ à 135° et 180° [8].
La courbe en trait plein indique la résolution en énergie des photons déduite de l'acceptance en impulsion du spectromètre, de la variation angulaire de la cinématique et des dimensions finies de la cible.

Le nouveau résultat de Hand et Schaerf [8] à Stanford (figure 4) concerne la photoproduction des π^+ aux grands angles. Ces auteurs trouvent un petit accident à 700 Mev qui est difficilement attribuable à la résonance B d'une part parce que le décalage en énergie de l'ordre de 50 Mev ne doit pas exister ici, d'autre part parce que sa largeur est inférieure à 15 Mev. Si cet accident est réel, il reste à expliquer.

Le second résultat est celui de Beneventano et al [9]. Ces auteurs ont présenté des résultats préliminaires sur la section efficace différentielle des π^+ aux petits angles à 600, 700 et 800 Mev qui ne semblent pas en accord avec les résultats de Boyden de la figure 3. Les sections efficaces trouvées semblent nettement plus grandes. Il se peut donc que l'aspect de la courbe dans cette région change bientôt. D'autre part, ces auteurs trouvent qu'à 700 Mev la dépendance angulaire de la section efficace aux petits angles est différente de celle trouvée à 600 et à 800 Mev.

Un autre résultat présenté par Querzoli [10] est la section efficace différentielle de photoproduction des π^0 à 90° dans le centre de masse (CM) (figure 5) qui a été mesurée avec une très bonne résolution en énergie en utilisant une chambre à étincelles pour détecter le proton de recul. Ils trouvent une résonance à E_γ = 740 ± 10 Mev avec une largeur de 60 Mev.

Une autre information importante sur la nature des résonances provient de la distribution angulaire dans la diffusion élastique π^- - p. Différents auteurs ont étudié cette distribution soit dans des chambres à bulles [11] [12] soit par des techniques de compteurs [13]. Ces distributions ont été approchées par des polynomes en $\cos^n \theta$ dont les coefficients sont indiqués sur la figure 6.

Les points de Meyer à Saclay et Bertanza à Pise n'existaient pas dans le rapport de Moyer.

Les points de Wood à Berkeley obtenus par la technique des compteurs sont les plus précis statistiquement ; mais ont une incertitude supplémentaire de ± 12 à 15 % sur la normalisation.

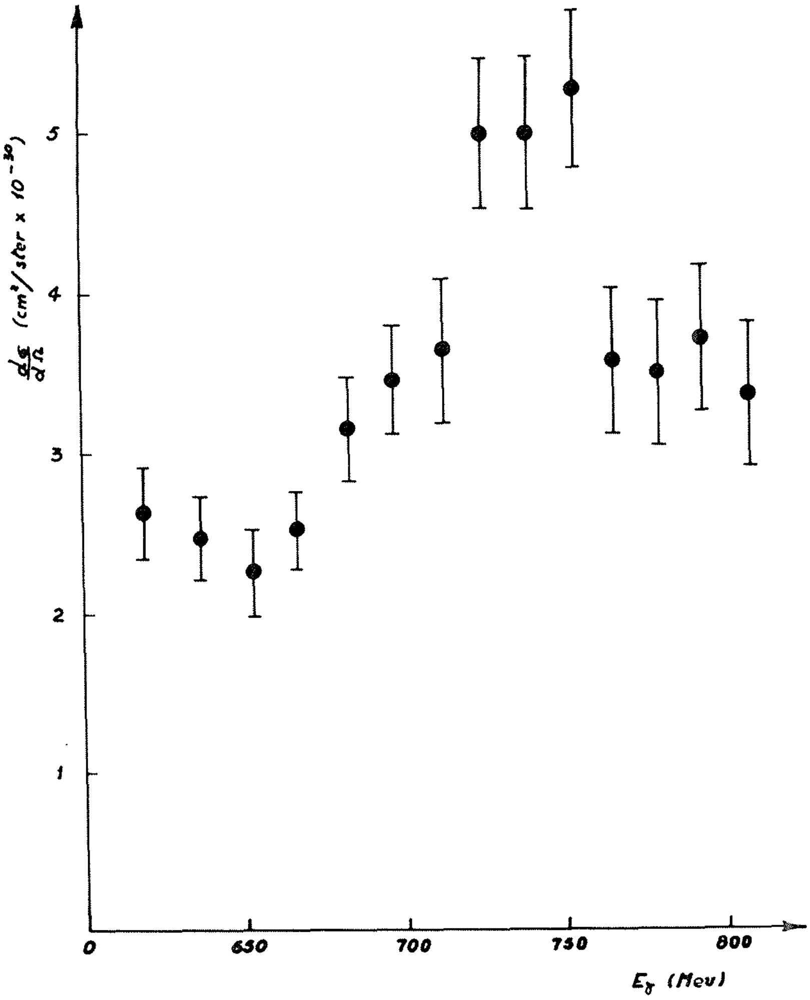


Figure 5 - Section efficace différentielle à 90° CM de la photoproduction des π^0 [10].

Le fait que a_3 soit grand à 450 Mev indique que l'on a besoin d'une interférence P D. Le fait que a_3 et a_4 soient petits à 600 Mev n'est pas en contradiction avec une résonance D car il s'agit d'une onde $D_{3/2}$, il n'y aura pas de terme supérieur à a_3 et la petitesse de a_3 peut résulter d'une interférence avec l'onde $P_{3/2}$.

En dépit du fait qu'aucun des arguments apportés ne soit décisif, on peut dire que l'attribution du maximum B à une onde résonnante $D_{3/2}$ dans l'état T = 1/2 à 600 Mev est l'hypothèse la plus simple qui n'est pas en contradiction avec les nombreux faits expérimentaux.

3 - CARACTERISTIQUES DE LA RESONANCE C

A 830 Mev il est possible d'évaluer les limites de la partie résonnante des sections efficaces totales, élastiques et inélastiques de la diffusion π-nucléon pour l'état I = 1/2 [2]. Ces limites sont compatibles avec un moment angulaire $J = \frac{5}{2}$ ou $J = \frac{7}{2}$. Le fait qu'à 900 Mev les coefficients supérieurs à a_5 (figure 6) ne soient pas nécessaires implique qu'il n'y a pas de terme en $J > 5/2$. La rapide croissance de a_5 entre 700 et 900 Mev implique une interférence $D_{5/2}$ $F_{5/2}$. Enfin, les résultats du groupe de Frascati sur la polarisation du proton dans la photoproduction à 800 Mev sont en faveur d'une parité opposée pour les résonances B et C c'est-à-dire d'une résonance $F_{5/2}$ pour C.

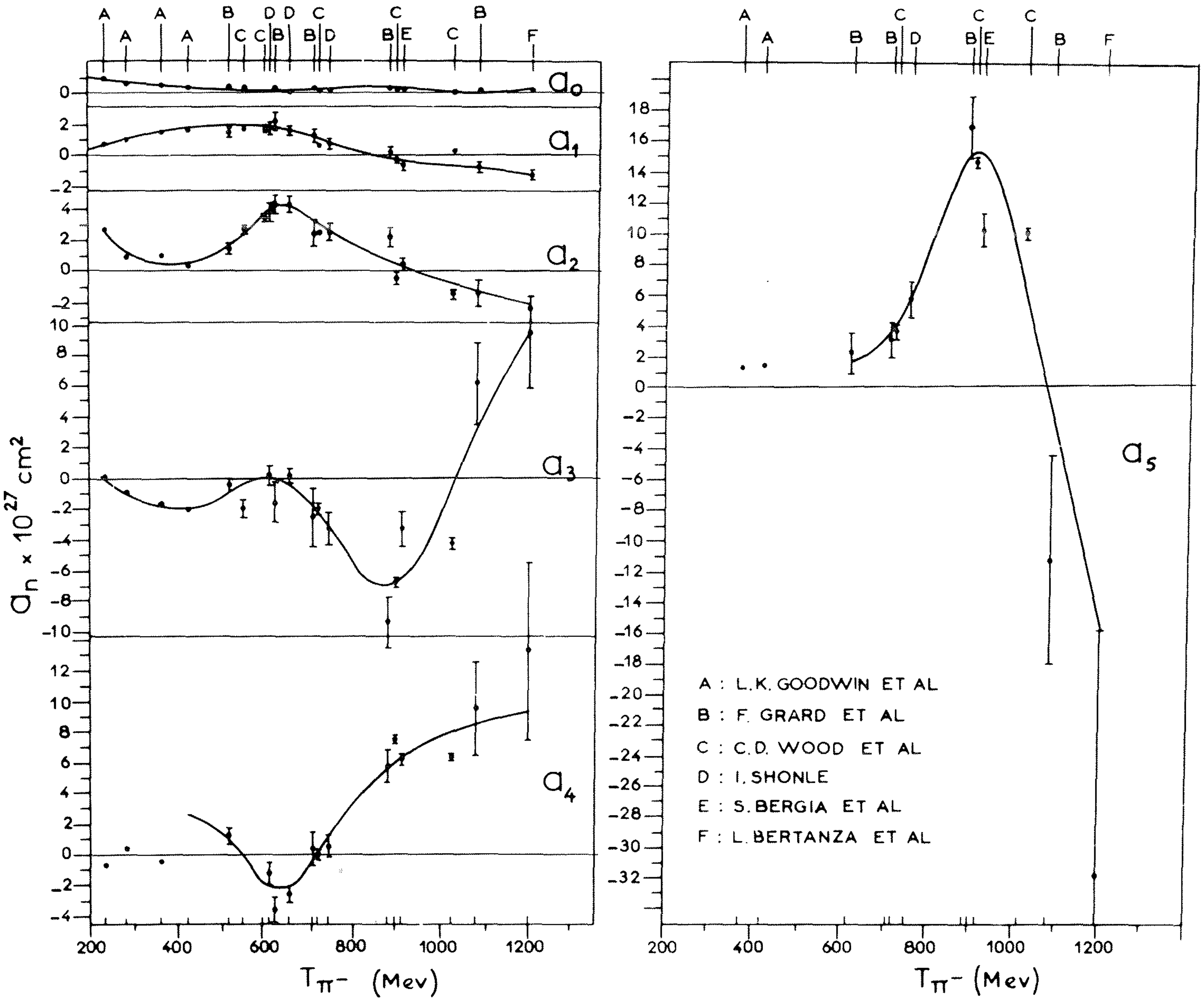


Figure 6 - Coefficients a_n de la distribution angulaire $\frac{d\sigma}{d\Omega} = \Sigma\, a_n \cos^n \theta$ de la diffusion élastique des π^- sur les protons [12].

Les données ne sont probablement pas suffisantes pour affirmer que l'onde $F_{5/2}$ résonne ou qu'elle est la seule onde résonnante à 900 Mev (*).

En ce qui concerne le maximum D, les données expérimentales sont encore trop restreintes pour que l'on puisse parler de résonance.

III - INFLUENCE DES ISOBARES π-NUCLEONS SUR LES REACTIONS DE PRODUCTION A 3 CORPS

L'existence des isobares π-nucléons et surtout celle de la résonance A domine les réactions de production à 3 corps

$$p + p \longrightarrow N + N + \pi$$

$$\gamma + p \longrightarrow N + 2\pi$$

$$\pi + p \longrightarrow N + 2\pi$$

1 - PRODUCTION DES π DANS LES CHOCS PROTON-PROTON

Les résultats expérimentaux sur la production simple des π dans les chocs p-p sont en accord assez remarquable avec le modèle isobarique de Lindenbaum et Sternheimer [14]. Dans ce

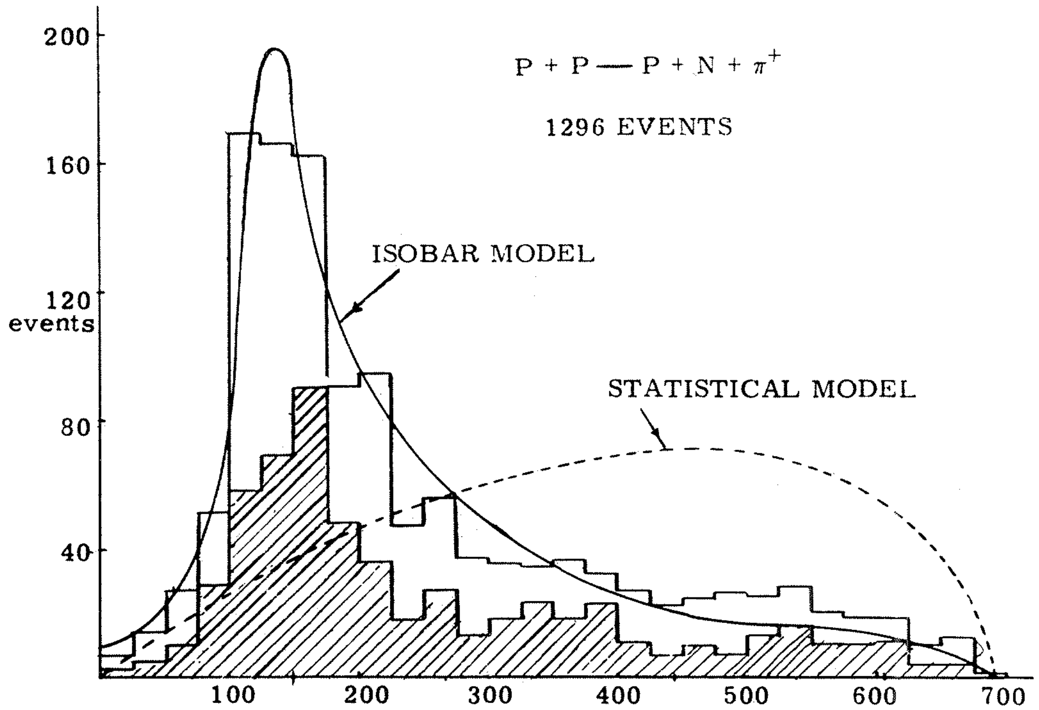


Figure 7 - Distribution des valeurs Q pour le système π proton dans la réaction $p + p \longrightarrow p + n + \pi^+$ à 2 GeV [15].

(*) La remarque faite par Höhler à la fin de la présentation de ce rapport et publiée d'autre part dans les comptes-rendus de la conférence (cf. Vol 1) est un argument très fort en faveur d'une onde résonnante $F_{5/2}$.

modèle qui suppose que toute la production simple des π se fait par l'intermédiaire de l'isobare A suivant le diagramme 1, la probabilité d'un état final donné est proportionnelle au produit de l'espace de phase à 2 corps par la probabilité de formation de l'isobare 3/2, 3/2 elle-même proportionnelle à la section efficace $\pi^+ p$.

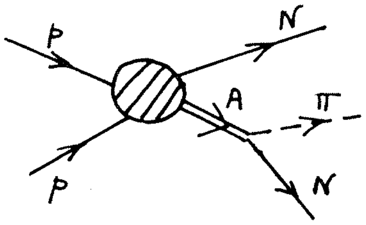


Diagramme 1

A 2 Gev et 2,85 Gev les résultats des groupes de Brookhaven et de Yale [15] [16] utilisant une chambre à bulles à hydrogène sont en très bon accord avec le modèle. La figure 6 présente par exemple le spectre de l'énergie cinétique totale $Q_{\pi p}$ du π^+ et du proton dans leur propre centre de masse pour la réaction $p + p \longrightarrow \pi^+ + p + n$ à 2 Gev.

La distribution angulaire des nucléons dans le centre de masse du système montrent que ceux-ci tendent à rester sur la ligne de vol initial, ce qui est une évidence de l'importance des interactions périphériques avec échange d'un méson π selon le diagramme 2 où N désigne un nucléon et A l'isobare 3/2 3/2. Les moments angulaires impliqués dans cette distribution sont de l'ordre de 5 à 2 Gev ce qui indiquerait une portée de 1,0 fermi en accord avec l'échange d'un méson π.

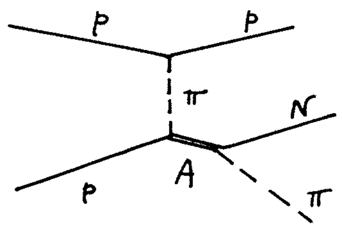


Diagramme 2

Les spectres de mésons π émis à différents angles dans la collision proton proton à 3 Gev mesurés par le groupe de Rochester et de Brookhaven et présentés par Yuan [17] montrent également que l'isobare est produit de façon très anisotrope dans le centre de masse (figure 8).

On a également observé l'influence des isobares B et C dans les collisions protons protons. La figure 9 représente à 2,9 GeV environ le spectre des protons émis à très petits angles, spectres obtenus par Duke et al [18].

Les flèches indiquent les positions des pics de protons que l'on devrait observer dans les réactions à deux corps $p + p \longrightarrow p + X$ où X est un des 3 isobares A,B,C. La figure 9 représente le spectre obtenu à $4,51°$ et la courbe en pointillé représente le calcul résultant de la contribution du seul diagramme 3, qui fait intervenir au vertex V la section efficace totale π°-p. Pour rendre

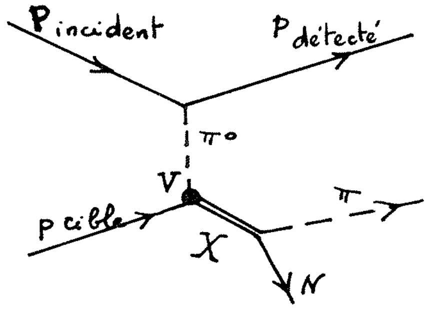


Diagramme 3

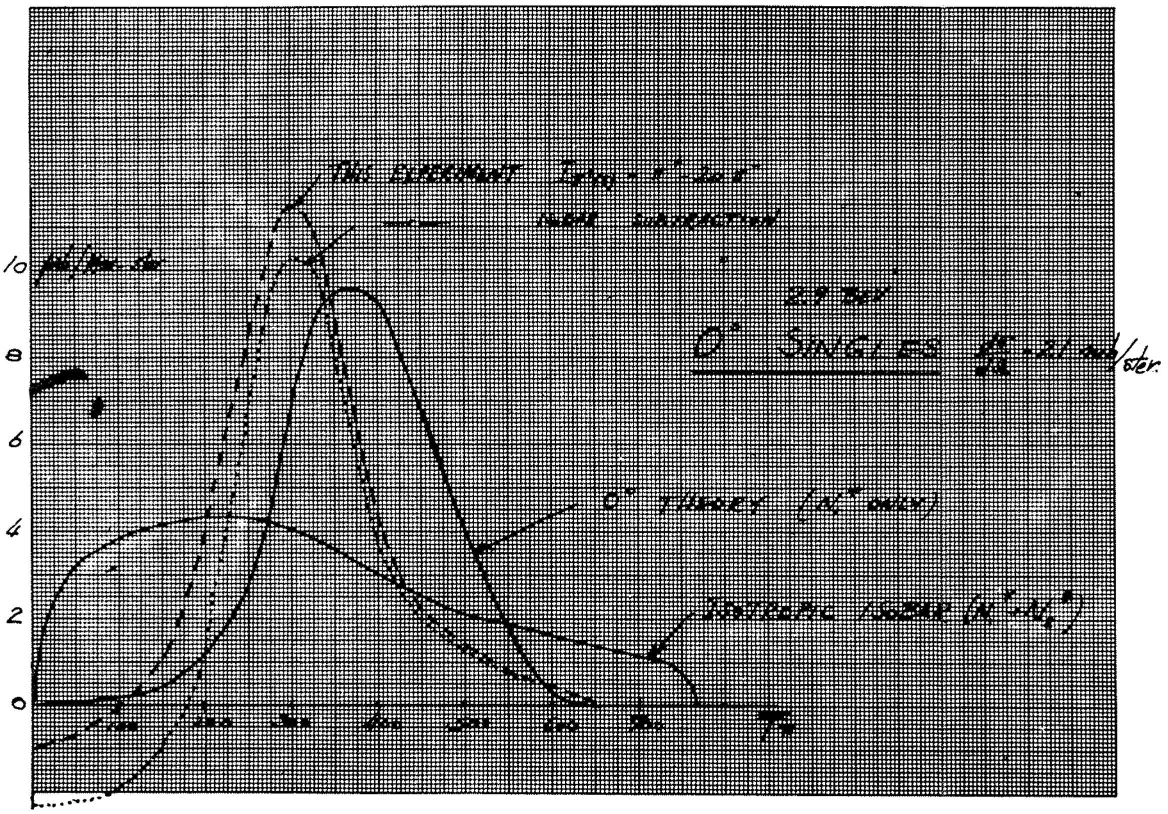


Figure 8 - Spectre des mésons π émis à 0° dans le choc p - p à 2,9 GeV.
Pour la signification exacte des différentes courbes voir référence [17].

compte des résultats expérimentaux, il faut ajouter à la contribution du diagramme précédent, entre autres, celle du diagramme 4 qui donne naissance également à l'émission de protons énergiques

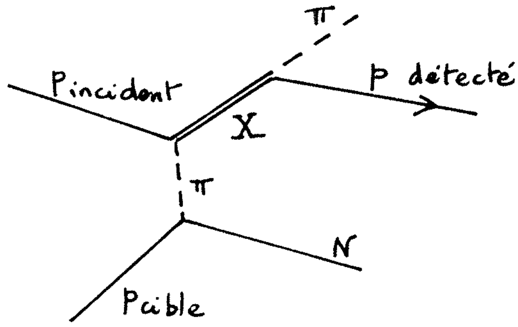


Diagramme 4

vers l'avant, mais la structure dûe à la présence de l'isobare est lavée par la désintégration de celui-ci qui est plus ou moins isotrope dans son centre de masse. On notera que dans la section efficace totale $\pi^\circ - p = \frac{1}{2}\ (\sigma_{\pi^+p} + \sigma_{\pi^-p})$ le maximum B est moins prononcé que les maxima B et C

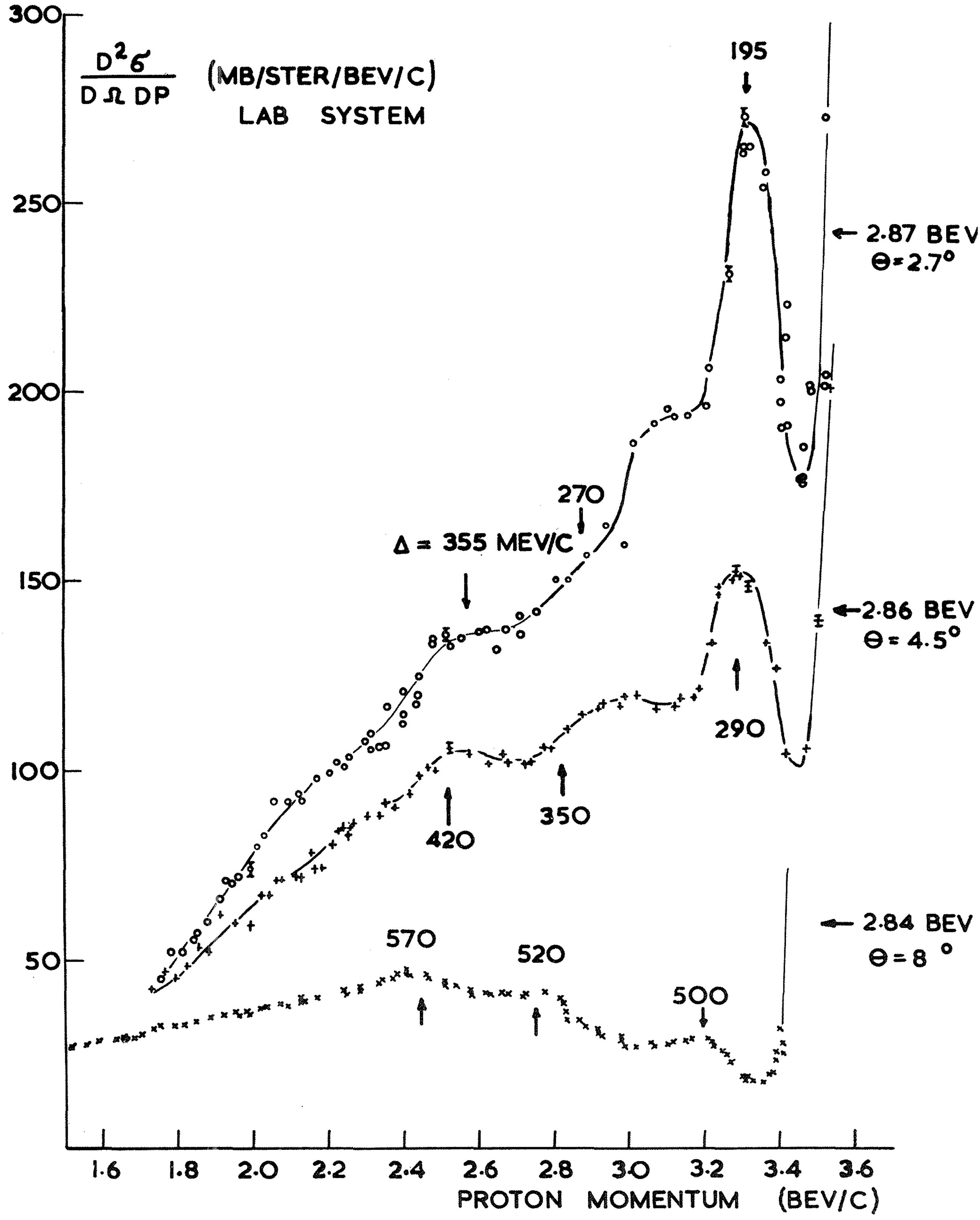


Figure 9 - Spectre des protons émis aux petits angles dans les chocs p - p [18].

ce qui explique qu'il n'apparaisse pas dans certaines distributions. Ce processus est probablement différent de celui observé à plus haute énergie [19] et qui implique une diffraction à l'un des vertex.

2 - PHOTOPRODUCTION DOUBLE

Kilner et al [20] à Pasadena ont étudié la réaction :

$$\gamma + p \longrightarrow \pi^- + \pi^+ + p \qquad \text{à} \qquad E_\gamma = 1\,230 \text{ Mev.}$$

Si les processus de Drell (diagramme 5) sont importants pour la production des π^- on doit observer un maximum de production aux petits angles.

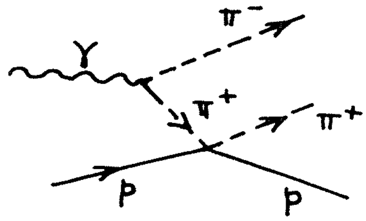


Diagramme 5

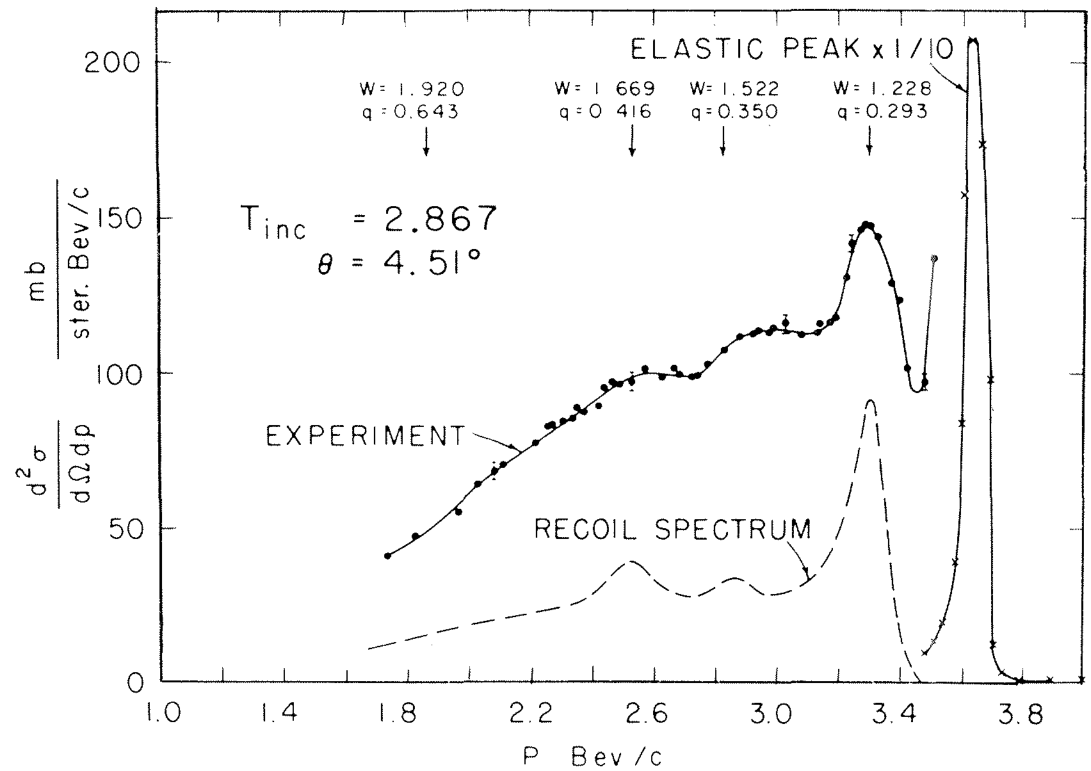


Figure 10 - Spectre des protons émis aux petits angles dans les chocs p - p [18]. La ligne pointillée correspond à la contribution calculée du diagramme 3.

Drell prédit :

$$\sigma(\vartheta, \omega) = \frac{\alpha}{8\pi^2} \frac{\sin^2\vartheta}{(1-\beta\cos\vartheta)^2} \frac{\omega(k-\omega)}{k^3} \sigma(\pi^+ + p)$$

où ω est l'énergie totale des π^- dans le CM ;

ϑ l'angle des π^- dans le CM ;

$Q = Q_{\pi^+p}$ la valeur Q du système π^+ proton ;

k l'énergie du photon incident.

La figure 11 montre quelques unes de ces distributions ainsi qu'à 17°2 la prédiction de Drell avec un maximum pour a = 145 Mev correspondant à la résonance A.

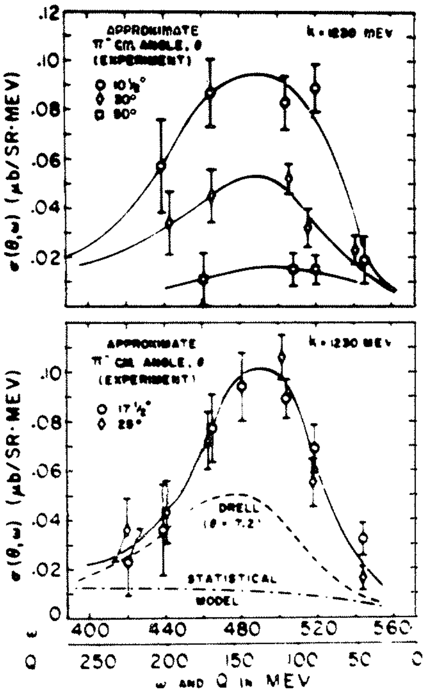


Figure 11 - Section efficace différentielle des π^- de la réaction $\gamma + p \longrightarrow \pi^- + \pi^+ + p$ à 1 230 Mev [20]. $\sigma(\theta, \omega)\, d\Omega\, d\omega$ est la section efficace de production d'un π- à un angle θ à l'intérieur de $d\Omega$ et avec une énergie ω à l'intérieur de $d\omega$, le tout dans le centre de masse. Les courbes pleines sont dessinées à travers les points expérimentaux. Les courbes pointillées sont calculées à partir du modèle de Drell (pour θ = 17,2°) et pour un modèle statistique isotrope.

La figure 12 montre $\sigma(\vartheta, \omega)$ pour $\omega \sim 490$ Mev.

Tout modèle isobarique doit donner la même dépendance en ω mais la rapide croissance pour ϑ petit est caractéristique des processus de Drell.

3 - PRODUCTION SIMPLE DES π DANS LES COLLISIONS π -NUCLEON

Les réactions du type $\pi + p \longrightarrow N + \pi + \pi$ sont déterminées à la fois par la forte interaction π-nucléon dans l'état 3/2 - 3/2 et par l'existence d'éventuelles interactions $\pi\pi$. Les évidences les plus claires de l'existence d'une forte interaction $\pi\pi$ dans l'état de spin isotopique I = 1 ont été obtenues cette année dans l'étude de cette réaction aux énergies supérieures à 1 GeV. Comme nous y reviendrons tout à l'heure, nous ne considérerons pour le moment que les résultats obtenus pour des énergies inférieures au GeV. La situation expérimentale a été résumée par Omnès [3]. Les principaux points sont les suivants : Pour expliquer la variation des coefficients a et les valeurs des sections efficaces inélastiques au voisinage des résonances π^--p on est amené à supposer que la production a lieu essentiellement par l'échange d'un méson π suivant le diagramme 6. Si c'est

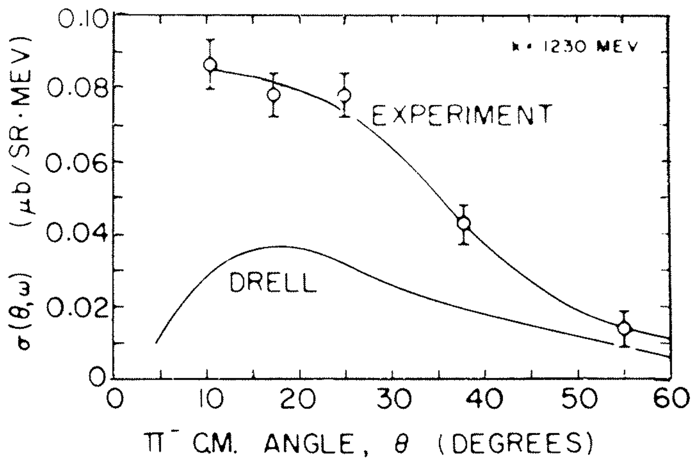


Figure 12 - Moyenne de la section efficace différentielle des π^- de la réaction $\gamma + p \longrightarrow \pi^- + \pi^+ + p$ à 1 230 MeV [20] pour les points à ω = 464, 505 et 619 MeV. La prédiction du modèle de Drell moyennée de la même manière en incluant la résolution expérimentale est donnée par la courbe du bas.

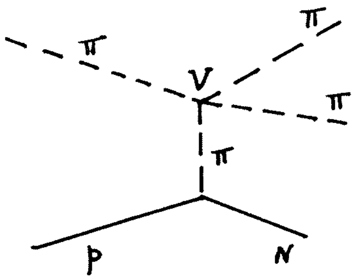


Diagramme 6

là le mécanisme primaire de production on peut prédire les rapports de branchement des différentes voies, selon que l'interaction $\pi\pi$ au vertex V a lieu dans un état de spin isotopique I = 0, 1 ou 2. Ces rapports de branchement sont donnés dans le Tableau I.

Tableau I

état initial		$\pi^- p$			$\pi^+ p$	
état final		$\pi^-\pi^+ n$	$\pi^-\pi^\circ p$	$\pi^\circ\pi^\circ n$	$\pi^+\pi^\circ p$	$\pi^+\pi^+ n$
pour une interaction $\pi\pi$ au vertex V dans l'état pur de spin isotopique	I = 0	2	0	1	0	0
	I = 1	2	1	0	1	0
	I = 2	2/9	0	4/9	2	8

En fait, dans ce domaine d'énergie l'un des 2 π aura souvent par rapport au nucléon une énergie voisine de celle de la résonance c'est-à-dire que la rediffusion du π par le nucléon dans l'état I = 3/2 aura une grande probabilité, ce qui conduit à accorder de l'importance au diagramme 7

Mais ceci ne changera pas beaucoup les rapports de branchement car l'échange de charge dans l'état I = 3/2 est petit.

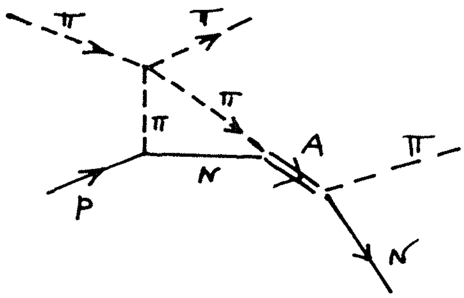


Diagramme 7

Le fait que la rediffusion est importante est apparent dans les figures 13 et 14. La première montre le spectre de π^+ de la réaction $\pi^- + p \longrightarrow \pi^- + \pi^+ + n$ obtenu par Barish et al [21] à 365 Mev. La seconde montre le spectre des π° obtenu par le groupe de Saclay pour la réaction $\pi^+ + p \longrightarrow \pi^+ + \pi^\circ + p$ à 900 Mev [22].

Mais à ces énergies, il existe déjà des évidences d'une interaction $\pi\pi$. Pour rendre compte de la forte section efficace $\pi^- + p \longrightarrow \pi^- + \pi^+ + n$ entre le seuil et 450 Mev, telle qu'elle avait été obtenue par Perkins et confirmée par Barish [21], Rodberg avait été conduit à admettre l'existence d'une forte interaction $\pi\pi$. Les calculs de Schnitzer ont donné un bon accord avec l'expérience en utilisant à la fois des ondes S dans l'état I = 0 et I = 2 et des ondes P dans l'état I = 1. La plus importante étant l'interaction I = 0 dans l'onde S l'accord étant obtenu pour une longueur de scattering de $a_\circ = 0,5 \dfrac{\hbar}{\mu c}$.

Le rapport de branchement des réactions $\pi^- p$ $r = \dfrac{\pi^\circ\ \pi^\circ}{\pi^-\pi^+ + \pi^-\pi^\circ}$ que l'on déduit des mesures de

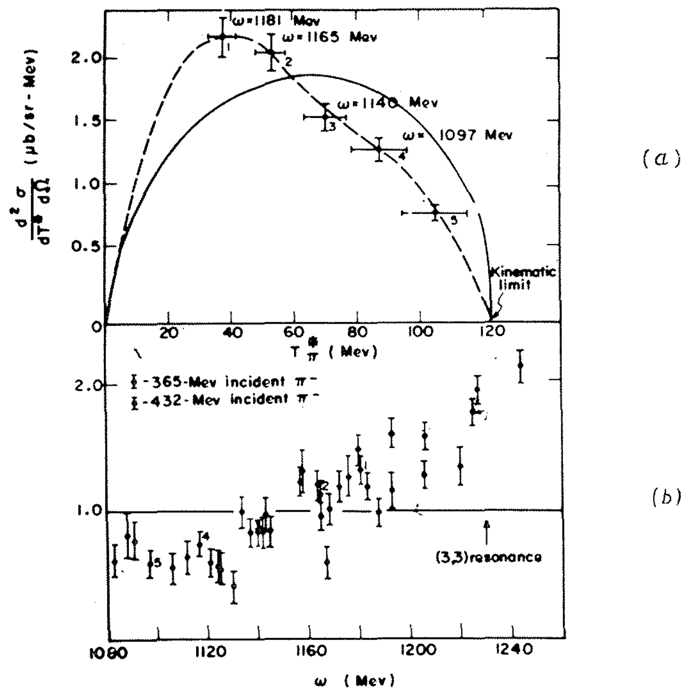


Figure 13 - (a) Spectre des π^+ de la réaction $\pi^- + p \longrightarrow \pi^+ + \pi^- + p$ à 365 Mev [21].
La courbe en trait plein est une distribution d'espace de phase normalisée à $d\sigma/d\Omega^*$.
(b) Rapport de $d^2\sigma/dT^*\, d\Omega^*$ à la valeur déduite de l'espace de phase en fonction de ω énergie totale du π^- et du neutron dans leur propre centre de masse.

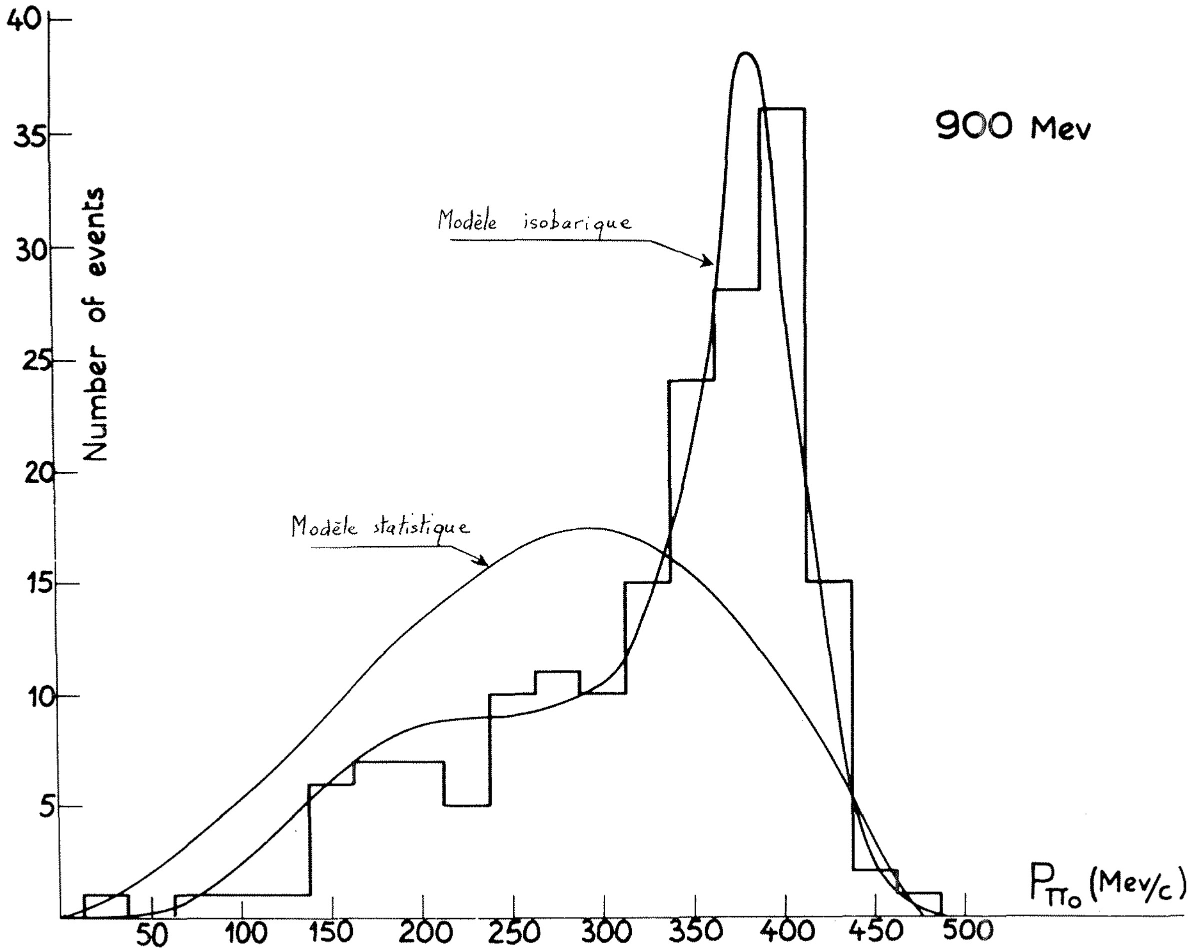


Figure 14 - Spectre des π^o de la réaction $\pi^+ + p \longrightarrow \pi^+ + \pi^o + p$ à 900 MeV [22].

Brisson et al [23] décroît continuement de 0,3 à 600 Mev à 0,15 à 1 100 Mev, indiquant que l'on a besoin d'interaction $\pi\pi$ dans l'état I = 0 à 600 Mev mais que l'importance de l'état I = 1 croît lorsqu'on monte en énergie. Le rapport de branchement dans la même réaction $R = \frac{\pi^-\pi^o}{\pi^-\pi^+}$ est de l'ordre de 0,5 à 1 Gev ce qui indique la nécessité d'interaction dans l'état I = 1.

Enfin la différence de comportement entre les réactions $\pi^+ + p$ et $\pi^- + p$ dans la diffusion élastique vers l'avant aux environs de 900 Mev telle qu'elle a été présentée par Meyer [12], [22] peut très bien être attribuée à l'existence d'une interaction $\pi\pi$ dans l'état T = 0 qui serait absente dans les réactions π^+p et présente dans les réactions π^-p.

Il reste encore beaucoup de travail à faire pour élucider le mécanisme exact des résonances et de la production au-dessous de 1 Gev. Certaines des expériences à faire ont été proposées par Omnès [3].

IV - LES RESONANCES DANS LES SYSTEMES DE PLUSIEURS π

1 - L'ETAT I = 1 DU DIPION DANS LA PRODUCTION DES π POUR LES CHOCS π NUCLEON

L'existence d'un état résonnant du système π-π dans l'état I = 1 ayant une masse d'environ 750 Mev a été mise en évidence de façon très claire pour la première fois cette année par le groupe de Wisconsin [24] dans la réaction $\pi + p \longrightarrow N + \pi + \pi$ à 1,9 Gev.

Le diagramme 6 correspondant à une interaction périphérique avec échange d'un pion doit être important, tout au moins pour les petites valeurs du transfert d'impulsion. Par suite de l'existence d'un pôle d'ordre 2 dans la matrice S pour $\Delta^2 = -\mu^2$ la distribution des évènements doit obéir à la relation [25]

$$\lim_{\Delta^2 \to -1} \frac{\partial^2 \sigma}{\partial \Delta^2 \, \partial \omega^2} = A f^2 \frac{\Delta^2}{(\Delta^2 + 1)^2} \frac{\omega}{q_{i_L}^2} K \sigma_{\pi\pi}(\omega) \tag{1}$$

où Δ est le quadritransfert d'impulsion, $\Delta^2 = 2 M T_N$ où T_N est l'énergie cinétique du nucléon de recul dans le laboratoire, M la masse du nucléon, ω est l'énergie totale des π dans leur propre centre de masse $\omega = 2\mu + Q_{\pi\pi}$, f est la constante de couplage π nucléon, q_{i_L} est l'impulsion du proton incident dans le laboratoire, $K = \sqrt{\frac{1}{4}\omega^2 - \mu^2}$ est l'impulsion de l'un ou l'autre des 2π dans leur centre de masse, $\sigma_{\pi\pi}$ une moyenne des sections efficaces $\pi\pi$ sur les différents états de spin isotopiques intervenant dans le processus particulier considéré, A un facteur numérique dépendant du processus particulier considéré.

Dans la formule, μ la masse du π a été prise égale à 1.

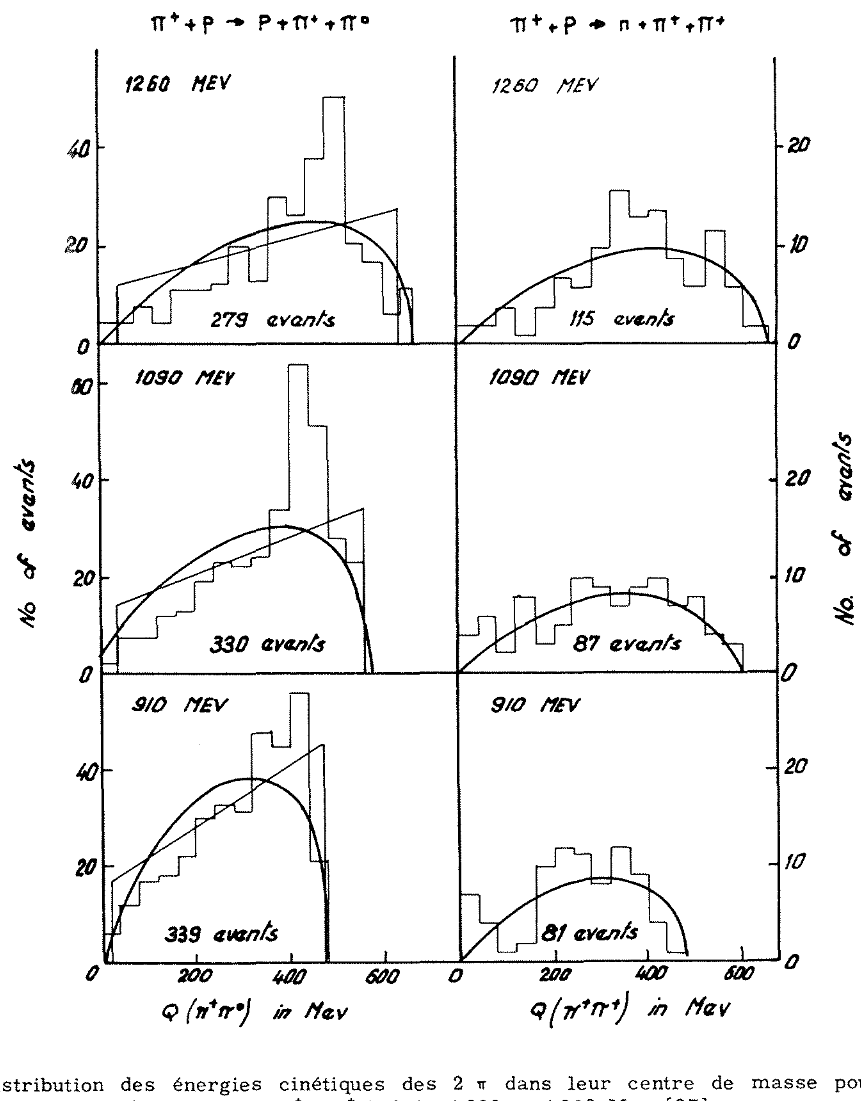


Figure 15 - Distribution des énergies cinétiques des 2π dans leur centre de masse pour les réactions $\pi^+ + p \longrightarrow \pi^+ + \pi^0 + p$ et $\pi^+ + p \longrightarrow n + \pi^+ + \pi^+$ à 910, 1090 et 1260 Mev [27].

Si de plus on est optimiste et si l'on a l'espoir que même dans la région physique seul le diagramme 6 est important, on peut laisser tomber le passage à la limite et le remplacer par les mots : "pour les petites valeurs de Δ^2". La relation

$$\frac{\partial^2 \sigma}{\partial \Delta^2 \, \partial \omega^2} = A f^2 \, \frac{\Delta^2}{(\Delta^2 + 1)^2} \, \frac{\omega}{q_{i_L}^2} \, K \, \sigma_{\pi\pi}(\omega)$$

pour les petites valeurs de Δ^2(2) permet alors après intégration sur Δ^2 de calculer $\sigma_{\pi\pi}$ lorsqu'on connaît la distribution des évènements.

Par suite, l'existence d'une forte interaction $\pi\pi$ est caractérisée d'une part par une forte proportion d'évènements avec petite impulsion de transfert, ce qui se voit sur la formule par le fort maximum de $\frac{\Delta^2}{(\Delta^2 + 1)^2}$ pour $\Delta^2 = + 1$; d'autre part par des maxima dans la distribution en ω correspondant aux maximas de $\sigma_{\pi\pi}$. Entre 0,9 et 1 Gev de telles anomalies avaient déjà été observées par différents auteurs [26].

La figure 15 montre le spectre de $Q_{\pi\pi}$ obtenu par le groupe de Yale [27] pour différentes énergies de π^+. On voit que l'anomalie qui apparaît dans $Q_{\pi^+\pi^0}$ et n'apparaît pas dans $Q_{\pi^+\pi^+}$, suggérant une interaction dans l'état I = 1, se déplace vers les grandes valeurs de Q lorsqu'on monte en énergie. A 1260 Mev le maximum est vers 730 Mev [28]. Ceci suggère qu'aux énergies inférieures à 1,2 Bev les mesures antérieures étaient entachées d'une sorte de biais physique lié à l'espace de phase.

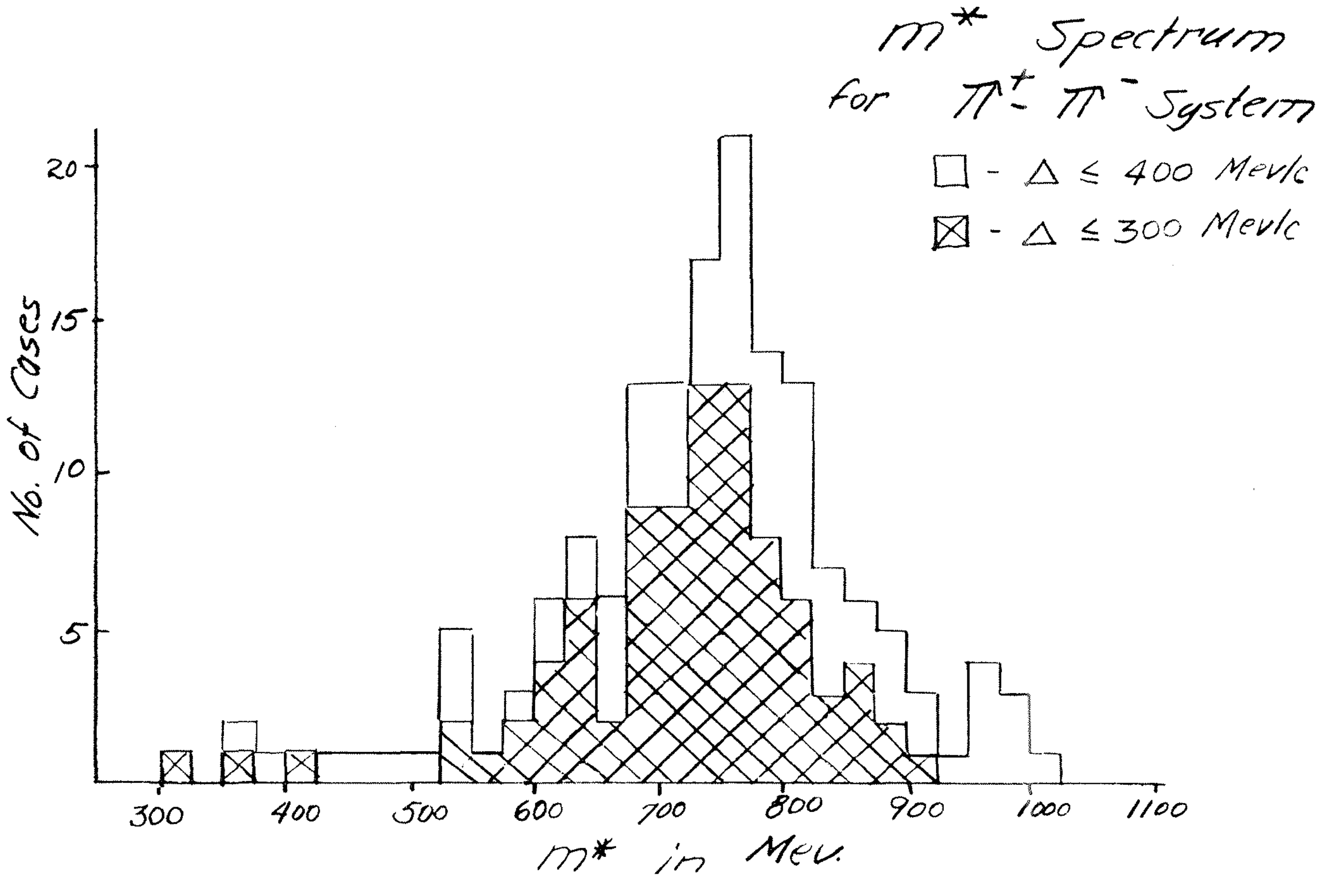


Figure 16 - Distribution en $m^* = \omega$ des évènements $\pi^- + p \longrightarrow \pi^- + \pi^+ + n$ à 1,9 GeV [24] pour les petits quadritransferts d'impulsion Δ .

Le groupe de Wisconsin et Brookhaven [24] étudiant les réactions

$$\pi^- + p \longrightarrow \pi^- + \pi^\circ + p \qquad (1)$$

et

$$\pi^- + p \longrightarrow \pi^- + \pi^+ + n \qquad (2)$$

à 1,9 Gev trouvèrent une distribution des nucléons piquée vers l'arrière dans le centre de masse. Les évènements pour lesquels le transfert d'impulsion était inférieur à 400 Mev/c montrait un fort maximum à 765 Mev. Les premiers résultats concernaient l'ensemble des 2 réactions (1) et (2). Les nouveaux résultats présentés sur la figure 16 se rapportent uniquement à la réaction (2).

Par intégration de la formule (2) ces auteurs obtiennent la section efficace $\sigma_{\pi\pi}$ de la figure 17 avec un maximum à ω = 750 Mev soit $\omega^2 = 29\,\mu^2$ et une largeur de 150 Mev environ. Malheureusement, en abandonnant le passage à la limite dans la formule (1) on a retiré toute rigueur à la formule et l'on n'est plus sûr de sa validité. On ignore l'influence des autres diagrammes en particulier celui de la rediffusion. Il est plus prudent de dire que ce qui est porté en ordonnée sur la figure 17 et sur les figures suivantes n'est pas $\sigma_{\pi\pi}$, mais le résultat de l'intégration de la formule (2) par rapport à Δ^2. C'est pour le moment la meilleure approximation de $\sigma_{\pi\pi}$ que nous ayons, mais nous ne savons pas jusqu'à quel point elle est mauvaise. Les évènements correspondant à 400 Mev/c sont dans le rapport $\frac{\pi^+\pi^-}{\pi^-\pi^\circ} = 1,7 \pm 0,3$ avec $\frac{\pi^\circ\pi^\circ}{\pi^-\pi^\circ} < 0,25$ la comparaison avec le tableau 1 montre que l'interaction a lieu dans l'état I = 1.

Les résultats de Erwin ont été confirmés à 1,25 Gev par Pickup [29] (figure 18) et par les

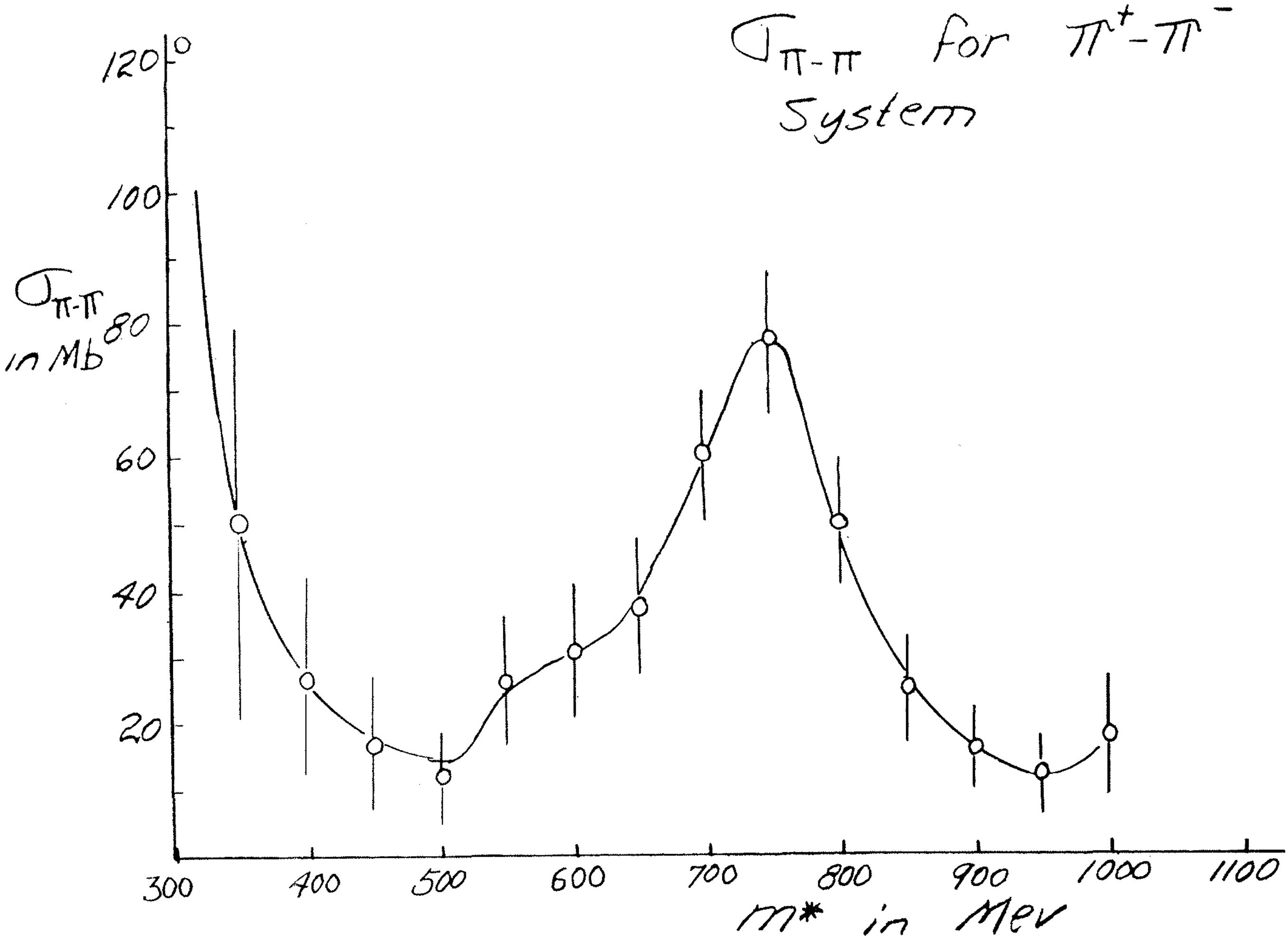


Figure 17 - Résultat de l'intégration de la formule (2) par rapport à Δ pour la réaction $\pi^- + p \longrightarrow \pi^- + \pi^+ + n$ à 1,9 GeV [24].

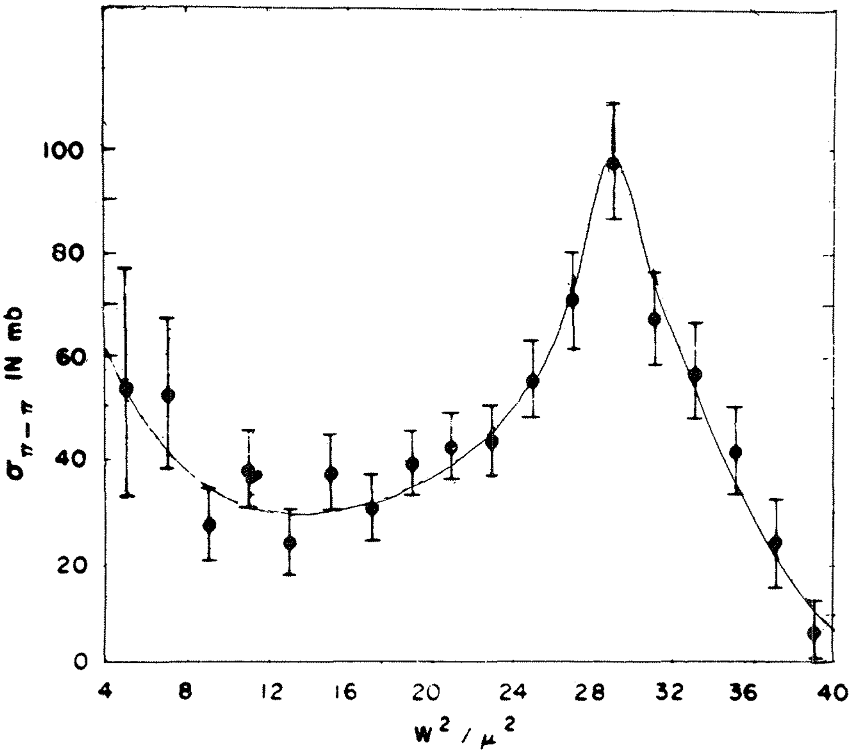


Figure 18 - Résultat de l'intégration de la formule (2) par rapport à Δ pour les réactions $\pi^- + p \longrightarrow \pi^- + \pi^+ + n$ et $\pi^- + p \longrightarrow \pi^- + \pi^\circ + p$ à 1,25 GeV [24] dans lesquels $\Delta^2 < 10\,\mu^2$.

groupes de Bari, Bologne, Orsay et Saclay (figure 19). Le maximum de $\sigma_{\pi\pi}$ est compris entre 80 et 100 mb alors que pour une résonance dans l'état I = 1 J = 1 il devrait être de $12\,\pi\,\lambda\!\!\!^{-2}$ = 120 mb. Compte tenu de l'incertitude dans l'interprétation de $\sigma_{\pi\pi}$ ceci est une indication du spin J = 1 de la résonance. Les différents auteurs ont essayé une analyse d'Adair pour obtenir une autre détermination du spin, la figure 20 montre les résultats du dernier groupe avec une distribution en $1 + 4\cos^2$ qui montre à côté de l'existence d'une résonance J = 1 l'existence d'un fond continu.

Des évidences de cette résonance I = 1, J = 1, $m^2 = 29\,\mu^2$ ont été également trouvées par Maglic et al. [31] dans l'annihilation en 5π des antiprotons en vol et par Cresti [32] dans l'annihilation des antiprotons à l'arrêt.

2 - L'ETAT I = 0 DANS LES COLLISIONS PROTONS DEUTERONS

C'est une technique complètement différente qui a amené Abashian, Booth et Crowe [33] [34] à postuler l'existence d'une interaction $\pi\pi$ dans l'état I = 0.

Ils étudient le spectre en impulsion des noyaux He^3 ou H_3 dans les réactions

$$p + d \longrightarrow He_3 + X \qquad (1)$$

et

$$p + d \longrightarrow H_3 + X \qquad (2)$$

où X est soit l'ensemble de deux π indépendants, soit une particule inconnue. Ils obtiennent la partie correspondant à l'état I = 0 de X en soustrayant le spectre de H_3 divisé par 2 de celui de He_3. La figure 21 montre le spectre de He_3 et H_3. Le continuum correspond aux réactions à 3 corps où X représente deux π indépendants. L'anomalie à 1 400 Mev/c dans le spectre de He_3 est attribué à une particule inconnue. La figure 22 représente la partie du spectre correspondant à l'état de

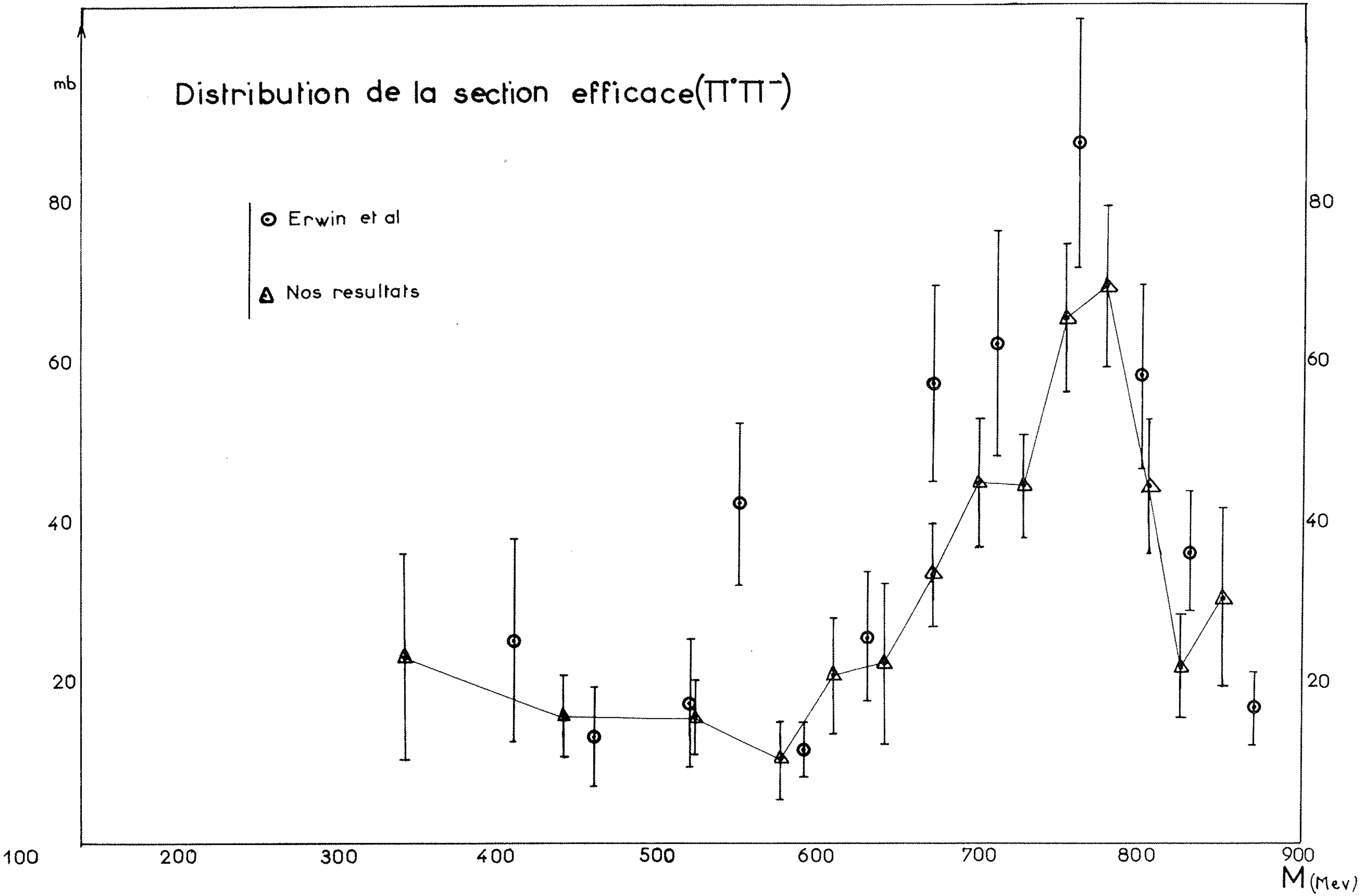


Figure 19 - Résultat de l'intégration de la formule (2) par rapport à Δ pour la réaction $\pi^- + p \longrightarrow \pi^- + \pi^\circ + p$ à 1,47 GeV [30].

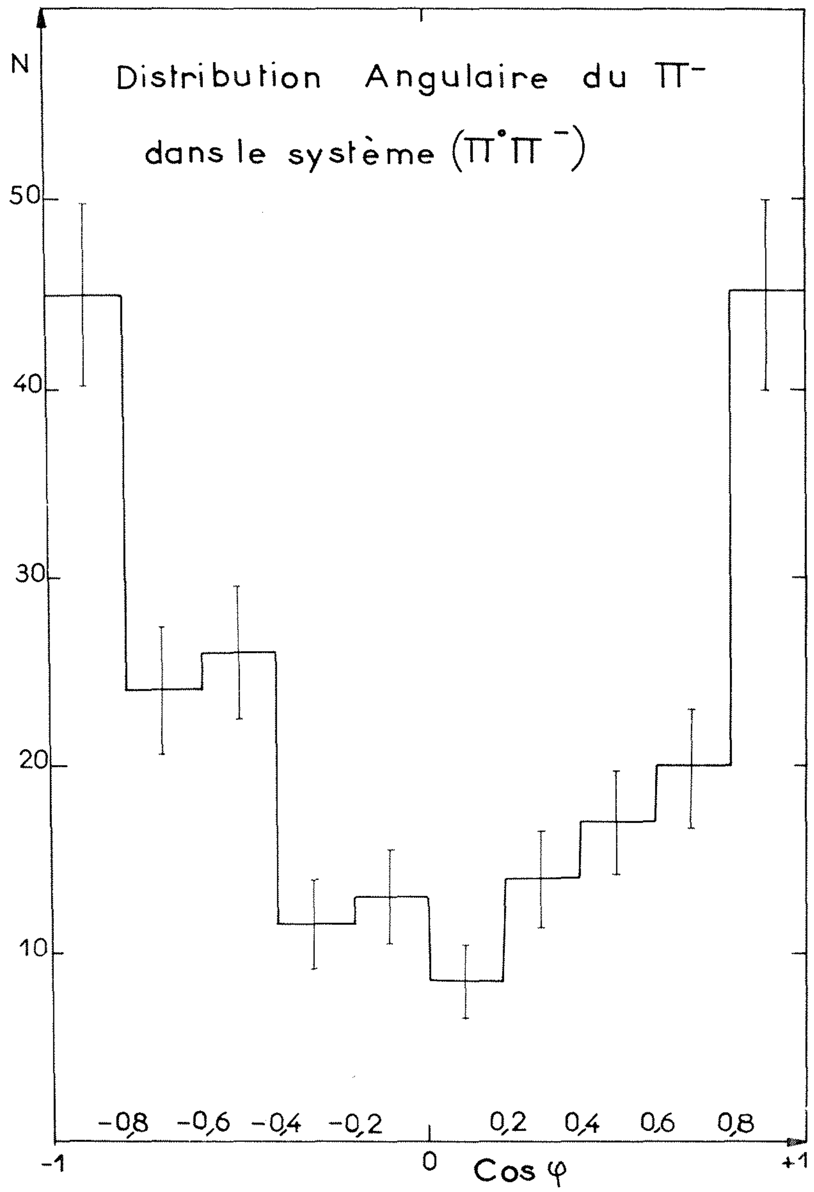
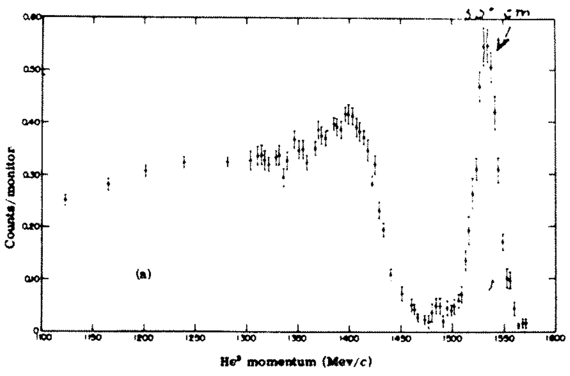


Figure 20 - Distribution angulaire dans le centre de masse $\pi^-\pi^o$ du π^- de la réaction $\pi^- + p \longrightarrow \pi^- + \pi^o + p$ à 1,47 GeV [30].



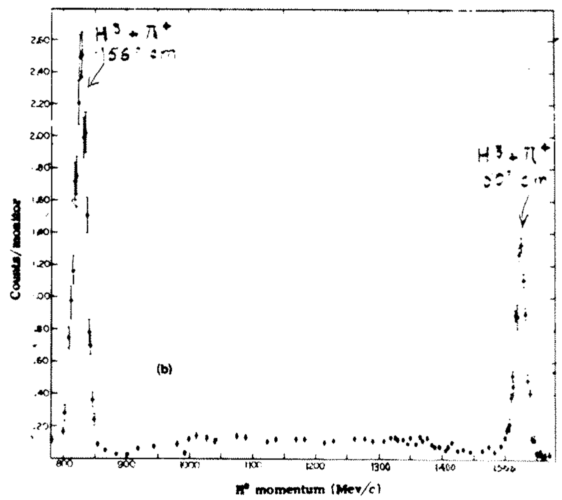


Figure 21 - Spectre des H_3 et He_3 des réactions (1) et (2) à 11,8° dans le laboratoire [34].

spin isotopique I = 0. La ligne pointillée représente l'espace de phase normalisé aux basses énergies. La courbe pleine représente le résultat d'un calcul d'amplification d'espace de phase dû à l'interaction des 2 π dans l'état final en supposant que cette interaction a lieu dans l'état S avec une longueur de scattering de $2,8\ \frac{\hbar}{\mu c}$. Comme Fubini [35] l'a indiqué, cette longueur de scattering est trop grande pour rendre compte de la production des π dans les diffusions $\pi^- - p$ au voisinage du seuil. Une autre interprétation de cette anomalie est qu'elle est due à l'existence d'une particule instable X de spin isotopique O. La figure 21 montre la différence entre l'espace de phase et les points expérimentaux de la figure 22. La masse obtenue pour cette particule serait m = 300 Mev, soit $m^2 = 4,6\ \mu^2$. La largeur serait de 25 Mev environ.

3 - L'ETAT I = 0 DU TRIPION DANS L'ANNIHILATION DES ANTIPROTONS EN 5 π ET DANS LES REACTIONS $\pi^+ + d$

Maglic et al [31] étudiant l'annihilation des antiprotons de 1,6 Gev/c en 5π suivant la réaction

$$\bar{p} + p \longrightarrow \pi^+ + \pi^+ + \pi^- + \pi^- + \pi^0$$

regardent la masse effective de chacun des triplets de π possible. Ces triplets sont classés suivant la valeur absolue de leur charge dans les 3 classes

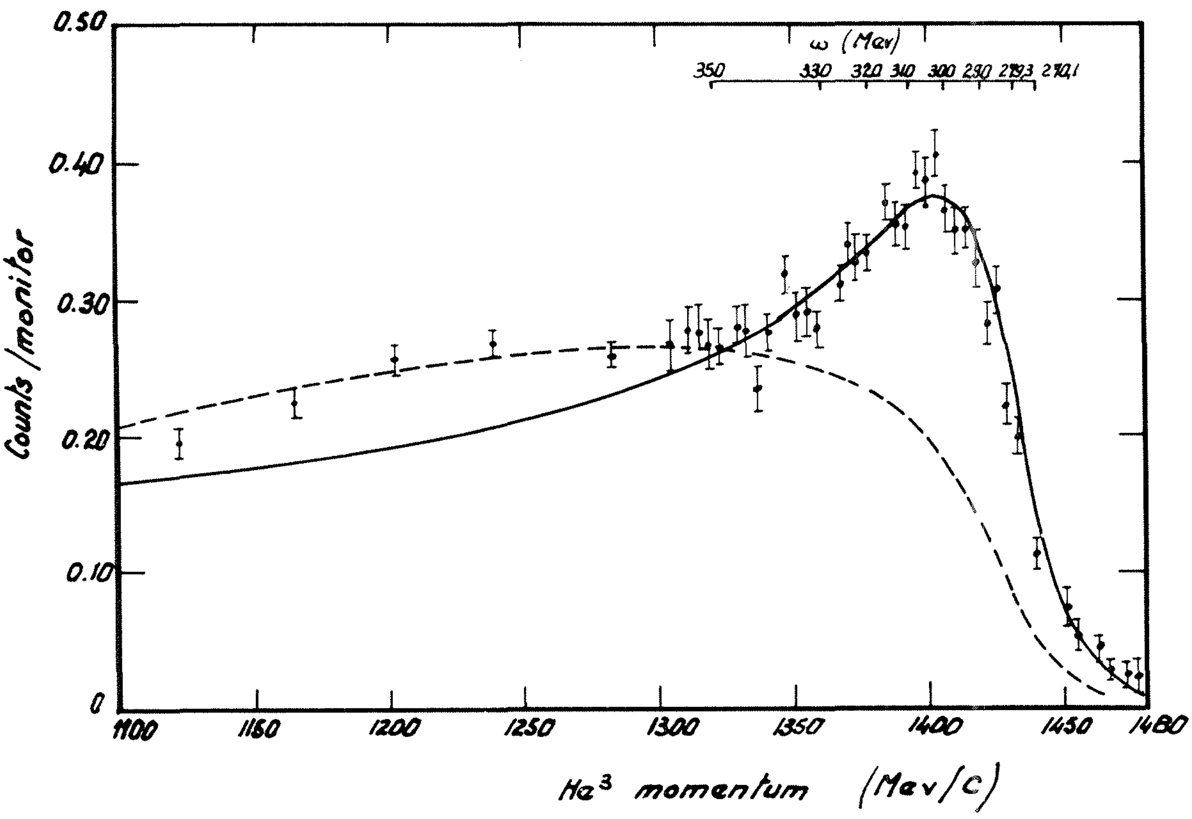


Figure 22 - Partie du spectre He_3 correspondant à l'état I = 0. La courbe pointillée est le volume d'espace de phase, normalisé sur les points inférieurs à 1 300 MeV/c. La courbe en trait plein est le volume d'espace de phase multiplié par le facteur d'amplification dû à l'intéraction $\pi - \pi$ pour une longueur de scattering $a_{so} = 2,8 \frac{\hbar}{\mu c}$. L'échelle W donne l'énergie totale des 2 π dans leur centre de masse.

A	$\lvert Q \rvert = 1$	+ + - - - +
B	$\lvert Q \rvert = 2$	+ + 0 - - 0
C	$\lvert Q \rvert = 0$	+ - 0

La figure 24 montre les spectres des évènements obtenus pour chacune des 3 classes. La courbe continue correspond approximativement à l'espace de phase. La distribution des évènements neutres montre un pic étroit à 787 Mev, sa largeur est inférieure à 30 Mev. Pour essayer de déterminer le spin de cette résonance les auteurs ont fait un diagramme de Dalitz analogue à celui utilisé pour la désintégration du τ. Ce diagramme est reporté sur la figure 25. Les diagrammes A et B correspondent à une bande témoin prise en dehors du pic. Les auteurs concluent à une valeur J = 1 et à une parité - du tripion, le raisonnement étant fondé sur l'existence d'une densité normale d'évènement le long des bords rectilignes du diagramme C et sur la dépopulation dans les coins inférieurs et le long du bord curviligne de ce diagramme.

L'existence de cette résonance à 3 π a été confirmée par le groupe de John Hopkins [36] dans les réactions $\pi^+ + D$.

La figure 26 représente le spectre des masses effectives des 3 π de la réaction

$$\pi^+ + D \longrightarrow p_{spec} + p + \pi^+ + \pi^- + \pi^0$$

où l'un des protons est spectateur, c'est-à-dire qu'il prend très peu de recul. Le pic à 780 Mev est nettement visible. La figure 27 montre le spectre de la masse effective des particules neutres dans la réaction

$$\pi^+ + D \longrightarrow p_{spec} + p + (\text{neutres})$$

6

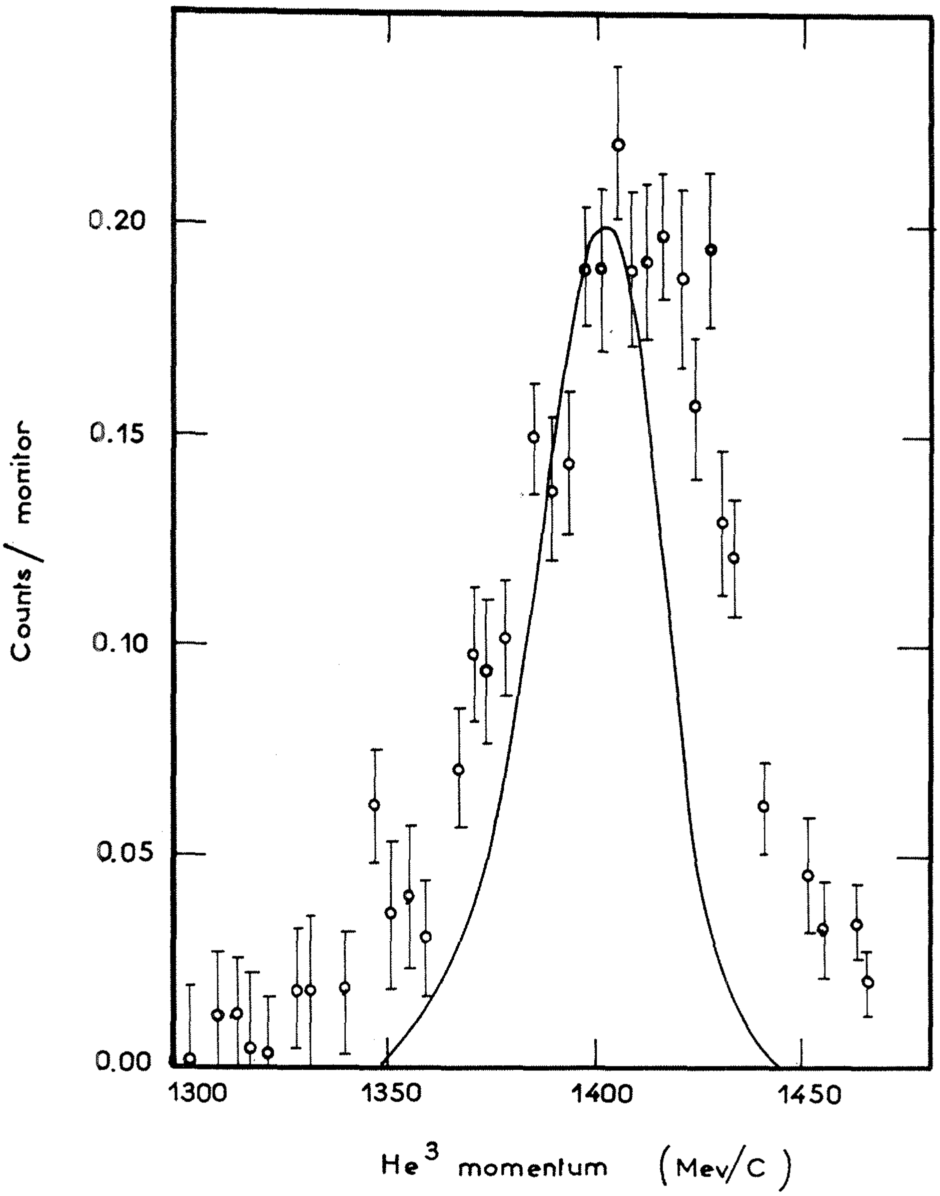


Figure 23 - Résultat de la soustraction du volume d'espace de phase des résultats expérimentaux. La courbe en trait plein est la résolution expérimentale calculée à 1 400 MeV/c.

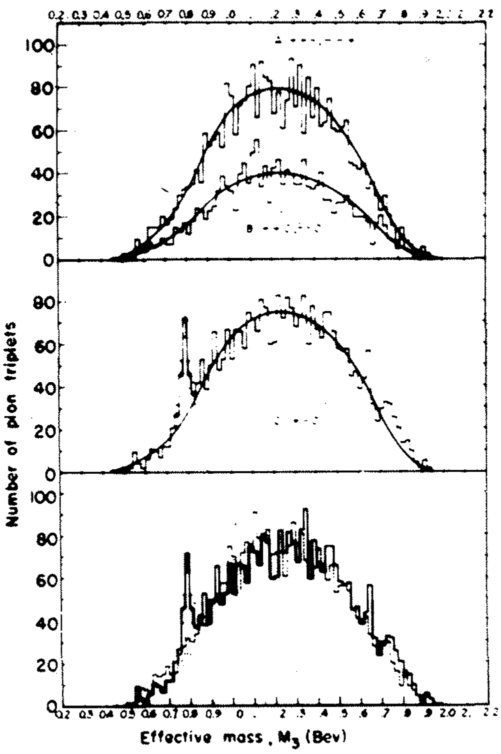


Figure 24 - Nombre de triplets de π en fonction de la masse effective de ces triplets pour la réaction $\bar{p} + p \longrightarrow 2\pi^+ + 2\pi^- + \pi^\circ$.
(A) distribution pour $|Q| = 1$.
(B) distribution pour $|Q| = 2$.
(C) distribution pour $|Q| = 0$.
Dans (D) les distributions combinées de (A) et (B) en pointillé sont comparées à la distribution (C) en trait plein.

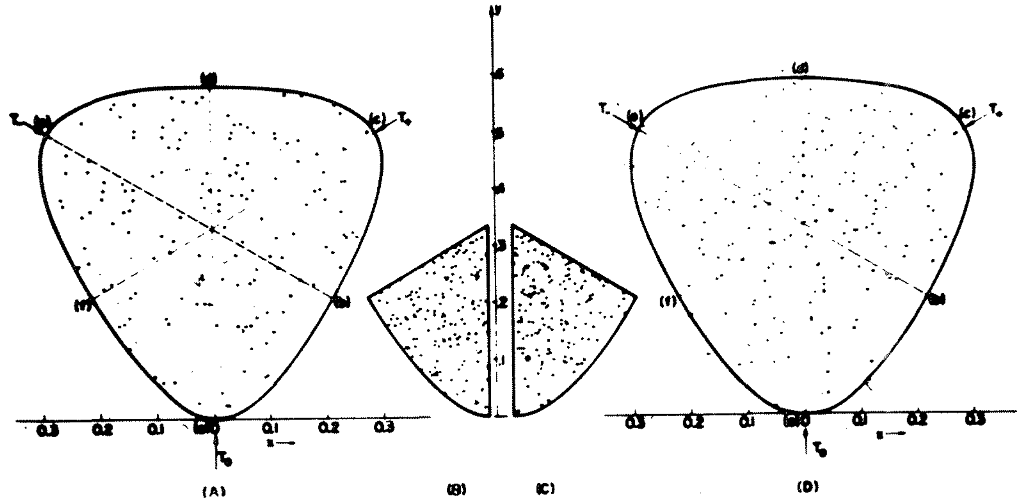


Figure 25 - (A) Diagramme de Dalitz pour 171 triplets de la région de contrôle ($820 \leq M_3 \leq 900$) ; (B) Diagramme replié de la région de contrôle, (C) Diagramme replié de la région du pic, (D) Diagramme de Dalitz pour 191 triplets dans la région du pic dont 43 ± 7 % sont dûs au tripion ω. T_+ T_- et T_0 sont les énergies cinétiques respectives du π^+, du π^- et du π°.

Ici le spectre à 750 Mev manque, ce qui s'explique s'il s'agit d'un état du tripion de spin isotopique O qui ne peut se désintégrer en 3 π° mais pour lequel la voie $\pi^+\pi^-\pi^\circ$ est très probable. L'accumulation d'évènements entre 500 et 600 Mev dans les 2 figures pourrait être due à un état à trois π pour lesquels les 2 modes de désintégration $\pi^+\pi^-\pi^\circ$ et $\pi^\circ + \gamma$ seraient compétitifs.

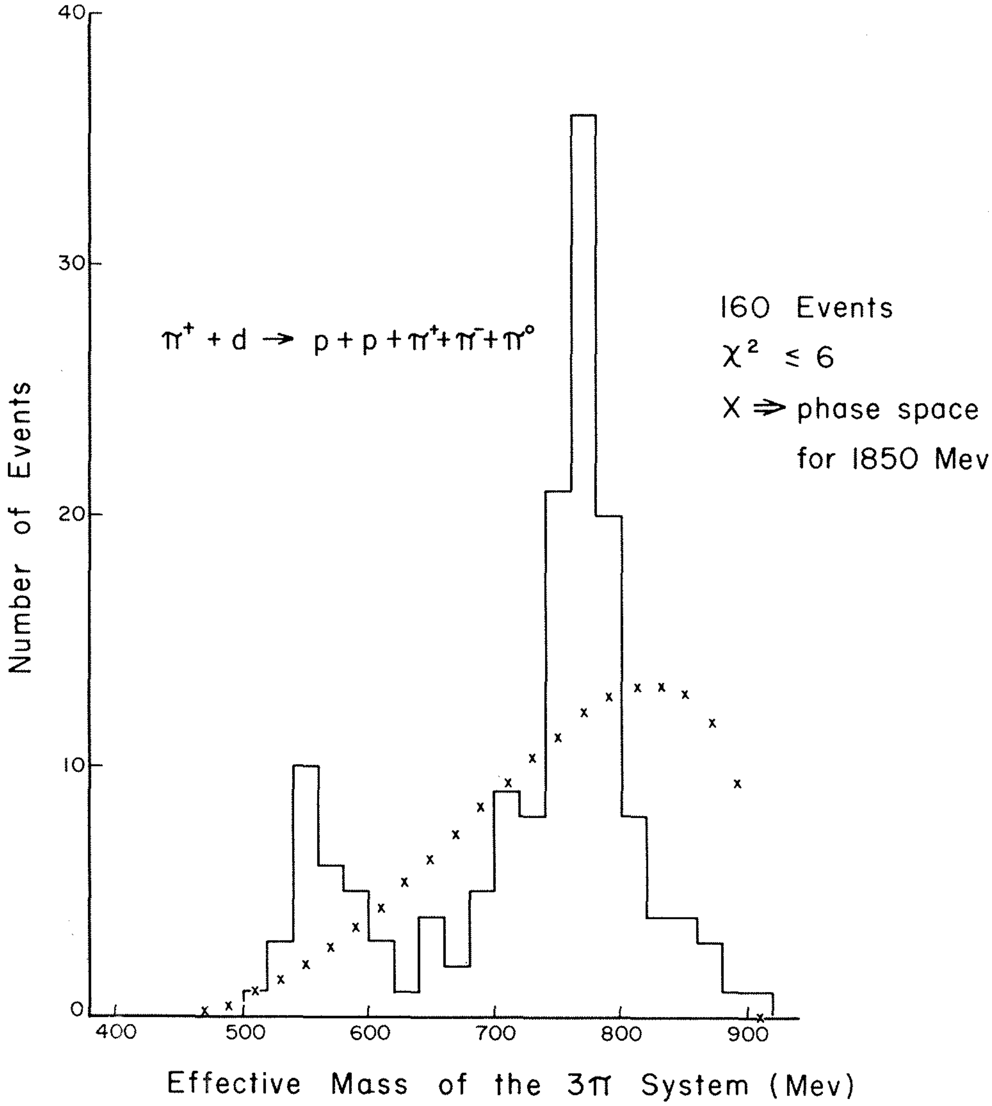


Figure 26 - Distribution des évènements $\pi^+ + d \longrightarrow p + p + \pi^+ + \pi^- + \pi^0$ pour lesquels le recul d'un des protons est faible [36] et classés d'après l'énergie totale des 3π dans leur centre de masse.

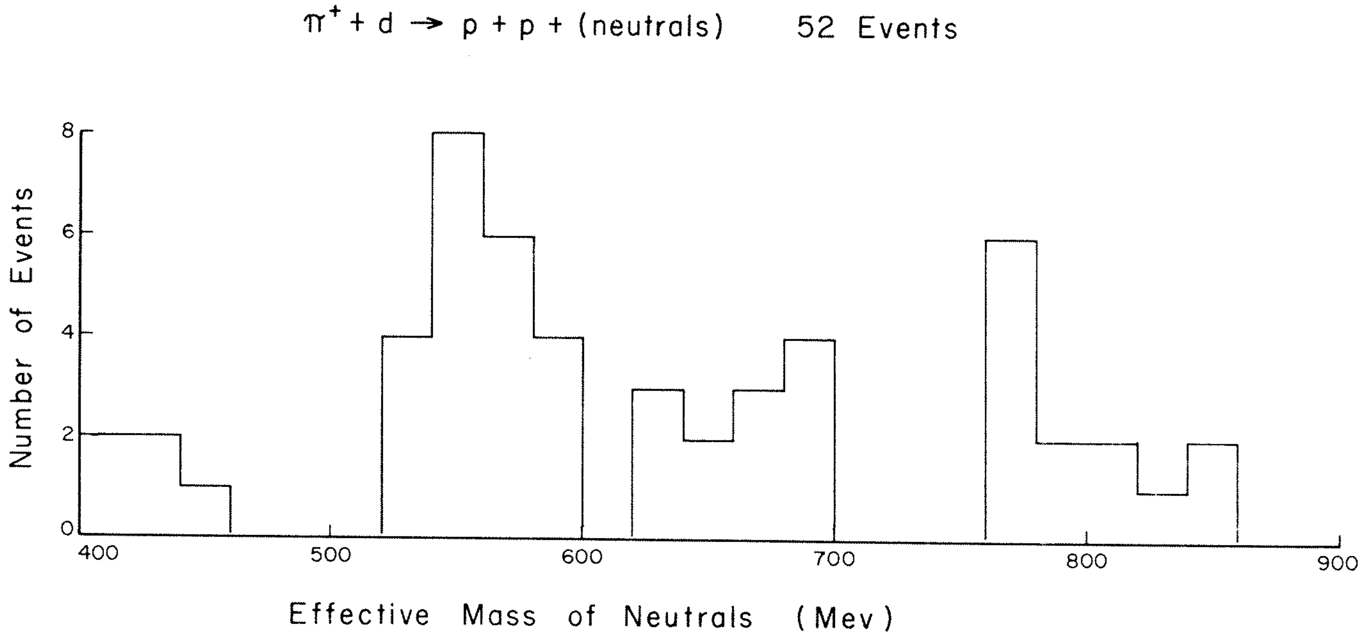


Figure 27 - Distribution des événements $\pi^+ + d \longrightarrow p + p +$ neutres pour lesquels le recul d'un des protons est faible [36] et classés d'après l'énergie totale des neutres dans leur centre de masse.

V - CONCLUSION

Les caractéristiques les plus probables des isobares π-nucléons (et des résonances des systèmes du π) sont résumées dans les tableaux suivants. Il se peut que dans un proche avenir ces tableaux s'allongent encore ; mais en ce qui concerne les résonances dont nous venons de parler il reste encore beaucoup de travail à faire, l'un des objectifs principaux étant d'obtenir la véritable section efficace π-π.

Tableau II

Isobares π-nucléon

	A	B	C	D
$T_{\pi\ \text{Mev}}$	200	605	890	1 300
I	3/2	1/2	1/2	(3/2)
J	3/2	3/2	5/2	
onde résonnante	P 3/2	D 3/2	F 5/2	

T_π énergie cinétique du π incident dans le laboratoire correspondant au maximum de la section efficace totale π-p.

Tableau III

Etats résonnants des systèmes de plusieurs π

multiplicité	2π	?	3π
I	1	0	0
J	1		(1)
Parité	+		(-)
m = masse en Mev	750	$(300)^*$	780
m^2	$29\mu^2$	$(46\mu^2)^*$	$31\mu^2$
Largeur à mi-hauteur	150 Mev	$(25 \text{ Mev})^*$	30 Mev

* Dans l'hypothèse où l'anomalie de Abashian, Booth et Crowe [34] n'est pas simplement due à une forte interaction π π dans l'état S.

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LOW ENERGY PION PHYSICS

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At the Geneva Conference in 1958 Chew reported on the Mandelstam [1] representation. This inaugurated a new era in pion physics, and there followed a large number of papers of great apparent complexity and erudition. By now the dust has settled and it is possible to see what is actually being done. In its simplest form this is rather simple, and it is perhaps an opportune moment to attempt a summary, and to draw attention to the types of experiment which would throw most light on the situation as we now see it.

There is of course still no theory of pion interactions, in the sense in which we have a theory of electrodynamics. All that is attempted is a semi-phenomenological analysis, which correlates the maximum amount of data with the minimum number of parameters. Apart from the obvious restrictions imposed by the conservation of angular momentum and isotopic spin, the main tools are the unitarity of the S-matrix, (conservation of probability), the relativistic covariance of the theory, and causality. Let us discuss unitarity first.

For a system with only one channel, the unitary condition shows that the elastic scattering, for a particular value of the angular momentum, J, and isotopic-spin, I, can be expressed in terms of a single real parameter, the phase shift, δ. The inverse scattering amplitude can be written as

$$T^{-1} = k\cot\delta - ik. \tag{1.1}$$

The imaginary part of the inverse amplitude is completely determined by unitarity. The dynamics of the system are described purely by the real part. For several two particle channels this result generalizes to

$$T^{-1}_{fi} = K^{-1}_{fi} - ik_f\,\delta_{fi} \tag{1.2}$$

where k_f is the centre of mass momentum in the fth channel and the imaginary part of the inverse matrix is again simply determined by unitarity. The matrix K^{-1} is the inverse K-matrix. It is real, and symmetric if the theory is invariant under time reversal. The elements of K are functions of the covariant $s = E_c^2$ (where E_c is the total centre of mass energy), and again completely describe the dynamics of the interaction.

If the system has a bound state, mass m_B (for example, the nucleon in pion-nucleon scattering), then K has a term which is formally identical with Born approximation for scattering via this intermediate state. We write this as

$$B = \frac{g^2}{s-s_B}, \tag{1.3}$$

where

$$s_B = m_B^2,$$

and g is the renormalized coupling constant between the bound state and the scattering particles. Note that B has a pole at the point s_B. This corresponds to an unphysical energy for scattering, but if it is near to the scattering region it may be expected to dominate the scattering in its neighbourhood. If we approximate K by this Born term, then multiplying through by B, Eqn. (1.2) can be written symbolically as

$$BT^{-1} = 1 - iBk. \tag{1.4}$$

This is precisely Heitler damping theory [2].

The momenta, k_f, when expressed in terms of the variable s involve square roots, whose arguments vanish at the thresholds for the different channels. Thus if we consider $T^{-1}(s)$ as a function in the complex s-plane, it can at best be analytic in a plane with a cut along the positive real axis, starting at the value of s corresponding to the lowest threshold of the system - the unitarity cut (figure 1). The Mandelstam conjecture, based on causality, is that the amplitude is in fact analytic in this cut plane (in this $K \simeq B$ approximation), and this enables one to modify (1.4) to the form of a dispersion relation [3]

$$BT^{-1}(s) = 1 - \frac{s-s_B}{\pi} \int_{s_T}^{\infty} \frac{B(s')k'}{(s'-s-i\varepsilon)(s'-s_B)} ds' \tag{1.5}$$

where s_T is the threshold for scattering, and we have incorporated the boundary condition that T(s) is given exactly by B at the pole (*). That is

$$BT^{-1}(s_B) = 1. \tag{1.6}$$

Note that if the position of the bound state and the strength of its coupling are known, (1.4) is an explicit expression for the scattering. This relation is valid at low energies (s slightly greater than s_T), provided s_B is reasonably close to s_T. We have so far merely reproduced in a covariant notation the well known result that, for example, low energy nucleon-nucleon scattering (for J = 1, I = 0) can be completely described by the position and coupling strength of the appropriate bound state - namely the deuteron (**). But this is a "shape independent" approximation and (1.5) must be modified to include some details of the shape of the potential.

More details of the form of the interaction appear in the theory in the following way. If we ignore spin, the complete amplitude for nucleon-nucleon scattering is a function J(s,t) of s and the invariant momentum transfer

$$t = -2k^2(1 - \cos\vartheta). \tag{1.7}$$

Covariance strongly suggests [4] that the amplitude for nucleon-antinucleon scattering is given by the same analytic function J(s,t), but with t now fixing the energy and s the momentum transfer, (see figure 2). Now the $n - \bar{n}$ system has a bound state - the single pion. The corresponding Born term, which in terms of the n - n is single pion exchange, (figure (3)a) should therefore be included in B, and accounts for the long range part of the n - n potential. The $n - \bar{n}$ system also has a unitarity cut (in the variable t) corresponding to the annihilation of $n - \bar{n}$ into two or more pions by the nucleon). After separating off the particular partial amplitude, T(s), this shows itself as an additional cut on the negative real axis of the s-plane-the lefthand cut. Thus figure 1 should be modified to figure 4 ; the lighter the mass of the particles exchanged, the closer the cut approaches the physical region. This lefthand cut gives rise to an additional integral in (1.5) which describes the effect of the 'potential' due to pion exchange. The integral here involves $mT^{-1}(s)$ for unphysical (negative) values of s ($=E_c^2$), where it is not determined by unitarity. This term is thus very far from being explicit, and it is in fact very difficult to evaluate its contributions even approximately in terms of the parameters already introduced (namely the renormalized coupling constants of the Born terms).

(*) This may be seen by evaluating

$$\oint \frac{B(s')T^{-1}(s')}{s'-s-i\varepsilon} ds'$$

around the contour illustrated in figure 1, using (1.2), (1.3) and (1.6). Since the sign of the imaginary part of T^{-1} is opposite above and below the unitarity cut : we have

$$B(s)T^{-1}(s) = \frac{1}{\pi} \int_{s_T}^{\infty} \frac{B(s')k'}{s'-s-i\varepsilon} ds' + C$$

where C is the contribution from the pole of B and the infinite circle of the contour. From this relation subtract the same relation evaluated at $s = s_B$ to obtain (1.5). From the relation

$$\frac{1}{s-i\varepsilon} = P\frac{1}{s} + i\pi\delta(s),$$

it can be seen that the imaginary part of T^{-1}, (uniquely determined by unitarity), is still given correctly, as in (1.4). The real part however is modified.

(**) The coupling strength of the deuteron is determined by the wave function normalization.
G.F. Chew and F.E. Low, Phys. Rev. 113, 1640, (1959).

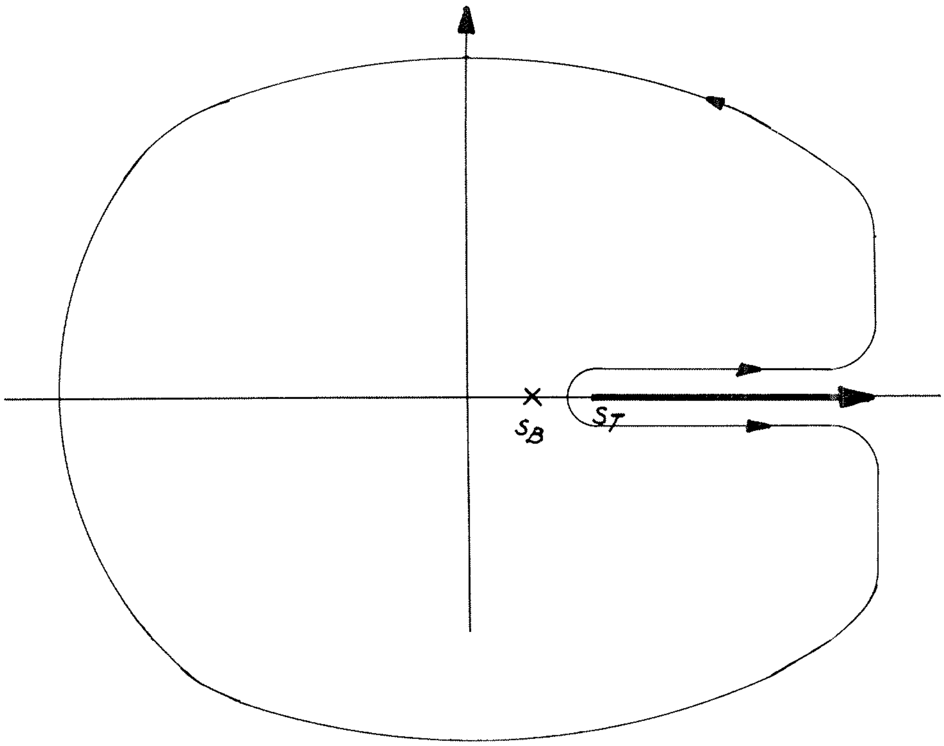


Figure 1 - Plot of the complex s-plane showing the unitarity cut, the Born pole at s_B and the contour which gives rise to the approximate dispersion equation (1.5).

However these terms can be included very simply if it is assumed that the 2π and 3π exchange takes place mainly through resonant states, which can be incorporated as pseudo particles and included, along with the single pion exchange term, in B (figure 4(b) and (c)). In this approximation we are back to the explicit formula (1.4) provided B, now includes the one pion and resonant 2π and 3π exchange terms. The price which has been paid is the introduction of several new parameters the positions of these resonances and their coupling strengths.

Let us emphasise again the simplicity of the model to which we have been led. The discussion above has been given in terms of the nucleon-nucleon interaction, but in fact applies to any low energy pion-nucleon system, where the long range interaction arises from pion exchange. All we have to do is to calculate lowest order Born approximation with the known stable particles and the 2π and 3π pseudo particles, and then make the result unitary be means of (1.5).

The real part of Eqn. (1.5) can be written

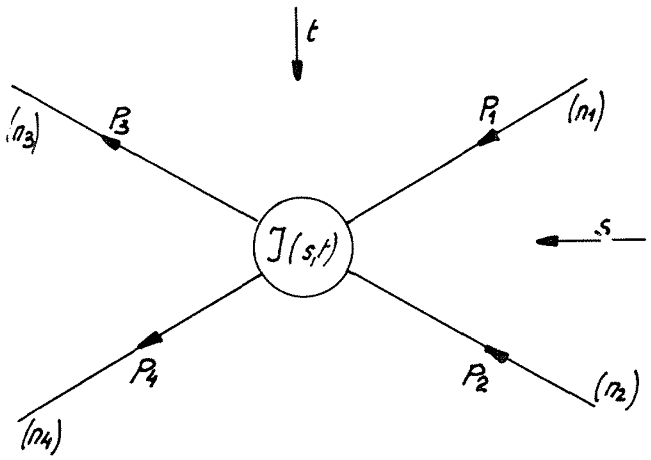
$$K^{-1} = B^{-1} - F \tag{1.8}$$

where F is an explicit integral. For a single channel this becomes

$$K = \frac{B}{1 - BF} \tag{1.9}$$

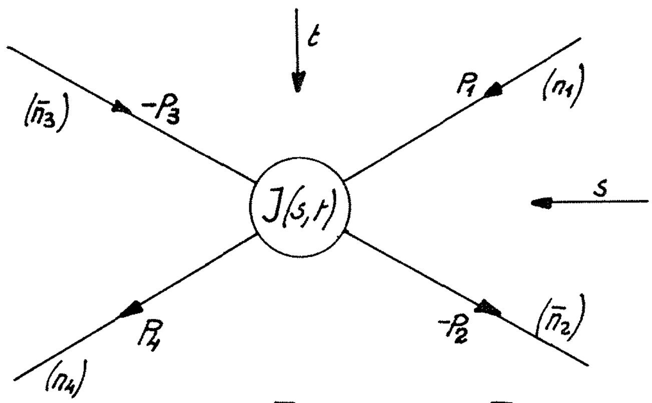
If B is very small the denominator can be neglected and we have simple Born approximation. However if B is large the denominator can be important, particularly when B is positive, (corresponding to an attractive potential), when the denominator can vanish, giving rise to a resonance.

The 2π and 3π pseudo particles can be in a great variety of states [5] specified by spin, i-spin, space parity and G-parity [6]. For reasons which will appear below, we will assume that the most important are two vector bosons with the properties given in Table 1.



$n_1 + n_2 \longrightarrow n_3 + n_4$

(a)



$n_1 + \bar{n}_3 \longrightarrow n_4 + \bar{n}_2$

(b)

Figure 2 - The matrix elements for (a) nucleon-nucleon, (b) nucleon anti-nucleon scattering in terms of the variables $s = (p_1 + p_2)^2$; $t = (p_1 - p_3)^2$.

Table 1

Particle.	J.	I.	P.	G.	Mass.	
V_μ	1	1	-	+	V	(2π)
S_μ	1	0	-	-	S	(3π)

Since our formula depends on the Born terms, we must specify the effective Lagrangians defining their interactions. These are

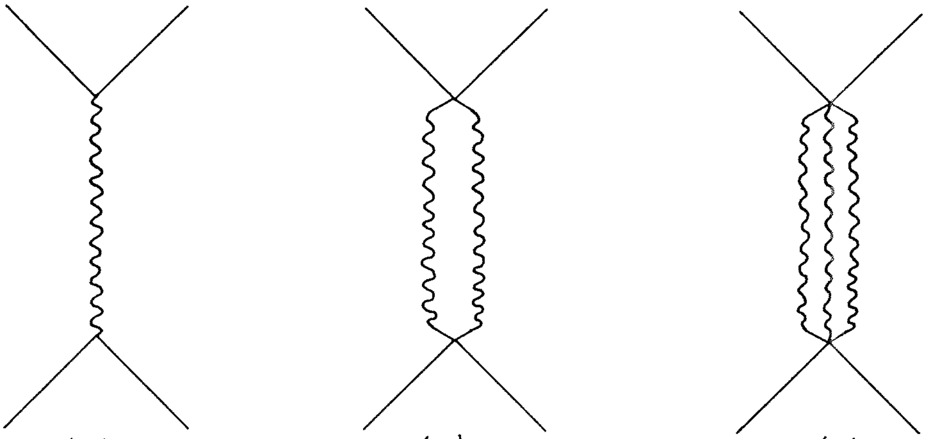


Figure 3 - The graphs which give rise to the nucleon-nucleon potential.
(a) Single pion exchange.
(b) Resonant 2π exchange.
(c) Resonant 3π exchange.

$$
\begin{aligned}
L_{SN} &= ig_1^s\, \bar{\psi}\, \gamma_\mu \psi\, S_\mu + \frac{g_2^s}{4M} \bar{\psi}\, \sigma_{\mu\nu} \psi\, (\partial_\mu S_\nu - \partial_\nu S_\mu), \\
L_{VN} &= ig_1^v\, \bar{\psi}\, \underset{\sim}{\tau}\, \gamma_\mu \psi\, \underset{\sim}{V}_\mu + \frac{g_2^v}{4M} \bar{\psi}\, \sigma_{\mu\nu} \underset{\sim}{\tau}\, \psi\, (\partial_\mu \underset{\sim}{V}_\nu - \partial_\nu \underset{\sim}{V}_\mu), \\
L_{V\pi} &= \frac{1}{2} g_\pi\, \varepsilon_{vst}\, V_{\mu v}\, \Phi_s \overleftrightarrow{\partial}_\mu \Phi_t , \\
L_A &= ef^s S_\mu A_\mu + ef^v V_\mu A_\mu ,
\end{aligned}
\tag{1.10}
$$

where ψ, Φ and A_μ are the nucleon, pion and photon fields, respectively. The parameters we have introduced are thus the two mass, S and V, and the seven constants

$$f^s,\ g_1^s,\ g_2^s$$

$$f^v,\ g_1^v,\ g_2^v \text{ and } g_\pi .$$

We will now show that most of these can be determined from the nucleon form factor data and our notation has been chosen with this in mind. We will then consider the effect of these resonances on other pion pheonomena.

2. NUCLEON FORM FACTORS -

The pion electromagnetic form factor $F_\pi(t)$ is defined by the relation

$$e F_\pi(t)\ (k + k')_\mu = \langle k' \mid j_\mu \mid k \rangle \tag{2.1}$$

where j_μ is the electromagnetic current, k and k' are the pion four momenta, and

$$t = (k + k')^2 . \tag{2.2}$$

A dispersion relation can be written for $F_\pi(t)$.

$$F_\pi(t) = \frac{1}{\pi} \int \frac{g(t')}{t' - t}\, dt' . \tag{2.3}$$

If it is assumed that the integral is dominated by 2π intermediate states represented by V_μ, this is equivalent to calculating $F_\pi(t)$ by perturbation theory from the Lagrangien (1.10), (figure 5). The result is

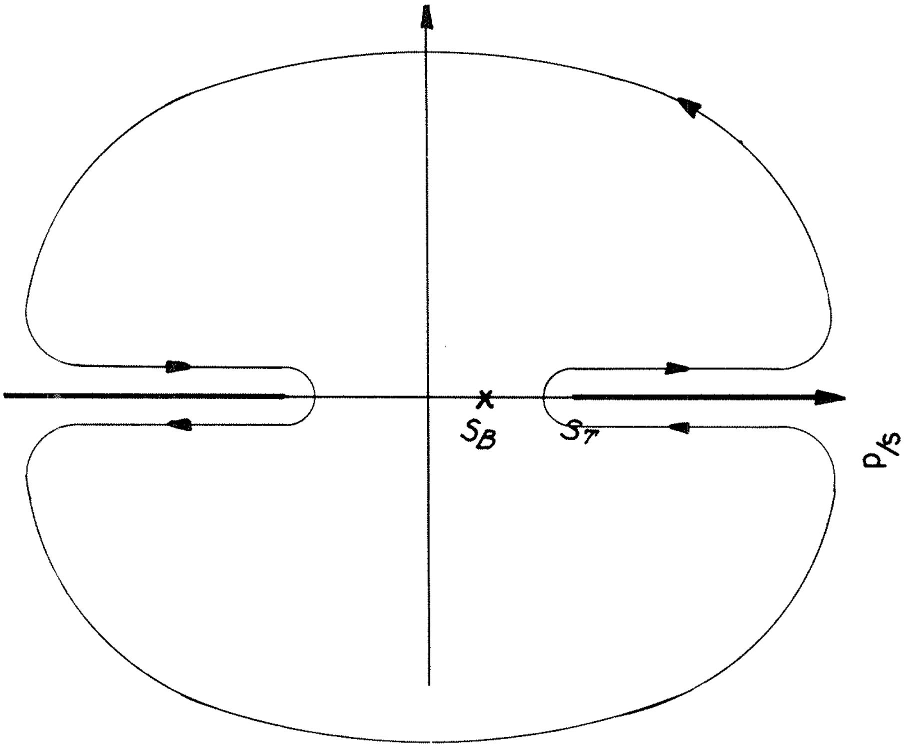


Figure 4 - The cut plane in which T(s) is assumed analytic, and the contour of integration for the correct dispersion relation.

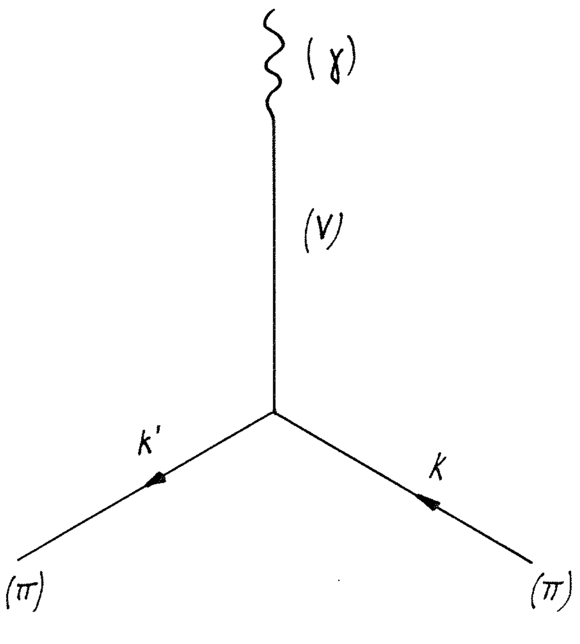


Figure 5 - The graph for the pion electromagnetic form factor, due to V-particle exchange.

$$F_\pi(t) = \frac{f^v g_\pi}{V^2 - t}. \qquad (2.4)$$

If we assume this is exact (i.e. no subtraction is necessary), then the requirement that the total charge in unity is that

$$F(o) \equiv \frac{f^v g_\pi}{V^2} = 1. \qquad (2.5)$$

The nucleon form factors are defined similarly :

$$< p' \mid j_\mu \mid p > = e\bar{u}(p') \, [F_1(t)\, i\, \gamma_\mu + F_2(t)\, i\, \sigma_{\mu\nu} \, (p'-p)_\nu] \; u(p) \qquad (2.6)$$

where

$$F_i = F_i^s + \tau_3 \, F_i^v. \qquad (2.7)$$

If these are calculated in the same approximation, then

$$F_i^w(t) = \frac{f^w g_i^w}{W^2 - t} + C_i^w, \quad \begin{matrix} (W = S, V.) \\ (i = 1,2.) \end{matrix} \qquad (2.8)$$

where the C_i^w are constants which represent the contributions from intermediate states of higher energy, (and possibly the subtractions constants, which may be necessary even in the exact dispersion relations). These constants can be eliminated by fitting to the charge and static moments at $t = 0$, and no attempt is then made to explain these quantities. In the case of the vector from factors, (W = V), f^v can be eliminated by (2.5). We thus arrive at the following expressions [7].

$$F_1^v = \frac{1}{2}\left[1 + \frac{2 g_1^v}{g_\pi} \frac{t}{V^2 - t}\right] \equiv \frac{1}{2}\left[1 + \frac{a_v t}{V^2 - t}\right], \qquad (2.9)$$

$$F_2^v = \mu_v \left[1 + \frac{1}{\mu_v} \frac{g_2^v}{g_\pi} \frac{t}{V^2 - t}\right] \equiv \mu_v \left[1 + \frac{b_v t}{V^2 - t}\right], \qquad (2.10)$$

$$F_1^s = \frac{1}{2}\left[1 + \frac{2 f^s g_1^s}{S^2} \frac{t}{S^2 - t}\right] \equiv \frac{1}{2}\left[1 + \frac{a_s t}{S^2 - t}\right], \qquad (2.11)$$

$$F_2^s = \mu_s \left[1 + \frac{1}{\mu_s} \frac{f^s g_2^s}{S^2} \frac{t}{S^2 - t}\right] \equiv \mu_s \left[1 + \frac{b_s t}{S^2 - t}\right] \qquad (2.12)$$

where μ_v and μ_s are the scalar and vector nucleon moments, in units of the nuclear magneton e/2M,

$$\mu_v = \frac{\mu_P - \mu_N}{2} = 1.85 \qquad \text{(experimental)} \qquad (2.13)$$

and

$$\mu_s = \frac{\mu_P + \mu_N}{2} = -0.06 \qquad \text{(experimental)} \qquad (2.14)$$

The final equalities in (2.9)-(2.12) show the relations between our parameters and the currently conventional parameters of the form factor problem. If S^2, V^2, a_s, a_v, b_s and b_v were known exactly, these, with (2.5), would give 7 relations between our nine parameters. (For details see the report of S. Fubini). In practice the situation is as follows :-

The proton data is much more reliable than that of the neutron. Since b_s appears always combined with the small factor μ_s, the proton moment distribution is dominated by F_2^v and determines [8]

$$V^2 \triangleq 28 \; m_\pi^2, \quad \text{(exp)} \qquad (2.15)$$

$$\frac{g_2^v}{g_\pi} \equiv \mu_v \, b_v \triangleq (1.8) \times (1.2), \qquad \text{(exp)}. \qquad (2.16)$$

The observed equality of vector, charge and magnetic moment, radii indicates

$$a \triangleq b \,, \tag{2.17}$$

and hence

$$2\mu_v g_1^v = g_2^v. \tag{2.18}$$

One more fact is then required to determine the vector parameters.

The situation with regard to the scalar parameters is less well determined. Since the S-particle cannot decay into two pions, there is no analogue to (2.5) to elimintate f^s. The marked difference in the observed shape of the proton charge,

$$F_1^p = F_1^s + F_1^v \,, \tag{2.19}$$

and moment distribution,

$$F_2^p = F_2^s + F_2^v, \tag{2.20}$$

indicates a lower mass for the iso-scalar resonance,

$$S \simeq 16 \; m_\pi^2 \quad (\text{exp}). \tag{2.21}$$

The vanishing of the neutron charge radius implies

$$\frac{a_s}{S^2} = \frac{a_v}{V^2} \,, \tag{2.22}$$

which in our notation is

$$f^s g_1^s = a_v \left(\frac{S^2}{V^2}\right) S^2. \tag{2.23}$$

The least well determined parameter is b_s since, as remarked above, it appears multiplied by a small factor. It is related to the coupling constants by

$$f^s g_2^s = \mu_s \, b_s S^2. \tag{2.24}$$

Combining the last two equations

$$\frac{g_1^s}{g_2^s} = \frac{a_v}{\mu_s b_s} \left(\frac{S^2}{V^2}\right). \tag{2.25}$$

This ratio proves to be very important in connection with nuclear forces (*).

The greatest success of this type of analysis so far has been the prediction of the V-resonance [9], which has recently been observed in pion production [10 a]. The width of the resonance is determined by the decay rate

$$\Gamma = \frac{1}{3}\left(\frac{g_\pi^2}{4\pi}\right) \; V \longrightarrow \pi + \pi \,.$$

By perturbation theory, this is

$$\Gamma = \frac{2}{3}\left(\frac{g_\pi^2}{4\pi}\right) \frac{q^3}{V^2} \equiv \gamma q^3 \tag{2.26}$$

(*) A slight variation of the avove argument is to assume that no subtraction is necessary in F_2^v. Then to give the difference of the static moments, μ_v correctly, by (2.5) and (2.8)

$$\frac{g_2^v}{g_\pi} = \mu_v \,.$$

Hence, as a prediction, by (2.16),

$$b_v = 1.$$

By (2.17) and (2.27)

$$a_v = 1 \text{ and } a_s = 0.5$$

These figures are not far from the Hofstadter-Herman [8] parameters.

where q is the momentum of the pions in the rest system of the V, and γ is the reduced width. Equating this to the observed value [10 b] gives

$$\frac{g_\pi^2}{4\pi} \simeq 1 . \qquad (2.27)$$

This with (2.15) - (2.18) roughly determines all the parameters of the V-particle (*).

3. PION-NUCLEON SCATTERING -

Since the S-particle cannot decay into two pions, it plays no role in pion-nucleon scattering. Born approximation in our model is given by the nuclear poles and the V-particle. (Figure 6(a) and (b).

If the S-matrix is written as (**)

$$S = 1 + i\,(2\pi)^4\,T \qquad (3.1)$$

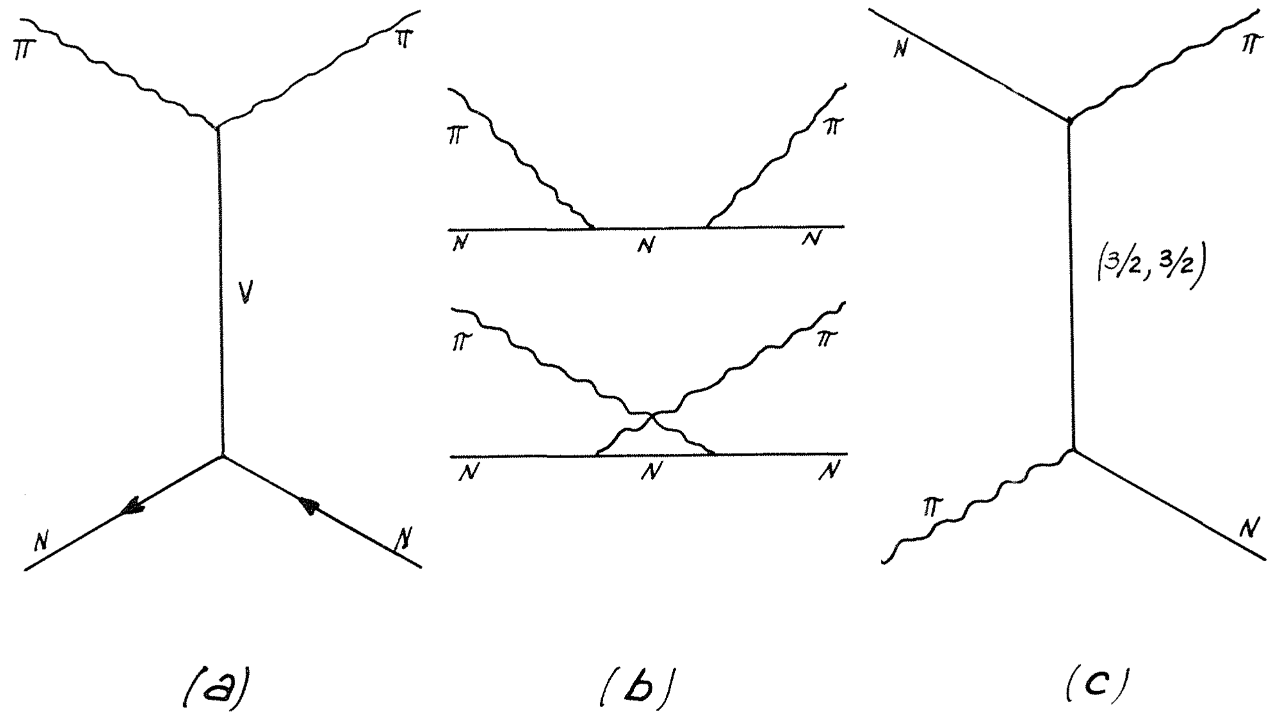


Figure 6 - Pole terms for pion-nucleon scattering : (a) the V particle, (b) the nucleon poles, (c) the (3/2, 3/2) resonance.

(*) If (2.26) is used to eliminate g_π in favour of γ in (2.9) and (2.10) we have

$$a_v = \frac{4}{3}\left(\frac{g_1^v g_\pi}{4\pi}\right)\frac{1}{V_\gamma} ,$$

(and a similar expression for b). This is identical with the expression founc by Cottingham Bowcock and Lurie [9] if we identify their parameters C_i by

$$C_i = -\frac{4}{3}\left(\frac{g_i^v g_\pi}{4\pi}\right)$$

(**) We use $\gamma_0 = -i\gamma_4$ anti-hermitian, $\gamma_{1,2,3}$ hermitian and a time-like metric $a_\mu b_\mu = a_0 b_0 - \underset{\sim}{a}\,\underset{\sim}{b}$

then

$$T = A + \frac{1}{2} i\gamma (k + k') B \tag{3.2}$$

where k and k' are the initial and final pion four momenta. The contribution to T from the g_1^v term in the V-particle pole (Figure 6(a)) is

$$A_v = 0 \tag{3.3}$$

$$B_v = g_1^v g_\pi (2 P_{1/2} - P_{3/2}) \frac{2}{V^2 + 2q^2(1 - \cos\vartheta)}, \tag{3.4}$$

where $P_{1/2}$ and $P_{3/2}$ are the i-spin projection operators, and q and ϑ are the center of mass momentum and scattering angle respectively. The contributions from the nucleon terms have been calculated repeatedly [11]. In the static approximation the scattering is purely p-wave, and the p-wave scattering is completely dominated by the attractive 'potential' in the (3/2, 3/2) state, coming from the nucleon terms (Figure 6(b)). Application of the theory outlined in § 1 then leads to a resonance in the (3/2, 3/2) state. Having found the resonance, a more refined theory is obtained by including the contribution of this resonance to the 'potential' by incorporating the appropriate pseudo-particle in the Born term, B. (Figure 6(c)). This incorporates 'crossing' symmetry and precisely reproduces the Chew-Low theory. The V-particle has no appreciable effect on this well known situation because it is so heavy.

Before considering the small p-waves we discuss the S-waves. If

$$f_s = \frac{e^{i\delta} \sin\delta}{q}, \tag{3.5}$$

then from the pole terms we have

$$f_s^{(B)} = \left(\frac{g_1^v g_\pi}{4\pi}\right) \begin{pmatrix} 2 \\ -1 \end{pmatrix} \frac{2\omega}{V^2 + 2q^2} - \frac{4Mf^2}{m_\pi^2} \left[1 - \frac{8}{q}\left(1 + \frac{2\omega_r}{M}\right)\right] \begin{pmatrix} 1 \\ 1 \end{pmatrix} \tag{3.6}$$

where ω is the meson energy, ω_r the (3/2, 3/2) resonance position, M the nucleon mass, and f^2 the dimensionless pion-nucleon coupling constant

$$f^2 = 0.08$$

The upper and lower numbers in brackets refer to i-spin 1/2 and 3/2 respectively. The first term in (3.6) comes from the V-particle, the second from the nucleon poles and the third from the (3/2, 3/2) resonance. The outstanding qualitative problem in pion scattering for many years has been to explain the strong i-spin dependence of the two S-wave scattering lengths, which is completely lacking in the nucleon terms. This is given directly by the V-particle term [12]. Since the S-wave is classically a low impact parameter effect we expect it to be particularly sensitive to high momentum transfer effects. These can be simulated by a subtraction constant, which it seems reasonable to take independent of i-spin (*). If it is assumed that this just cancels the nucleon terms, we have

$$f_s = G_1 \begin{pmatrix} 2 \\ -1 \end{pmatrix} \frac{2\omega}{V^2 + 2q^2}, \quad G_1 = \frac{g_1^v g_\pi}{4\pi} \tag{3.7}$$

The scattering lengths are

$$a_s = G_1 \begin{pmatrix} 2 \\ -1 \end{pmatrix} \frac{2m_\pi}{V^2}. \tag{3.8}$$

To fit the experimental values [13],

(*) That the high momentum transfer effects will be much more effective for s-waves than p-waves can be seen by representing them by the exchange of a heavy J = 0 I = 0 boson, mass M_σ, and coupling strength Λ. The contribution to the s and p wave scattering lengths are proportional to Λ/M_σ and Λ/M_σ^3, respectively. This is essentially a dimensional argument, and therefore quite general.

$$a_s = \binom{+\ 0.178}{-\ 0.087}\ 1/m_\pi\ , \tag{3.9}$$

we require

$$G_1 \simeq 0.9, \tag{3.10}$$

which, using (2.16) implies

$$\frac{g_\pi^2}{4\pi} \simeq \frac{1}{2.2} \tag{3.11}$$

and is not in serious disagreement with the estimate from the width ot the π-π interaction (*).

The small p-waves are now given by

$$f_1^1 = \frac{8}{3} G_1 \frac{\omega q^2}{(V^2 + 2 q^2)} + \frac{2 q^2}{M} G_1' \frac{1}{(V^2 + 2 q^2)^2} - \frac{8}{3} \frac{f^2 q^2}{\omega} \frac{\omega_r}{\omega + \omega_r},$$

$$f_3^1 = \frac{8}{3} G_1 \frac{\omega q^2}{(V^2 + 2 q^2)} - \frac{2}{3} \frac{f^2 q^2}{\omega} \frac{\omega_r}{\omega + \omega_r}, \tag{3.12}$$

$$f_1^3 = -\frac{4}{3} G_1 \frac{\omega q^2}{V^2 + 2 q^2} - \frac{q^2}{M} G_1' \frac{1}{(V^2 + 2 q^2)^2} - \frac{2}{3} \frac{f^2 q^2}{\omega} \frac{\omega_r}{\omega + \omega_r},$$

where

$$G_1' = \frac{(g_1^v + g_2^v)\, g_\pi}{4\pi} \tag{3.13}$$

The final term in each expression being the well known nucleon contribution (see for example Cottingham, Bowcock and Lurie [9], Eqn. (5.4)). The remaining terms arise from the V-particle, pole.

Substituing the numerical values which have been used above one obtains for the scattering lengths

$$C_{2J}^{2I} = f_{2J}^{2I}/q^2 \Big]_{q^2=0} \tag{3.14}$$

the values given in Table 2.

Table 2

The theoretical values of small p-wave phase shifts compared with experimental values of Hamilton [13]

	Exp.
$C_1^1 = -0.14 + G_1 \times (0.06)$	- 0.02
$C_3^1 = -0.035$	- 0.07
$C_1^3 = -0.035 - G_1(0.03)$	- 0.04

These quantities are not very well established experimentally, but it can be seen that reasonable values are obtained by taking the value of G_1 determined by the S-waves. The main function of the V-particle contribution is to reduce the value of C_1^1, which appears to be considerably too large from the nucleon contribution alone.

(*) A less optimistic programme is to use the scattering lengths to determine the high energy parameters, and employ the theory to predict the energy dependence of the s-waves [9].

4. NUCLEON-NUCLEON INTERACTION -

In the nucleon-nucleon interaction both the S and V particles contribute to the 'potential', representing, it is hoped, the main contribution from three pion and two pion exchange. (The graphs are shown in figure 4). Owing to the exceptionally large mass of the V-particle, it appears, from the nucleon form factors, that 3π exchange (S particles) has a longer range, and is consequently more important, than 2π exchange.

If T is defined as in (3.1), we introduce $T(\gamma)$ by the relation

$$T = \bar{u}(p_1')\bar{u}(p_2')\,T(\gamma^{(1)}, \gamma^{(2)})\,u(p_1)u(p_2), \tag{4.1}$$

where p_1, p_2 and p_1', p_2' are the four momenta of the initial and final nucleons respectively.

Let T_π and T_s be the contributions to the Born approximation graphs coming from single pion, and single S exchange respectively, (excluding exchange effects). Then if the result is expressed in terms of the 5 'β-decay' invariants. (See the appendix to Amati, Leader and Vitale [14]).

$$T_\pi = g^2 \frac{\gamma_5^{(1)}\gamma_5^{(2)}}{t - m_\pi^2} \tag{4.2}$$

$$\begin{aligned} T_s = \Bigg[& 1^{(1)}.1^{(2)}\, g_1 g_2 \left(\frac{4M^2 - 2s - t}{4M^2}\right) - \gamma_\mu^{(1)}\gamma_\mu^{(2)}\, g_1(g_1 + g_2) + \sigma_{\mu\nu}^{(1)}\sigma_{\mu\nu}^{(2)}\, g_2(g_1 + g_2)\frac{t}{4M^2} \\ & + \gamma_5^{(1)}\gamma_5^{(2)}\, g_2(g_1 + g_2)\left(\frac{4M^2 - 2s - t}{4M^2}\right)\Bigg]\frac{1}{t - S^2} \end{aligned} \tag{4.3}$$

$$T_v = T_s \left(\frac{t - S^2}{t - V^2}\right) \tau^{(1)}.\tau^{(2)}. \tag{4.4}$$

The coupling constants in the formula for T_s should, of course, all be g^s. In the expression for T_v they should be replaced by g^v. In the expression for T_π, g is the π-N coupling constant,

$$\frac{g^2}{4\pi} = \left(\frac{2M}{m_\pi}\right)^2 f^2. \tag{4.5}$$

To make contact with the conventional discussion of low energy nucleon-nucleon interaction, this can be further reduced to a matrix between two component Pauli-spinors in the c.m. frame. The corresponding amplitude is

$$\begin{aligned} t(\sigma) = & \frac{(\underset{\sim}{\sigma}^{(1)}.\underset{\sim}{\Delta})(\underset{\sim}{\sigma}^{(2)}.\underset{\sim}{\Delta})}{4M^2}\frac{g^2}{\Delta^2 + m_\pi^2} + \Bigg[g_1^2 + \frac{i(\underset{\sim}{\sigma}^{(1)} + \underset{\sim}{\sigma}^{(2)}).\underset{\sim}{n}}{4M^2}(3g_1^2 + 4g_1g_2) \\ & + \frac{S_{12}(\Delta)}{4M^2}(g_1 + g_2)^2 - \frac{2}{3} - \frac{\underset{\sim}{\sigma}^{(1)}.\underset{\sim}{\sigma}^{(2)}}{4M^2}\Delta^2(g_1 + g_2)^2 \Bigg]\frac{1}{\Delta^2 + S^2}, \end{aligned} \tag{4.6}$$

where

$$\begin{aligned} \underset{\sim}{\Delta} &= \underset{\sim}{p}_f - \underset{\sim}{p}_i, \\ S_{12}(\Delta) &= (\sigma^{(1)}.\Delta)(\sigma^{(2)}.\Delta) - \frac{1}{3}(\sigma^{(1)}.\sigma^{(2)})\Delta^2, \\ \underset{\sim}{n} &= \underset{\sim}{p}_f \wedge \underset{\sim}{p}_i. \end{aligned} \tag{4.7}$$

If this matrix element is now interpreted as an effective potential, the first term is the well known one pion exchange term, (OPEP). In the square bracket the term g_1^2 is a repulsive core of approximately the required magnitude as remarked by Nambu [7]. For nucleon anti-nucleon scattering this term changes sign and becomes a deep attractive well.

The next term is a spin-orbit interaction of the type introduced so effectively by Marshak and Signell [15]. Its coefficient $(3g_1^2 + 4g_1 g_2)$ is not positive definite. The term $3g_1^2$ is of the required sign as has been stressed by Breit [16] and Sakural [17], but for the whole term to be of the correct sign, we must have, $(g_1 > 0)$,

$$g_2 > -\frac{3}{4}g_1. \tag{4.8}$$

In terms of the parameters of the nucleon electromagnetic form factors, this implies, by (2.11) and (2.12),

$$\frac{b_s}{a_s} < \frac{3}{8\,|\mu_s|} \simeq 6 \tag{4.9}$$

This condition is satisfied by the Hofstadter [8] parameters, but the parameters of Bergia and Stanghellini [8] lead to a spin-orbit potential of the wrong sign. This plausible relationship betwenn the spin-orbit potential and the nucleon form factors gives special interest to an accurate determination of the difficult parameter b_s.

The two pion exchange described by the V-particle gives rise to a term of precisely the same form as from the S-particle, with the obvious change to the V-particle parameters, and an i-spin dependence of $\underset{\sim}{\tau}^{(1)}.\underset{\sim}{\tau}^{(2)}$. The best one can hope for in terms of the theory presented here, would be an explanation of nucleon-nucleon phenomena in terms of some refinement of this potential, with the parameters determined by the nucleon form factors. (See the work of Amati, Cottingham and Vinh Mau, report to Session H, p.347, Vol.1, of this Conference).

5. CONCLUSIONS AND π - π SCATTERING -

To summarise the situation, a semi-phenomenological synthesis of low energy pion physics has been attempted in terms of a dipion (I = 1, J = 1) and a tripion (I = 0, J = 1) resonance, involving nine parameters, one of which, f^s, we eliminate. The parameters of the nucleon form factors provide, in principle, seven relations between the remaining eight, but one of these parameters (b_s) is very poorly determined.

The most successful prediction [9] of the theory is that there should be a 2 π resonance at 20 - 30 m_π^2. This has since been suggested by a number of experiments and finally beautifully confirmed by Walker [10] et al. who found a strong correlation between the final pions in $\pi^- + p \longrightarrow N + \pi + \pi$ at 28 m_π^2. In pion-nucleon scattering the Chew Low theory is unaffected by the two pion resonance because it is so massive. If one assumes that the nucleon pole contributions are cancelled by high energy effects, then the strong i-spin dependence of the s-waves is naturally explained by the dipion pole [12]. The magnitude of the coupling constant (the last remaining free parameter) can be chose to give the observed value of the phase shifts. This then determines the width of the dipion resonance in reasonable agreement with experiment [10b]. The three small p-wave phase shifts are now determined and are in reasonable agreement with not very well determined experimental values.

The tripion resonance makes a contribution to the nucleon-nucleon potential which satisfactorily explains the repulsive core [7]. It also contribute a spin-orbit term, which may be of the correct sign but this depends critically on the parameter b , which is not accurately determined by the form factor data.

There has also been work done on photoproduction [17] and the higher pion-nucleon resonances, [18] but this is still at a very preliminary stage.

By far the most specific unobserved prediction of the theory is the existence of the tripion resonance at about 16 m_π^2, and certainly the most interesting experimental problem in pion physics at the moment is the direct observation of this state. The particle can decay by strong interactions into 3 pions, but for masses up to $5m_\pi$, it is strongly inhibited by phase space, and is more likely to decay electromagnetically (*)

(*) Assuming a point interaction

$$H_{nt}^{f} = \frac{e\lambda}{2m} F_{\mu\nu} (S_\mu \, \partial_\nu \, \Phi - S_\nu \, \partial_\mu \, \Phi)$$

for electromagnetic decay,

$$\frac{1}{\tau_\gamma} = \frac{e^2\lambda^2}{12\pi\, m^2} \left(\frac{S^2 - m_\pi^2}{2S}\right)^3 .$$

For

$$S_\mu \longrightarrow 3\pi ,$$

we assume

$$S_\mu \longrightarrow \pi^\circ + \gamma$$

with a life time comparable to that of Σ°.

In this connection there have been two experiments with positive, but puzzling results. The first is that of Abashian, Boothe and Crowe [19] who have studied the recoil of He^3 and H^3 in the reactions

$$\begin{aligned} p + d &\longrightarrow He^3 + ? \\ &\longrightarrow H^3 \;\; + ? \end{aligned}$$

and found evidence for a resonance in the I = 0 state with mass 2.2 m_π. This is very hard to understand, since if it is a J = 0, 2π -state, it should have a strong effect in τ decay which is not observed. On the other hand if it is J = 1, 3π - state (S_μ particle), one would expect

$$K^+ \longrightarrow \pi^+ + S_\mu \longrightarrow \pi^\circ + \gamma$$

which would have been observed experimentally as anomalous τ' decays [20] ($\tau' \longrightarrow \pi^+ + 2\,\pi^\circ$). This consideration puts a lower limit of 2,5 m_π on the mass of any such particle, unless one invokes the fortuitous vanishing of the K^+ decay matrix element. Further, the search for a resonance in the neizhbourhood of 2 or 2 1/2 m_π range in the reaction $\gamma + p \longrightarrow p + ?$ at Frascati has yielded negative results [21].

On the other hand three-pion correlations have been studied directly by Maglic, Alvarez and Rosenfeld [22] in the 5-pion annihilation of $p + \bar{p}$ and by Pevsner et al [22] in the reaction $\pi^+ + d \longrightarrow p + p + \pi^+ + \pi^- + \pi^\circ$. Both groups have found a resonance with mass 5 1/2 m_π. Their very preliminary evidence favours an iso-scalar, 1-, particle (S_μ). This has the quantum numbers required for the isoscalar nucleon form factor, but is heavier than expected on the present form factor analysis. In this connection, therefore, it is very interesting that Pevsner has preliminary evidence for another 3π resonance at $4\,m_\pi$. As remarked above such a state would decay more readily into $\rightarrow \pi^\circ + \gamma$ and not show up so distinctly as the 5 1/2 m_π state in these experiments.

Our choice of quantum numbers for the dipion and tripion resonances has been strongly influenced by the nucleon form factors. The two pion states, which directly effect π-N scattering, are restricted by Bose statistics to satisfy

$$(-1)^{J+1} = 1$$

Apart from V_μ, one could have an iso-scalar, 0+, (σ particle) or, 2+, and the phenomenological analysis of the influence pion-pion interaction on pion-nucleon scattering by Hamilton et al. [23] suggests such an interaction. The possible three pion states are much more various. However one is reluctant to introduce them until one knows they are there, since each one involves a considerable increase in the number of parameters. Thus the initiative at the moment seems to be with the experimenters, and the key questions are the dominant features of the dipion and tripion interactions. These can be studied in any interaction in which two or more pions are produced and once the main facts are known, it may be possible to piece together a 'theory' in which at least the number of quantitative predictions exceeds the number of parameters - even so the number of parameters will still be distressingly large.

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suite de la note de la page précédente.

$$H^{3\pi} = \frac{\lambda'}{m^3}\, \varepsilon_{\mu\nu\pi\rho}\; \partial_\pi\, S_\rho\, \varepsilon_{rst}\; \partial_\mu\, \Phi_r\; \partial_\nu\; \Phi_s\, \Phi_t,$$

and derive

$$\frac{1}{\tau_{3\pi}} = \left(\frac{\lambda'^2}{4\pi}\right) \left(\frac{S}{m}\right) \left(\frac{m_\pi}{m}\right) \left(\frac{\varepsilon^4}{m^4}\right) \frac{1}{2^6\, 3\, \sqrt{3}\,\pi}\; m_\pi,$$

where

$$S = 3\,m_\pi + \varepsilon\,,$$

and (1/m) is the range of the interaction ($\simeq 1/m_\pi$).

This arbitrariness can be considerably reduced if the pion resonances can be shown to form a self-consistent picture. Thus the 2π resonances should appear in pion-pion scattering in the same way as for example the (3/2, 3/2) resonance appears in pion-nucleon scattering. A calculation very much in the spirit of this paper has recently been published by Zachariasen [24]. He puts in a V-particle, with mass V, and coupling constant, g_π, as a pseudo-particle and determines these parameters by requiring that the position and width of the consequent resonance in π - π scattering are self-consistent. Reasonable values are obtained. This approach also implies strong s-wave scattering in at least one i-spin state [25]. This is a feature of more bearned calculations [26] and may be confirmed by experiment [19].

Given the V-particle, one may then expect an S-particle since each pair in the 3π system, then sits in the T = 1, J = 1 state [27]. Alternatively one can consider V- π scattering according to the theory presented here and show that the most attractive 'potential', and hence the strongest 3π resonance, should be expected in this state [25].

An alternative approach, which also considerably reduces the arbitrariness of the model, is to explain the occurence of 2π, 3π and $k\pi$ resonances in terms of higher symmetries of the strong interactions. This will be discussed in a wider context by Salam.

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HIGH ENERGY PHYSICS ABOVE 10 GeV

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CERN

The title of this talk and the arbitrary limit of 10 GeV shows clearly how badly specified is the topic. It seems to me that the talk should be devoted to the interactions of elementary particles at high energy and in particular to production processes in high energy collisions.

One would have thought that this is the main topic of the physics done with high energy accelerators, indeed the great actual problem of physics seems to be the strong interactions, and as pointed out long ago by Heisenberg, strong interactions will manifest themselves in phenomena of multiple production if enough energy is available. Therefore the study of such production events should be a way towards the understanding of strong interactions. In fact this approach has been, till now, somehow disappointing. This came probably from the very complexity of the problem either on the theoretical side where the sole possible approach was the statistical theory, or on the experimental side where one did not know what special parameters were most meaningful.

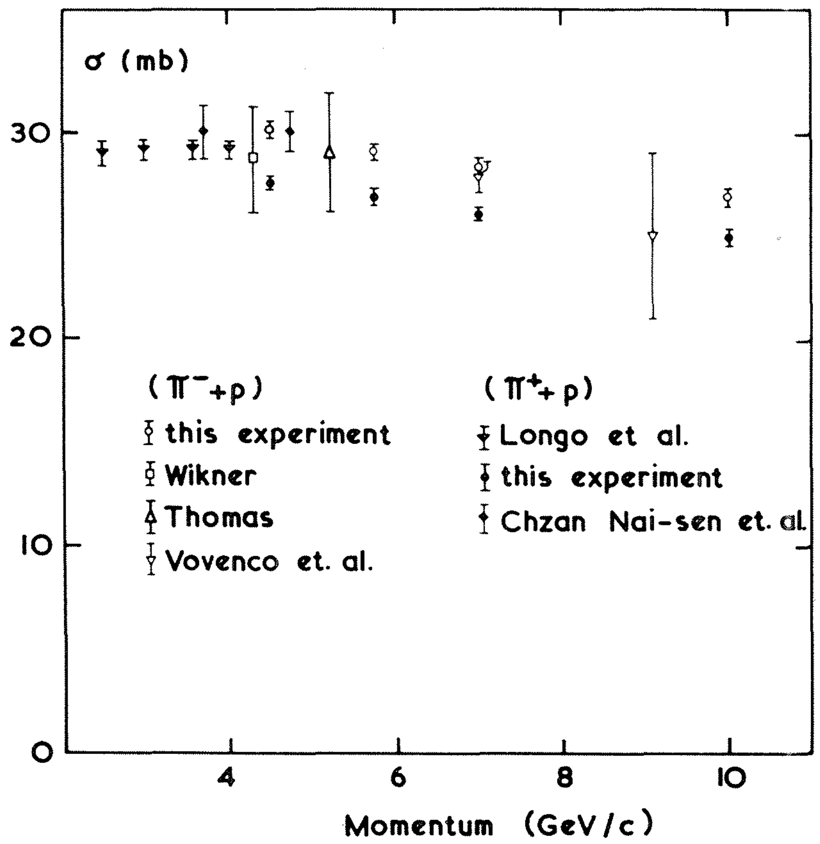


Figure 1 - $\pi^{\pm}$-p cross sections in function of energy.

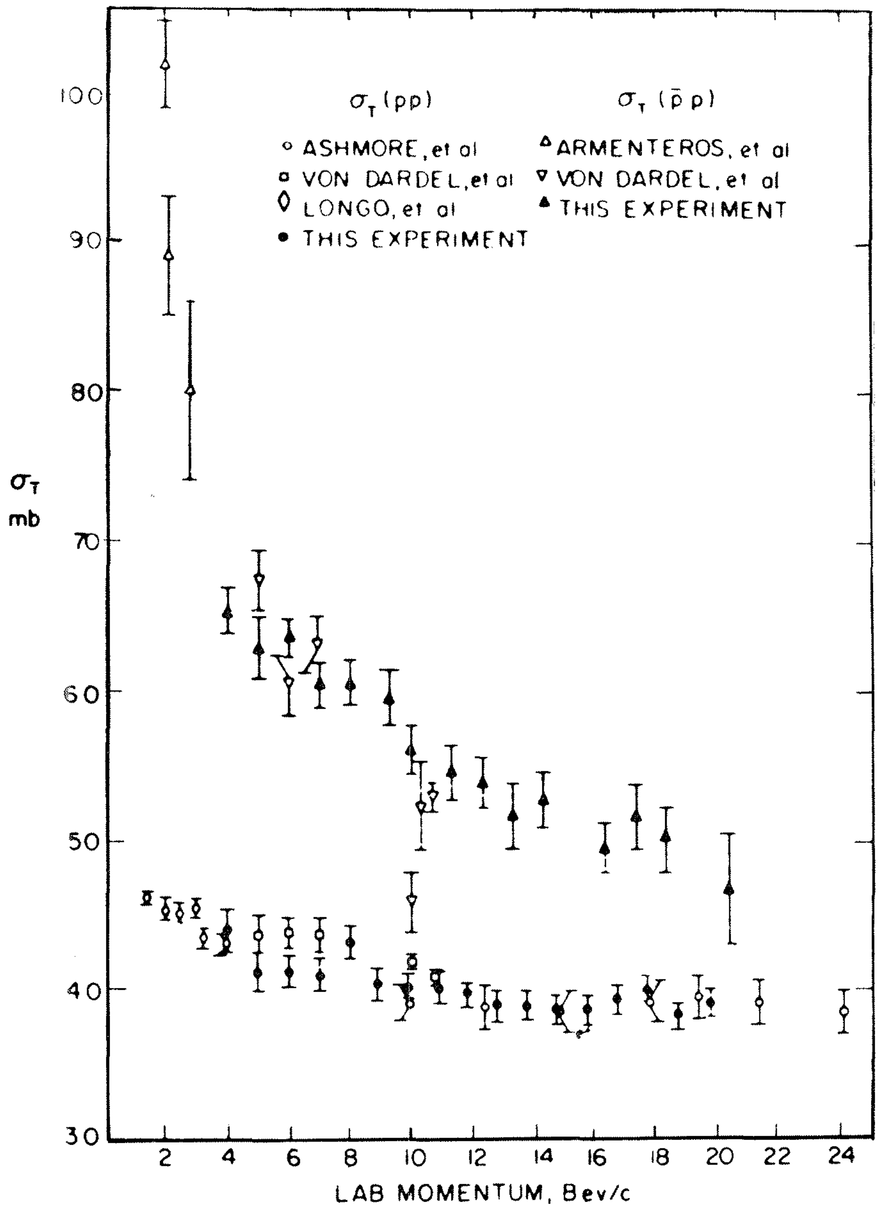


Figure 2 - p-p and $\bar{p}$-p cross sections in function of energy.

I wish to show, in this talk, what results have been obtained till now and how, very slowly, methods are being found by which it should be possible to make the study of high energy interactions contribute in an important manner to our knowledge of the elementary particles. Lack of time has forced me to restrict myself to experiments done with high energy machines. I will, therefore, not speak of the pioneer work done in cosmic rays on high energy collisions at energies of $10^{12} \longrightarrow 10^{14}$ ev and even higher. An excellent resumé of the results obtained in this field has been given by Perkins [1] at the CERN conference on high energy phenomena. It must be understood that, in this talk, the length of time devoted to a particular experiment is not proportional to its importance in physics. The precise value of a cross section can be reported in few seconds, whereas qualitative results, whose meaning is not immediately obvious, may request a long time of explanation.

TOTAL CROSS SECTIONS -

The first thing which can be measured at high energy is the total cross section. There exist the predictions of Pomerantchuk [2] to which experimental results can be compared. The difficult problem of identification of particles at very high energy has been solved in an elegant way by the use of large gas Cerenkov counters. The work in this field has been done in CERN essentially by von Dardel [1] and his group. Their latest and most precise result presented by Vivargent at this conference concerns $\pi^{\pm}$ - p cross sections and is shown in figure 1. The latest results on p-p and $\bar{p}$-p cross sections, Lindenbaum et al. [2] figure 2 and on $K^{\pm}$ - p cross sections [3] figure 3 come now from Brookhaven and have been presented at this conference by Yuan [4] and by Cool [5] respectively. In no case are the cross sections equal for the positive and negative particles, therefore, from the point of view of Pomerantchuk, we are not yet at high energies. From the work in cosmic rays we know that the total cross sections have at 10^3 - 10^4 GeV roughly the same value as at 10 - 20 GeV.

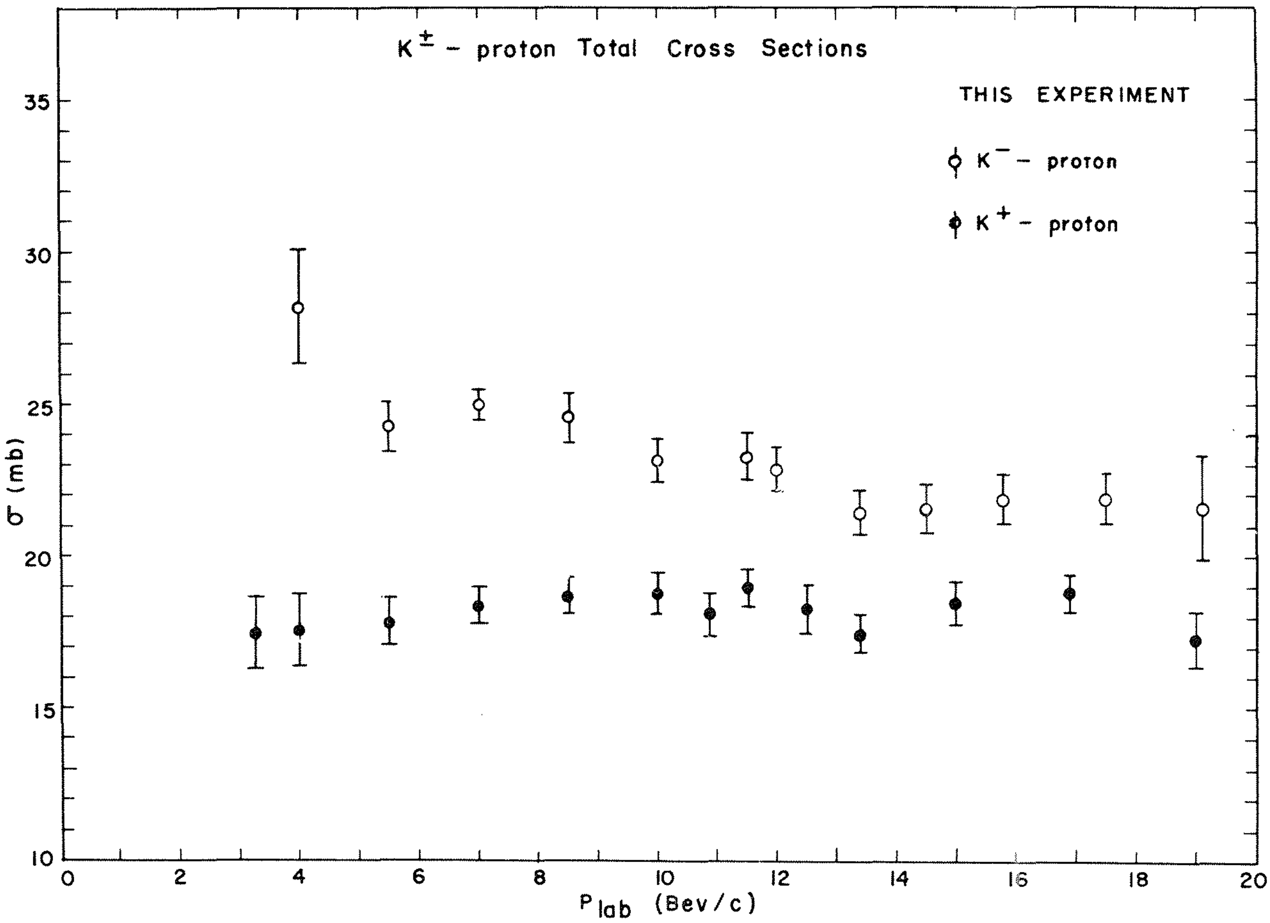


Figure 3 - $K^{\pm}$-p cross sections in function of energy.

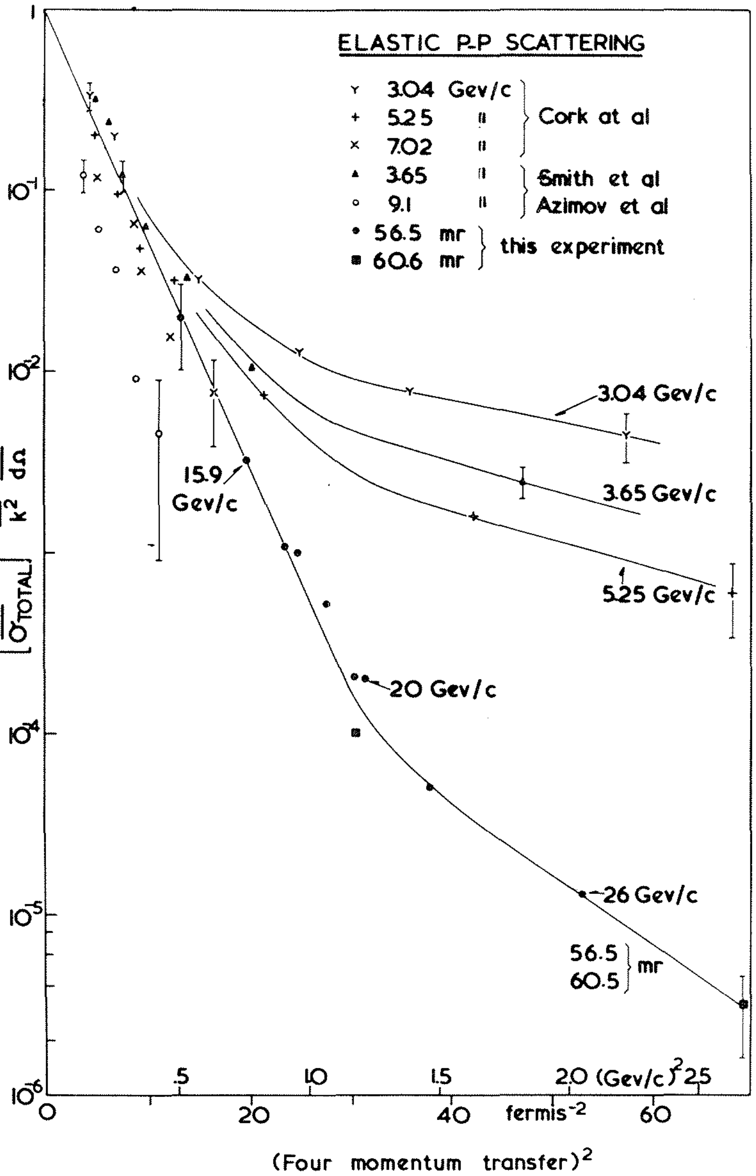


Figure 4 - Universal p-p elastic scattering plot.

ELASTIC SCATTERING -

At energies larger than 10 GeV the largest angular range scattering has been covered by Cocconi and his coworkers [6]. They have essentially worked at a fixed angle and varied the energy of the primary proton. Their results and those obtained previously by other workers at lower energies are shown on figure 4 which represents a sort of universal plot for elastic scattering. The abscissa is q^2 the square of the four-momentum transfer, the ordinate is $\left(\frac{4\pi}{\sigma_{tot}}\right)^2 \frac{f}{k^2} \frac{d\sigma}{d\omega}$ (on a logarithmic scale). By application of the optical theorem this quantity is 1 for $q^2 = 0$, if the forward scattering amplitude is purely imaginary. It is seen that this is the case at all the energies recorded on the plot. Furthermore for small values of q^2 the experimental points fall well on a straight line which represents a gaussian law. The authors think that this gaussian function can be an approximation of the Bessel function representing the diffraction scattering of a black disc, the slope of the straight line gives for the radius of the disc : R ~ 1 fermi.

However, at larger values of q there is a definite change of slope, i.e. there are more high energy transfers than expected from an extrapolation of the gaussian law found for small q's. It is remarkable that the q at which the change in slope occurs is not the same for different energies of the primary protons. In other words $\frac{1}{k^2}\frac{d\sigma}{d\omega}$ is no more energy independent.

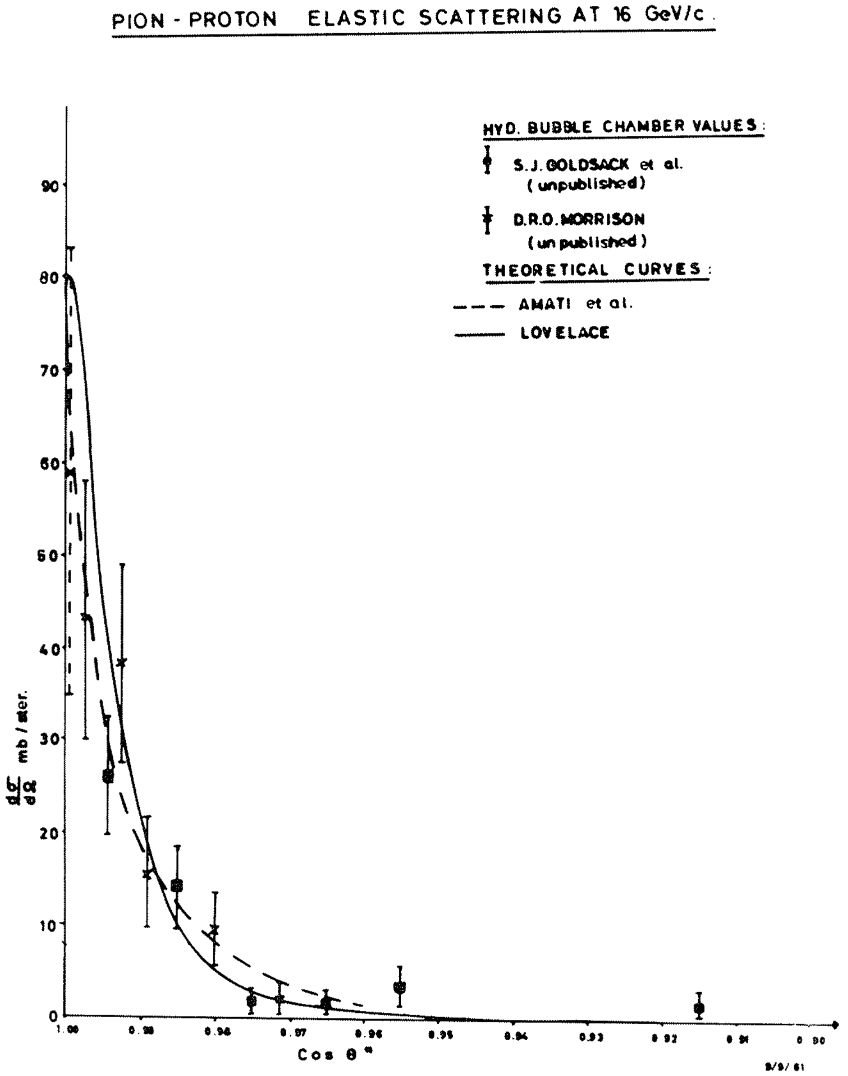


Figure 5 - Angular distribution of elastic scattering (π^--p) with π^- of 16 GeV.

Other experiments on p-p scattering have been reported at this conference [7, 8] ; as they were obtained in emulsion or in bubble chambers they concern essentially small momentum transfer and confirm the extrapolation to the optical theorem point. The value obtained for the cross sections are all in agreement within experimental limits. One can say that 8 mb < σ_{el} (p-p) < 10 mb.

On π-p elastic scattering, there is, at the time being, no data at high energies obtained by counter techniques. The results obtained by the British university collaboration and CERN, for the elastic scattering of 16 GeV π^- mesons on protons, were extracted from pictures taken by the CERN 30 cm hydrogen bubble chamber, and are both represented on figure 5. The theoretical curves are from Amati et al. [9] and Lovelace [10]. The cross section for π^--p elastic scattering is found to be 4 mb ± 0,5. In the same work of the British universities and of CERN, a careful scanning was done for zero prong events, i.e. for interactions of the type : $\pi^- + P \longrightarrow N + n\pi^\circ$, the cross section for such events which is an upper limit for charge exchange scattering ($\pi^- + P \longrightarrow N + \pi^\circ$) is found to be 0,25 mb, that is more than 10 times smaller than the cross section for elastic scattering.

Before concluding this chapter it must be noted that elastic scattering with high momentum transfer, with a primary of high energy, is an important method of investigating the structure of elementary particles, free of the complexity of inelastic phenomena and that it is still practically an open field.

INELASTIC HIGH ENERGY PHENOMENA -

I would like to divide this chapter in three parts, (a) general study of multiple π production, (b) glancing collisions, (c) production of strange particles.

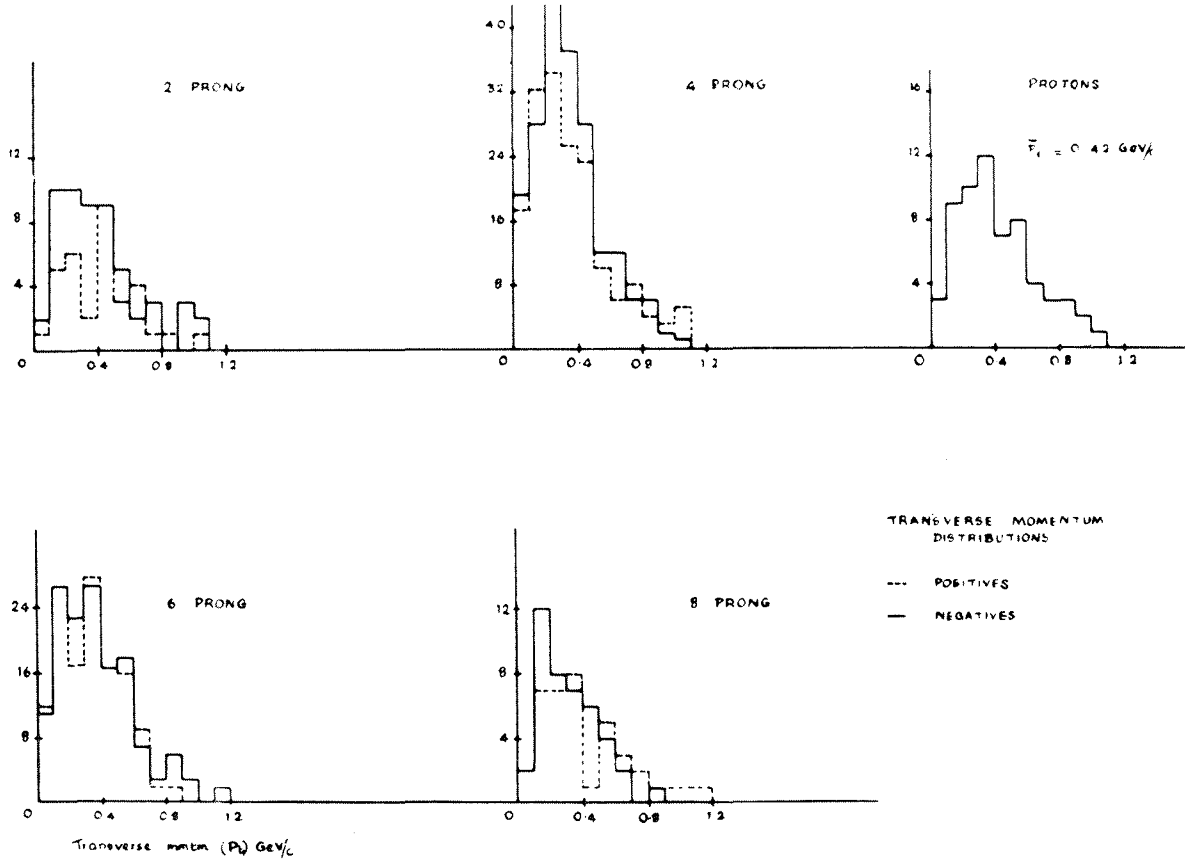


Figure 6 - Transverse momentum distribution of secondaries from π^--p interactions.

- Multiple π production (jets) - The study that I am reporting here has been made by a collaboration of 3 British universities : Imperial College-London, Birmingham and Oxford universities [11, 12] using photographs taken by the CERN 30 cm hydrogen bubble chamber. The primaries were 16 GeV π^- mesons and 24 GeV protons.

The transverse momentum distribution figure 6 of the secondaries hardly extends beyond 1 GeV/c, the average transverse momentum being 0,4 GeV/c. This is a very general picture of high energy interactions and it is true whether the primary is a proton or a π. It is true up to the highest energy (10^{14} ev) as shown by cosmic ray work. It is true for hyperons or K meson secondaries, as well as for π mesons as will be shown later.

The angular distribution of secondaries in CM system figures 7 and 8 is strongly anisotropic (peaked forwards and backwards) for events with 2 or 4 charged secondaries but tends to become isotropic for higher multiplicities. In other words the spectra of momenta and transverse momenta are quite different for low multiplicity jets and very similar for large ones.

In the case where the centre of mass is not a centre of symmetry of the system (π^-, P), there is, for low multiplicities jets (2 and 4 prongs) a marked asymmetry of the secondaries. The negative prongs go mostly forwards as was the incident π^-, the positive go mostly backwards as was the proton.

In many cases a proton could be identified by its greater ionization. Since this was done by eye judgement without any refined ionization measurement, the recognition was only possible for small momenta up to 750 MeV/c. The spectrum of these protons in CM system is shown in figure 9 for p-p collisions and compared to the spectrum foreseen by the statistical theory. One could object that the discrepancy is due to a bias, the only recognized protons being slow in the laboratory system, they must have a large momentum (backwards) in the CM system. However, the authors point out that they have identified such a large number of protons that the unfound ones, cannot possibly bring the two spectra into agreement. The results on protons of π-P collisions is very similar. This has been established by the study of hydrogen like events on the pictures taken in Ecole Polytechnique propane bubble chamber exposed to π^- of 6-11-18 GeV [13]. Here the great size of the chamber and the good stopping power of propane gives an easier identification of protons.

The British groups find that the secondary π mesons, contrary to protons, have a momentum spectrum in good agreement with the prediction of the statistical theory. This contradiction does not appear very serious to me (this opinion is rather personal). Indeed in the π spectrum the events of large multiplicities carry a larger weight (more mesons produced) than in the proton spectrum. This effect is even stronger if we think that to produce many mesons a proton must lose almost all its energy in the CM system, i.e. be fast in the lab. system. Furthermore it is clear that the statistical theory makes better predictions on the spectrum of the produced particles (π mesons) than on the inelasticity of the collision which is directly correlated to the proton spectrum.

What is now the conclusion we can draw from that work ? 2 and 4 prong jets constitute more than half of all the inelastic events, they show a strong anisotropy and even asymmetry. The spectrum of the protons after collision show that they retain a large amount of their energy in the CM system. All these points out to the fact that a large fraction of high energy collisions have the character of glancing or peripheral collisions, which bring us to our next point.

Glancing collisions - I have chosen this word to represent collisions in which the energy spent in the production of particles is a small fraction of the primary energy (small inelasticity). In other words the primary particles (incident proton or π, target proton) retain most of the energy and momenta after the collision. I have avoided "quasi elastic" which seems ambiguous, and peripheral which seems to have a definite theoretical meaning (one meson exchange).

The most interesting results on this type come from Cocconi et al. [14] and are already relatively well known. They are summarized in figure 10 which represents the spectrum of protons scattered (elastically or inelastically) by hydrogen, the angle of scattering is fixed 56,5 mrad, the energy of the incident protons is changed from a curve to another. One sees a large maximum of the spectrum, 1 GeV below the elastic peak. In fact this maximum is split in two. The position of the two maxima vs the elastic peak correspond very well to the excitation of the two isobars of higher energies (600 and 900 MeV) observed in π-P scattering. So that everything happens as if the incident proton had scattered on the target proton leaving it in an excited state. However, one of the most striking results of the work is the total absence of a peak corresponding to the excitation of the 3/2, 3/2 resonance.

Before going further into the discussion of these results and of the connected work done in the bubble chamber I must venture into the field of theoretical explanations.

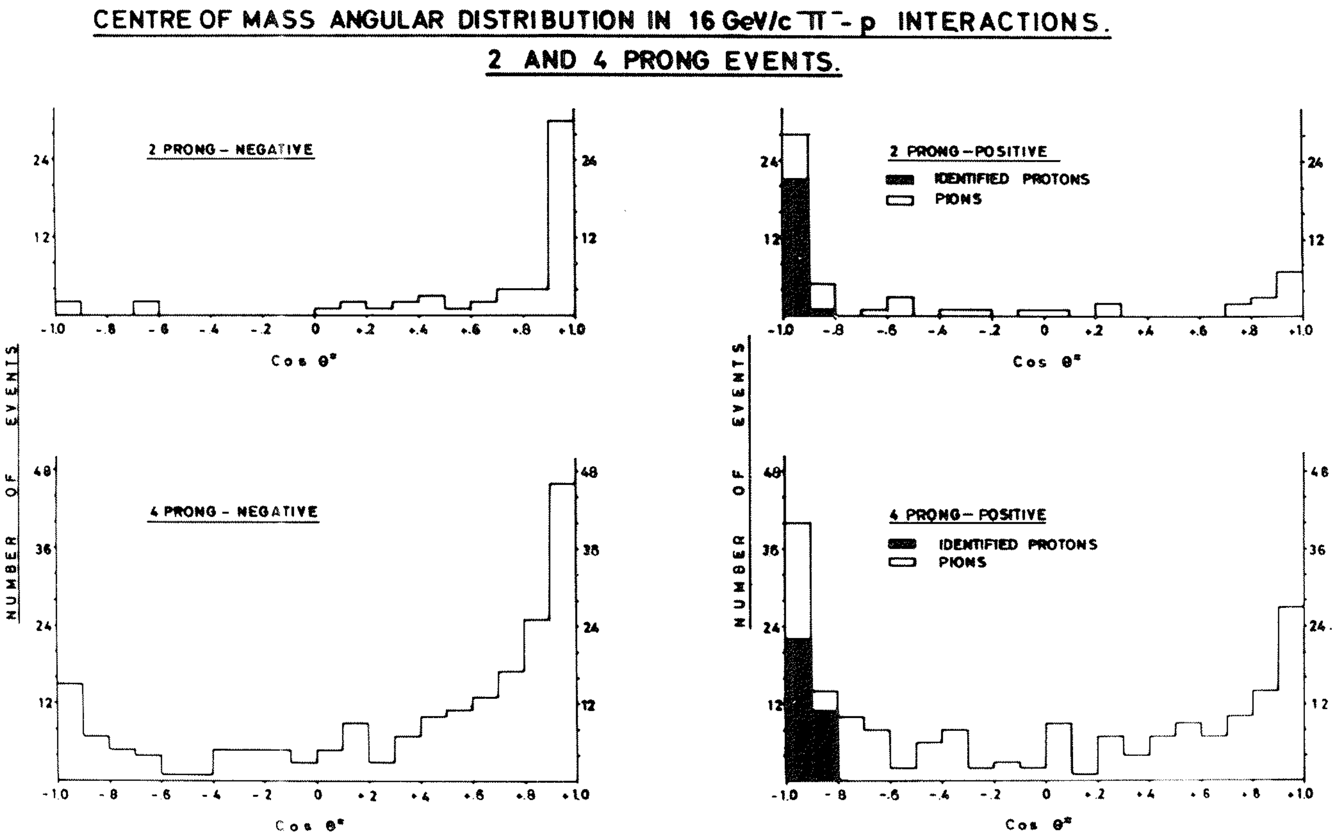


Figure 7 - CM system angular distribution of secondaries of π^--p interactions - (2 and 4 prongs jets).

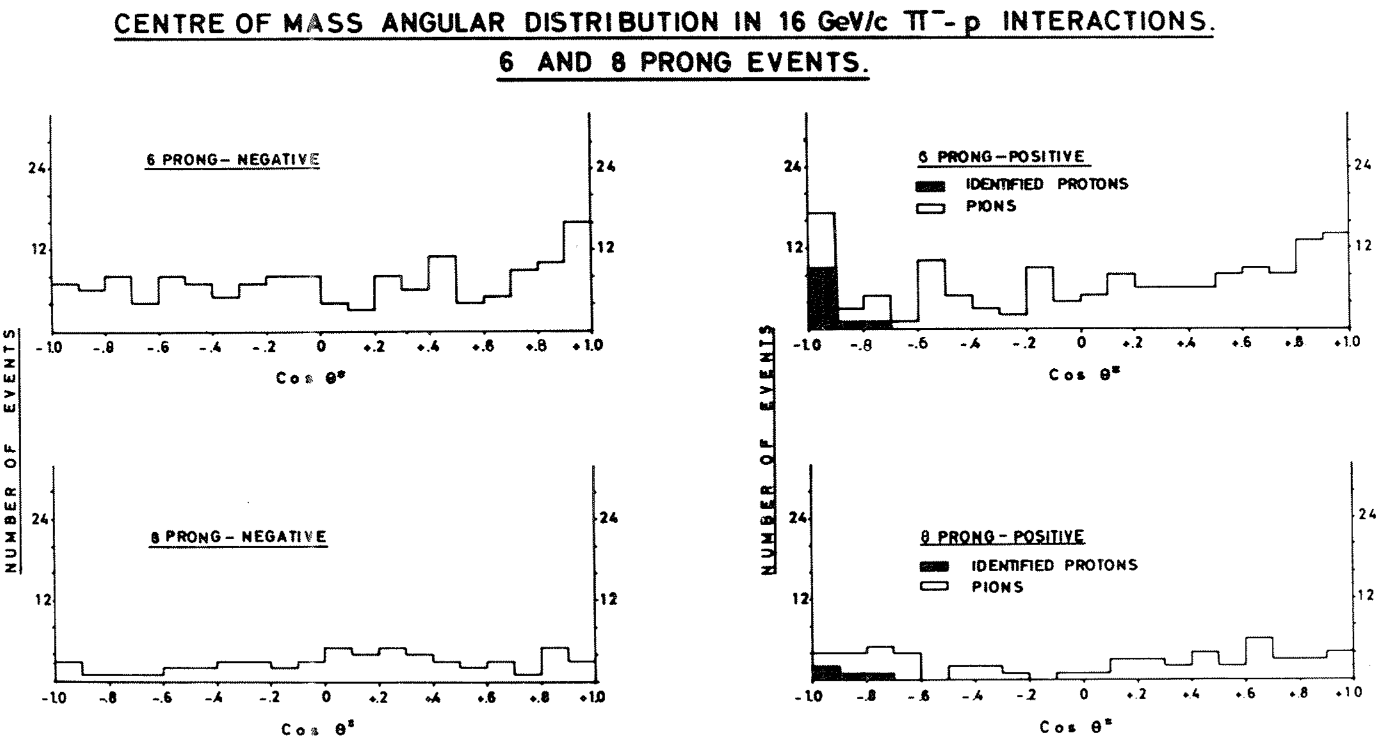


Figure 8 - CM system angular distribution of secondaries of π^--p interactions - (6 and 8 prongs jets).

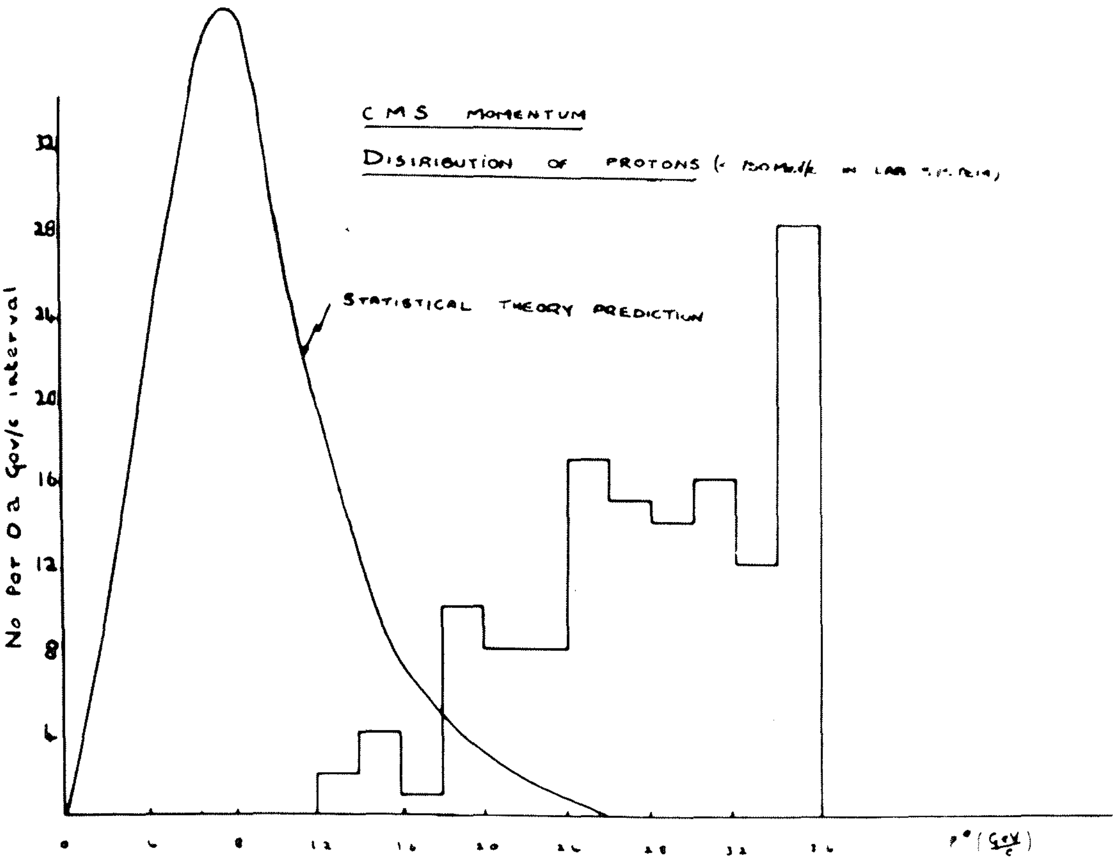


Figure 9 - CM system spectrum of secondary protons in p-p collisions.

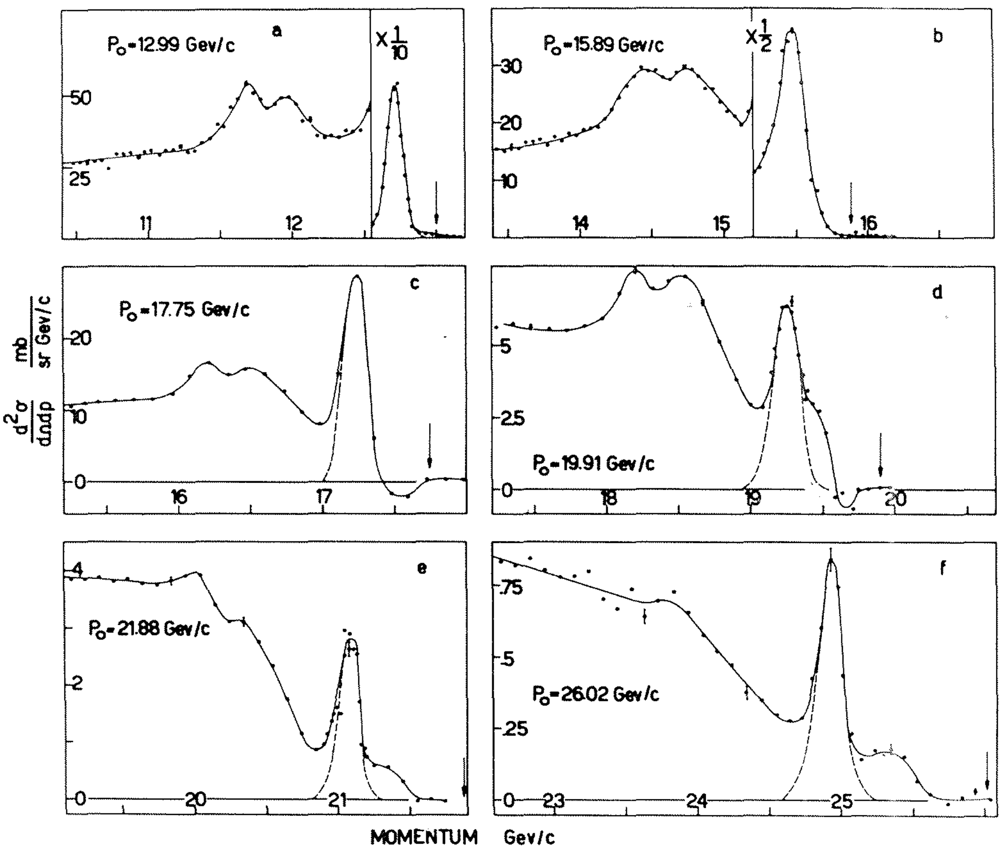


Figure 10 - Spectrum of protons scattered by hydrogen at a fixed angle and different primary energies.

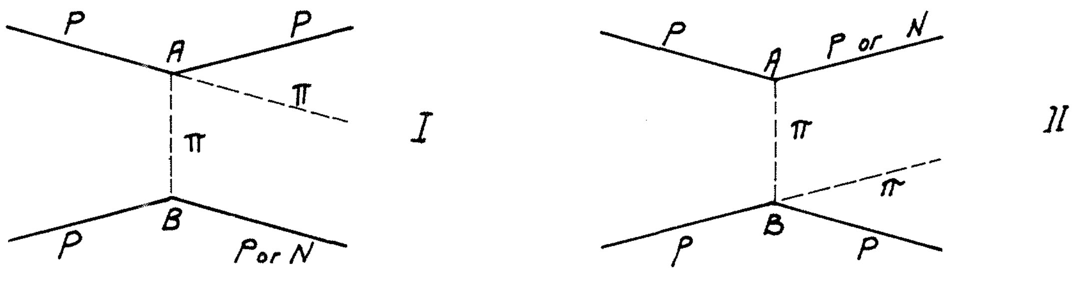
Let us consider the two graphs I and II, by convention I consider that the top proton is the incident proton, the bottom one the target proton. The graphs I and II are of course entirely symmetrical, and if one integrates on all possible final states they must equally contribute to the p-p inelastic events. However, we are considering a given final state : the one studied by the apparatus of Cocconi et al (a given scattering angle, a small energy loss) and, therefore, graphs I and II can give very different contributions to the cross section for arriving at this given final state. A priori one would like graph II to be the dominant one. Indeed the scattering in B could be dominated by the resonances and explain the double maximum. However, Drell and Hiida (*) [15] have calculated that the contribution of this graph is small. This is due to the relatively large scattering angle which corresponds to a relatively large momentum transfer which has to be carried by the virtual meson A-B, this one is therefore very far from the energy shell and the contribution of the graph is small. In graph I however the virtual π can be as close to real as possible ; a further enhancement comes from the large known cross section for diffraction scattering π-P which can be used in A. The character of diffraction scattering for the vertex A explains very well the small energy loss of the protons.

The Drell explanation can be expressed in very naîve terms for experimentalists use (**). Protons have a meson cloud. Field theory, propagators, and all that, favours virtual mesons as close to reality as possible, so let us suppose that they are completely real. Then we should observe an elastic scattering of the proton on a real π. As usual if a particle scatters at a given angle on a light particle it will suffer a larger energy loss as if it was making a scattering of the same angle on a heavy particle. This explains the difference in energy between these protons (elastically scattered on π) and the one elastically scattered on protons. One understands also, in this way, the angular correlation between the π and the scattered proton predicted by the Drell theory.

Of course this theory does not explain the double maximum but it is suggested that this could be understood by a final state interaction. This in fact does not seem contradictory to the experimental facts. A look at figure 10 shows indeed that the main experimental fact is the existence of an important group of protons which have lost a small amount of energy, the double maximum looking more like a superimposed feature.

As we have already mentioned if we consider graph II as the symmetric of I it must give an equally important contribution but the final state will be different. In I the final state was characterized by a π meson of small energy in the lab system i.e. vs the target proton. In II the final state must be a π with small energy vs the incident proton, i.e. large in the lab system. Approximate calculations indicate that this energy is very roughly 6 GeV, therefore, the incident proton must have lost about the same amount. Furthermore in I the incident proton remains a proton because in A we have a diffraction scattering, but the target proton can become a neutron in 2/3 of the events (by charge independence) in II the situation is symmetric. Therefore to the group of nucleons having lost $\sim$ 1 GeV must correspond an equally important group of nucleons having lost $\sim$ 6 GeV the distribution in energy being much broader than the first one. If one considers only protons the ratio of the numbers in the two groups must be 1/3.

We can now bring our attention to the work done by Morrison [16] in the hydrogen bubble chamber. The idea was to see if it was possible to identify events of the type found by Cocconi et



(*) The explanation of the phenomenon by graph I with a diffraction scattering in A had been made by Amati and Prentki independently of the works of Drell and Hiida.

(**) When I, or other experimentalists of my type, speak of Feynman graphs, I have always the impression of hearing a very young, pure girl, of strict Victorian education, discussing the Kinsey report.

al in a bubble chamber and what they looked like, and possibly to extend that work to π-P events. From the pictures of the 30 cm chamber, one selected the events with only two charged secondaries (two prong jets). The elastic scatterings were separated (this can be done accurately since it requires only angle measurements and momentum measurements on the low energy tracks). Therefore the work was concentrated on inelastic events of the type :

$$1 \quad \left\{ \begin{array}{l} P + P \longrightarrow P + P + \pi^{\circ} + x\pi^{\circ} \\ \text{or} \\ \pi^{-} + P \longrightarrow \pi^{-} + P + \pi^{\circ} + x\pi^{\circ} \end{array} \right.$$

$$2 \quad \left\{ \begin{array}{l} P + P \longrightarrow P + N + \pi^{+} + x\pi^{\circ} \\ \text{or} \\ \pi^{-} + P \longrightarrow \pi^{-} + N + \pi^{+} + x\pi^{\circ} \end{array} \right.$$

Most of the glancing collisions should be contained in this sample. One also hopes that a relatively large proportion of the events are true cases of single π production (either π^+ or π°) i.e. that x is 0. Strictly speaking it is impossible to prove it. The precision in momentum measurement is much too poor to make sure of the non-production of extra π°'s. In the course of the discussion it will be clear that for the most interesting events x was probably 0, and that if not, the extra π°'s donot affect the main features of the conclusions. The validity of this hypothesis has been very well confirmed in the work reported by Fiorini [13]. They have made, in the E.P. propane bubble chamber, a detailed study of events of the type 1 $\pi^- + P \longrightarrow \pi^- + P + \pi_0 + x\pi_0$ the π_0 are identified by materialization of the γ rays, and they find, among other things, that in 70% of such inelastic events there is one and only one π° produced.

In his work, Morrison selected of course, events close to the beam entrance in order to have the best possible precision for the measurements of the momentum of the secondaries. Even so, because of the small size of the chamber, the precision is not too good for high energy tracks and in particular for secondaries of momentum larger than 10 GeV/c no signification should be attributed to the exact value given, but for the fact that it is larger than 10 GeV. Therefore, it is impossible to identify "Cocconi-like events" as such (the primary has lost $\sim$ 1 GeV). Nevertheless some significant results were found just by measuring the momentum of the low energy tracks, and the angles. Also an essential part of the work was the measurement of ionization made by the mean gap length technique which made an identification possible between protons and π^+'s up to a momentum of 1.5 GeV/c.

In the proton-proton experiment, the centre of mass is a centre of symmetry therefore we can concentrate our attention on what happens in the backwards hemisphere (of CM system) which correspond to low energy particles in the lab system which are then easy to measure. The results are exhibited in figure 11 ; figure 11 is a so-called $P_L^* - P_T$ diagram, each particle is represented by a point, the co-ordinates are P_L^* (longitudinal momentum in CM) and P_T transverse momentum. In other words each point is the terminal point of the momentum vector of the particle starting from the origin of the co-ordinates (centre of mass). The advantage of such plots is that they give a more complete picture of the situation than separated angular distribution and momentum distribution histograms. Their disadvantage is that they are sometimes a little confusing. But if one remembers that the transverse momenta are small (as can be checked on the plot) one can use, instead of angular distribution plots, a histogram of the distribution in P_L^* which still does not introduce any confusion between two particles having both, for instance, $\cos\vartheta^* = -.95$ but very diférent momenta. Figure 12 shows such histograms for identified protons and identified π^+.

As can be seen on the figures all but three particles of the backward hemisphere have been identified (protons or π^+ mesons). This is not in fact perfectly true, the unidentified particles which are forwards as π^+ would be backwards as protons. They have been considered as π's for reasons of symmetry and anyhow they would not distort the results very much.

The π^+'s are clustered close to zero P_L^* and their spectrum does not extend very far in the region of high P_L^*. On the contrary the protons have most of the time a high negative P_L^*, close to the values they had before the collision. This proves that most of the events selected have indeed the character of glancing collisions.

A direct comparison with the Cocconi results can be obtained by the use of the mirror system (Dobrotin and Slavatinsky) [17] i.e. a system in which the incident proton was at rest and the target proton had a total energy of 25 GeV. The determination of the high energies in the system can be

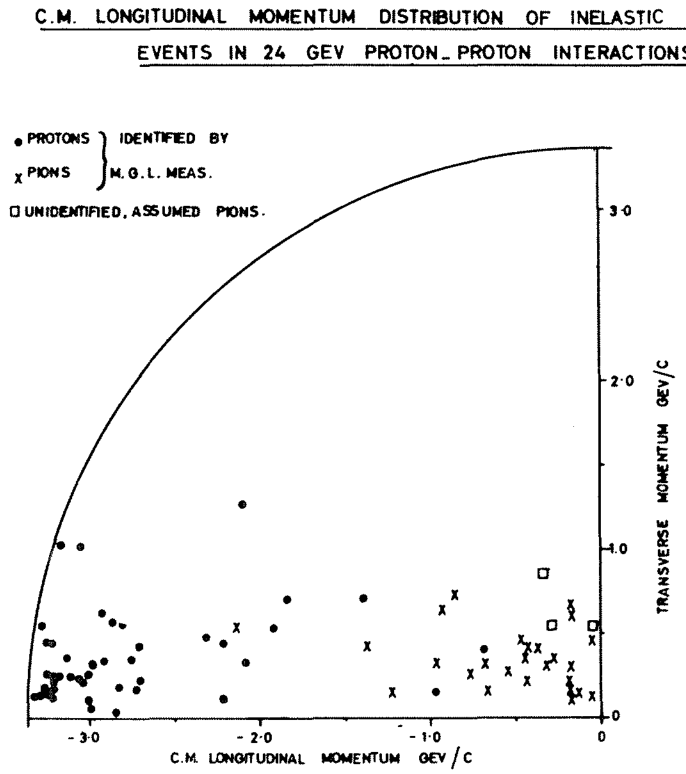


Figure 11 - $P_L^* P_T$ plot of two prongs p-p collisions (backward hemisphere).

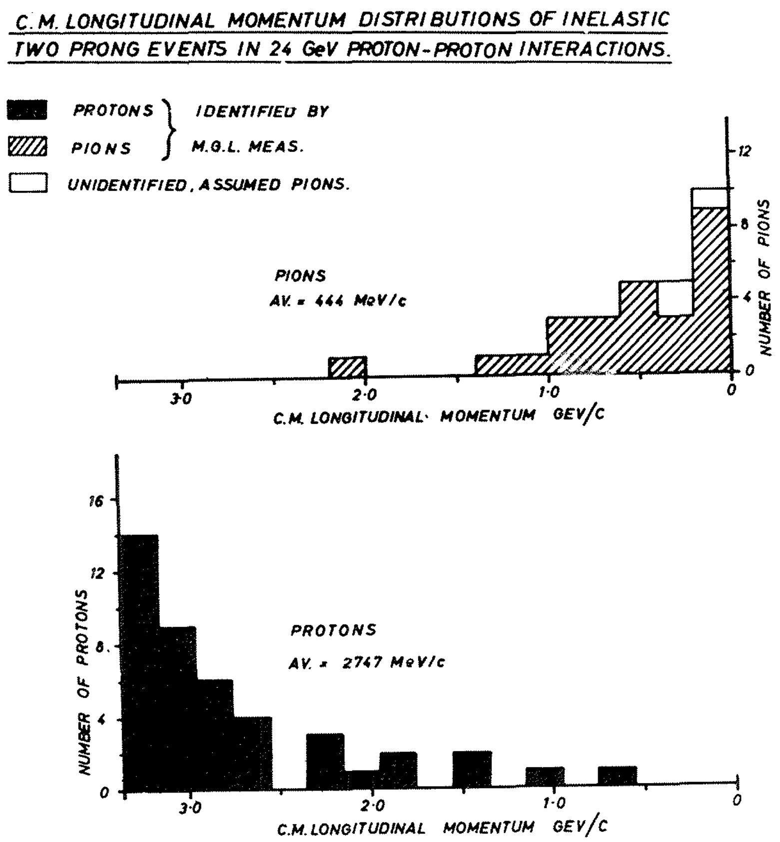


Figure 12 - P_L^* histograms for secondaries of two prong p-p collisions (backward hemisphere).

done with good precision since it depends only on angle low momentum measurements in the lab system. 25 GeV/c can be measured with an accuracy of ± 200 MeV/c.

The results are given in figure 13 which is a combination of the data obtained by Morrison and by the British university collaboration. The figure represents a histogram of the energy distribution of the high energy protons (remember that elastically scattered protons are already eliminated). The essential features of glancing collisions are clearly exhibited. Most of the protons have lost less than 5 GeV. The most probable energy loss is ~ 1 GeV. Because of lack of statistics, all these events are grouped in the interval 25-24 GeV. No double bump is found in such a small sample, but the precision would allow one to hunt for it with a larger statistic.

CALCULATED MIRROR ENERGY FOR "RECOIL" PROTON

FOR INELASTIC TWO PRONG EVENTS IN 24 GeV p-p INTERACTIONS.

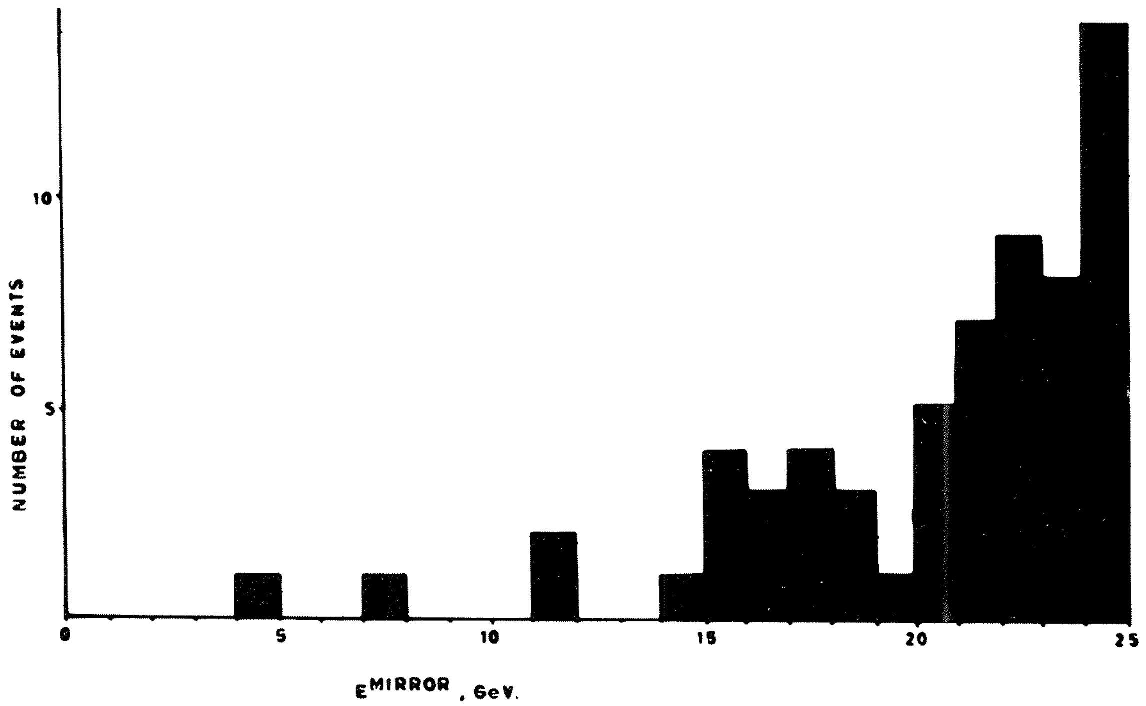


Figure 13 - Energy distribution of protons in the mirror system.

Of course one must remember that the choice of events constitutes a preselection of glancing collisions, therefore it is useful to state that the number of events in the first interval (energy loss ~ 1 GeV) correspond to a cross section of 1 mb. The total cross section for inelastic events is 30 mb.

With some imagination one can even guess the existence of a group of protons having lost ~ 6 GeV or more which would correspond to the graph II discussed previously. Their number seems even to be about 1/3 of the other.

The results found in the π^--P collisions are exhibited in figure 14 P_L^* - P_T plot and P_L^* histogram for negative particles, and in figure 15 P_L^* histograms for identified or assumed π^+ and for identified protons.

The distribution of π^+ and protons resemble very much the one found in proton-proton collision. Here no argument of symmetry permits the classification of unidentified particles in π^+ or protons. It is thought that they are π^+ but it is not proved. If they were protons the spectrum of protons would have a long tail in the region of small P_L^* values.

The distribution of negative particles (π^-) is quite different. There are two groups, one of small P_L^* with a distribution rather resembling the one found for π^+, another of large positive P_L^*.

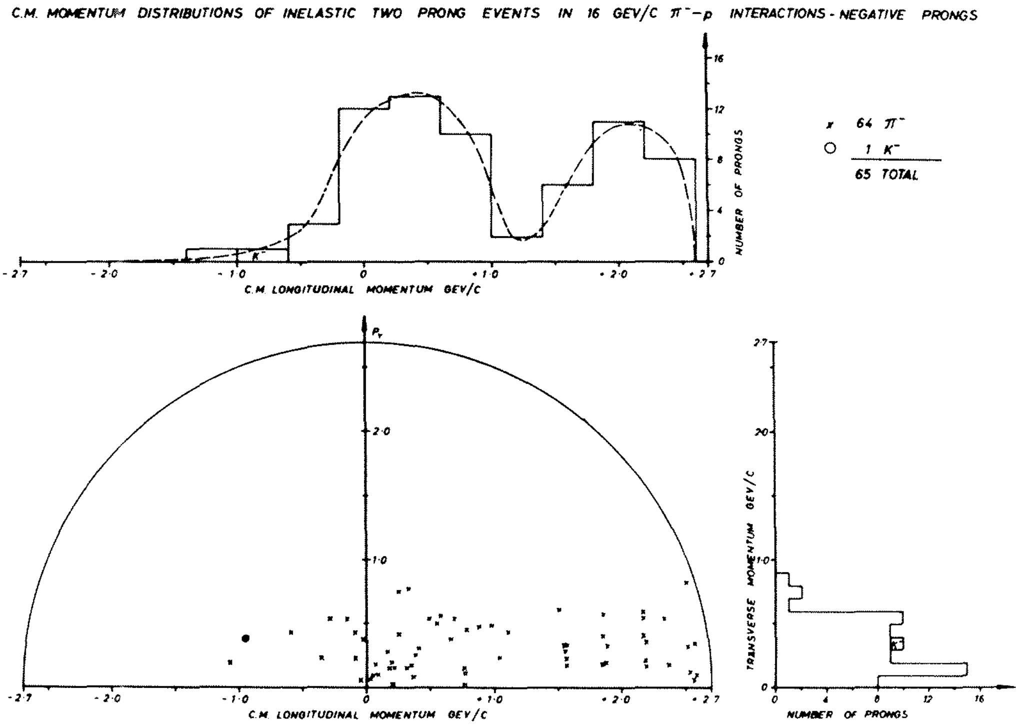


Figure 14 - $P_L^* \, P_T$ plot and P_L^* histograms for negative prongs from two prong π^--P collisions.

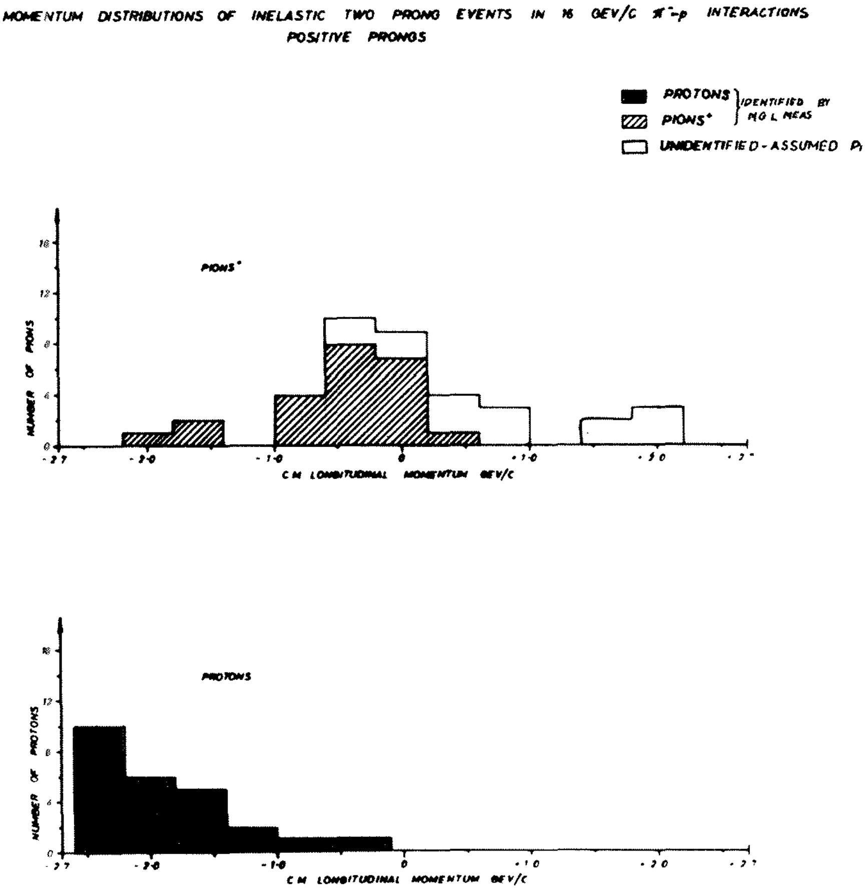


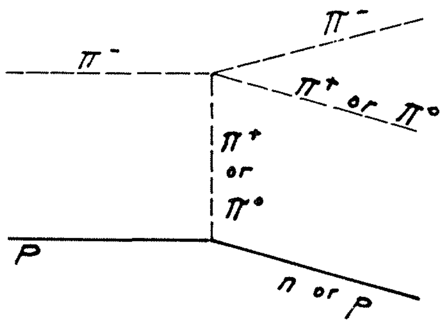
Figure 15 - P_L^* histograms for proton secondaries from other prongs π^--P collisions.

Of course the separation into two groups can be accidental, but there are certainly more than 20 events with a P_L^* larger than 1.5 GeV/c, and forwards. One must remember that the precision of momenta measurements for such particles is poor so that they can in fact be better grouped than they appear. So it seems that there are, here, quite a number of good examples of glancing collisions in which the π^- keeps most of its momentum and even its identity. Finally, we come to the correlation plot of figure 16. Here the P_L^* of positive particles is plotted vs the P_L^* of the negative particle, event for event. The shading of the squares corresponds to the nature of the positive particle (proton, π^+, unidentified). The group of high P_L^* π^-'s appears on the right. One sees that in these events there was either a very backwards proton, or a π^+ emitted backwards with small momentum. This encourages us to identify the events with one of the two reactions :

a) $\pi^- + P \longrightarrow \pi^- + P + \pi^\circ$

b) $\pi^- + P \longrightarrow \pi^+ + N + \pi^-$

and to consider they can be represented by a graph similar to the one mentioned for P-P interactions :



where the important process is the scattering of a π^- on a virtual (but almost real) π^+ or π°. The fact that the π^- retains most of the energy indicates that this process has the character of a diffraction scattering. In principle a cross section for π-π scattering can be deduced from the cross section for this type of event. It must be noted that if the π^-, π° and π^-, π^+ cross sections were equal, charge independence will predict twice as many π^+ as π° emitted. At this stage of the game there are 16 π^+ for 8 π° in the group considered.

Of course, not everything is as clear cut as we make it appear. At the top of the plot there are five events which represent the emission of a π^+ with large positive P_L^* where a π^- remain with a smaller but also positive P_L^*. If these events are also interpreted as π-π scattering, they are some sort of head-on collisions or if one prefers to call it that way, double charge exchange scatterings.

Finally the majority of events are still the ones in which the π^- and the π^+ have a small P_l^* (or the proton a large negative P_L^*), therefore in these events most of the momentum in the lab system has been given to a neutral particle, probably not a neutron, but one or several π°'s. It is possible, even probable, but not proved (see the original paper) that they correspond to events with multiple π production :

$$\pi^- + P \longrightarrow \pi^- + P + \pi^\circ + \pi^\circ$$

or

$$\pi^- + P \longrightarrow \pi^- + N + \pi^+ + \pi^\circ$$

and are therefore more complex than the one described above.

It might appear that I have given very long explanations for a work which concerned only two times 70 events. But it seemed to me that this is a good example of what is going to happen to the physics of high energy collisions. One selects special kinds of events, one tries to find meaningful correlations between several parameters and one hopes in this way to select from the swamp of all high energy collisions pure physical processes which can be interpreted more or less correctly by a simple theory. The hunt for resonances represents in fact the same sort of attitude. And so, step by step, like an onion is peeled, a clearer total picture of high energy collisions will appear.

Strange particles - We come now to our last topics : strange particle production in high energy collisions. There are at least three reasons for studying such processes :

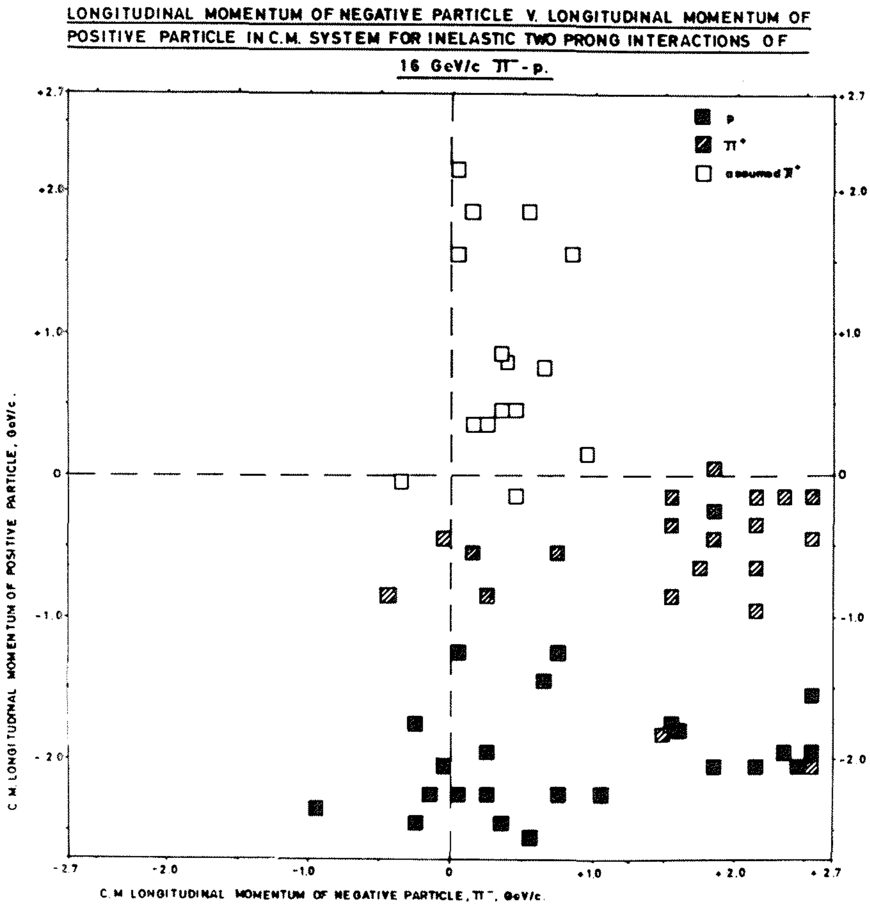


Figure 16 - P_L^{*+}, P_L^{*-} correlation plot.

1/ Strange particles are a fact of life, they are produced in high energy collisions, so we have to study this production.

2/ Since the production of hyperon pairs is very rare, hyperons are made from the initial nucleon or nucleons. Therefore, by looking at hyperons, one might have an idea of the destiny of a baryon in a high energy collision. This idea will be less biased than if one tries to identify directly a proton. This is only a working hypothesis, based on ignorance, in later developments the difference of behaviour of a baryon when it remains a nucleon, or becomes a hyperon will be the interesting topic.

3/ Since K interactions have a short range, since nucleons have this interesting property of becoming hyperons, it is possible that strange particles production is the thing to study if one wishes to arrive at a deep understanding of the nature and of the possible structure of a nucleon ; more so than π mesons which seem to be like the always present and the always annoying mosquitoes of the jungle.

The results I am reporting are essentially based on the work done in the 30 cm CERN hydrogen bubble chamber (primaries π^- of 16 GeV, protons 24.5 GeV). Groups : (CERN-Pisa-Trieste) [18] and the E.P. propane bubble chamber (primaries π^- of 6, 11 and 18 GeV) group (E.P.-CERN-Milano-Torino-Padova) [19].

CROSS SECTIONS -

In the hydrogen bubble chamber the cross section for hyperon production (Λ° + $\Sigma^{\pm}_{\circ}$) is 1,3 ± 0,2 mb, for K° production 2,9 ± 0,3 mb. If one assumes that in associated production there are equal chances of producing K^+ and K°, and in $K\bar{K}$ pairs equal chances for all the possible type of pairs, one find $\sigma_{YK} = 1,3$ mb, $\sigma_{K\bar{K}} = 2,2$ mb. So as already noted by Soloviev [20] (using π^- of 7 GeV) σ_{YK} does not change very much, but $\sigma_{K\bar{K}}$ is increasing with energy.

The most striking feature of hyperon production is the backwards peak of the angular distribution in the CM system (already reported in Rochester 1960). This is illustrated in figure 17 $P_L^* P_T$ plot for Λ°'s. The conventions of the plot are quite sophisticated. Each Λ is represented by a circle. Since each Λ has a probability, smaller than one, to decay in the chamber each event found has a weight which is the inverse of this probability, the area of each circle is proportional to the weight. Each circle is more or less black, this is correlated to the number of charged prongs of the interaction which produced the Λ, blanks mean no charged prong, total black means 8 prongs or more. Finally there are circles whose black is not black but shaded, they correspond to V° events which could be either K's or Λ's, one sees that their presence does not distort the distribution. Anyhow, statistical sampling, indicates that most of them are Λ's.

In spite of the complication one sees on the figure the advantages of a P_L^*, P_T plot. Clearly a Λ with $P_L^* \sim - 200$ MeV/c and $P_T < 100$ MeV/c is physically not very different from a Λ with $P_L^* \sim + 200$ MeV/c, $P_T < 100$ MeV/c but on a cos ϑ^* plot they will however appear as extremely different. Whereas a Λ with a very high negative P_L^* will be on a cos ϑ^* plot confused with one of them, and probably represent something very different for a physical interpretation.

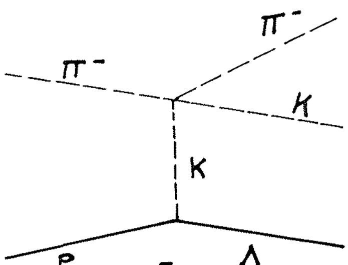
It has been often argued, especially by the authors themselves, that this distribution could be just the consequence of a bias. Indeed forward Λ's will be high energy Λ's in the lab system, they will normally leave the chamber before decaying and even if they decay inside the chamber they will escape detection because of their high energy. The answer is the following, if the Λ distribution were in fact symmetric, from all Λ's emitted forwards with $P_L^* > +500$ MeV/c, 13 will decay inside the chamber and we detected one. In P-P collisions where everything is symmetric we should have found 7 Λ's of such high energy and we found 6. So it seems that the physicists who scanned the pictures are not entirely blind to such high energy Λ's.

Now, how can the feeling conveyed by the P_L^* - P_T plot be worded ? My suggestion is the following. In a Λ producing collision, the Λ has equal probability to appear with any longitudinal momentum from the maximum possible to O provided the longitudinal momentum is negative. This extends also to a small region $P_L^* < 500$ MeV/c of the positive P_L^*. But it is highly unlikely that the nucleon suffers a big longitudinal momentum reversal, i.e. a great momentum change without loss of energy. This is probably connected to the narrow transverse momentum distribution which is also valid here as a glance on the plot will show.

This result is valid for any sort of collision in contrast to the proton distribution, reported in the preceding chapter for glancing collisions. The Σ distribution is essentially similar, possible differences are not established well enough to be discussed.

Therefore the next question is the following : What is the K° distribution ? It is shown in figure 18 also a P_L^*, P_T plot, which a first glance looks like a statistical cloud. The fact that the distribution does not extend to very high P_L^* is of course trivial. K mesons are heavy and their birth cost a lot of energy. Therefore they cannot get a very large momentum especially if the primary particles do not like to lose all their capital (in energy) as shown previously.

But it is more interesting to notice that the cloud is not centered on the zero P_L^* line, rather it is displaced into the positive P_L^* region (forwards emission). Since the transverse momentum, is as usual limited, this feature appears in a striking way on a cos ϑ^* plot figure 19 where a forwards peak is found (after correction for detection probability). Now of course comes the question what type of production contributes to the forward peak : $K\bar{K}$ or YK or both. This can only be answered by the study of pairs of V events (either $K\bar{K}$ or YK). Pairs of V events can possibly give information on the detailed mechanism of strange particle production. For instance the question has often been asked : does associated production go via π exchange or K exchange as sketched in the two following graphs ?



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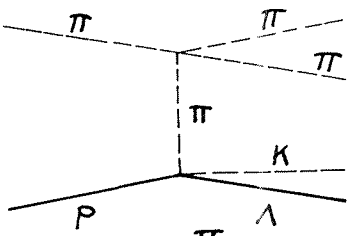

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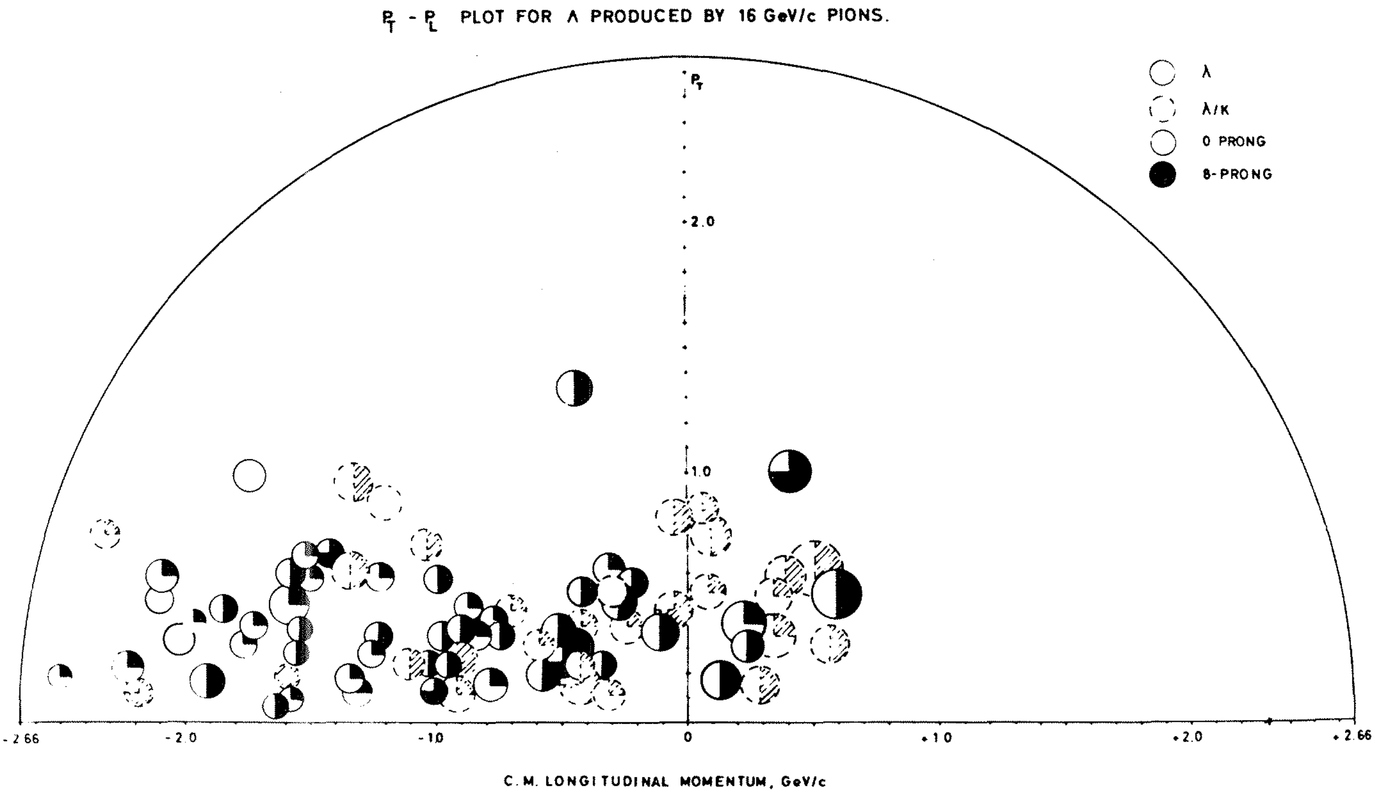

Figure 17 - P_L^*, P_T plot for Λ^0's produced in π^--P collisions.

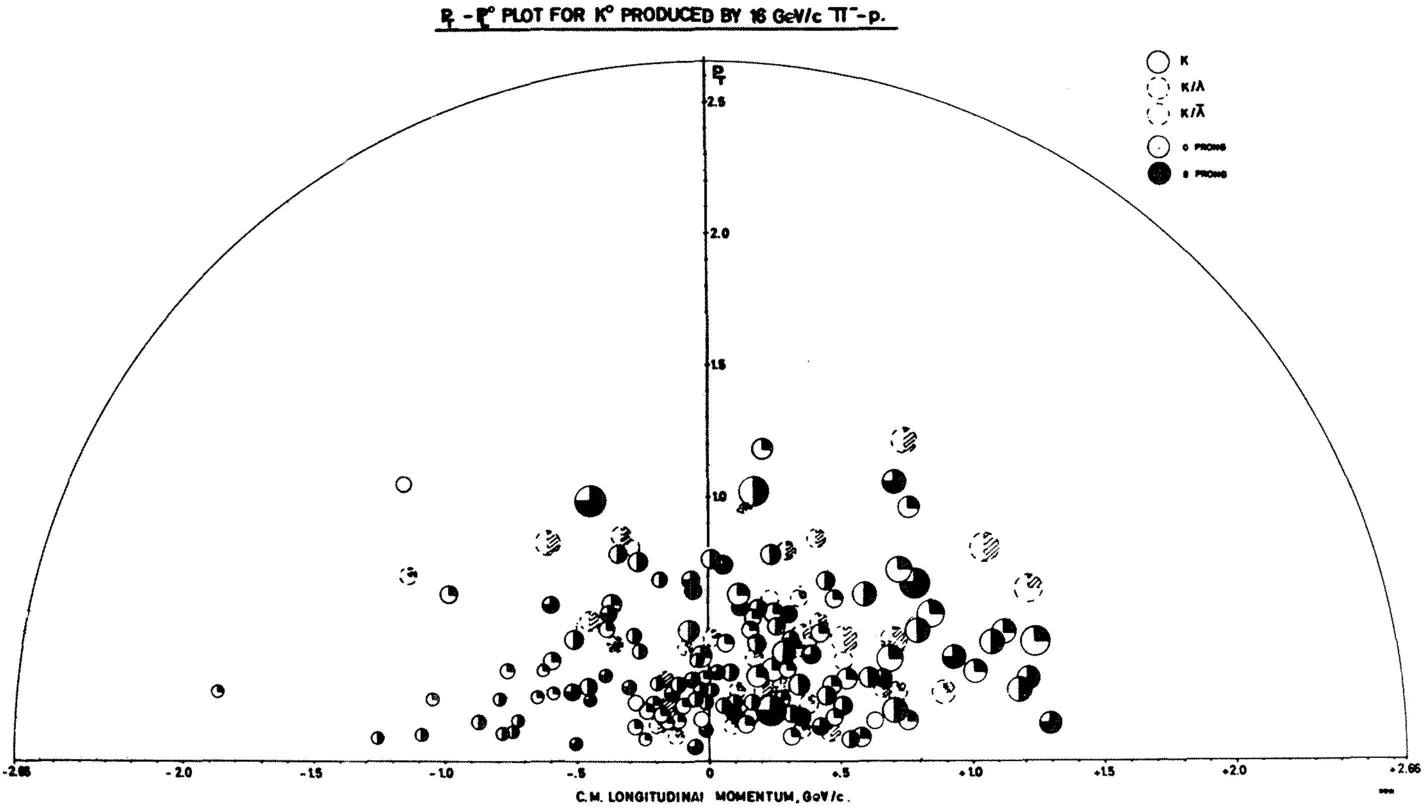

Figure 18 - P_L^*, P_T plot for K^0's produced in π^--P collisions.

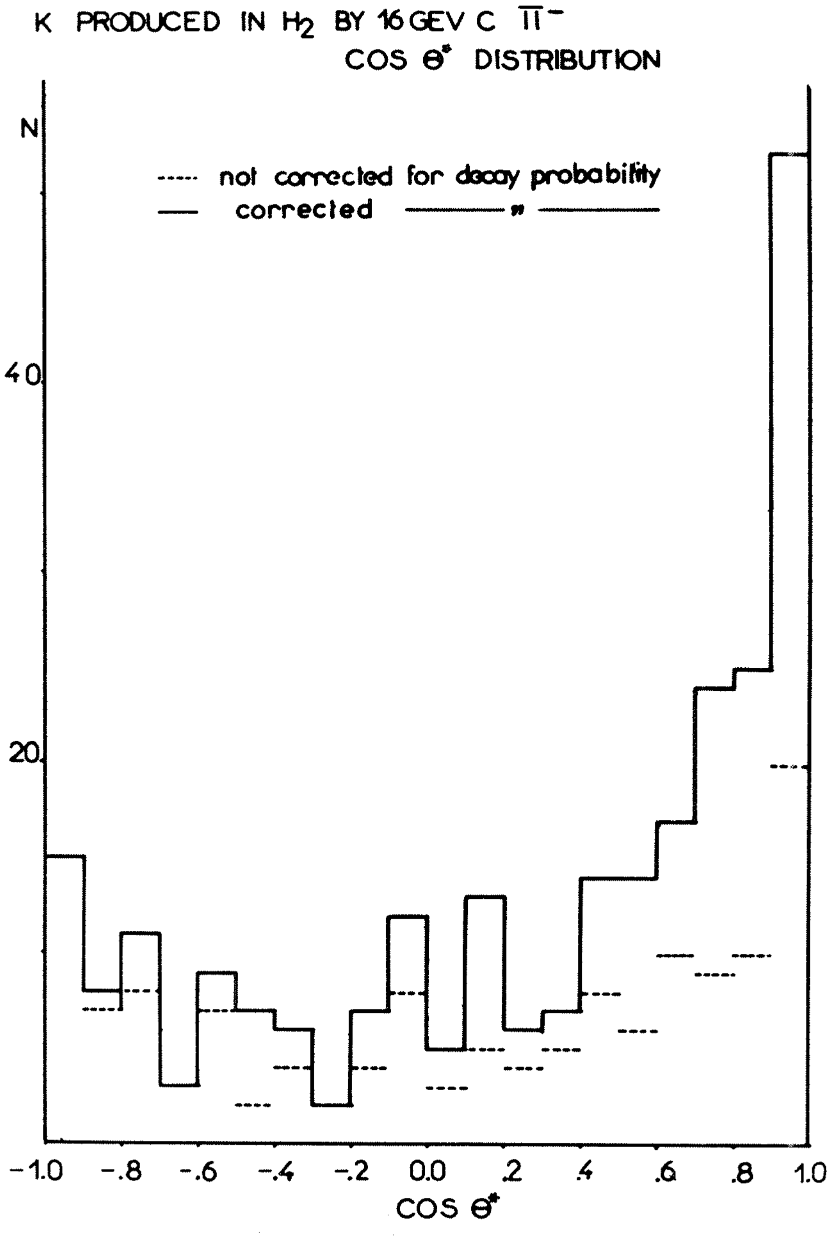

Figure 19 - CM angular distributions of K° produced in π^--P collisions.

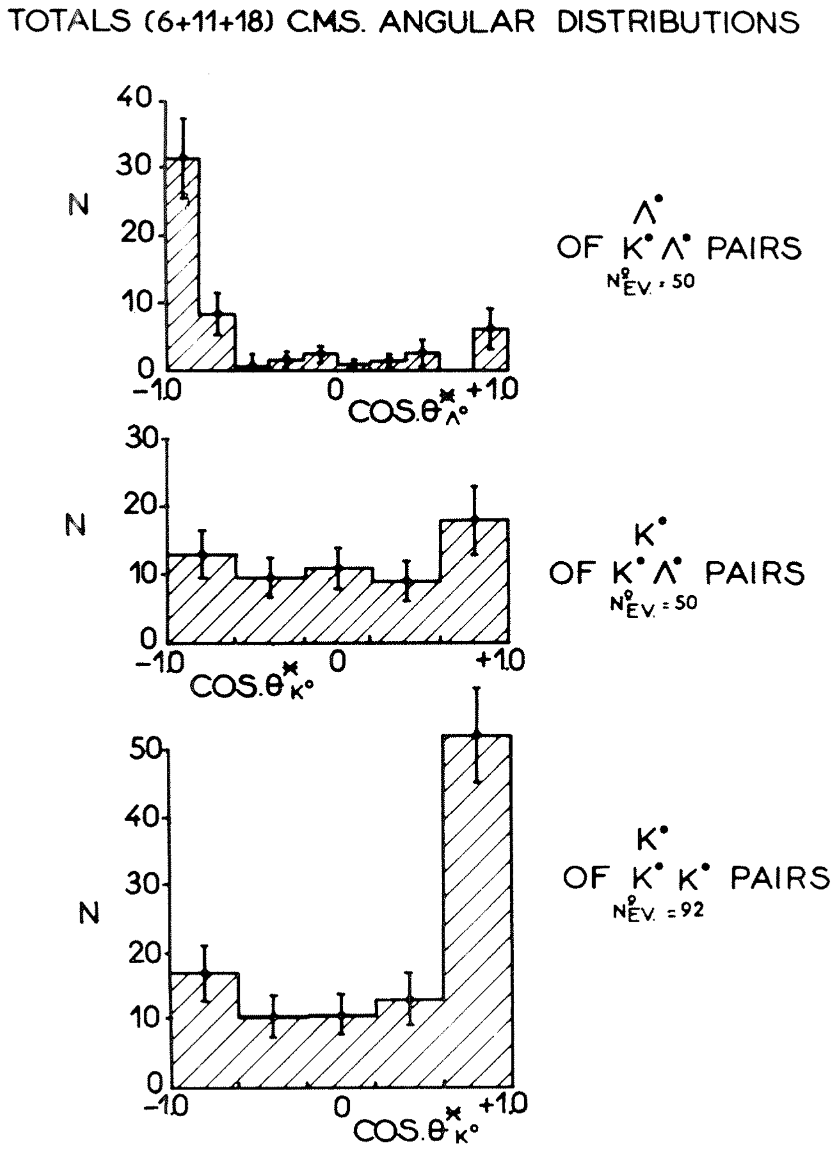

Figure 20 - CM angular distribution of Λ's and K's from ΛK and K$\bar{K}$ pairs.

Of course extra π's can be added to any of the vertices. One could expect, following Salzmann [21] that the particles produced at the top vertex would go forwards, and those from the bottom vertex would go backwards. Therefore K exchange will give K forwards, Λ backwards, π exchange will give K and Λ both backwards. These kind of conclusions are probably a little too simple. In particular in graph I the process at the top vertex can be a sort of diffraction scattering which will give the K still going backwards in the CM system. The experiment could not be done in the 30 cm chamber because of its small size and a special run was done in the 80 cm Saclay bubble chamber, the pictures are not yet analyzed. However the group working on the pictures of the E.P. 1 m propane chamber have already found very interesting results which are illustrated in figures 20 and 21. Figure 20 shows the angular distributions of Λ's and K°'s produced in ΛK events, and in $K^{\circ}\bar{K}^{\circ}$ pairs. One finds the usual backwards peak of the Λ's, a flat distribution for the Ks produced with Λ's, but a strong peak forwards is observed for the Ks of K$\bar{K}$ pairs. Figure 21 shows the distribution of the angle in the CM system, between the two K of the same pairs. Here again a strong correlation is observed, that is the two K's of a pair are both going forwards. This is a really exciting

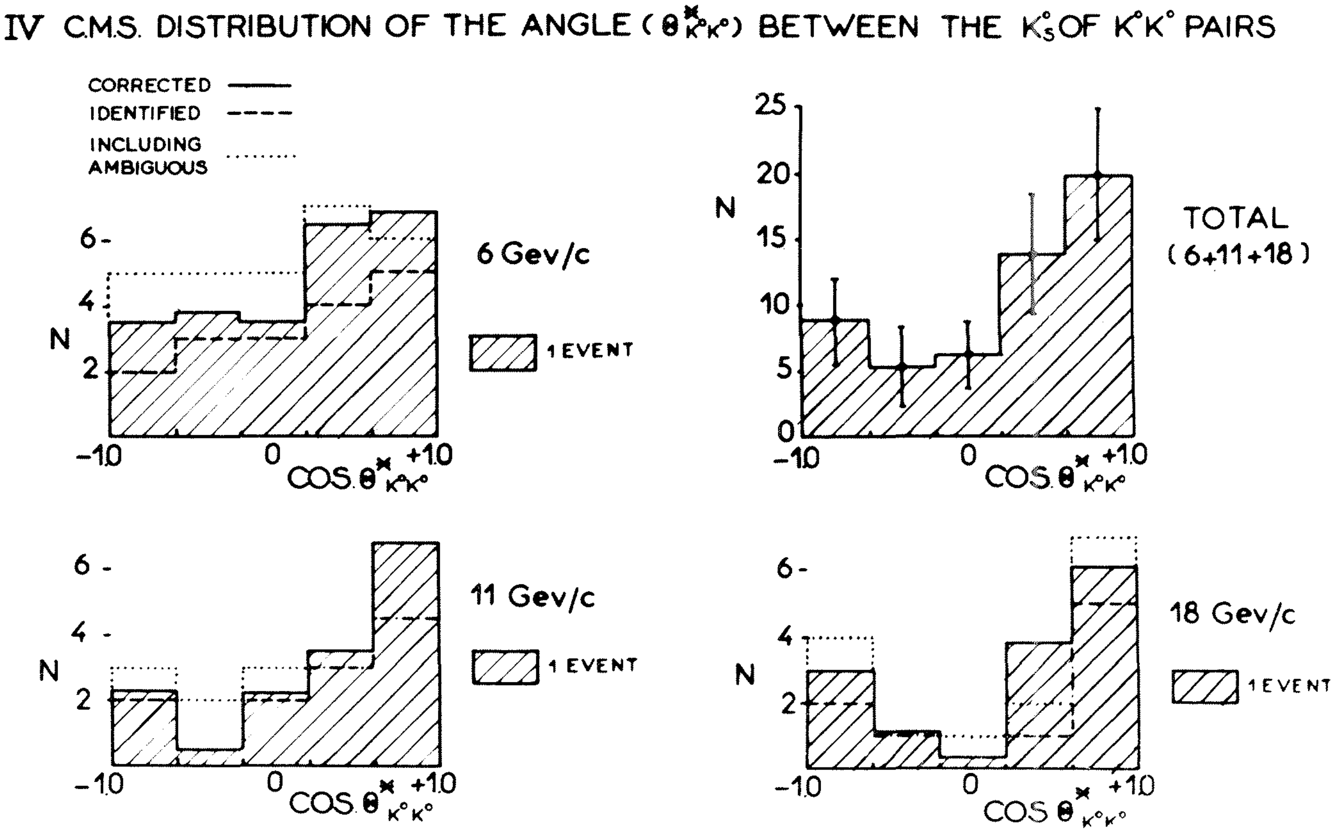

Figure 21 - Angular correlations between K° of $K^\circ\bar{K}^\circ$ pairs.

result. It is too early to decide what is its signification. Is it just the result of a peripheral collision with the particles of the top vertex going forwards ? Does it mean a strong $\pi\pi$, $K\bar{K}$ interaction going maybe through a resonant channel ? This will be known soon, but again preliminary as it is, the result shows how progress is made towards a somewhat deeper understanding of high energy collisions.

Before concluding this talk I would like to add a few words. The actual theoretical situation of high energy collisions has some resemblance to the one which existed more than 10 years ago in cosmic rays. There was a great controversy between the plural and the multiple production. The partisans of the plural production (Heitler) thought that only one meson was produced in elementary proton-proton collision. Meson showers were then thought to be the result of several collisions inside of a complex nucleus. Whereas Heisenberg thought that there was multiple production of π mesons in an elementary collision.

We know, now, that Heisenberg was right, there is multiple π production. But the ideas on peripheral collisions, which, to me, look like a sort of improved perturbation approach, bear in fact a strong resemblance to the basic attitude of the plural production theory. Now, if we are studying peripheral processes, it is not only because of their intrinsic interest, it is also because we want to isolate them, eliminate them, to be left with something that I will call catastrophic events. We hope that such events will bring information on the inside, the core of elementary particles. Of course two dangers are waiting for us on this path. It is possible that we will go on peeling the onion indefinitely, going from peripheral events to others a little less so, but without encountering any drastic change and at the end there will be no core. It can also be that the core collisions will always manifest themselves as obeying the statistical theory. But we should hope for the best.

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HIGH ENERGY PHYSICS ABOVE 10 Gev

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I - INTRODUCTION -

We shall be concerned with the collisions of strongly interacting particles at very high energy, which, for the purposes of this talk, has been defined to mean >10 GeV in the laboratory system. The richest part of physics for study in this very high-energy region is the large momentum transfer collisions which probe deep into the mysterious cores of the interacting particles, as is done, for example, in the electron-proton scattering studies discussed earlier in this conference [1]. However, in contrast with electrodynamic interactions, we are not using known (•) probes, such as electrons and photons, but pions and protons for which it is hoped to gain an understanding of the detailed character of their interactions at large momentum transfers from studies of their "cores" in such collisions. It still remains a task for the future in strong interaction theory to grab this problem by the horns. Cocconi [2] has already emphasized at the CERN conference in June that with energies in the laboratory of >10 GeV one can experimentally probe with very large momentum transfers of > 5 GeV/c, corresponding to lengths smaller than several hundredths of a fermi. This is an attractive possibility. However, in the present state of theory, it is preferable to concentrate on a more limited class of processes such as can be analyzed with dispersion methods, which provide us with some rule statements with the aid of the optical theorem, or with phenomenological comparisons to other measured parameters at lower energies.

The discussion of this paper is aimed at such processes, which are generally characterized as high-energy, low momentum transfer collisions. In all cases the analyses which carry beyond the stage of a pure phenomenological comparison of data can be taken to be accurate only to the extent to which the enigmatic core can be ignored.

II - RELEVANT DATA (••) -

The total proton-proton cross-section is constant, $\sigma_{pp,\,total} \approx 40$ mb from 10 to 24 GeV as measured by accelerators, and according to the Perkins [3] report to the CERN conference in June, appears to remain so, up beyond 10^4 GeV in the cosmic ray measurements.

The total elastic p-p cross-section is ≈ 9 mb at 24 GeV though its trend with energy above 10 GeV is not well known. It appears, however, that the diffraction peak in elastic scattering becomes more narrow with increasing energy ; for example the elastic cross-section at a momentum transfer of 1 GeV/c has fallen to less than 2×10^{-4} of its forward value, for an incident proton laboratory energy of 24 GeV, whereas it has fallen only to $\approx 2 \times 10^{-3}$ for this momentum transfer at a lower incident energy of 6 GeV.

The total anti-proton-proton cross-section is still slowly decreasing with energy between 13 GeV, where $\sigma_{\bar{p}p,\,total} \approx 52$ mb, to 20 GeV, where $\sigma_{\bar{p}p,\,total} \approx 46$ mb, and does not appear to have reached its asymptotic limit (if in fact such a limit exists for $\sigma_{\bar{p}p}$ at very high energy).

The total pion-proton cross-sections are likewise still very slowly falling with increasing energy in this region and appear to be still approaching but not yet to have attained their asymptotic values, as shown by the following table.

(•) That is "known" to the limits of present tests of quantum electrodynamics.

(••) Reviewed in preceding report of C. Peyrou.

E_π(GeV)	$\sigma_{\pi^-p,\ total}$ (mb)	$\sigma_{\pi^+p,\ total}$ (mb)
6	29	27
10	27	25
16	25	

The difference ($\sigma_{\pi^-p,\ total} - \sigma_{\pi^+p,total}$) is very roughly a constant 2 mb at these energies. The elastic σ_{π^-p} cross-section drops from 5 mb in the 5 GeV energy region to 4 mb at 16 GeV and the angular distribution, as in the p-p case, resembles diffraction scattering with the forward peak narrowing as the energy rises.

The strange particle production cross-sections are a small fraction (< 3 mb) of the total interaction cross-sections, which may be discussed initially in terms of pion and nucleon interactions alone.

In particular classes of inelastic cross-sections, corresponding to processes in which the incident projectile has made a low-momentum transfer, or peripheral, collision and has retained all but a small fraction of its initial energy, particle groups, or bumps, in the cross-section, have been found.

III - HIGH-ENERGY THEOREMS -

From the observed general features, physics above 10 GeV has a more or less uniform appearance. No new resonances are found and total cross-sections suggest a slow approach to constant asymptotic values, whereas elastic cross-sections continue to decrease slowly with energy apparently less rapidly than (energy)$^{-1}$, as their forward diffraction peaks become narrower. That both pion and nucleon cross-sections share this behaviour is what may be anticipated if these interactions involve identical intermediate states.

It appears that higher laboratory energies are necessary, perhaps > 50 GeV so that there is > 10 $Mc^2 \approx 10$ GeV available in the centre of mass, before the total cross-sections settle into their asymptotic limits. In this limit Pomeranchuk [4] has shown that if the total cross-sections for pion-nucleon and nucleon-nucleon scattering do approach constants, the following relations hold :

$$\sigma_{\pi^+p} \longrightarrow \sigma_{\pi^-p} \tag{1}$$

and

$$\sigma_{\bar{p}p} \longrightarrow \sigma_{pp}$$

as $\omega_L \longrightarrow \infty$.

The Pomeranchuk theorem was first proved in 1958 on the basis of the forward scattering dispersion relations together with simple and very reasonable physical arguments ; (the original assumptions have since been sharpened by others [5]). In the pion-nucleon case, for example, the dispersion relation for the forward charge exchange scattering amplitude shows that :

$$\mathrm{ReA}_{ch.\,ex.}\,(\omega_L, 0^\circ) \longrightarrow \Delta\omega_L \ln \omega_L + 0\,(\omega_L)$$

where $\Delta \equiv \sigma_{\pi^+p,t} - \sigma_{\pi^-p,t}$ is assumed constant. This behaviour contradicts the limit which we arrive at on the basis of a simple physical picture for interactions of finite range, according to which :

$$\sigma_t\,(\omega_L) > \pi\,\vartheta_o^2(\mathrm{ReA}_{ch.ex.}\,(\omega_L, 0))^2 \approx \frac{\pi}{(\omega_L R)^2}\,(\Delta\,\omega_L \ln \omega_L)^2 \approx (\text{constant})\ \Delta^2\ (\ln \omega_L)^2.$$

Since this inequality is violated by the observed constancy of σ_t we conclude that $\Delta = 0$; the alternative is that the interaction range increase at least as fast as $\ln \omega_L$.

Evidently we must wait quite some time before passing judgment on the prophecy of Pomeranchuk, as a difference of between 5-10 % in the corresponding cross-sections in (1) still persists at present energies.

There are other high-energy theorems based on analytic properties of the scattering amplitudes alone and free of any injection of experimental information or of physical assumptions. These lead to less severe, if somewhat more compelling, restrictions on the cross-sections and seem to be well obeyed if present trends continue. Thus Froissart [6] has been able to put the limit $\sigma_t \lesssim (\ln \omega_L)^2$ as $\omega_L \longrightarrow \infty$ by using the analyticity properties inferred from the Mandelstam representation to help terminate the partial wave series - the essential point in this is that only a finite number of subtractions need be made - and unitarity to bound each term of this series. Unitarity together with analyticity of the scattering amplitude in the cosine of the scattering angle within the Lehmann ellipse are themselves sufficient to provide the weaker limit of $\sigma_t \lesssim \omega_L (\ln \omega_L)^2$ as deduced by Greenberg and Low [7].

IV - GENERAL FEATURES -

Returning to present energies, the general features observed suggest that certain simplifying assumptions, which are physically attractive, are also useful for an approximate analysis of the diffraction elastic and total cross-sections. We assume that the elastic scattering is due entirely to diffraction of the incident particle wave accompanying absorption into the numerous open, strongly coupled inelastic channels as in figure 1.

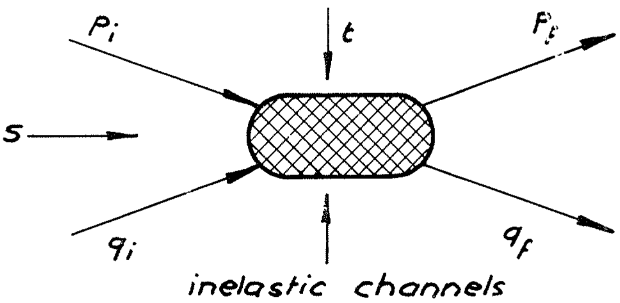

Figure 1

The scattering amplitude is given by its absorptive part :

$$A(s,t) = iA_i(s,t) \tag{2}$$

in this approximation. A(s,t) is a function of the total mass $s = (p_i + q_i)^2 = 4E_{cm}^2 \approx 2m_t \omega_L$, where m_t is the mass of the target particle and ω_L, the incident laboratory energy, and of the invariant momentum transfer $t = (p_i - p_f)^2 = -2p_{cm}^2(1 - \cos\vartheta)$. By the optical theorem the forward amplitude $A_i(s,0) = (s/8\pi m_t)\,\sigma_t(s)$, and therefore the forward differential elastic cross section is :

$$\left(\frac{d\sigma_{el}}{d\Omega}\right)_{0^\circ,\,lab} \simeq |A_i(s,0)|^2 = \left[\frac{s}{8\pi m_t}\,\sigma_t(s)\right]^2,$$

a relation well-satisfied experimentally above several GeV. Writing for non-forward angles :

$$A(s,t) \simeq iA_i(s,0)\,g(st), \quad \text{where} \quad g(s,0) = L, \tag{3}$$

we find :

$$\left(\frac{d\sigma_{el}}{d\Omega}\right)_{s,t} \approx \left[\frac{s}{8\pi m_t}\,\sigma_t(s)\right]^2 |g(s,t)|^2 \tag{4}$$

and :

$$\sigma_{el}(s) = \int \frac{d\sigma_{el}}{d\Omega}\,d\Omega \simeq \frac{\sigma_t^2(s)}{16\pi}\int_0^\infty |g(s,t)|^2\,dt \tag{5}$$

where, for large energies, extending the upper limit in the momentum transfer integral to ∞ introduces negligible error (*). Equation (5) correlates the observed narrowing of the diffraction peak $|g(s,t)|^2$ with the observed decrease in $\sigma_{el}(s)$ as energy increases.

(*) When considering scattering of identical particles as in the case of p-p scattering, a factor of 2 must be removed here.

9

Inserting experimental numbers into (5) we find for 10 GeV < ω_L < 30 GeV :

$$\int_0^\infty |g(s,t)|^2 \, dt = \frac{16\pi\,\sigma_{el}(s)}{\sigma_t^2(s)} \leqslant 6\, m_\pi^2 .$$

This has the simple and appealing consequence that low momentum transfer, $\sqrt{-t}$ < 350 MeV/c, and large impact parameter collisions exhaust most of the observed cross-sections. This encourages us to neglect the core as a useful first approximation in the high-energy studies and to concentrate on the outer regions of the pion clouds in studying σ_t and $d\sigma_{el}$ at large s. Also the large values of σ_t together with the small and decreasing ratio of $\sigma_{el}/\sigma_t \lesssim 1/4$ lead to a qualitative picture of nuclear particles surrounded by a large "grey" and fuzzy cloud of pions.

Translated into a mathematical approximation for σ_t and $d\sigma_{el}/d\Omega$, this becomes the "strip approximation" of Chew and Frautschi [8]. In the language of diagrams this means approximating figure 1 by figure 2 :

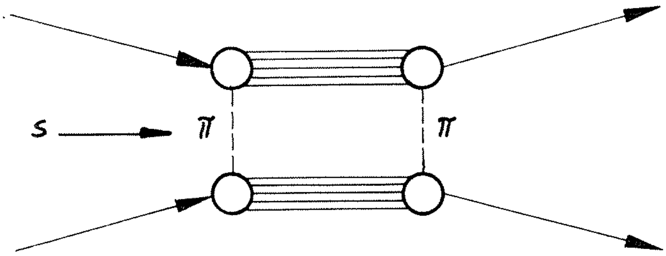

Figure 2

or by one pion exchange amplitudes in calculations of inelastic processes, and via the unitarity relations, by two pion exchange amplitudes for the diffraction scattering.

V - FORMAL DEVELOPMENTS -

Formal reasons supporting this approximation are found in the study of the analytic properties of scattering amplitudes. There is a singularity in the momentum transfer dependence of A (s,t) when it is continued as a function of t from the physical region t < 0 to t > 0 values corresponding to the vertical exchange of any real physical particles in figure 2. The singularity nearest to the physical region comes from two pion exchange and we assume that the main contributions come from graphs of type 2 which exhibit this singularity which starts near 4 m_π^2 for large s values. Then following the Mandelstam program we write a dispersion relation for the scattering amplitude in t :

$$A_i(s,t) = \frac{1}{\pi} \int_{4m_\pi^2 + t_0(s)}^{\infty} \frac{\alpha_{2\pi}(s,t')\,dt'}{t'-t} \tag{6}$$

where the spectral region runs from the 2π exchange threshold, $4m_\pi^2 + t_0(s) \approx 4m_\pi^2(1 + 4m_\pi^2/s)$ and the contribution to the spectral function $\alpha_{2\pi}$ is to be computed by including only 2π exchange contributions. This is the "strip approximation" which constitutes the new look in attempts to calculate the high-energy cross-sections. Implementing this program by calculating $\alpha_{2\pi}$ still presents a formidable task and before outlining an approach three points may be made :

1/ Since we rely on the proximity to the t singularities we may not carry this program too far out in the t variable in the study of the diffraction cross-sections - i.e., the "core" and observed shoulders in the angular patterns are beyond the realm of present remarks.

2/ The success of this approach is yet to be established ; it is still largely untried.

3/ Lovelace [9] of Imperial College of London has recently communicated an ingenious method of using (6) to find the Mandelstam spectral function for the πN diffraction cross-section directly by extrapolation from experiment. He goes about it in this way. Combining (2) and (6) we write :

$$\frac{d\sigma_{el}}{d\Omega}(s,t) = \frac{1}{\pi} \int_{4m_\pi^2 + t_0(s)}^{\infty} \frac{\rho(s,t')\,dt'}{t'-t} \tag{7}$$

since the square of an analytic function satisfies a dispersion relation in t with the same cuts ; spin and isotopic complications all disappear at large s. Now instead of computing ρ in the unphysical $t > 0$ region to find $\frac{d\sigma}{d\Omega}$ in the physical one, we reverse the procedure and from a measured $\frac{d\sigma}{d\Omega}$ at one s_1, we determine $\rho(s_1, t)$ by extrapolation to $t > 0$.

The s dependence is then constructed from analytic arguments developed by Regge [10] in potential scattering theory, and $\frac{d\sigma}{d\Omega}$ computed with no free parameters at all energies. The fit he obtains is unbelievably good in all cases. The new trick making Lovelace's extrapolation possible is a mapping first discussed by Frazer and by Giulli and Fischer [11] of the complex t plane into the interior of the unit circle in the η plane, figure 3.

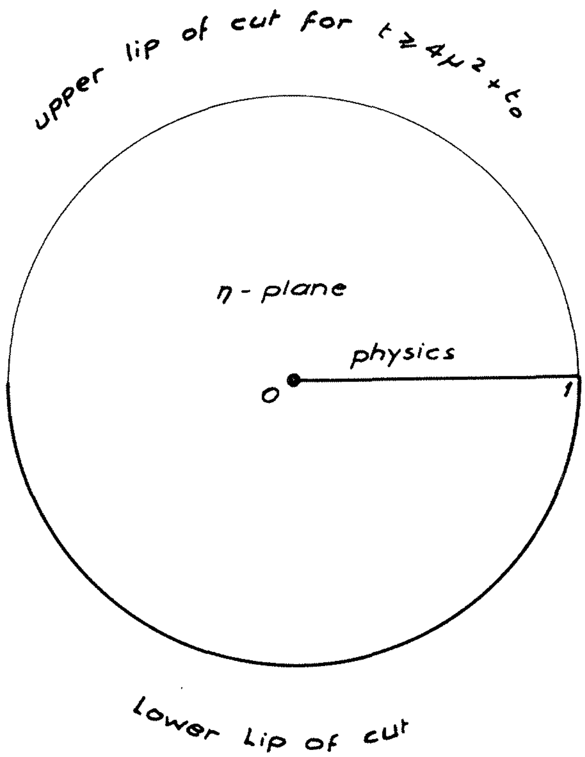

Figure 3

Dispersion relations assure the analyticity of $\frac{d\sigma}{d\Omega}$ within this circle and so Lovelace fits $\frac{d\sigma}{d\Omega}$ by a simple $f(\eta)$ at a given energy s_1 and computes directly $\rho = \mathrm{Im} f(e^{i\rho})$. To give an idea of how it goes he uses very accurate measurements by Thomas [12] at 5.17 GeV/c, which fit a Gaussian in η, $e^{-b\eta^2}$ with $\frac{\Delta b}{b} < 5\ \%$. The resulting ρ is a violently oscillating function of t, a behaviour clearly demanded by (7) if $\frac{d\sigma}{d\Omega}$ is to decrease as rapidly as observed in the experiments. Such behaviour was already found to be required by the observed constancy of total cross-sections by Goebel [13] and by the formal analysis of Regge [10] in potential scattering. Regge found this behaviour in s at very large momentum transfers t and low energies s where it is plausible that field theory is similar to potential theory. From this Chew, Frautschi and Mandelstam [14] have argued that crossing symmetry suggests the similar behaviour in t at large energy s, low momentum transfer t, the region of present interest. Combining crossing with the Regge result and the experimentally observed approximate constancy of $\sigma_{\pi p, t}$ in order to fix the s dependence of his parameters has led Lovelace to several interesting results :

(•) This implies a radius increasing with energy as $R \propto (\ln s/m_\pi^2)^{\frac{1}{4}}$ which is not fast enough to satisfy the Pomeranchuk theorem, as discussed earlier.

a)

$$\left(\frac{d\sigma_{el}}{d\Omega}\right)_{cm} = 0.2663 \frac{p_{cm}^2}{m_\pi^2} \left(\frac{s}{m_\pi^2}\right)^{-2.7\,\tau^2} \text{mb}$$

which agrees marvelously with all data above 2.5 GeV ;

b) the width of the diffraction peak narrows at higher energies leading to a :

$$\sigma_{el} \sim \left(1/\ln \frac{s}{m_\pi^2}\right)^{\frac{1}{2}};$$

c) the radius defined by $R^{-2} = \int |g(s,t)|^2 dt$ increases (*) from ~ 1 f at 1 GeV to 1.3 f at 100 GeV ;

d) according to experiment, ρ violently oscillates in t, and attempted approximate calculations must take this into account. The amplitude of the oscillations is sharply peaked in the "strip region" between the 2π and 4π thresholds (the peak is also below the threshold for 3π exchange, which is possible in N-N scattering). This is encouraging for the strip approximation.

It is of course true that other functions which are arbitrarily small in "physics" and arbitrarily large at the cut on the perimeter in figure 3 can always be constructed with hard enough effort, but I think this Lovelace application is of interest because of its simplicity, success, and suggestiveness as to the character of the spectral function.

Turning now to the problem of calculating $\alpha_{2\pi}$ in (6) our main hope as always in the strong coupling calculations lies in relating it to other physical observables. How to go about doing this is suggested by the following diagram in the s-t plane, figure 4 :

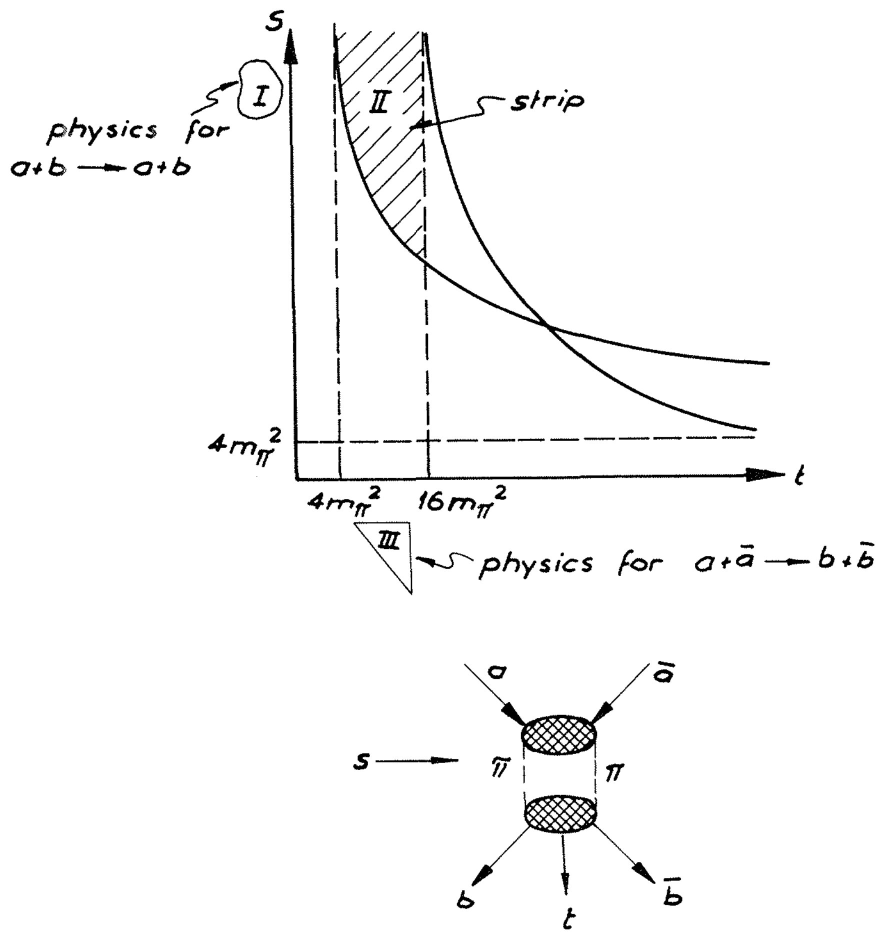

Figure 4

To calculate cross-sections for $a + b \longrightarrow a + b$ in region I, we have been led to compute $\alpha_{2\pi}$ in region II. That is an unphysical region, but down in III we find the physical region for the crossed reaction $a + \bar{a} \longrightarrow b + \bar{b}$ at energy $t > 0$ and momentum transfer $s < 0$. Let us then compute $\alpha_{2\pi}$ down in III. It is the absorptive part of A in the t channel and is given there by the physical unitarity condition in terms of the amplitudes $a + \bar{a} \longrightarrow 2\pi$ and $2\pi \longrightarrow b + \bar{b}$, in the strip approximation. This is a useful step to make in this framework because we are always dealing with two particle scattering amplitudes which we can now continue back to $s > 0$ because of their fine analyticity properties. Hence we continue the amplitudes from $s < 0$ to $s > 0$ in region II by writing dispersion relations in the s channel for fixed t. If particles a and b are pions we are led to an integral equation for $\pi\pi$ scattering ; if a = pion and b = nucleon we find an integral equation for π-p scattering in terms of a $\pi\pi$ kernel and so on. These equations were originally written by Mandelstam [15] in 1958.

Pictures may help convey the underlying idea. With the assumption of purely absorptive scattering we approximate $A(s,t)$ to its contribution from figure 1 with only real intermediate states included ; the dispersive, or off-the-mass shell part is neglected. With the strip assumption we further simplify to figure 2 and then use the dispersion relations to express the contribution for all masses of the two exchanged virtual pions in terms of the absorptive amplitudes for the two pions real. This gives :

$$A_{ab \to ab}(s,t) = \int ds_1 \int ds_2 \int \frac{dt'}{t' - t} K(s, s_1, s_2 t') \, A_{a\pi}(s_1, t') \, A^*_{b\pi}(s_2, t') + A^r_{ab \to ab}(s,t) \qquad (8)$$

where K is a complicated kernel, but $A_{a\pi}$ and $A_{b\pi}$ are absorptive amplitudes for physical processes for real particles π and a (or b) scattering through real states of mass s_1 and s_2. We can thus introduce physics for these amplitudes. The added term on the right $A^r(s,t)$ includes the important low energy resonance contributions not contained in the high-energy diffraction approximation.

Equation (8) has been written by Amati, Fubini, Stanghellini and Tonin [16] and is the starting point of their very interesting study of high-energy processes. Study of the kernel K has led to progress in the solution of (8). They observe that the energies s_1 and s_2 are much smaller than s when K is limited to the strip region in t and therefore lower energy parameters on the right hand side lead to predictions on the high-energy behaviour. As s increases so do the limits on s_1 and s_2 until they increase to the point that the $a + \pi$ process is also in the high-energy diffraction region and itself must be opened up. In this way a chain develops as in figure 5.

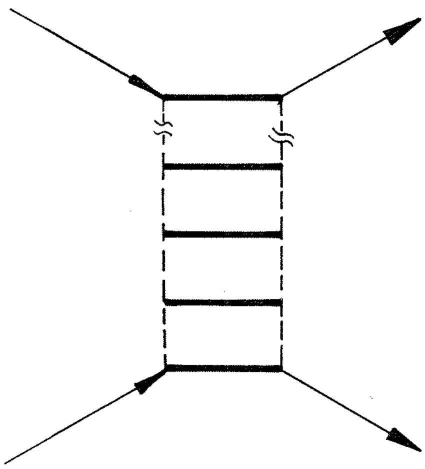

Figure 5

Successive terms in the chain arise from successive iterations in solving (8) and as energy s increases, the higher terms become increasingly important. Thus the length of the chain increases with energy and leads in calculation to a narrowing of the diffraction pattern, as observed. In the several GeV region, the first term in the iteration solution of (8) is a good approximation and the t dependence can be directly computed from $\int \frac{dt'}{t' - t} K(s, s_1, s_2 t')$ if we simplify $A^r(s,t)$ to $A^r(s)$ due to its smooth low energy behaviour. Preliminary calculations according to this program give good fits to observed πp and pp diffraction peaks.

Computing the total cross-section by the optical theorem from $A(s,o)$ leads back to the peripheral formulae of Dremin and Chernavskii [17] and of Salzman and Salzman [18] if in (8) $A(s,t) \longrightarrow A(s,o)$ on the right hand side. This will not be a valid approximation at very high energies which allow large values of s_1 and s_2 in the strip. For large energies the t dependence of the scattering amplitudes on the right becomes important as these processes also develop diffraction peaks. They too must then be opened up as in figure 5. In N-N scattering for example, for incident energies of $\geqslant 4$ GeV, the contributions to scattering in the strip region of the integrand come from 2 GeV. Further applications of this approach are now in progress and a model for the inelastic processes has been given by Amati et al. [16].

In another approach to an analysis of the diffraction cross-sections, Blankenbecler and Goldberger [19] have studied the Fourier-Bessel transform of the scattering amplitude at high energies showing that it also satisfies a Mandelstam representation. This approach is modelled after the classical impact parameter approximation in high-energy potential scattering and offers a convenient new point of departure for calculations which are now in progress at Princeton.

VI - PERIPHERAL MODEL FOR INELASTIC PROCESSES -

Finally we come then to the peripheral model for inelastic processes where our formal methods are considerably more primitive as we have no Mandelstam representation to help us in dealing with production amplitudes. The approach here is to look for particular processes and unusual kinematic conditions so that some particular Feynman graph has a very small energy denominator for the exchange of one pion between two vertices. In this limited phase space region such a one-pion exchange graph may dominate over the myriad of all other uncalculable ones if at the same time there are enhancement factors at the vertices into which the pion line is absorbed. Proceeding with optimism we consider a diagram such as shown in figure 6, for $t \sim m_\pi^2$:

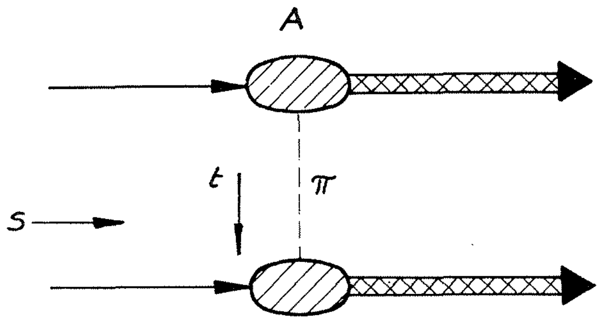

Figure 6

and write for the amplitude according to the usual Feynman rules :

$$\frac{V_A V_B}{t - m_\pi^2}$$

where corrections to the pion propagator which vanish at $t = m_\pi^2$ are neglected. If we also neglect corrections at the two vertices which vanish at the one-pion exchange pole $t = m_\pi^2$, we can insert physically observed cross-sections for the processes initiated by the pions at A and B and in this way correlate different experimental amplitudes. The accuracy and regions of validity of this procedure can therefore be directly checked by experiment. If for example two nucleons are incident in figure 6, a correlation between π-N and N-N cross-sections is predicted and can be tested directly (*).

By itself, this does not yed teach us any new physics. However, if we find a domain of validity for this approximation we can stay there with the kinematics and by changing one of the incident or final particles, learn new parameters of physical interest which cannot at present be

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(*) This procedure is an extension of the original suggestions of Chew and Low and of Goebel who proposed actually extrapolating the measurements from the physical region to the pole at $t = m_\pi^2$, and takes advantage of enhancement factors in the physical region. See the report on this subject to the Berkeley Strong Interactions Conference for a more detailed discussion of this point with applications to various problems, and for a bibliography of the numerous contributions of many authors (Rev. Mod. Phys. 33, 458 (1961).

studied more directly by any other means. The π-π cross-section is a case in point and we have heard earlier from Profs. W.D. Walker and G. Puppi [20] of their application of this method to find the $\pi\pi$ resonance. Also this past winter, Ferrari, Selleri and Da Prato [21] in a series of calculations have obtained extremely good fits to one-pion production events in p-p collisions in the several GeV range, and work reported to this conference by Chadwick showed clearly the three π-N scattering resonances in the cross-section for pion production in nucleon-nucleon scattering in the low momentum transfer collisions. These arise from the graph in figure 7 according to the peripheral model :

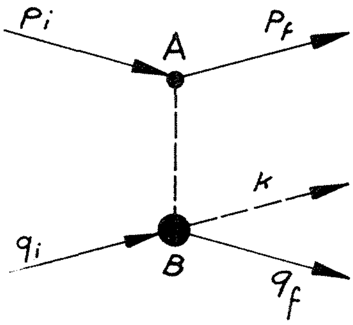

Figure 7

where p_i and p_f represent the incident and scattered nucleon in the several GeV energy and small scattering angle region, and at vertex B there takes place π-N scattering at an energy $E_i - E_f$.

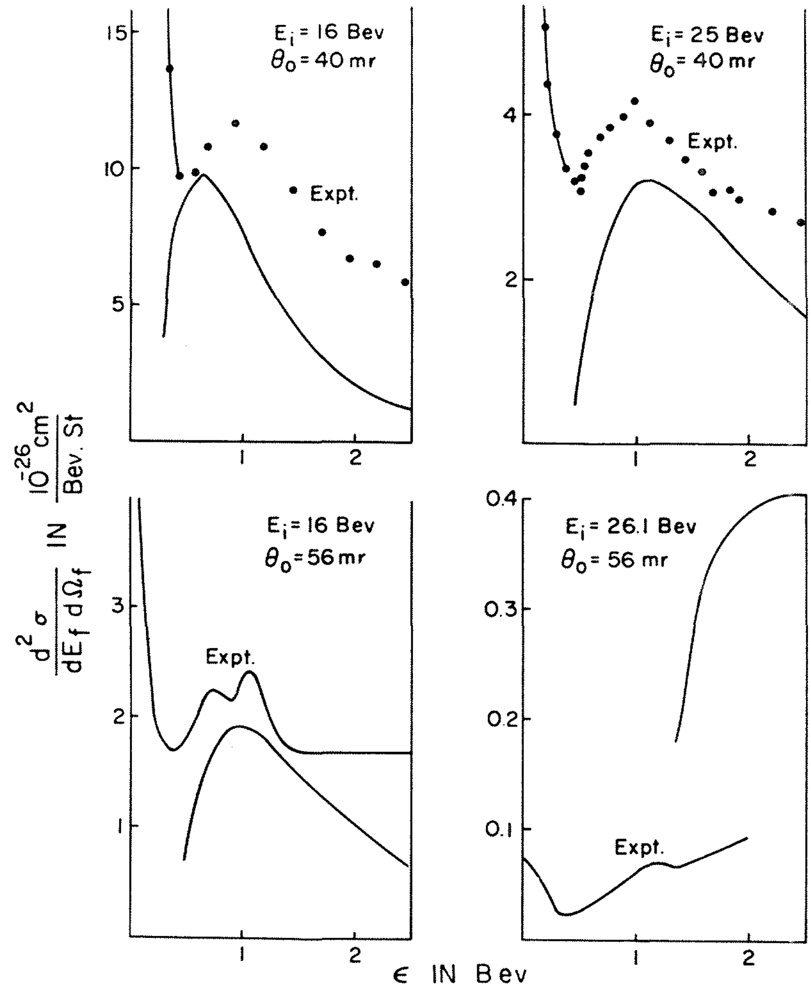

Figure 8

Returning above our 10 GeV boundary condition, we come to some recent CERN experiments that have turned up an amusing inelastic bump which the peripheral theory has managed to explain. Some of the data of the CERN [22] group showing typical parameters and bumps are shown in figure 8. Diddens presented many newer results in his talk on Friday morning. The general characteristics of this bump are that it lies ~ 1 GeV below the elastic peak and that its amplitude, like that of the elastic peak, falls off with momentum transfer in a manner similar to diffraction scattering. Since we are dealing here with high energy, low momentum transfer inelastic processes, we try to find a mechanism leading to the bump in the peripheral approximation. This leads us to consider two diagrams of figure 7 and 9 both of which contribute to one-pion production. In this region of almost elastic collisions it is easy to show that production of more than one pion is of little importance due to phase space limitations [23].

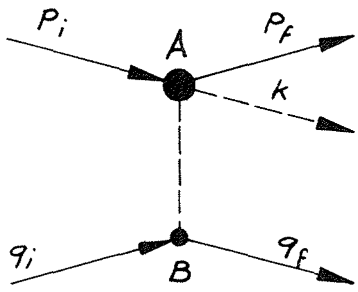

Figure 9

Some of the non-peripheral contributions, corresponding to pion bremsstrahlung in nucleon-nucleon diffraction scattering, as illustrated in figure 10 :

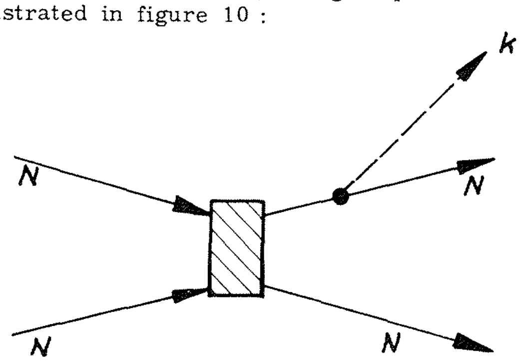

Figure 10

(there are four such graphs) can be estimated. In the present analysis this contribution is relatively small [23]. This leaves it up to the two graphs of figures 7 and 9. Figure 7 will lead to structure as we already noted in the inelastic spectrum due to the π-p scattering resonances at B. However, this is a weak candidate for a peripheral calculation in the present experimental conditions for three reasons. The pion emerges from vertex A with $t \sim -$ (1 GeV/c)^2 and is therefore rather remote from the peripheral region near the pole at $t = +m_\pi^2$. Also the 3-3 resonance peak is not found experimentally, although the camel's hump structure in the bumps coincide with the 2nd and 3rd resonances in π-N scattering quite accurately in all the recent experiments which have resolved this structure [24]. The 3-3 peak should be seen if present at least to 10 % of the probability of the higher resonances. Finally the over-all magnitude of the calculated results is smaller than the experimental number when the peripheral formula is applied.

A these high energies figure 9, dominates because it has a greater enhancement factor at vertex A corresponding to forward diffraction scattering of the very high energy incident proton from a slow pion in the cloud of the target nucleon. This I would like to add is my understanding of the diffraction dissociation mechanism proposed by Good and Walker [25] last year, and shows the relation of their discussion to the peripheral models. In the integration over the undetected final pion $\underset{\sim}{k}$ the dominant contribution comes when $\underset{\sim}{k}$ is parallel to $\underset{\sim}{p}_i - \underset{\sim}{p}_f$ and therefore the exchanged pion between A and B can closely satisfy the peripheral conditions. It is in fact a test of this model to show this angular correlation of $\underset{\sim}{k}$. The observed bump emerges in this mechanism because the nucleon sails by the pion, retaining its energy in a small angle diffraction scattering. This increases the cross section as the emerging nucleon retains a larger fraction of its initial energy until the amplitude runs out of phase space for producing the pion. The detailed calculations [23] give the predicted behaviour rather well as shown in figure 8. The bump exhibits the diffraction character of the π-p scattering at A and in the experimentally probed region is observed to stay between 0.8 and 1.3 GeV below the elastic peak. The camel's hump structure is taken to be an evidence of final state interaction which has not been included, but can be if one wants to, in this simple model. Since the diffraction scattering at A does not change the quantum nombers, the system (k,q) remains in the $I = \frac{1}{2}$ state as required for the higher resonances. The normalization of the experimental curves may be uncertain by as much as 50 % ; this will affect one parameter in this calculation - namely the cut-off in the integral over the mass of the exchanged pion which is introduced at $\sim 4\,m_\pi$. Any cut-off in the range of $4\,m_\pi$ - $5\,m_\pi$ will do. The need for a cut-off shows the weakness of peripheral approaches which in themselves provide no clue of how calculate off-the-mass shell corrections ; and for the production amplitudes we are unable to move everything on to mass shells by analytic continuation as in the Mandelstam program.

There must of course be an identical process to figure 9 in which diffraction scattering occurs at vertex B since evidently the two protons p_i and q_i are equivalent, especially when viewed from the centre-of-mass system. This is just figure 7 when the incident nucleon at A loses several GeV of its energy so that the exchanged pion can then have a large diffraction cross section for forward scattering at B. I mention this because we are also led to expect rather intense and well-collimated hyperon (Λ, Σ°, Σ^+) beams in the multi-GeV region emerging from such interactions with single K exchange and serving as useful secondary beams for high-energy hyperon-nucleon scattering studies ; the production cross-section is roughly (*) 0.1 barn/ster-GeV. The accuracy of this calculation for one K meson exchange is questionable but the qualitative prediction may be of practical value.

We have dwelled on this point to show that the mechanism is well enough understood to encourage us to look for the analogous bump in π-p collisions in the hope of measuring π-π diffraction scattering at high energies at vertex A. Morrison [26] has done this as reported earlier to the conference, and found this process which now gives us a first hint at the high-energy π-π cross-section. Integration of the theoretical formula for his results coupled with the assumption that the angular width of the π-π diffraction scattering at A is comparable to that in π-p scattering of the same energy has yielded through the connection in Eq. (5) the very reasonable result that (**) $\sigma_{\pi\pi,\,\mathrm{total}} \approx 20$ mb. Interestingly and encouragingly for this interpretation, the ratio of events in which a π^+ is emitted to those in which a π° is emitted at B to scatter the incoming π^- is 2 : 1 as one would expect from the observed smallness of the charge exchange scattering and the Pomeranchuk theorem for π-π scattering.

These peripheral events represent only a small fraction of the inelastic cross-section since we have so restricted the final phase space that there is an enhancement factor at one end of the pion line only. With greater energy losses diagrams such as figure 11 become very important :

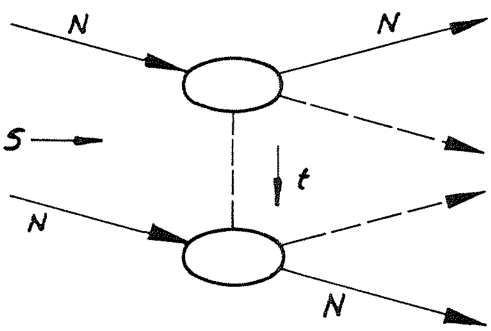

Figure 11

(*) This is computed from Eq. (3) in Phys. Rev. Letters 5, 342 (1960) using an incident E_i = 25 GeV, a final E_f = 12 GeV, forward angles, and $f_k^2 = 0.1$.

(**) A somewhat larger number of 30 mb was quoted at the conference on the basis of a rough approximate integration of the formula in Ref. [23] ; since then Dr. Hiida has performed a more accurate calculation and finds the value quoted above (to only one significant figure in view of the small number of events observed so far experimentally).

An interesting result first obtained by Berestetskii and Pomeranchuk [27] from application of the peripheral model to figure 11 is that :

$$\sigma_{NN,t} = \frac{3\, m_\pi^2}{8\pi^3}\, \sigma_N(1)\, \sigma_N(2)\, \ln s/M^2$$

if the pion mass is restricted to a constant $|t| < m_\pi^2$. This shows that the elastic cross-sections must decrease at least as fast as $(\ln s)^{-\frac{1}{2}}$ if σ_t is not to increase with energy. This result is of interest in connection with the similar finding of Lovelace on the decrease of $\sigma_{\pi N} \sim (\ln s)^{-\frac{1}{2}}$.

Finally, in closing there are several experiments which the present optimistic climate for peripheralism indicates to be interesting as well as possible.

It was pointed out earlier [28] that the electromagnetic form factor of the pion could be measured by the following sequence of experiments. First high energy small angle photoproduction of a charged pion is measured in the reaction $\gamma + p \longrightarrow \pi^{\pm} + (n)$, where (n) denotes all other strongly coupled particles which may emerge from the interaction ; only the $\pi^{\pm}$ is detected. Comparison of the measurements with the predictions of the pole formula for the exchange of one real pion corresponding to figure 12 provides the necessary information on the validity of the peripheral approximation as a function of scattering energies and angles.

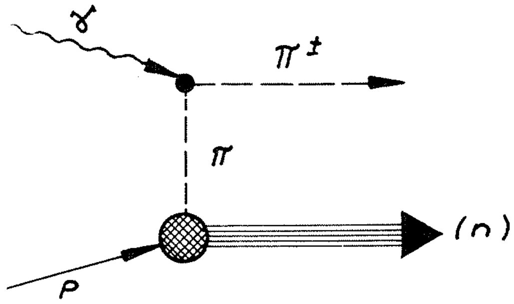

Figure 12

Once a region of phase space in which this approximation is accurate is found, the kinematics are specified so as to keep the "mass" of the virtual exchanged pion and the total energy of the π-p interaction unchanged and the experiment is repeated using an incident electron beam directly and detecting the scattered electron and the high energy $\pi^{\pm}$ in coincidence. The relevant graph is shown in figure 13 and the new information obtained is the electromagnetic form factor of the pion for space-like momentum transfers $q^2 < 0$.

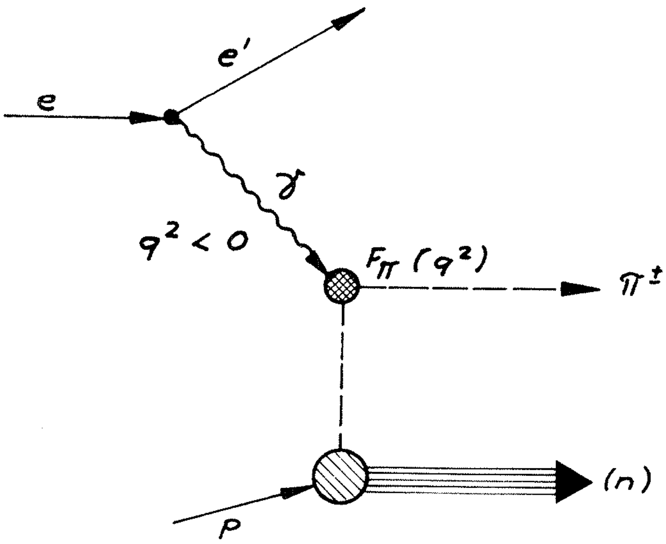

Figure 13

In order to measure the pion form factor for time-like momentum transfers $q^2 > 0$ we want to turn these diagrams around as in figures 14 and 15 :

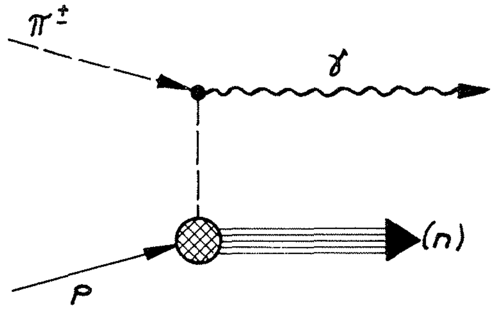

Figure 14

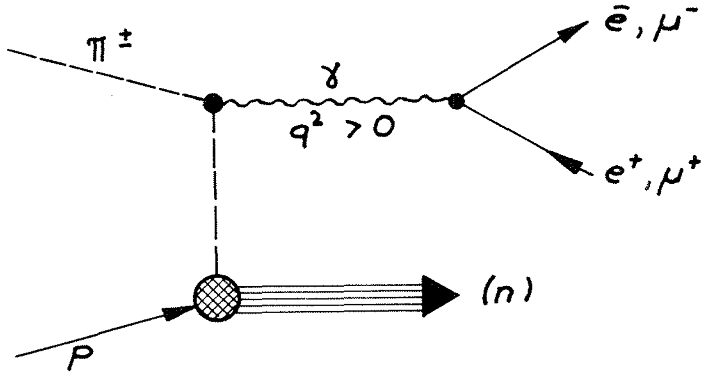

Figure 15

and produce an e^--e^+ pair. Pion production of a high-energy photon at small angles is the control experiment in this case to probe for regions of quantitative validity of the pole approximation (*). It serves the same function as photon production in figure 12. Detection of the e^--e^+ pair in the region of validity thereby established for this approximation then measures $|F_\pi(q^2)|^2$ for $q^2 > 0$. Detailed calculations by Hadjioannou [29] are encouraging with regard to the counting rate for this process for q^2 in the neighbourhood of the suspected resonance in the pion electromagnetic form factor [30] at $\approx (780\ \text{MeV})^2$. A high-energy pion beam (6 - 10 GeV) is required in order to satisfy the clashing requirements in this process of high mass q^2 for the virtual photon and low mass near the pole for the virtual exchanged pion.

If this measurement proves possible, it raises one very exciting new possibility - that of measuring the mu-meson electromagnetic vertex for very large time-like momenta simply by detecting a μ^- - μ^+ pair in place of the e^--e^+ pair in figure 15. The ratio of the $\mu^\mp$-pair to the $e^\mp$-pair cross-section should be unity to within negligible corrections of $\approx (m_\pi/\mu\text{-energy})^2$ and any deviations must be attributed to non-electromagnetic structure corrections to the μ-meson interaction for time-like momenta of ≈ 800 MeV at the vertex (**).

ACKNOWLEDGEMENT -

This report was prepared at CERN which I would like to thank along with the Guggenheim Foundation for financial support. Discussions with both experimental and theoretical colleagues here have been of great value.

(*) This experimental possibility of detecting a high-energy photon from an incident charged pion by the inverse of the reaction in figure 12 has been pointed out by Prof. B. Richter of Stanford and more recently by F. Salzman and G. Salzman in the CERN conference loc. cit. p.283.

(**) These remarks assume the absence of such corrections for the electron interaction.

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THE BOTANY OF STRANGE PARTICLES

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The importance and diversity of reports concerning the interactions of strange particles which have been presented at this meeting make it practical and desirable to confine this report largely to an exposition and summary of that work. Originally the title of this meeting was, in English and in French, "The Zoology of Strange Particles". Both Dalitz and myself, separately, wrote the organizing committee asking for a more precise definition. The committee presumably assumed that this implied a criticism and obligingly changed the English title to the "Botany of Strange Particles". I wish to take this opportunity to thank the committee for their consideration.

By the process of elimination, Botany or Zoology, in this context, seems to refer to all Strange Particle Physics, excepting resonances. As usual, it is convenient to classify the contributions, and organize this discussion in terms of the weak interactions of strange particles divided further into the leptonic decays and non leptonic decays ; and the strong interactions. Again, as usual, there are contributions which bear directly upon precisely defined problems, and there are contributions which are more nearly programmatic in nature. It is much easier to discuss the former, and in common with most reporters I shall unfairly neglect the less specific experiments which are so essential in providing us with the firm bases on which we progress.

A basic postulate, or better, faith, in our concern with the weak interactions, is our belief in "The Universal Fermi Interaction". I shall loosely define this as the hypothesis that all weak interactions can be described in terms of one basic interaction. According to a simple interpretation of the U.F.I., and a simple picture of strange particles, we should expect the Σ^-, and the Λ^0 to undergo β-decay and μ-decay with a probability essentially the same as the neutron, modified, of course, by straight forward considerations such as volume of phase space. Such a calculation predicts branching ratios for β- and μ-decay of the Λ^0 of a few percent. It has been known for some time that the actual probability or branching ratio is much smaller, a conclusion, however, generally based on events noticed during experiments designed for other purposes. Particularly the summaries prepared from the results of several such measurements may suffer from discovery biases ; those that notice an event say so, those who do not, keep quiet. A most interesting and specific experiment has been performed by a group at the Ecole Polytechnique searching for lambda β-decay. A beam of 1.2 Bev π^--mesons from Saturne at Saclay, was directed into a propane-freon chamber, 50 % freon, with dimensions of 1 meter × 50 cm × 50 cm, in a magnetic field of 17.5 Kilogauss. From 12,000 photographs, 3 000 Λ^0-decays were noted, of which 8 were observed to β-decay. Of these 8, one was identified as an electron from the kinematics of a δ-ray, 2 were established by analysis of the variation of curvature with range, and 5 stopped. Stopped electrons curl uniquely at the end of their path. The probability of identifying an electron decay is estimated as 86 %. These numbers result in a branching ratio of $0.30 \pm {}^{.15}_{.12}$ %, an order of magnitude less than that predicted by the elementary theory. An example of Λ^0-β decay is shown in the proceedings.

In the course of these measurements 254 Σ^- decays were observed, none of which β-decayed, a result which is again in contradiction to the theoretical result, but not so dramatically.

Presumably we would prefer to think of this, not of the breakdown of the U.F.I., but evidence that the hyperons, or perhaps all strange particles are in some way different. We recall in this context that the matrix element for the decay $K \longrightarrow \mu + \gamma$ is smaller than that for $\pi \longrightarrow \mu + \gamma$.

During the discussion of this result it was emphasized that, if the μ and e are as similar as we believe, the ratio of β and μ-decay should be that predicted by the theory. The Λ^0 μ-decay seems to be quite difficult to measure ; however, a Columbia group found a particularly striking example in a hydrogen bubble chamber run.

Some time ago, it was noticed that the Ξ_i-hyperon did not seem to decay into a nucleon and a pi-meson, suggesting that decays in which $\Delta S = 2$ were forbidden to weak interactions. There are limits to our ability to test this directly by the non observance of Ξ_i-nuclear decays. However, the possibility of $\Delta S = 2$ transitions should play, an important part in the mass difference between the K_2° and K_1°, where, as usual $|K_1^\circ\rangle = (1/\sqrt{2})\,(|K^\circ\rangle + |\overline{K}^\circ\rangle)$ and $|K_2^\circ\rangle = (1/\sqrt{2})\,(|K^\circ\rangle - |\overline{K}^\circ\rangle)$. As we know, these different linear combinations are respectively even and odd under CP, and as result of this different symmetry, they are linked to different real states and decay with different lifetimes. We can say that the imaginary part of their mass is different. Since they have different symmetry properties they are also linked to different virtual states and have then different self energies or different masses. The difference in mass comes from transformations of K° to $\overline{K}^\circ$ and vice versa as these states have different relative signs for K_2° and K_1°. We can estimate the mass change as in second order perturbation theory from a transition such as $K^\circ \xrightarrow[\Delta S=1]{} n\,\bar{n} \xrightarrow[\Delta S=1]{} \overline{K}^\circ$ and the inverse, then : $\Delta M = \dfrac{g/\sqrt{\hbar c}\,\Gamma\ .\ g/\sqrt{\hbar c}\,\Gamma}{M_n}$ where g/hc the weak interaction coupling constant is about 10^{-13} and Γ is an energy of the order of the K-mass. Then $\Delta M \simeq 10^{-6}$ volts $\simeq \hbar/\tau_1$, where τ_1 is the lifetime for K_1 decay. If $\Delta S = 2$ transitions, of the appropriate symmetry, are allowed a much different situation will obtain : we have then transitions such as $K^\circ \xrightarrow[\Delta S=0]{} n\Lambda^\circ \xrightarrow[\Delta S=2]{} K^\circ$ where the $\Delta S = 0$ transition will occur through the strong interactions. Then $\Delta M = \dfrac{G/\sqrt{\hbar c}\,\Gamma\ \ g/\sqrt{\hbar c}\,\Gamma}{M_n}$ where $\dfrac{G^2}{hc} \simeq 1$ and $\Delta M \simeq 1$ volt. A measurement which would differentiate between these two very small but vastly different values would then test the $\Delta S = 2$ rule. Such a measurement with, however, limited statistical validity, was performed some time ago and suggested that the mass difference was small, of the order of $\hbar/\tau_1$, and hence that $\Delta S \neq 2$.

Several recent measurements establish this more firmly. Of particular interest is a measurement reported by a Berkeley group based on an unusual effect of considerable intrinsic interest, the coherent regeneration of K_1°-mesons from a beam of K_2°-mesons in matter.

Consider the passage of K_2°-mesons through matter as the passage of the linear combination $(1/\sqrt{2})\,(|K^\circ\rangle - |\overline{K}^\circ\rangle)$. The interactions of these states will be somewhat different as they have different strangeness quantum numbers. The effect of matter will result in a modification of the incident plane wave for each of these states in a manner anologous to tne effects of an index of refraction : $e^{ikz} \longrightarrow e^{ikz}\,e^{i(a+ib)z}$. The attenuation of the beam intensity here e^{-2bz}, must be equal to $e^{-n\sigma}$, where n is the number of nuclei/cm^3 and σ is the total nucleus cross section for the K° or $\overline{K}^\circ$: therefore, $b = 1/2\,n\sigma$. From the optical theorem, $\sigma = (4\pi/k)\,\mathrm{Im}\,A(0)$, where $A(0)$ is the nuclear forward scattering amplitude ; hence $b = 2\pi n\,\mathrm{Im}\,A(0)/k$, here k is the K-nucleus wave number. This relation between the imaginary parts can be extended to the real parts and $a = 2\pi n\,\mathrm{Re}\,A(0)/k$. We can relate a to a potential, the optical potential, which would induce the same phase changes. If $\mathrm{Re}\,A(0)$ is of the order of the nuclear radius, the potential depth will be about 10^{-6} volts.

This refraction will be different for the K° and $\overline{K}^\circ$ since, having different strangeness numbers, they interact differently with matter. In particular the phase between K° and $\overline{K}^\circ$ will change in passing through matter and the odd linear combination $K_2^\circ = (1/\sqrt{2})\,(|K_\circ\rangle - |\overline{K}_\circ\rangle)$ will have induced a part of the even combination $|K_1^\circ\rangle = (1/\sqrt{2})\,(|K_\circ\rangle + |\overline{K}_\circ\rangle)$; K_1° mesons will be regenerated in the beam. The beam, unchanged in direction will contain K_1° and K_2°-mesons.

A classical optical analogue exists. Consider the passage of right circularly polarized light, which can be considered, of course, as a linear combination, with appropriate phase, of light plane polarized in two perpendicular directions, X and Y, through a medium in which the index of refraction for light plane polarized in the two directions, X and Y, is different. The phase relation between the two components will change and the resulting wave can be described as a mixture of right and left circularly polarized light.

Further qualitative remarks may be made concerning this simple picture of K_1° regeneration, a description which neglects, in particular, the decay of the K_1° and the mass difference between the K_1° and K_2°. The generation of K_1° occurs along the K_2° beam, every increment of length dl_1 adds coherently to the amplitude of the K_1° state. The finite life of the K_1° state results in a continuous decay of the K_1° amplitude and limits the magnitude of K_1° production. A large $K_1^\circ - K_2^\circ$ mass difference results in a more crucial effect. If this mass difference is large, the K_1° amplitude will fall out of phase with the K_2° amplitude and contributions from different increments of length dl will be out of phase, and on the average, incoherent. The resultant K_1° intensity will result from a sum of infinitesimal intensities instead of a sum of amplitudes in phase. The sum of squares will be small compared to the square of sums, and if the mass difference is large, the regeneration effect will be negligible.

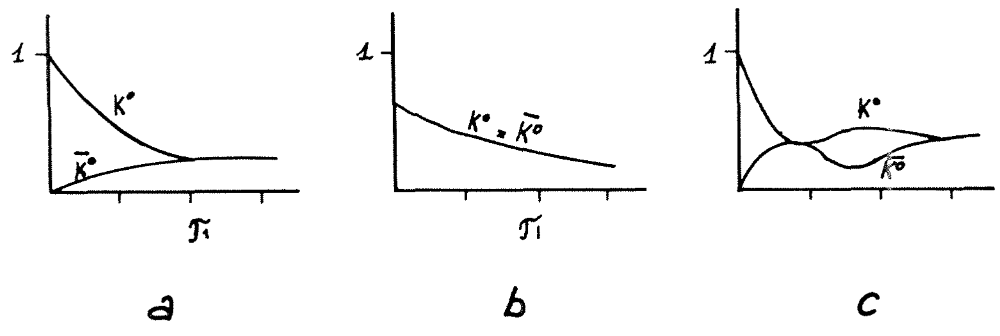
The existence of this effect is then the basis of this measurement of the mass difference by the Berkeley group. They constructed a beam of K_2°-mesons with a mean momentum of about 600 Mev/c. This beam passed through a propane bubble chamber fitted with a steel or lead plate. On the exit side of the plate, K_1° mesons were seen to decay, which had been proceeding almost exactly in the direction of the beam. The slide shows vividly the existence of this effect. The sharp peak, almost at $\cos\vartheta = 1$ results from the coherent regeneration, the broader distribution exhibits the diffraction production of K_1° mesons by the individual nuclei in the iron plate. A quantitative analysis by this group results in the conclusion that the mass difference is about $\hbar/\tau_1$ more precisely $0.84\,{}^{+.24}_{-.22}\,\hbar/\tau_1$, thus excluding an allowed $\Delta S = 2$ transition.

I will take this opportunity to make an observation, which is, I believe, due to M. Good, who developed the preceeding arguments. If the antiparticles, K° and $\bar{K}^\circ$ are oppositely effected by gravity, their energies in the earth's gravitational field would differ by about one volt, the phase between K_0 and K_0 would change rapidly, and the $K_2^\circ \longrightarrow K_1^\circ$ regeneration of the type noticed here, would occur even in the absence of matter with an intensity which would rapidly deplete the K_2° state. The existence of the long lived K_2° is then immediate experimental evidence that the gravitational effect on at least some matter and anti-matter is the same.

Two other recent measurements of the mass difference are in agreement with the conclusion that the mass difference is small. These are both based on a somewhat different principle. Since strangeness is conserved in strong interactions such interactions invariably produce either K° or $\bar{K}^\circ$ states. For the sake of definiteness I will discuss the Wisconsin-Padua-Berkeley work from which both a $K_1^\circ - K_2^\circ$ mass difference and the very interesting information concerning the $\Delta S = \Delta Q$ rule, have been derived.

In this experiment a K^+ beam is introduced into a propane chamber and observations are made on the reaction chain $K^+ + X \longrightarrow K^\circ + X''$, $K^\circ + X''' \longrightarrow \Sigma/\Lambda + X''''$ where the X represent, as initiated states, carbon or hydrogen nuclei, and as final states any of several reaction products. Since the K^+ has $S = +1$, the K-meson produced in the first reaction will be the K°. Since the hyperon has $S = -1$, the initial K-meson in the second reaction must be a $\bar{K}^\circ$.

The initial K° state can be written as $|K^\circ\rangle = (1/\sqrt{2})\,(|K_1^\circ\rangle + |K_2^\circ\rangle)$. The $|K_2^\circ\rangle$ state will decay leaving the state : $K_2^\circ = \frac{1}{\sqrt{2}}\,(|K^\circ\rangle - |\bar{K}^\circ\rangle)$, which is half $\bar{K}^\circ$. In the absence of any mass difference between K_1° and K_2° this leads to a complete description of the intensities of K° and $\bar{K}^\circ$ as a function of time, which follows a pattern as in a)

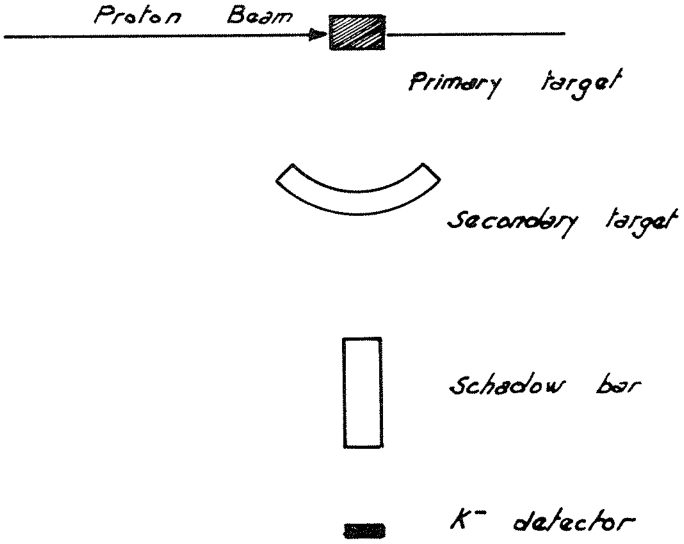


If however, a mass difference exists the relative phase of the K_1° and K_2° will vary with time and the composition of the neutral K will oscillate between K° and $\bar{K}^\circ$ very rapidly. When the mass difference is very large these oscillations will occur within any experimental resolution and the K° and $\bar{K}^\circ$ will effectively be equally present every-where. Figure b represents that situation, and figure c is a badly drawn estimate of an intermediate solution with $\Delta M \simeq \hbar/\tau_1$.

The reaction $\bar{K}^\circ + X \longrightarrow \Sigma/\Lambda + X'$ is a probe or measure of the intensity of the $\bar{K}^\circ$ state and the distribution of such events is used to determine the mass difference. In particular, the paucity of such events, initiated very close in space and time to the K-charge exchange, excludes the large mass difference required by an allowed $\Delta S = 2$ transition. More precisely, the data suggests a value of ΔM of about $\hbar/\tau_1$ consistent with the value discussed previously.

A Princeton group working at Berkeley has measured the small mass difference in a way which is conceptually similar but experimentally much different. They observe with counters

K-mesons produced in a secondary target by K-zero mesons produced by a proton beam. The set-up is shown schematically in the sketch.



The reaction chain is $P + X \longrightarrow K^{\circ} + X'$, $\bar{K}^{\circ} + X'' \longrightarrow K^{-} + X'''$. The geometry and proton beam energy is chosen such that it is unlikely that any negative strangeness hyperons or $\bar{K}^{\circ}$ or K^{-} can be produced in the primary target and interact in the secondary target. This is possible because of the different kinematics of the reactions producing these particles. They tend to go more nearly forward than the K°. The counting rate as a function of secondary counter distance is the basis for the conclusions that $\Delta M \simeq 1.8\ h/\tau_1$, which, within the quoted errors, is in reasonable agreement with the other values, though larger. More important $\Delta S = 2$ is again excluded.

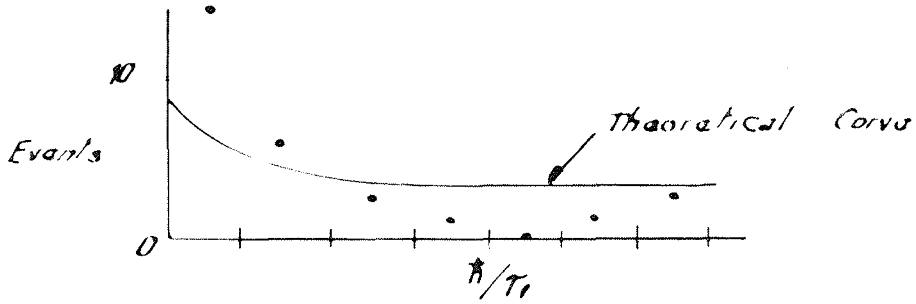
The results of the Wisconsin-Padua-Berkeley group concerning the $\Delta S = \Delta Q$ rule have been presented twice at this meeting already. The presentations have been of a critical nature, designed to present the evidence in a manner which would allow a careful appraisal of these very important results. The importance of this work is such that it would still be useful to those of us, who are, like myself, not experts.

Prof. Lee has explained the important consequences and aesthetic desirability of the relation. This rule is usually stated as a requirement that, in the strangeness changing leptonic decays of mesons and baryons to other mesons and baryons the change in strangeness of the meson or baryon must be equal to the change in charge. Reactions forbidden by this relation are :

$$\Sigma^{+} \longrightarrow n + e^{+} + \nu, \quad K^{\circ} \longrightarrow \pi^{+} + e^{-} + \bar{\nu}, \quad \bar{K}_{0} \longrightarrow \pi^{-} + e^{+} + \bar{\nu}.$$

In each of these cases $\Delta S/\Delta Q = -1$. Though the β-decay of the Σ^{+} has not been observed and is known to be, at least, rare, all hyperon β-decays are uncommon and it has not been established that this decay is unusally rare. The best experimental access to this hypothesis is then by means of K-zero decays. If the rule is correct, we should expect the number of e^{+} decays measured as a function of time to follow the intensity of K° and the number of e^{-} decays to be proportional to the intensity of $\bar{K}_0$ as shown in figure c. Furthermore, the number of decays in a specific time interval should be proportional to the number of hyperon producing neutral K-meson interactions even as these also should be proportional to the intensity of $\bar{K}^{\circ}$ present. Further, the total number of electron decays of both sign should be proportional to the total intensity of K-zero present, a quantity which will fall to 1/2 the initial intensity with the decay of the K_1°-mesons. The experimental data is in accord with none of these relations ; most surprising it is not even in agreement with the last relation concerning the variation of the total intensity with time.

The most serious contradiction to the expected e^{-}/e^{+} ratio is provided by the existence of 3 e^{-} out of a total of but 36 events, in the interval of $0.\hbar/\tau_1 - 2\hbar/\tau_1$, a highly improbable result if the $\Delta S = \Delta Q$ rule is valid. In a first interval of time there are 5 e to 3 $\bar{K}_0$ interactions, in a later interval there are 3 e to 28 interactions illustrating the strong violation of the expected constancy of this ratio. The following figure illustrates the variation of total intensity.



Quite preliminary data from Berkeley concerning similar measurements in the 72" hydrogen chamber where K°-mesons are produced primarily by the reaction $\pi^- + p \longrightarrow K^{\circ} + \Lambda^{\circ}$, is, on the basis of but 10 events, in better agreement with theory on the charge ratio, but without what may be important corrections, shows the same puzzling total intensity variation.

Two results have been presented concerning K_2° decay branching ratios, a measurement by the Ecole Polytechnique and by a Brookhaven group. The Ecole Polytechnique group used an expansion chamber with metal plates parallel to the beam in a field of 8,500 gauss. The chamber was placed 15 meters from a target situated in a curved section of Saturne at Saclay at angle of 65° with the beam. Of 455 V° seen decaying in the chamber 329 were analyzable K_2° decays.

The Brookhaven group used a 50 cm long liquid hydrogen bubble chamber to examine the decays and interactions of K_2° mesons produced by π-mesons striking a polyethylene target. The angle of detection was chosen so as to obtain rather slow K_2°-mesons, from 200 to 400 Mev/c.

The branching ratios of these decays is of considerable interest in as much as the ratios are almost completely predictable from the K^+ branching ratio, if the $\Delta I = 1/2$ rule for non leptonic decays, and the $\Delta S = \Delta Q$ rule for leptonic decays, is valid. Conversely agreement between the observed branching ratios and the predicted ones lend support to these hypotheses.

Below is a table presenting calculated and observed branching ratios. The calculated ratios are taken essentially from calculations of the Brookhaven group.

	Theoretical	Brookhaven	Ecole Poly.
π	32 ± 4	34 ± 4	26 ± 8
$\pi\mu$	33 ± 4	35 ± 6	37 ± 8
$\pi^+\pi^-\pi^{\circ}$	12 ± 1	9 ± 2	14.5 ± 2
$\pi^{\circ}\pi^{\circ}\pi^{\circ}$	22 %	Not measured	Not measured

The agreement between the two experiments, and with the theory is gratifying, and provides some support for $\Delta S = \Delta Q$, though, of course the agreement could be accidental. Furthermore, the ratio of decays in the BNL experiment is found to be $.91 \pm .18$ in agreement with the theoretical value of exactly 1.0 as the K_2° is composed of equal parts K° and $\bar{K}^{\circ}$.

The momentum distributions and angular correlations were also determined by the Brookhaven group. These distributions depend upon the form of the weak interaction coupling and the structure of the π^- meson, or the pi form factor. The results clearly excluded a tensor coupling and with a reasonable pion form factor favored the vector coupling as is to be expected from the V-A description of weak interactions.

It has been shown experimentally in the last few years that the decays of the lambda and sigma do not conserve parity. Originally this conclusion was reached by observing the up-down asymmetry of the hyperon decay products with respect to a plane of production. In particular, the reactions $\pi + p \longrightarrow \left\{ \begin{matrix} \Lambda^{\circ} \\ \Sigma^+ \end{matrix} \right. + K$ result in such an effect. We may say that we observe a pseudoscalar $(\vec{p}_{\pi} \times \vec{p}_{\Lambda}) . \vec{p}_p$.

Two interactions are involved, a production reaction conserving parity, resulting in a polarization p of the spin parallel to the direction $(\bar{p}_{\pi} \times \bar{p}_{\Lambda})$ and a non parity conserving decay interaction, resulting in the decay of the lambda such that the proton is emitted preferentially with respect to the spin direction proportional to $1 - \alpha \cos \vartheta$, where ϑ is the angle between the direction of polarization p and the direction of decay. Experimentally one measures the product αp ; since we do not know enough of elementary particle dynamics to know the sign or value of p, we learn only a lower limit to the value of α ; since $|p| \leq 1$, $|\alpha| \geq |\alpha p|$.

In this non parity conserving decay, the nucleon will, in general, be polarized in the direction of its emission. The pseudoscalar $(\vec{p}_p \cdot \vec{\sigma}_p)$ will be observable. The value of this polarization cherality is just equal to $-\alpha$ where $\alpha = \frac{S \cdot P}{S^2 + P^2} \cos \delta$ where S and P represent the magnitude of the S and P amplitude and δ represents the difference in phase between these amplitudes. As we believe the interaction is invariant under time reversal, this δ is just the difference between the two π-nuclear phase shifts at the same energy, which we know to be negligible i.e. $\cos \delta \approx 1$. In general the measurements of such longitudinal polarizations are difficult as, for example, the angular distributions of scattering or reactions are independent of longitudinal polarization, and can only be dependent on the polarization, P, perpendicular to the beam. Furthermore, the analyzing reaction, as in optics, must be a polarizing reaction. One detects a polarization of the incident particle by noting a left-right asymmetry at an angle. This ratio $\frac{N(L) - N(R)}{N(L) + N(R)} = P.A$, but the analyzing power A is just equal to the polarization produced by scattering of a polarized particle.

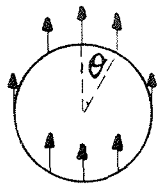
The translation of useless longitudinal polarization to useful transverse polarization is produced simply by the change from the C.M. system to the laboratory system, as below.

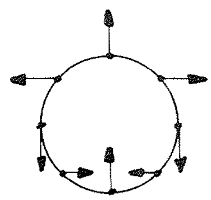
We emphasize that all of this refers to unpolarized Λ^0 though the results are independent of the polarization. Favorite analyzers are carbon and metal plates which are known to be good proton polarizers at small angles and laboratory proton energies of 100 Mev to 300 Mev. Fortunately protons emitted from Λ and Σ hyperons from the interactions $\pi + p \longrightarrow Y + K$ at π energies near a Bev characteristically have momenta in this range.

The next figure shows the experimental set up used by a Berkeley group to measure α for the decay of Σ^+ and for the decay of Λ^0. The spark chamber is fitted with carbon plates. The left-right asymmetry of the scattering of decay protons in the plates establishes the value of $\bar{\alpha}$. The spark chamber is triggered by the detection of K^+ mesons stopped in a water Cerenkov counter. These results are $\alpha(\Sigma_+) = + .75 \pm .17$. Previously there had been two measurements of the α for Λ^0 decays. Assuming that $|\alpha| \geq 0.7$ from measurements of αP, a very early experiment by an MIT group using interactions in a large multiplate expansion chamber found a positive α with a probability of about 12 : 1, a later measurement by a Berkeley group, measuring scatterings in a propane chamber found that α is negative with a probability of about 1/100 to one. The Berkeley measurement just discussed also finds the value of α is negative. Recently two quite accurate results on α, one by a Brookhaven group, again using the spark chamber, finds a value of α of $- .64 \pm .24$ which excludes positive value of 0.7 by a factor of about 1/4 000. A Syracuse-Duke-Johns Hopkins group has looked at the scattering of protons from lambda decays in a helium bubble chamber. The reaction chain here is $K^- + He \longrightarrow \Lambda^0 + ---$, $\Lambda^0 \longrightarrow P + \pi^-$, $P + He^4 \longrightarrow P + He^4$. The α particle is known to be a good analyzer of proton polarization. This group has results, again of great statistical weight, in agreement with a negative value of α for the lambda decay.

There has been considerable theoretical interest in these values, particularly in the signs of α. Remembering that α is the negative chirality, $= -(\sigma_p \cdot p_p)$. The standard (V-A) weak interaction used with global symmetry predicts that α_Λ and α_Σ are opposite, in agreement with these results. A very simple model based on the idea of the Σ and Λ having opposite parity the Σ being essentially a bound Λ and S-wave π predicts these two have the same sign. This result is now contradicted by these measurements.

The part of this experiment which concerns the lambda polarization provides another quantity of great interest. The measurement of α does not tell us whether the S-wave or P-wave intensity dominates in the decay. If the value of α is -0.7, from our relation $\alpha = 2S.P/(S^2 + P^2)$, we know that $|S|^2/|P|^2 \approx 10/1$ or $1/10$. One can determine which of these two possibilities is true by measuring the correlations between the polarization of the lambda and the proton. Since we know that the lambdas are highly polarized at most angles and production energies in the reaction $\pi + p \longrightarrow \Lambda^\circ + K^\circ$, such a measurement is feasible. Consider lambdas polarized upwards, which decay via pure S-wave or pure P 1/2-wave ; the final state is then $Y_0^\circ \uparrow$ or $\sqrt{1/3}\ Y_1^\circ\uparrow + \sqrt{2/3}\ Y_1^1\downarrow$ which gave the following polarizations vs. angle.





s-wave   p-wave

While the true situation, being a mixture of these, is more complicated, the general features are illustrated thus. Then a measurement of the in-out polarization of protons emitted by polarized Λ° will determine the relative S and P wave contributions. The results of this Berkeley experiment show that the decay intensity is predominantly S-wave, with a probability of about 20 : 1. This is quite important, as through a chain of reasoning due to Dalitz, this suggests strongly that the Λ-K parity is odd. Let me review this quite briefly.

The binding energy of a Λ in light nuclei is to a first approximation similar to the binding of a particle in a square well with Radius $A^{1/3}$ and a depth of something like 10 Mev. There are deviations from this curve, in particular, a Λ is bound to He^3 more strongly, relatively, than to He^4. This indicates that the Λ°-nucleon interaction is spin dependent ; He^4 is symmetric so there can be no preferred Λ-N alignment while there can be such in the binding of a Λ to He^3. But we do not know whether the Λ is preferentially aligned or antialigned.

eg.	we may have	$\uparrow n \uparrow p$	$\uparrow n \uparrow p$
		$\downarrow p \downarrow \Lambda$	$\uparrow p \uparrow \Lambda$

If the Λ is antialigned the proton will be left in a symmetric state and the first reaction will dominate. If the Λ is aligned the proton will be aligned. The Pauli principle will not allow it to be found in a low state and it will escape. The second reaction will dominate. Since experimentally the first reaction is more probable, we know the Λ is preferentially antialigned and then that the ${}_\Lambda He^4$ ground state spin zero.

The existence of the reaction :

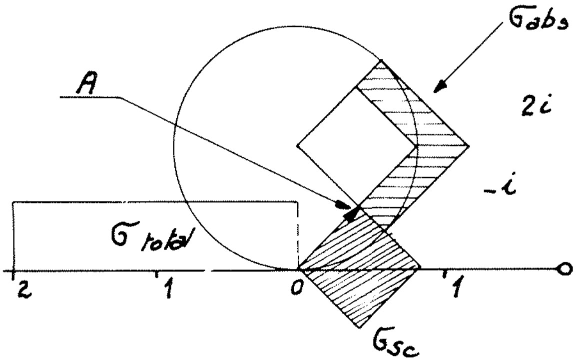
$K^- + He^4 \longrightarrow {}_\Lambda He^4 + \pi^-$ which has been observed can only conserve both angular momentum and parity if the K-Λ parity is odd. More properly this is the (K, n, Λ) parity but by convention we choose n and Λ parity as even. The reliability of such a conclusion depends on careful analysis of these statements, of course. I believe Dalitz will say more about this.

Most of the material concerning the strong interactions of strange particles has concerned resonances, a subject which has been reserved for this afternoon's session. There is, however, an interesting result which does not seem to show resonances, that is a measurement of the total cross section for the interaction of K -mesons with protons and neutrons ; a measurement performed by a Berkeley group. The results are presented in the proceedings, the truly anomalous character of these results is immediately evident ; there are no resonances.

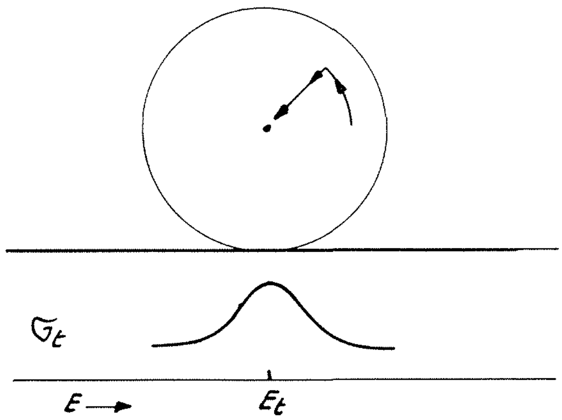
There is, however, some evidence of structure, perhaps a discontinuity in slope of the cross section curve near 1 100 Mev/c. This is of particular interest now since Ball and Frazer have suggested that the resonance like peak in the K^--p cross section near 1 100 Mev/c results, not from an ordinary resonance, but from the existence of the threshold for production of the K' at this momentum. If this were indeed the case we might suspect structure in the ($K^+ p$) cross section at the same energy as this is the threshold for the strangeness + 1 K' which should exist.

I have had little chance to see the paper of Ball and Frazer and I am not sure I could understand it completely if had the opportunity to study it, however, some features of this problem may be understood in a simple form.

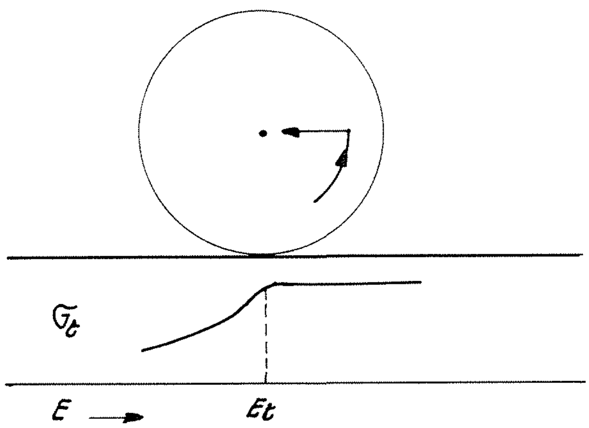
a)



b)



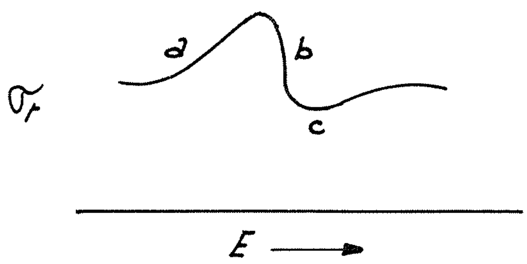
c)



At the threshold for a new process, or a new channel, other processes are disturbed or effected and this effect can be calculated to a first approximation using solely the unitarity and analiticity of the S-matrix. The behavior of the scattering amplitude is easily calculated and can be shown most simply graphically. Figure a shows the demensionless scattering amplitude $A = i(1-S)$ plotted in the complex plane. The shaded areas represent various cross sections in units of $\pi\lambda\!\!\!^{-2}$ times appropriate statistical factors.

On such a diagram, at the S-wave threshold for a new channel, the scattering amplitude makes a left hand turn. Figures b) and c) illustrate two interesting conditions. If the new state has a finite width the turn will be rounded and the width the turn will be rounded and the cross sections might be about as illustrated. Note that with approximately the same change in absorption cross section (or increase in K' production) the shapes simulate approximately the (K^- p) resonance, and the (K^+, p) structure. Incidentally, size of the discontinuities are such that in order to explain them, the K'-S-wave production must quickly approach unitary, and the K' spin must be at least one.

This suggestion, that the structure in these cross section is related, and the result of the threshold, is interesting and plausible and must be investigated further. However, I am disinclined to believe that this is actually the case. The K^--p resonance shows a very particular slope sketched here : In particular the slow rise (a) the sharp fall (b) and the dip (c) seem to be definitely indicated by the data. All of this fits quite precisely a simple Breit-Wigner resonance representation, where a background phase of about 60° exists. Neglecting absorption the relation $\delta = \tan^{-1} \dfrac{\Gamma/2}{E_\lambda - E} + 60°$ fits the shape rather precisely. While such a shape is not derivable simply from the threshold theory, it is probable that certain accidents of variation could simulate it. However, while a beast which looks like a cow might be a malformed horse, there is much to be said for assuming that it is a cow.



Aside from anomalies, this general program of precise measurements of the K^+-p cross sections together with the angular distributions also measured by this group are essential in efforts to learn parities of the Λ and Σ and the coupling constants relevant to the virtual processes $\Lambda \longrightarrow K + n$, $\Sigma \longrightarrow K + n$, by using the zero momentum transfer, or forward direction dispersion relations.

These have a general form such as :

$$\mathrm{Re}\, A(\omega) = \frac{1}{\pi}\, P \int_{-\infty}^{+\infty} \frac{\mathrm{Im}\, A(\omega')}{\omega - \omega'}\, d\omega'$$

where ω is the total energy.

If we consider this as the relation for the scattering of K -mesons, the negative energy part of the integral is related to K -p scattering, however there is a non physical part $-M_K \leq \omega \leq M_K$ which is not accessible but contains contributions from singularities corresponding to the virtual transitions $K + n \longrightarrow \Lambda$, and $K + n \longrightarrow \Sigma$, as well as contributions from the virtual transitions $K + n \longrightarrow \Lambda + \pi$, $K + n \longrightarrow \Sigma + \pi$. The magnitude and sign of the contributions from the singularities are dependent upon the coupling constants and parities respectively. To the extent that the Y^* dominates the rest of the unphysical region its parity and coupling constant determine the contribution from that part. The imaginary part of the forward scattering amplitude is proportional to the total cross section. Then the differential cross section in the forward direction is proportional to the $(\mathrm{Re}\, A)^2 + (\mathrm{Im}\, A)^2$. Sufficient measurements of these quantities can then help establish the values in the unphysical region.

SOME TOPICS IN STRANGE PARTICLE PHYSICS (*)

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1 - THE PARITIES OF STRANGE PARTICLE STATES -

At the present stage in elementary particle physics, our most important problem is the determination of the intrinsic parameters - spin, parity and isotopic spin - for each of the elementary particle eigenstates, both for the states whose decay is by weak interactions and electromagnetic processes (the so-called "elementary particle" states) as well as for the more transitory resonant states which lie in the continuum and whose decay is through strong interactions. A knowledge of these parameters will be essential in any attempt to recognize relationships between these states which reflect further symmetry principles for the strong interactions. The determination of the coupling strengths for their mutual interactions will also be important, for their values may suggest such relationships or may provide direct and quantitative tests for such symmetry principles. For most of these states, the isotopic spin is known from their multiplicity. From studies of their decay processes, the Λ and Σ hyperons have been shown to have spin $\frac{1}{2}$, and the K-meson has been shown to be spinless. The determination of their parities has provided a more difficult problem, mainly because of the failure of parity conservation in the weak decay processes. The determination of parities therefore depends on the analysis of strong and electromagnetic processes, since these interactions conserve parity. Owing to the strangeness selection rule $\Delta S = 0$ for these interactions, an absolute parity assignment is not possible for states of odd strangeness : the convention we shall follow is to assign even parity to the Λ hyperon, so that the K-parity or the Σ-parity (or the K^*-parity, or the Y_1^* parity, etc.) will be specified relative to that of the Λ hyperon. The parity of the Ξ hyperon (or of the Ξ^* state, if such exists) does not depend on this convention, of course. In this section, then, we shall discuss the status of the experimental indications on the K-parity and Σ-parity ; some discussion of the Y_o^* and Y_1^* states will be given in Section 3 below.

K-meson. The outstanding evidence on the K-parity comes from the existence of the reactions,

$$K^- + He^4 \longrightarrow {}_\Lambda He^4 + \pi^-, \tag{1.1a}$$

$$\qquad\qquad\quad {}_\Lambda H^4 + \pi^\circ, \tag{1.1b}$$

observed in a helium bubble chamber by Block et al. [1]. The rate observed for these reactions is quite high, amounting to 3 % per K^- stop (the ${}_\Lambda He^4/{}_\Lambda H^4$ ratio being 2, as required by charge independence). The conclusion from (1.1) that the K-parity is odd depends on the fact that $J = 0$ holds for the $({}_\Lambda He^4, {}_\Lambda H^4)$ doublet, since these reactions are then strictly forbidden for even K-parity. We emphasize here that this is a strict selection rule and that this conclusion does not depend upon the validity of the conclusion by Day [2] that K^- capture in liquid helium takes place predominantly from s-orbitals.

Our belief that $J = 0$ holds for this doublet stems from the high branching ratio observed by Ammar et al. [3] for the two-body decay mode :

$$ {}_\Lambda H^4 \longrightarrow \pi^- + He^4, \tag{1.2}$$

which represents a fraction $R_4 = 0.67^{+0.06}_{-0.05}$ of all ${}_\Lambda H^4$ decay modes leading to π^- emission. In fi-

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(*) The preparation of this report was carried out under the program of the U.S. Atomic Energy Commission at the University of Chicago.

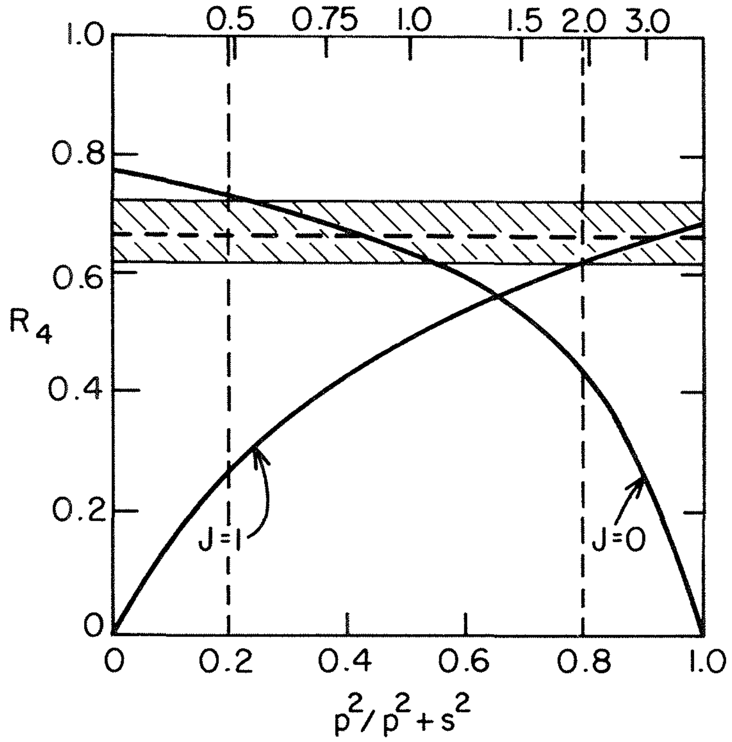

Figure 1 - The fraction R_4 of all ${}_\Lambda H^4$ decays giving a π^- meson, which are in the two-body mode $\pi^- + He^4$, is plotted as function of $p^2/(p^2 + s^2)$ for J = 0 and J = 1, and is compared with the experimental value (shaded region).

gure 1, this fraction R_4 is compared with the values calculated by Dalitz and Liu [4] as function of the relative strengths of the s and p channels of Λ decay, defined such that for free Λ decay :

$$M(\Lambda \longrightarrow p + \pi^-) = s + p\, \sigma \cdot q/q_\Lambda. \qquad (1.3)$$

Here q denotes the pion momentum, with q_Λ its magnitude for free Λ decay. Time-reversal invariance requires that the coefficients s and p should be essentially real ; s and p are now known to have opposite signs [5]. Until recently, our only knowledge of their relative magnitudes came from the value [6],

$$-\alpha_\Lambda = -2ps/(p^2 + s^2) \geq 0.78 \pm 0.08\,, \qquad (1.4)$$

for the asymmetry coefficient in Λ decay ; this required that $p^2/(p^2 + s^2)$ lie between the limits 0.2 and 0.8. From a measurement of the polarization of protons resulting from polarized Λ decay, Beall et al. [7] have now obtained a preliminary value $p^2/(p^2 + s^2) = 0.17^{+0.16}_{-0.07}$, in general accord with expectation from the Karplus-Ruderman argument [8] on the rate of non-mesic hypernuclear decay. The evaluation of the error on this value is not yet complete ; the error may increase somewhat but the mean value is not expected to change appreciably. From figure 1, this value clearly requires J = 0 for ${}_\Lambda H^4$; qualitatively, this conclusion simply depends on the fact that, with J = 1, the pion must be emitted into the p-wave in the two-body mode (1.2) and that, with a decay interaction (1.3) giving predominantly s-wave pions, the possibility J = 1 would therefore allow the two-body mode only with low frequency, contrary to observation. It then follows that the K-meson is pseudoscalar.

This conclusion could be avoided only through one remote possibility. If the K-meson happens to be scalar, so that reaction (1.1) is forbidden, ${}_\Lambda H^4$ and ${}_\Lambda He^4$ decay events could still be observed following K^--He^4 capture if there existed a bound excited state (${}_\Lambda H^{4*}$, ${}_\Lambda He^{4*}$) with J = 1 and if all the observed events resulted from the sequence :

$$K^- + He^4 \longrightarrow {}_\Lambda He^{4*} + \pi^-, \qquad {}_\Lambda He^{4*} \longrightarrow {}_\Lambda He^4 + \gamma\,. \qquad (1.5)$$

At the moment, it is not clear whether there should exist an excited state for this hypernucleus. The rather crude calculations of Dalitz and Downs [9] on this system suggest that there should exist an excited state with Λ binding energy B_Λ^* of about 0.5 Mev (i.e. with excitation energy about 1.7 Mev), whereas the more elaborate calculations of Dietrich et al. [10] on ${}_\Lambda He^5$ and ${}_\Lambda He^4$ indicate that no excited state is to be expected (*). Further calculations on this question are now very desirable (**). On the other hand, it would seem feasible to attempt a direct experimental test of the possibility (1.5), since it would require that a γ-ray of between 1 and 2 Mev be emitted with every hypernucleus formed, that is for 3 % of K^- stops in helium : further, in 2 % of K^- stops, this γ ray would be emitted in coincidence with a fast π^- meson from the first step in this sequence (***). Finally, the rate observed for the reactions (1.1) is already about twice that estimated on the basis of the impulse approximation [9]. For $B_\Lambda^* = 0.5$ Mev, the estimated rate would be about 0.7 %, so that the high rate observed provides a weak argument against the possibility that all the observed reactions follows the sequence (1.5).

The value of the coupling constant $G_{K\Lambda}$ is not yet known. The most hopeful method for its determination is from the observation of the photo-electric term in the process :

$$\gamma + p \longrightarrow \Lambda + K^+, \tag{1.6}$$

by an extrapolation to the K^+ pole in the momentum transfer or angular distribution. As pointed out by Moravcsik [14], this procedure may also provide a confirmation of the $K\Lambda$ parity. To date, however, the cross sections observed [15] for this reaction are essentially isotropic, and it is not reasonable to attempt such an extrapolation unless the forward peaking characteristic of this term is apparent in the physical data on the angular distribution. Studies of the angular distribution of (1.6) at forward angles at the highest possible photon energies will be of great interest. Another reaction of interest from this point of view is :

$$\pi + N \longrightarrow K^* + \Lambda, \tag{1.7}$$

since the pole term associated with K exchange may be expected to give a forward peaking in the K^* production, which may be identified and interpreted in terms of the K-parity and $G_{K\Lambda}$ (assuming the K^* spin is established and that the width for $K^* \longrightarrow K + \pi$ establishes the $(K^* K \pi)$ coupling constant) - on the other hand it is possible that K^* exchange may also contribute strongly to such a forward peaking and it may not be easy to distinguish between these two contributions.

For the threshold process (1.6), the Kroll-Ruderman theorem [16] is not available to justify interpretation of the s-wave production cross section in terms of $G_{K\Lambda}$. However Mc Daniel et al. [17] have compared their data with perturbation-theory calculations by Capps [18]. This involves uncertain assumptions about hyperon magnetic moments, but the comparison made suggested $G_{K\Lambda}^2/4\pi \approx 2.2$.

Σ-hyperon. The Σ-parity is a crucial question for all symmetry principles proposed for the strange particles. Is the Σ triplet closely related with the Λ singlet, as envisaged by the hypothesis of global symmetry [19] or of the doublet approximation [20] which require even Σ-parity, or do they have opposite parities and belong to different representations of larger symmetry groups ? Should the

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(*) The calculations of Dietrich et al. [10] use an extension of the method developed by Mang and Wild [11] for calculating the binding energies of light nuclei. They use Λ-N and N-N potentials whose form consists of an attractive square well potential outside a repulsive hard core of radius 0.2f. We note that the accuracy of their method is probably least for the case of a lightly-bound Λ particle (as for ${}_\Lambda He^{4*}$ with $B_\Lambda^* = 0$). More extensive calculations on this question, following the same lines with other potential shapes and hard core radii, are very desirable.

(**) Uncertainties in deciding this question because of the possibility of three-body hypernuclear forces have been stressed previously [12]. However there are qualitative theoretical arguments [13] for expecting such three-body forces to be predominantly non-central and, although strong, relatively ineffective in binding of the Λ particle. Also, there are no indications yet that such three-body forces need be invoked for the interpretation of hypernuclear binding energies. Although the possibility of three-body forces should be borne in mind, the most reasonable viewpoint at this stage is to interpret hypernuclear binding energies in terms of predominantly two-body forces.

(***) The formation of the excited state ${}_\Lambda He^{4*}$ is also possible for a pseudoscalar K^- meson, but then only for a small fraction of the reactions leading to a hypernuclear decay. As Day [2] has argued, it is probable that the K^--He^4 capture processes occur predominantly from s-orbitals, and the formation of $J = 1$ ${}_\Lambda He^{4*}$ is then forbidden for a pseudoscalar K-meson. However Day's calculations show that 10 % of K^--He^4 capture events could occur from p-orbitals and the fraction might well be considerably larger than this : a small fraction of these p-orbital captures could lead to ${}_\Lambda He^{4*}$ and to the subsequent emission of a 1-2 Mev γ-ray.

Σ triplet be regarded as a composite system, a bound s-state of the Λ-π system as envisaged by Sakurai and Nambu [21], which requires odd Σ-parity ? At this stage, the resolution of this question is overdue, in view of its great importance for future developments, and merits a massive experimental effort. Let us review briefly the experimental situations which have been considered to bear on this question.

1 - The hyperon-nucleon and pion-hyperon interactions. The data on Λ hypernuclei now suggest [13] that the Λ-N interaction has strong spin-dependence, the ^1S force being quite strong (well-depth parameter ≈ 0.7), the ^3S force being relatively weak and including a substantial fraction of K-exchange terms. This spin dependence follows naturally with even $\Lambda\Sigma$ parity, and the strength of the ^1S force corresponds to $G^2_{\Lambda\Sigma\pi} \approx G^2_{NN\pi}$. The dependence of this comparison on $G_{\Sigma\Sigma\pi}$ appears quite weak and is being investigated further at present. The hypothesis of global symmetry, $G_{\Sigma\Sigma\pi} \approx + G_{NN\pi}$, also accounts well [13] for the Σ-N interactions observed for Σ^- hyperons coming to rest in hydrogen and for the final Σ-N states following the K^--d capture reactions.

With odd Σ parity, this spin-dependence of the Λ-N potential does not arise naturally and must be attributed to $K\Lambda$ and $K\Sigma$ couplings of suitably chosen strength [22].

As pointed out by Gell-Mann [19], global symmetry would require the existence of two $j = \frac{3}{2}$ π-hyperon isobars with I = 1 and I = 2, direct analogues to the (3.3) π-N isobar. This interpretation has been proposed by Amati et al [23] for the Y_1^* resonance observed in the π-Λ system, and these authors point out that the existence of these isobars does not depend critically on an exact global symmetry. In view of the above remarks, with even Σ parity and a strong coupling $\pi + \Lambda \longrightarrow \Sigma$, this interpretation of the Y_1^* resonance is rather natural (especially as there is now no other I = 1 π-Λ resonance which could be identified naturally with this isobar state) and does not yet disagree with any of the experimental data (cf. Section 3). Indeed, as pointed out by Wentzel [24] using strong coupling theory and by Franklin [25] using the one-meson approximation, the situation with $G_{\Sigma\Sigma\pi} \approx 0$ allows three $j = \frac{3}{2}$ pion-hyperon isobars, an I = 1 isobar lying lowest and the I = 0 and I = 2 isobars lying close together at a higher total mass. This picture may well be closely parallel to the physical situation. In addition to the I = 1 Y_1^* resonance at 1 385 Mev, there has recently been established by Ferro-Luzzi et al. [26] the existence of an I = 0, $j = \frac{3}{2}$ resonance at 1 525 Mev. Further, there are indications [27] of a π-Σ resonance with I = 1 or 2 at about 1 580 Mev ; since there is no evidence for an I = 1 resonance in the K^--p system in this region, this resonance (if confirmed) might well represent an I = 2 isobar. On the other hand, it is not yet clear whether the Σ-N interactions could be well accounted for with a small value of $G_{\Sigma\Sigma\pi}$; the extent to which these two situations may be compatible is at present under investigation. To sum up, the existence of a $j = \frac{3}{2}$, I = 1, π-Λ isobar resonance appears a rather natural consequence of even Σ parity and a large $G_{\Lambda\Sigma\pi}$, so that the determination of the Y_1^* spin will be quite crucial on this point. The further resonances mentioned above may also be very relevant ; the confirmation and determination of the I-spin of the 1 580 Mev resonance are particularly urgent.

2 - The photoproduction process :

$$\gamma + p \longrightarrow \Sigma^0 + K^+ \tag{1.7}$$

has been studied at 1 140 Mev (100 Mev above threshold) at Cornell by Edwards et al. [15,28]. The angular distribution consists of three points, compatible with an isotropic distribution [28], so that no conclusion can be drawn from these data at present concerning the Σ parity. At such low energies, where the photoelectric term does not have a major influence on the angular distribution, an extrapolation to the K^+ pole is not at all meaningful, but at much higher energies, a knowledge of the angular distribution will be very useful in this respect.

3 - Following the suggestion of Adair [29] and of Baz and Okun [30], there has been some hope of determining the Σ parity from a study of the cusp behavior in the reaction :

$$\pi^- + p \longrightarrow \Lambda + K^0, \tag{1.8}$$

expected to occur at the threshold for the competing reactions to Σ + K states. At this threshold, strong s-wave production of the Σ + K system has been observed by Wolf et al. [31], and a determination whether this cusp occurs in the $s_{\frac{1}{2}}$ or the $p_{\frac{1}{2}}$ channel of the $\Lambda + K^0$ system would establish the Σ parity as even or odd, respectively. Nauenberg and Pais [32] have pointed out, however, that there is a Minami ambiguity in the analysis of $d\sigma/d\Omega$ and $P(\vartheta)$ for reaction (1.8), so that for every solution for the partial wave reaction amplitudes for (1.8) which assigns the $\Lambda + K^\circ$ cusp state to $s_{\frac{1}{2}}$ there exists a corresponding solution which assigns the $\Lambda + K^\circ$ cusp to the $p_{\frac{1}{2}}$ state. It is conceivable that these two possibilities could be distinguished by a study of $d\sigma/d\Omega$ and $P(\vartheta)$ for (1.8) as function of π^- energy down to the $\Lambda + K^\circ$ threshold, appealing to the continuity of the reaction amplitudes as function of energy in this region and distinguishing the $s_{\frac{1}{2}}$ and $p_{\frac{1}{2}}$ states of the $\Lambda + K^\circ$ system by the requirement that the $p_{\frac{1}{2}}$ amplitude vanish linearly with momentum $p_{\Lambda K}$ at this threshold, while the $s_{\frac{1}{2}}$ amplitude remains finite at threshold. Eisler et al. [33] have reported at this meeting a study of the reaction (1.8) below the $\Sigma + K$ threshold, which shows that the amplitude for $s_{\frac{1}{2}}$ production is rather small at threshold and that p- and d-wave production becomes of importance already quite close to the threshold. As a result, the possibility of resolving the Minami ambiguity in this way appears an exceedingly difficult proposition, and a determination of the Σ parity in this way has become much less hopeful. Also, the observation of the cusp effect has also proved quite difficult. Although both the Columbia and Berkeley groups [31, 33] agree that the angular distribution $d\sigma/d\Omega$ for reaction (1.8) appears to vary with unusual rapidity as the π^- incident energy varies from below to above the $\Sigma + K$ threshold, the nature of this change is not yet clearly established, and it is not at all clear yet just which spherical harmonics in $d\sigma/d\Omega$ and $P(\vartheta)$ show this cusp effect most strongly.

4 - The angular distributions for hyperon production in pion-nucleon collisions show a striking difference between the cases of Λ and of Σ hyperons. The results reported at this conference by Alles-Borelli et al. [34] at 1.6 Gev/c and by Erwin et al. [35] at 1.9 Gev/c illustrate this difference : the Λ production in the reaction (1.8) is very strongly peaked toward backward angles (relative to the incident π^- direction), whereas the Σ^- production is just as strongly peaked forward in the reaction :

$$\pi^- + p \longrightarrow \Sigma^- + K^+. \tag{1.9}$$

In the range 1100 to 1400 Mev/c, the cross section $\sigma(\pi^+ + p \longrightarrow \Sigma^+ + K^+)$ also shows [36, 37] some forward peaking for the Σ^+ hyperons. On the other hand, Crawford et al. [38] have found $\sigma(\pi^- + p \longrightarrow \Sigma^\circ + K^\circ)$ to be roughly symmetrical about 90° at 1.2 Gev/c, with perhaps some peaking in the forward and backward directions.

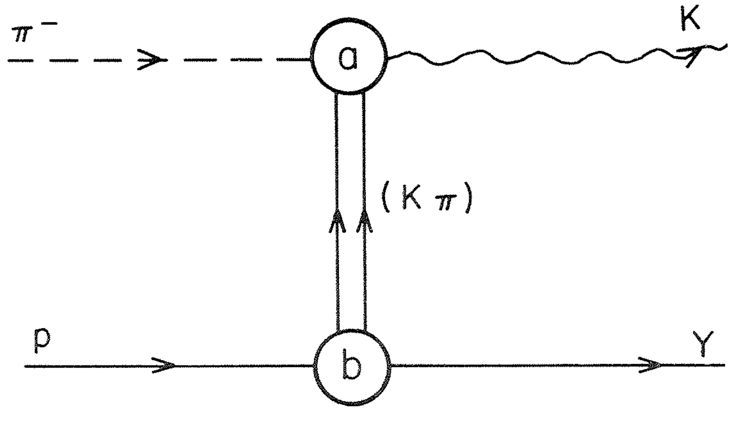

Figure 2 - Peripheral graph for the process $\pi^- + p \longrightarrow K + Y$.

It has frequently been suggested that this difference may reflect opposite parity for Λ and Σ hyperon. Tiomno et al. [39] have pointed out that the backward Λ peaking may be due to the exchange of a K-π resonance state, as shown in figure 2, which leads to the denominator factor :

$$\{(E_Y - E_N)^2 - (\underline{p}_Y - \underline{p}_N)^2 - m^2_{K\pi}\}^2 \tag{1.10}$$

favoring backward hyperon production. It is necessary to consider also the vertex (b) for $N + (K\pi) \longrightarrow Y$. If the $(K\pi)$ resonant state has zero spin, this vertex is essentially a constant if the KY parity is odd, whereas for even KY parity the vertex has (non-relativistically for the baryons) the form $\sigma.(\underline{p}_Y - \underline{p}_N)$ which contributes to the cross section the factor $(\underline{p}_Y - \underline{p}_N)^2$ which favors forward production for the hyperon. Thus, with odd KΛ parity and even $\bar{K}\Sigma$ parity, this mechanism leads directly to backward Λ production and forward Σ production. Beg et al. [40] have pointed out that at sufficiently high energies (for which the relativistic form of the vertex (b) must be used) this mechanism then leads again to predominantly backward Σ production. Studies of the energy dependence of $d\sigma/d\Omega$ for all these hyperon production reactions for still higher energies would be of much interest.

However it seems doubtful whether the hypothesis of odd parity is itself sufficient to account for the behavior of these distributions, without quite a number of further special assumptions. First, it appears probable that the K-π resonance may have spin 1. In this case the vertex (b) has the form $\gamma_\mu(p_K + p_\pi)_\mu \approx (E_K + E_\pi)$ for odd KY parity, and $\gamma_5\gamma_\mu(p_K + p_\pi)_\mu \approx \underline{\sigma}.(\underline{p}_K + \underline{p}_\pi)/2M_B$ for even KY parity ; neither of these factors leads to forward peaking for hyperon production. Secondly, the exchange of a K-π system can contribute to Σ^- production according to figure 2 only if it has $I = \frac{3}{2}$, whereas it can contribute to Λ production only if it has $I = \frac{1}{2}$, so that different isotopic spin states of the K-π system would necessarily be involved in the two processes ; the known resonant state (K^*) has $I = \frac{1}{2}$. Also, Beg et al. [40] point out that the observation of strong hyperon polarization in these reactions shows that the graph of figure 2 can only represent a portion of the production amplitude since alone it would predict zero hyperon polarizations. It is clear that the assumption of opposite parity for Λ and Σ does not lead directly to a simple explanation for the observed behavior of these hyperon production cross sections. The striking difference between Λ and Σ production therefore does not really provide any convincing evidence for odd Σ parity although its interpretation does pose an intriguing and important problem for the strange-particle physicist.

5 - The use of forward-scattering dispersion relations should now be mentioned briefly, as the data on forward and total cross sections for K^+N and K^-N scattering are rapidly becoming more accurate and more complete. As illustrated in figure 3, the K^+p and K^-p data for physical energies is to be used to determine the strengths of four poles occurring for unphysical energy values (as well as some background in the unphysical energy range given by $(M_x^2 - M^2 - m_K^2)/2M$ for $M_\Lambda + m_\pi \leqslant M_x \leqslant M_N + m_K$). These poles correspond to strange particle eigenstates X contributing to the forward scattering amplitude through the mechanism :

$$K^- + p \longrightarrow X \longrightarrow K^- + p \,. \tag{1.11}$$

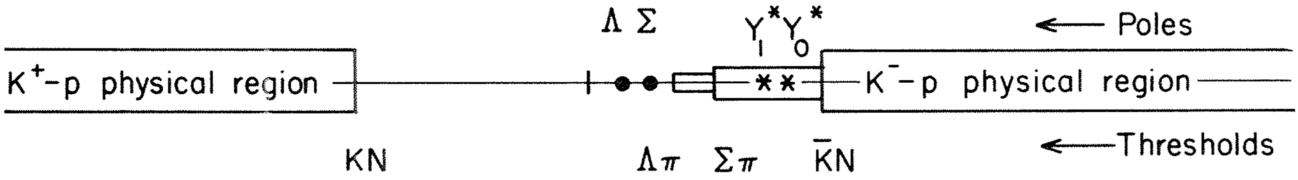

Figure 3 - The poles and threshold branch cuts for the $\bar{K}$-N forward scattering amplitude as function of the $\bar{K}$ laboratory energy E.

These eigenstates are the Λ and Σ hyperons, and the Y_1^* and Y_0^* resonant states. They each contribute a pole term :

$$R_x/(E - (M_x^2 - M^2 - m_K^2)/2M), \tag{1.12}$$

where the magnitude of the residue R_x measures the strength of the vertex $\bar{K} + N \longrightarrow X$ for all three particles on the mass shell (i.e., the coupling constant G_{KNX}^2) and the sign of R_x is positive or negative according as the $\bar{K}N$ orbital angular momentum l in (1.11) is even or odd. For the Λ pole, odd K-parity requires l = 1 and a negative residue R_Λ. For the Y_0^* and Y_1^* states, the sign of this residue will depend on their nature. If they are $j = \frac{1}{2}$ $\bar{K}$-N bound s-states, then l = 0 holds and the residues R_1 and R_0 will be positive : in this case the effective residues may be estimated by an extrapolation from the low-energy $\bar{K}$-N scattering data, based on a reaction-matrix with a suitable energy dependence. If Y_1^* is a $j = \frac{3}{2}$ π-Λ isobar, then odd K-parity requires $l = 1$ and this residue R_1 will be negative. Karplus et al. [41] have pointed out that the present data are best for energies well away from the threshold regions and that, in this case, it is difficult to determine more than an average of these four residues, $\bar{R} = (R_\Lambda + R_\Sigma + R_1 + R_0)$. Investigations by Kerth [42] have shown that this average $\bar{R}$ is probably negative, and at least this is consistent with what now appears the most likely situation (K-parity odd, Σ-parity even, Y_1^* an isobar state and Y_0^* either absent or a $\bar{K}$-N virtual bound state).

6 - The process which appears most directly available for the determination of the Σ parity at present is the Σ° decay mode :

$$\Sigma^\circ \longrightarrow \Lambda + e^+ + e^- , \tag{1.13}$$

which represents internal pair conversion for the usual electromagnetic transition :

$$\Sigma^\circ \longrightarrow \Lambda + \gamma . \tag{1.14}$$

For (1.14), the effective interaction has the form $\lambda \underline{\sigma}\cdot\underline{H}$ for even Σ parity, $\nu\underline{\sigma}\cdot\underline{E}$ for odd Σ parity, where $\underline{E}$, $\underline{H}$ denote the electric and magnetic vectors of the electromagnetic field. These interractions correspond to effective currents $\underline{J} = \langle \Sigma | \underline{j} | \Lambda \rangle$ for this transition, given by :

a) even Σ parity :

$$\underline{J} = \lambda \underline{\sigma} \times \underline{k} , \tag{1.15a}$$

b) odd Σ parity :

$$\underline{J} = \nu \underline{\sigma} . \tag{1.15b}$$

The process (1.13) corresponds to the graph of figure 4. The energy and momentum transferred to the electron-positron pair by the electromagnetic field is given by :

$$k_0 = E_+ + E_- \approx m_\Sigma - m_\Lambda = \Delta , \tag{1.16a}$$

$$\underline{k} = \underline{p}_+ + \underline{p}_- , \tag{1.16b}$$

where m_Σ, m_Λ denote the Σ and Λ mass values. In the expression for the decay probability the only part not calculable from electrodynamics is the square $\underline{J}^2$ of the vertex representing the transition

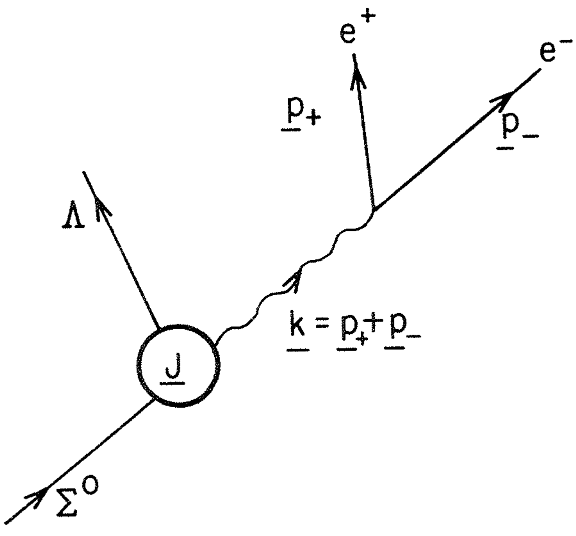

Figure 4 - Feynman graph for the decay process $\Sigma^\circ \longrightarrow \Lambda + e^+ + e^-$.

current : for odd Σ parity, this is essentially a constant, but for even Σ parity, it is proportional to $(\underline{p}_+ + \underline{p}_-)^2/\Delta^2$. The branching ratio for small angle pairs (for which $|\underline{k}| \approx k_0 = \Delta$) is determined from electrodynamics alone, since the value of $\underline{J}^2$ which is then appropriate is just that effective in the normal decay (1.14). The branching ratio for wide angle pairs (for which $|\underline{k}| < k_0 \approx \Delta$) will be markedly less for even Σ parity than for odd, since $\underline{J}^2$ falls rapidly with decreasing $\underline{k}^2$ in the former case. This results in a higher pair conversion ratio for odd Σ parity than for even Σ parity : the calculation has been made by Feinberg [43] and by Feldman and Fulton [44] who find a branching ratio $(\Lambda + e^+ + e^-)/(\Lambda + \gamma)$ of 1/161 for odd Σ parity compared with 1/182 for even Σ parity, a difference of 12 %. However, it is not really the absolute rate, nor the branching ratio, which is of interest to us here for, as remarked above, the branching ratio for the dominant small-angle pairs is determined essentially by electrodynamics. For the Σ parity question, what is of crucial interest is the distribution of $\underline{k}^2$, or more conveniently the distribution of the covariant combination :

$$\begin{aligned} x^2 &= (E_+ + E_-)^2 - (\underline{p}_+ + \underline{p}_-)^2 , \\ &= 4(m^2 + q^2), \end{aligned} \tag{1.16}$$

where q denotes the electron momentum in the electron-positron barycentric frame, for it is this distribution which distinguishes directly between the two possibilities (1.15) and therefore bears most directly on the Σ parity. For example, the graph given by Feinberg shows that the fraction of electron-positron pairs with $x \geqslant 10\ m_e$ (comprising about 40 % of all pairs) is about 20 % larger for odd Σ parity than for even Σ parity. The study of this distribution appears a very promising procedure for the determination of the Σ parity since a high yield of (unpolarized) Σ° hyperons may readily be obtained by stopping K^- mesons in a hydrogen chamber.

Several minor complications in the interpretation of such data should be mentioned briefly here. For virtual electromagnetic fields, the general form of the current $\underline{J}$ may be more complicated than (1.15), as follows (•) :

a) even Σ parity :

$$\underline{J} = \lambda(\underline{k}^2)\ \underline{\sigma} \times \underline{k} + L(\underline{k}^2)\,\underline{k} \tag{1.17a}$$

b) odd Σ parity :

$$\underline{J} = \nu(\underline{k}^2)\ \underline{\sigma} + N(\underline{k}^2)\,\underline{k}\ \underline{\sigma}\cdot\underline{k}/M^2 , \tag{1.17b}$$

where $M = \frac{1}{2}(M_\Sigma + M_\Lambda)$. The terms L and N do not contribute to the γ-decay rate, since they vanish for a transverse electromagnetic field (i.e. with $m = \pm 1$ relative to direction $\underline{k}$). However, for the pion decay (1.13), the intermediate electromagnetic field can be longitudinal (i.e. with $m = 0$ relative to $\underline{k}$), and these additional terms will contribute (incoherently) to the rate for this process. The general expression for the distribution of these pairs is [45],

$$R = \frac{\alpha}{4\pi} \int_{2m}^{\Delta} \frac{dx}{x} \int_{-1}^{+1} d(\cos\vartheta) \left(1 - \frac{x^2}{\Delta^2}\left(1 + \frac{\Delta^2}{4M^2}\right) + \frac{x^4}{4M^2\Delta^2}\right)^{\frac{1}{2}} \left(1 - \frac{4m^2}{x^2}\right)^{\frac{1}{2}} \left(1 - \frac{2x^2}{4M^2 + \Delta^2}\right) \times$$
$$\left\{ \left(1 + \frac{4m^2}{x^2} + \left(1 - \frac{4m^2}{x^2}\right)\cos^2\vartheta\right) R_T(x) + \left(\sin^2\vartheta + \frac{4m^2}{x^2}\cos^2\vartheta\right) \frac{2(2M+\Delta)^2 x^2}{(2M\Delta + x^2)^2} R_L(x) \right\} . \tag{1.18}$$

(•) The general form for this current has been given in relativistic notation by Feldman and Fulton [50], as follows :

a) $J_\mu^e = f_1(k^2)\,(k^2\gamma_\mu - \Delta k_\mu)/M^2 + f_2(k^2)\,\sigma_{\mu\nu}k_\nu/M$,

b) $J_\mu^\circ = F_1(k^2)\,(k^2\gamma_5\gamma_\mu + (M_\Sigma + M_\Lambda)\gamma_5 k_\mu)/M^2 + F_2(k^2)\,\gamma_5\sigma_{\mu\nu}k_\nu/M$,

where f_1, f_2, F_1 and F_2 are appropriate form factors, and $M = \frac{1}{2}(M_\Sigma + M_\Lambda)$. It may be noted that the terms of f_1 and F_1 are at most of order of magnitude $\Delta/M \approx 1/10$ relative to the terms f_2, F_2. The k^2 dependence of these form factors is expected to be quite unimportant since k^2 varies only between 0 and Δ^2. For example, with global symmetry for case a), $k^2 f_1(k^2)$ and $f_2(k^2)$ would equal the electric and the magnetic form factors of the neutron, respectively ; in this case $f_1(k^2)$ is known to be negligible, and $f_2(k^2)$ varies by only 1 %, over this range of k^2. The spatial parts of J_μ° and J_μ^e reduce to the forms (1.17) in non-covariant notation, for suitable λ, ν, L and N.

Here ϑ denotes the angle between the momentum $\underline{q}$ in the electron-positron barycentric frame and their total momentum $\underline{k}$ in the Σ° rest frame, given in terms of the electron and given positron energies in this latter frame by :

$$E_+ - E_- = |\underline{k}| \left(1 - \frac{4\,m^2}{x^2}\right)^{\frac{1}{2}} \cos\vartheta, \tag{1.19}$$

and $R_T(x)$, $R_T(x)$ specify the squares $\underline{J}_T^2$, $\underline{J}_L^2$ of the transverse and longitudinal currents normalized to the value of $\underline{J}_T^2$ for real photons. It will be noticed in (1.18) that the transverse and longitudinal contributions to this distribution each have characteristic distributions in ϑ by means of which they may be distinguished (•). These additional contributions therefore do not provide any essential difficulty for this experiment although it is desirable to check that they are not significantly large. Actually, as remarked in footnote (page), there is no reason to expect these additional contributions to be appreciable, for their amplitude is of order $\Delta/M \approx 1/10$ relative to the first amplitude ; also, because of the smallness of the momentum transfers $\underline{k}^2$ involved in this process, the possibility of form-factor variations for λ and ν introduces only a negligible uncertainty in the distribution. Although these complications are logically possible, there is no reason to expect them to be of significance here and the first investigations of these processes could well ignore them.

If a source of strongly polarized Σ° particles becomes available, the examination of polarization correlations in the pair decay (1.13) will allow an independent means for determining the Σ parity. This has been investigated by a number of authors [47], who have pointed out that the Λ polarization is related to the incident Σ polarization by the expression :

$$\underline{P}_\Lambda = -\hat{\underline{k}}\,\underline{P}_\Sigma \cdot \hat{\underline{k}} \mp \alpha\{\underline{n}\,\underline{P}_\Sigma \cdot \underline{n} - \hat{\underline{k}}\mathrm{x}\underline{n}\;\underline{P}_\Sigma \cdot (\hat{\underline{k}}\mathrm{x}\underline{n})\}, \tag{1.20}$$

where $\underline{n}$ is the unit vector normal to the plane of the electron-positron pair in the Σ° rest frame. For given x^2 (averaging over $\cos\vartheta$), the coefficient α is given by $\alpha = 2(1 - 4\,m^2/x^2)(1 + 2\,m^2/x^2)$ with the x-distribution obtained from (1.18). Averaging over all pairs, the mean value $\bar{\alpha}$ is 0.43. In (1.20), the sign before α is the sign of the Σ parity. The presence of strong longitudinal currents L or N would tend to depress the average value $\bar{\alpha}$ but only very weakly. More sensitive determinations of the sign may be obtained by including the x-dependence of α in the analysis, but it must then be remembered that the relative contribution of the unknown longitudinal terms will be much larger in the region of large x than is their relative contribution to the mean value $\bar{\alpha}$. Since this experiment offers a particularly clean-cut method for the determination of the Σ parity, it merits a massive effort at present, and it is therefore of great importance to determine circumstances in which Σ° particles are produced with a high degree of polarization ; perhaps counter methods can be devised for carrying out such a search quickly, although it seems clear that the polarization correlation experiment itself will necessarily require the bubble chamber technique.

7 - Since recent developments make it appear probable that targets containing polarized protons may become available before many more years, it is worth mentioning here the importance this would have for strange particle physics. This would allow experiments which would bear very directly on the question of parities for strange particle states, as pointed out by Bilenky and Ryndin [48] and others. For example, consider the reaction :

$$\pi^+ + p \longrightarrow \Sigma^+ + K^+ \tag{1.21}$$

on protons with polarization $\underline{P}_p$, for Σ^+ particles produced at 0°. In this configuration, the polarization $\underline{P}_\Sigma$ is given by :

a) $K\Sigma$ parity odd :

$$\underline{P}_\Sigma = \underline{P}_p\,, \tag{1.22a}$$

b) $K\Sigma$ parity even :

$$\underline{P}_\Sigma = 2\hat{\underline{k}}\;\hat{\underline{k}}\cdot\underline{P}_p - \underline{P}_p\,, \tag{1.22b}$$

where k is the incident pion momentum. In particular, for a transversely polarized target, the relative sign of P_Σ and P_p is opposite to that of the $K\Sigma$ parity.

(•) For the process $\pi^- + p \longrightarrow n + e^+ + e^-$, the identification of the longitudinal contributions has been achieved in this way by Kobrak [52]. We note here also that, for odd Σ parity, the current $\nu\underline{\sigma}$ already gives rise to a longitudinal contribution in (1.18), with $R_T = 0.5$.

Ξ hyperons. At present nothing is established concerning the spin and parity of the Ξ hyperon. Their determination is obviously of the greatest importance as concerns their relationship with the other hyperons. Similarly, the possibility of π-Ξ resonant states is also a matter of deep interest and, as pointed out by Lundby [49], may be explored by the method of Dowell et al. [27], by the analysis of K^+ following the reaction :

$$K^- + p \longrightarrow \Xi^* + K^+. \tag{1.23}$$

The Y_0^* and Y_1^* resonant states will be discussed in Section 3.

2 - THE $\bar{K}$-N INTERACTION IN THE LOW ENERGY REGION -

In the past year, the data from the Berkeley hydrogen bubble chamber on K^--p interactions at low energies have been re-evaluated and analyzed by Ross [50] and Humphrey [51]. Ross has obtained cross sections for elastic and charge-exchange scattering, and Humphrey for the reactions leading to Σ^+ and Σ^- hyperons (as well as for Σ^0 and Λ hyperons, although with large uncertainties for these neutral reactions), as function of K^- momentum between 100 and 275 Mev/c laboratory momentum. Ross has obtained an elastic angular distribution at several momenta, for angles including the interesting region of Coulomb-nuclear interference. Humphreys has reevaluated the reaction rates for K^- coming to rest in hydrogen, with the results Σ^- : Σ^+ : $(\Sigma^0 + \Lambda)$ = 1 553 : 722 : 1 200 and $\Lambda/(\Sigma^0 + \Lambda) = 0.186 \pm 0.017$. Much of this data is shown on figures 5 through 8. The main differences between this data and that used in previous analyses [52] are (i) the larger values obtained for the total absorption cross sections leading to $(\Sigma^+ + \Sigma^-)$ hyperons, (ii) the availability of the energy dependence of the elastic cross section, (iii) the inclusion of new Coulomb-nuclear interference data, and (iv) the inclusion of the energy-dependence of the Σ^-/Σ^+ ratio.

The analysis of this data may be considered at various levels of sophistication. The simplest approach is in terms of constant scattering lengths, A = a + ib - "the zero-range approximation" - such that the (complex) $\bar{K}$-N scattering phases are given by :

$$k \cot \delta_I = 1/A_I, \tag{2.1}$$

for the two isotopic spin channels I = 0 and I = 1. In this analysis, it is essential to take into account the deviations from charge independence due to the K-meson and nucleon mass differences, as well as the influence of the K^--p Coulomb interaction. Further constant parameters must be introduced for the description of the reaction processes : ε, which denotes the ratio $\Lambda/(\Sigma + \Lambda)$ for the I = 1 channel, and ϕ, a constant phase angle such that the ratio of the I = 0 and I = 1 amplitudes for the reaction $\bar{K} + N \longrightarrow \Sigma + \pi$ is given by (in the absence of the charge-dependent corrections)

$$M_0(\Sigma)/M_1(\Sigma) = (b_0/b_1(1-\varepsilon))^{\frac{1}{2}} e^{i\phi}. \tag{2.2}$$

This angle ϕ is actually the phase difference between the I = 0 and I = 1 reaction amplitudes, including these charge-dependent corrections, at the charge-exchange threshold. Complete expressions for all the scattering and reaction cross-sections in terms of these parameters are given in reference [53].

These six parameters a_0, b_0, a_1, b_1, ε and ϕ have been determined by Ross and Humphrey by a least squares search procedure starting from the four parameter sets, (a+), (a-), (b+) and (b-), given in reference [52]. This search led to the two parameter sets listed in Table I. Both sets have acceptable χ^2 values, for the number of degrees of freedom was 58. We note that these solutions are not yet closely determined (*), especially not the real parts of A_0 and A_1. In solution I, a_0 has a very large error and even a_1 is poorly determined, not even the sign being determined. In solution II, a_1 is rather well determined but a_0 can lie anywhere between quite wide limits. The imaginary parts b_0 and b_1 have not varied widely from their starting values and are relatively well determined. The fits to the experimental data obtained with the best values of these parameters

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(*) On figures 5 through 8, the fit to the present experimental data given by the (a-) parameter set of reference [52] is also shown. The χ^2 for the fit given by this set, which led to the solution I, was 92, to be compared with χ^2 = 74 obtained for the final solution II. We mention this here to emphasize that the limits of error quoted in Table I should be interpreted rather generously - parameter sets lying outside these limits may still give quite acceptable χ^2 values. When further data is obtained and analyzed, it is not at all improbable that the final parameter sets may be outside these limits.

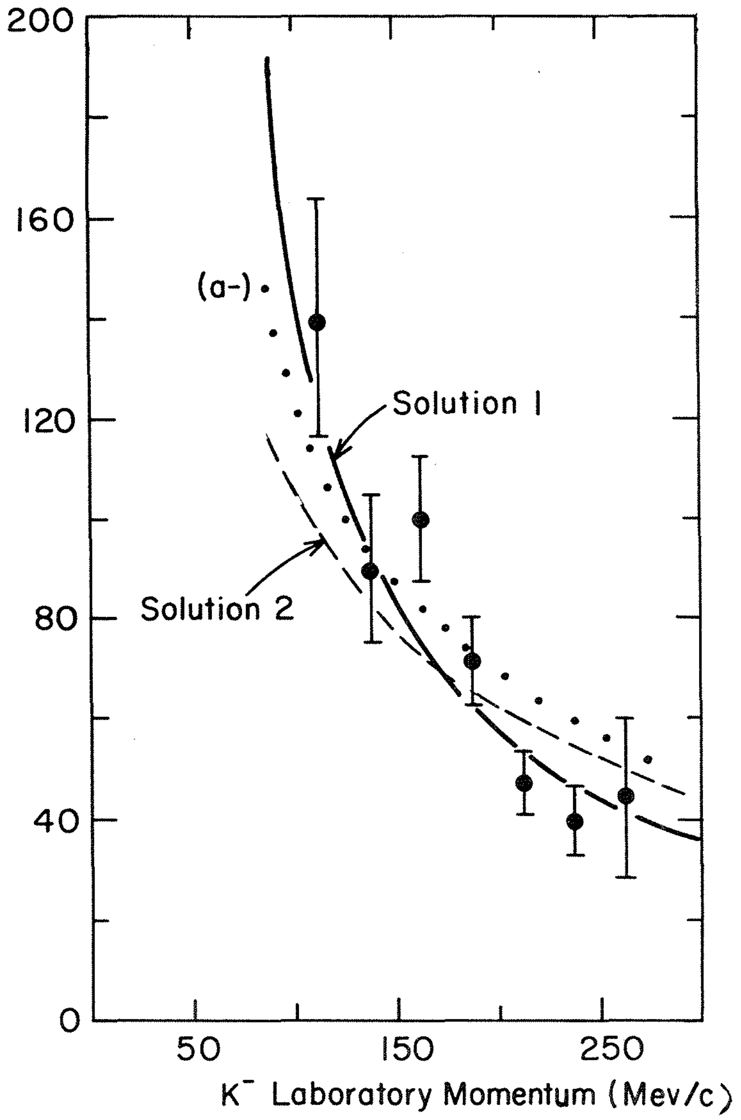

Figure 5 - The elastic cross sections measured by Ross [50] are compared with those obtained in the constant scattering length approximation, for the parameters corresponding to solutions I and II, as well as for the old (a-) parameters.

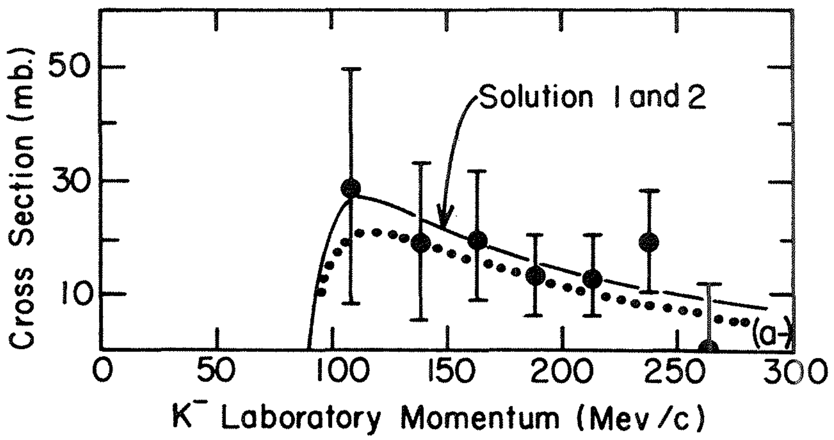

Figure 6 - The charge-exchange cross sections measured by Ross [50] are compared with the curve given by the constant scattering length approximation, for solutions I and II, as well as for the old (a-) parameters.

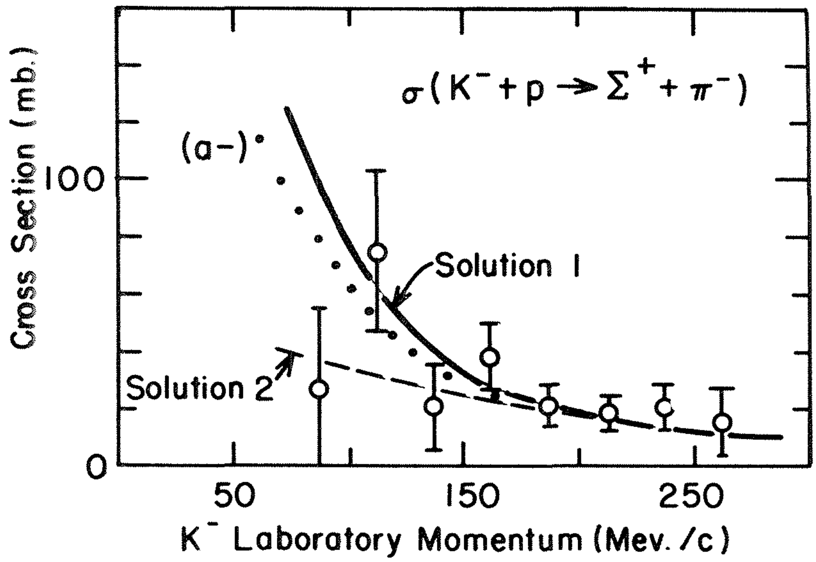

Figure 7 - The energy-dependence of the cross sections measured by Humphrey [51] for the reaction $K^- + p \longrightarrow \Sigma^+ + \pi^-$ are compared with that given by the constant scattering length approximation with the parameters of solutions I and II, as well as with the (a-) parameters.

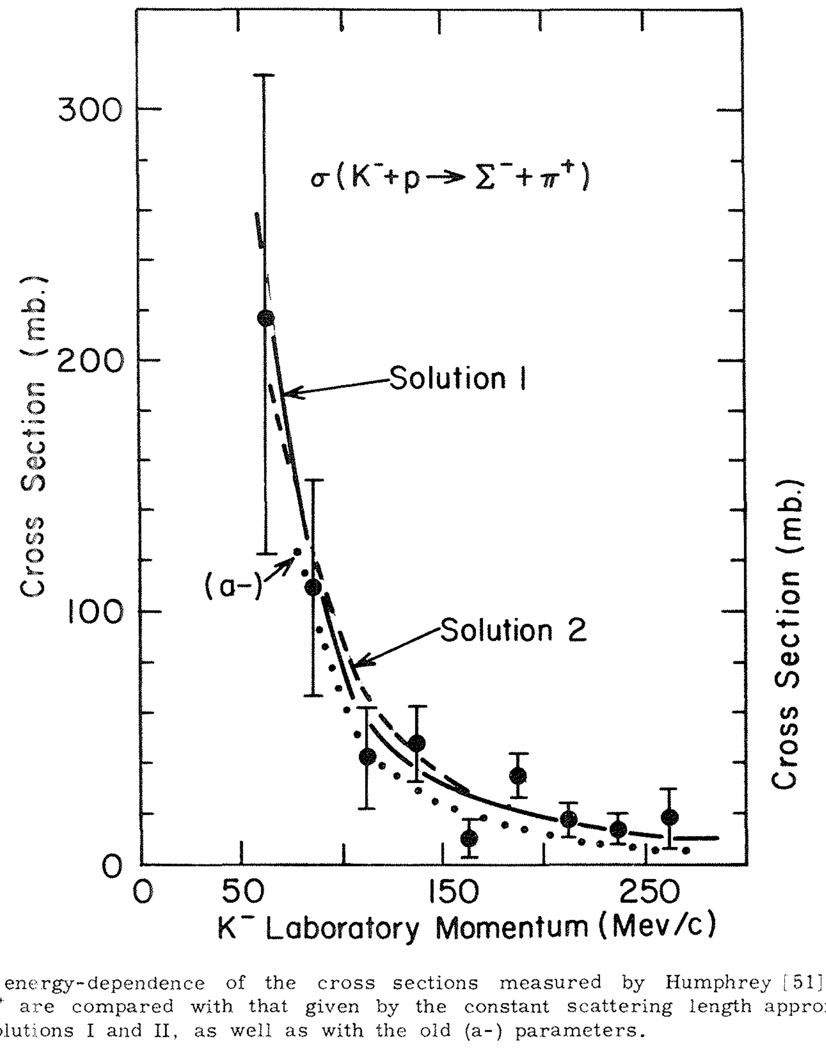

Figure 8 - The energy-dependence of the cross sections measured by Humphrey [51] for the reaction $K^- + p \longrightarrow \Sigma^- + \pi^+$ are compared with that given by the constant scattering length approximation with the parameters of solutions I and II, as well as with the old (a-) parameters.

are shown in figures 5 through 8, where it will be seen that these scattering lengths do give a good representation (*) of the data from 100 Mev/c to 275 Mev/c.

Table I

	a_o	b_o	a_1	b_1	ε	γ	ϕ	χ^2	Starting Set
Solution I	-0.22	2.74	0.02	0.38	0.40	2.15	96°	57.9	(a-), (b+).
	±1.07	±0.31	±0.33	±0.08	±0.03	±0.16			
Solution II	-0.59	0.96	1.20	0.56	0.39	2.04	-50°	73.5	(a+), (b-).
	±0.46	±0.17	±0.06	±0.15	±0.02	±0.18			

The two sets of scattering parameters determined by Ross and Humphrey. The parameters a_o, b_o, a_1 and b_1 are given in unit 10^{-13}cm. The parameter γ here denotes the Σ^-/Σ^+ ratio for K^--p capture at rest and is directly related with ϕ ; the angle ϕ given corresponds to the best values for all the other parameters.

It is next of interest to attempt to distinguish between these two parameter sets by comparing their predictions with the data for other situations involving $\bar{K}$ interactions.

1 - K_2^o-p interactions have been studied by Luers et al. [54] for K_2^o momenta in the momentum range 200-500 Mev/c. The hyperon production reactions are then those due to the I = 1 interactions alone. Values for ε may then be obtained very directly from the ratio of the reaction rates for $\Lambda + \pi^+$ and $\Sigma^o + \pi^+$ (equal to the $\Sigma^+ + \pi^o$ rate), for these two processes may very readily be distinguished in this situation in contrast to the difficulty experienced in separating the final states $\Lambda + \pi^o$ and $\Sigma^o + \pi^o$ for $K^- + p$ reactions at low momenta. Luers et al. find $\varepsilon = 0.35 \pm 0.15$ at 220 Mev/c and 0.2 ± 0.1 at 350 Mev/c, to be compared with the Ross-Humphrey value of 0.4 ± 0.03 for the best overall fit for the range 0-250 Mev/c. The analysis of Ferro-Luzzi et al. of their data [26] on K^--p reactions also gives a mean value $\varepsilon = 0.4 \pm 0.15$ over the range 300-400 Mev/c. These values are all essentially in agreement, so that there is no evidence at present for any rapid energy variation of ε over this wide momentum range.

The cross section for the reaction :

$$K_2^o + p \longrightarrow K_1^o + p \tag{2.3}$$

is given by the expression [55] :

$$\pi\left|\left(\frac{1}{2}\left(\frac{\alpha_o}{1 - ik\alpha_o} + \frac{\alpha_1}{1 - ik\alpha_1}\right) - \frac{A_1}{1 - ikA_1}\right)\right|^2 \tag{2.4}$$

where α_o, α_1 are the (energy-dependent) real scattering lengths for the I = 0 and I = 1 K-N systems, which are known from K^+ scattering data in this range to correspond to repulsive interactions. The interference between the K-N and the $\bar{K}$-N terms in (2.4) then allows the possibility of distinguishing between solutions I and II, since a_1 is large (and attractive) for solution I and gives constructive interference for the K_1^o reaction whereas a_1 for solution II is small and will give rise to very little interference. Because uncertainties in the K_2^o flux make the measurement of absolute cross sections uncertain, it is most convenient to compare the observed ratio K_1^o/(all hyperons) with the ratio of expression (2.4) to the total absorption cross section $2\pi b_1/(k|(1 - ikA_1)|^2)$. The results reported by Luers et al. are shown in figure 9. At 230 Mev/c, they favor solution I ; at higher energies (a very large extrapolation for these zero range solutions!) they require something like solution I again rather more strongly.

2 - K^--d interactions at rest have been studied in considerable detail by Miller et al. [56] at Berkeley. It is well known [57] that these data show marked effects due to final state interactions, especially of the reaction $\Sigma + N \longrightarrow \Lambda + N$. However Capps and Schult [58] have pointed out that, with the reasonable assumption that the final pion does not interact strongly with the spectator nucleon, there exist a number of sum rules which remain valid irrespective of the nature of these

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(*) We note that the calculated curves shown in figures 5, 7 and 8 are not quite correct in the region of the charge-exchange threshold at 90 Mev/c, since the form of the small cusps occuring at this threshold have not yet been computed.

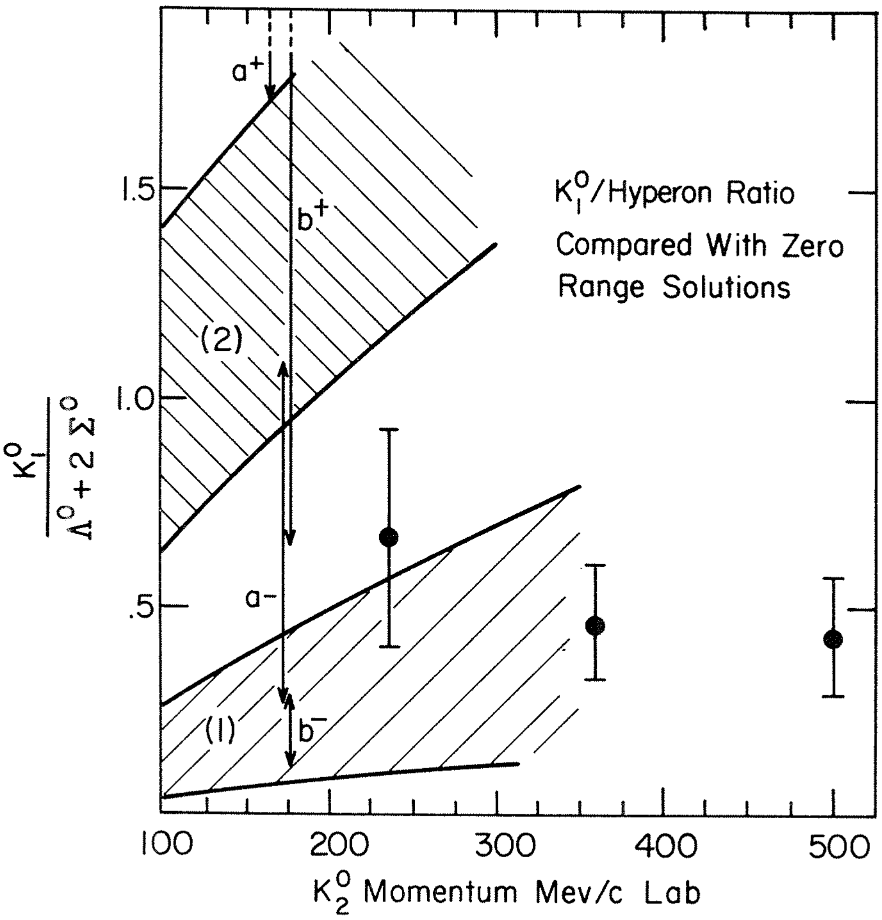

Figure 9 - The ratios K_1^0/(hyperons) observed in K_2^0-p collisions by Luers et al. [54] is compared with the values expected in the constant scattering length approximation with the parameters of solutions I and II. The values expected with the old (a ±) and (b ±) parameter sets are also indicated.

hyperon-nucleon final state interactions. For the reactions $K^- + d \longrightarrow Y + N + \pi$, there are three such sum-rule rates :

a) R_3, for $I = \frac{3}{2}$ Σ-N states. In terms of the I = 0 and I = 1 charge-independent reaction amplitudes M_0 and M_1 mentioned above, the total rate is $R_3 = \left|\left(M_1 + \sqrt{\frac{2}{3}}\, M_0\right)\right|^2$. Empirically, R_3 is given by $2\,n(\pi^+)$, appart from a common normalizing factor.

b) R_1, for $I = \frac{1}{2}$ Σ-N states taken together with the Λ-N states which have resulted from Σ conversion (the latter are recognizable from the pion momentum, which is characteristic of the $\bar{K} + N \longrightarrow \Sigma + \pi$ reaction). This rate is $R_1 = 2\left|\left(M_1 - \frac{1}{\sqrt{6}}\, M_0\right)\right|^2$; empirically, R_1 is given by $(3n(\pi^-) - n(\pi^+))/2$.

c) R_Λ, for the directly produced Λ particles. This rate is $R_1 = 3N_1^2$, where N_1 is the I = 1 reaction amplitude for $\bar{K} + N \longrightarrow \Lambda + \pi$. N_1^2 is related to the above parameters by the equations :

$$b_1/b_0 = (M_1^2 + N_1^2)/M_0^2, \qquad \varepsilon = N_1^2/(M_1^2 + N_1^2). \tag{2.5}$$

This rate R_Λ does not include those Λ particles which result from the production and decay of Y_1^*, which amounts to about 10 % of the total K^--d absorption rate and is to be reckoned as an additional competing channel.

The experimental and theoretical estimates of the relative values for these rates are compared in Table II (•). The agreement is excellent for solution I, but solution II gives rather strong disagreement with the data. This comparison reflects both the relative magnitudes of b_0 and b_1, and the phase angle ϕ. Since :

$$(R_1 + R_3)/R_\Lambda = (b_0 + 3b_1(1-\varepsilon))/3b_1\varepsilon, \tag{2.6}$$

agreement for R_Λ requires a large value for b_0/b_1, which is not the case for solution II. Agreement for R_3/R_1 then requires a small negative value for cos ϕ, as is the case (•) for solution I but not for solution II. It must be emphasized here that there are corrections arising from multiple scattering in the initial state and from the K^--$\bar{K}^0$ mass difference which have not yet been included in the discussion of the K^--d reaction rates. These corrections are not likely to upset the comparison (2.6) but they could well modify the theoretical estimate for R_3/R_1.

Table II

Rate	Experimental Values		Calculated values (I)	(II)
R_3	$2n(\pi^+)$	= 0.49	0.47	0.60
R_1	$\frac{3}{2}n(\pi^-) - \frac{1}{2}n(\pi^+)$	= 0.40	0.41	0.15
R_Λ	$n_d(\Lambda)$	= 0.11	0.12	0.25

Capps and Schult used the old (a-) parameters for which $b_0/b_1 \approx 10$, and found agreement for R_Λ but not for R_3/R_1. They pointed out that, if there existed an I = 0 resonance close to the $\bar{K}N$ threshold (at about -10 Mev), this would invalidate the assumption of the impulse approximation that the reaction parameters could be taken as constants since the relative phase of M_0 and M_1 would then vary rapidly as the $\bar{K}$-N energy became negative. The variation in ϕ which they estimated using the old (b-) parameters was then sufficient to give agreement for R_3/R_1. This remark was the first suggestion that a Y_0^* resonance may exist ; some tentative experimental evidence was also presented by Alston et al. [59] at about the same time. We see here that it may not be necessary to invoke this resonance to account for the K^--d data. In fact, we may ask whether this agreement casts doubt on the existence of the Y_0^* resonance ; we think not, since the empirical evidence [59] for this resonance puts it at about -30 Mev, relatively distant from the energies appropriate to the $\bar{K}$-N systems in K^--d capture.

K -d cross sections in flight have been measured [57] for the sum of elastic (K^--d) and inelastic (K^-np) scattering at several momenta. Calculations of the total cross sections ($\sigma_{el.} + \sigma_{inel.}$) have been carried out by Chand [60] with the inclusion of multiple scattering : these are compared with the data in Table III and clearly favor solution I.

(•) Note added subsequent to the Conference. The discrepancy between these conclusions and those of Schult and Capps [58] lies in the treatment of the corrections which arise from the mass-differences. If these corrections are neglected in the analysis of the zero-energy K^--p data, we have $M_1^2/M_0^2 = 0.137$, and cos $\phi = +0.7$, together with $N_1^2/(N_1^2 + M_1^2) = 0.4$. From these follow the values $R_3 : R_1 : R_\Lambda = 0.65 : 0.19 : 0.16$, in considerable disagreement with the observed values, as was pointed out by Schult and Capps. In the text, we have used the charge-independent amplitudes M_0 and M_1, and phase ϕ, obtained after allowing for the charge-dependent corrections. Clearly, in discussing K^--d capture, the charge-dependent corrections appropriate for this situation should also be computed and included before comparison with experiment is made. In particular, it is apparent that, for very large n-p separations, these will modify the K^--p capture in just the same way as they do for free K^--p capture, whereas there will be no modification for K^--n capture. For intermediate n-p separations, however, the modifications due to multiple scattering of the $\bar{K}$ meson before capture can be large and can affect the value of the interference term cos ϕ which enters so sensitively into the expressions R_3 and R_1. These corrections are still to be calculated an could seriously affect the values given in Table II ; probably the final result will be intermediate between those and the estimate above. Finally, we should recall that the phase ϕ arises from final state π-Y scattering. In K^--d capture, the final hyperon is in strong interaction with the spectator nucleon in many configurations, and it is quite possible that this could modify this relative phase ϕ sufficiently to affect the ratio R_3/R_1 quite significantly.

Table III

p_L Mev/c	125	135	175	200	210
Experiment (mb.)	145 ± 35	-	55 ± 15	-	95 ± 25
Solution I	-	105	-	70	-
Solution II	-	175	-	130	-

The total cross sections for the reactions $K^- + d \longrightarrow K^- + d$ and $K^- + n + p$ calculated by Chand [60] for solutions I and II are compared with the experimental data [57].

c) The nuclear optical model potential U for K^- mesons is known to be attractive [61]. Hetherington and Ravenhall [62] have recently made careful calculations of U based on zero-range scattering amplitudes and the methods of many-body theory. These calculations were made for the former sets of parameters, but they suggest that solution II will lead to an attractive U. For solution I, the large error in a_0 and the uncertain sign of a_1 make it unclear what sign is to be expected for U ; calculations of U in this way may lead to further conditions on A_0 and A_1 which will reduce this uncertainty in a_0 and a_1.

From the above remarks, we conclude that solution I leads to an excellent fit to quite a wide variety of data and is the solution to be preferred.

The relationship of the Y_0^* and Y_1^* resonant states to these sets of scattering parameters is quite unclear at present. The extrapolation to the Y_1^* energy (50 Mev below the $\bar{K}$-N threshold) is a long one, especially as the zero range theory is likely to be less reliable below this threshold than above (because of the existence of singularities expected in the $\bar{K}$-N scattering amplitudes in the unphysical region). Discussion of the possible relationship of Y_1^* with the low energy $\bar{K}$-N data will probably require the use of an energy-dependent K-matrix (see below) : the analysis of the $\bar{K}$-N interactions should then be carried out both (a) using the Y_1^* parameters as a further constraint on the $\bar{K}$-N K-matrix (assuming the Y_1^* to be a $\bar{K}$-N s-wave bound state), and (b) ignoring the Y_1^* state, as will be appropriate if it is a $j = \frac{3}{2}$ π-Λ isobar. Similarly, the relationship of the Y_0^* state with these solutions is far from clear. If desired, a fit could be obtained to Y_0^* using solution II and some energy-dependence in the scattering lengths. With solution I, b_0 is so large that it is difficult to see how such a narrow resonance ($\Gamma/2 \approx 10$ Mev) could be compatible with such a strongly absorptive interaction.

For discussion of these questions, the simple scattering length theory needs to be extended to the use of a K-matrix with elements relating all the channels, $\bar{K}N$, $\pi\Sigma$ and $\pi\Lambda$. To illustrate, consider the two-channel situation appropriate to I = 0. Taking :

$$K = \begin{pmatrix} \alpha & \beta \\ \beta & \gamma \end{pmatrix} \tag{2.7}$$

we then have the scattering matrix T given by :

$$T = K(1 - ipK)^{-1}, \tag{2.8}$$

where the momentum matrix (p) is diagonal, with elements k and q, when the base states are ($\bar{K}N$) and ($\pi\Sigma$) respectively. For the $\bar{K}N$ scattering, the phase shift δ_K is then given by $k\cot\delta_K = 1/A_K$ where :

$$A_K = \alpha + i\beta q \left(\frac{1}{1 - i\gamma q}\right)\beta,$$

$$= (\alpha - \gamma b) + ib, \tag{2.9}$$

where $b = \beta^2 q/(1+\gamma^2 q^2)$. For the $\pi\Sigma$ system, the phase shift δ_Σ is given by $q\cot\delta_\Sigma = 1/A_\Sigma$, where :

$$A_\Sigma = \gamma + i\beta k \left(\frac{1}{1 - i\alpha k}\right)\beta. \tag{2.10}$$

Below the $\bar{K}N$ threshold, we must take $k = +i\varkappa$, so :

$$A = \gamma - \beta\varkappa \frac{1}{1+\alpha\varkappa}\beta . \tag{2.11}$$

The existence of a $\pi\Sigma$ resonance of the "virtual $\bar{K}N$ bound state" type then corresponds to the vanishing of the denominator in A_Σ, $1 + \alpha\varkappa = 0$. The location of the resonance energy $E^* = (M_N^2 - \varkappa^2)^{\frac{1}{2}} + (m_K^2 - \varkappa^2)^{\frac{1}{2}}$ is therefore determined by $\alpha = a + \gamma b$, so that even a knowledge of $A = a + ib$ as function of energy is not sufficient for a unique determination of E^*. The difference between α and a can be quite large, especially if b is large. Further, the width of the resonance depends primarily on the value of β^2/α at $E = E^*$, since this governs the rate at which δ_Σ passes through $90°$ with increasing energy, and this combination also cannot be expressed in terms of a, b without a knowledge of γ.

An effective range approach has been developed for the multichannel K-matrix by Ross and Shaw [63]. This formalism closely parallels the usual effective range formalism for one-channel processes, where the quantity p^{2l}/K is expanded in powers of the energy. The natural quantity to consider proves to be the matrix $(p^l)K^{-1}(p^l)$, where (p^l) denotes the matrix which is diagonal in the channel representation and then has elements $k_i^{l_i}$, where k_i, l_i are the momentum and orbital angular momentum in channel i. The effective range approximation then consists of the linear approximation,

$$(p^l)K(E)^{-1}(p^l) = (p_o^l)K(E_o)^{-1}(p_o^l) + B(E-E_o). \tag{2.12}$$

Ross and Shaw point out that, in the channel representation, the off-diagonal elements of the matrix B are generally small relative to the diagonal elements, and suggest the use of a series of effective ranges R_{ii} for each channel, such that :

$$B_{ii} = E_o C_{l_i} R_{ii}^{1-2l_i}, \tag{2.13}$$

where $C_o = 1$, and $C_1 = -3$. We should remark here that the approximation (2.12) is likely to become inadequate quite quickly as the energy E falls below the $\bar{K}$-N threshold ($M_N + m_K$) : if there are strong $\bar{K}$-N potential interactions arising from the exchange of systems of pions (for example, from the exchange of two pions with a strong I = 0 s-wave attraction at low relative energy), these K-matrix elements will have dynamical singularities in the unphysical region (for example, a branch cut beginning at $E = (M_N^2 - m_\pi^2)^{\frac{1}{2}} + (m_K^2 - m_\pi^2)^{\frac{1}{2}}$, for the two-pion exchange mechanism). The existence of these singularities is likely to cause the expression (2.12) to deviate rapidly from the linear approximation as E approaches the upper limits of these branch cuts : we note that explicit tests of the expansion (2.12) for multichannel situations have been made only for potentials of square-well form and that no such dynamical singularities occur for potentials which fall off more rapidly with increasing r than any exponential.

A correct zero-range theory may be obtained by taking B = 0 in the expression (2.12) : with $l_i = 0$ for all channels, this corresponds to the assumption of a constant K-matrix, but even this still allows quite appreciable energy dependence for A_K^{-1} in many cases, because of the energy dependence of q. In general, moderate values for R_{ii} will lead to quite strong energy dependence for A_K. Although an effective range expansion $A_K^{-1} = A_K^{-1}(0) + \frac{1}{2} Rk^2$ is always possible for $k\cot\delta_K$, the complex "effective range" R may be very much larger than would be suggested by the interaction ranges appropriate to the physical mechanisms giving rise to the $\bar{K}$-N scattering and this expansion may be valid only over a very limited energy range.

This effective range expansion (2.12) for the K-matrix has recently been used by Shaw and Ross [64] in an attempt to find a set of K-matrix elements giving a fit to the Y_1^* parameters and to the earlier data on K^--p interactions at rest and at 175 Mev/c, as well as giving reasonable behavior for s-wave $\bar{K}$-N scattering in the 300-400 Mev/c region. All the effective ranges R_{ii} were taken equal to R, in order to reduce the number of parameters. They then found that sets of parameters $K_o(E_o)$ consistent with all of this data could be obtained only for even $K\Sigma$ parity (irrespective of the $K\Lambda$ parity), and then only if R were taken of the order of 0.4×10^{-13} cm. This parity requirement appears to stem from the high value observed for ε from Y_1^* decay, in comparison with the value $\varepsilon \approx 0.4$ required in the low energy K^--p interactions ; all of their parameter sets indicated that ε should increase with increasing energy beyond the $\bar{K}N$ threshold, with a value close to unity for the s-wave interactions at 300 Mev/c. None of the parameter sets obtained led to an I = 0 resonance. More elaborate investigations of this type, with and without the interpretation of one or both of the Y^* resonances as a $j = \frac{1}{2}$ state, will become desirable (with more general assumptions on the matrix B) as the $\bar{K}$-N scattering data increase in variety and accuracy.

More sophisticated phenomenological approaches to these multichannel situations have been based on the use of partial wave dispersion relations, with the inclusion of parameters which measure the strengths of the many interaction mechanisms which may contribute to these processes and their energy dependence. Simplified calculations of this kind have been made by Ferrari et al. [65], who have emphasized the relationship between the K-N and $\bar{K}$-N potential interactions which arise from the exchange of resonant pion systems between meson and nucleon, and by Wali et al. [66], who stressed the importance of the Yukawa interaction poles in the energy variable (due to $\bar{K} + N \longrightarrow Y \longrightarrow \bar{K} + N$, for example) for the case of odd Σ parity. As Salam [66] will be discussing the application of partial wave dispersion relations to these problems at this meeting, we shall not discuss this approach in further detail.

3 - THE Y_1^* AND Y_0^* RESONANCE STATES -

The Y_1^* resonance appears with particular prominence in the reactions :

$$\bar{K} + p \longrightarrow Y_1^* + \pi \longrightarrow \Lambda + \pi + \pi \tag{3.1}$$

which have been studied for K^- mesons at 750 and 850 Mev/c [68] and at 1 150 Mev/c [69], and for K_2^0 mesons at about 975 Mev/c [70]. The threshold momentum for this reaction is about 450 Mev/c, so that it is reasonable to expect that, at 750 and 850 Mev/c, the primary pion is emitted predominantly into low angular momentum states.

In the K^--p reaction, the final state can be reached through two channels :

$$K^- + p \longrightarrow \left\{ \begin{matrix} Y_1^{*+} + \pi^- \\ \\ Y_1^{*-} + \pi^+ \end{matrix} \right\} \longrightarrow \Lambda + \pi^+ + \pi^-, \tag{3.2}$$

and the amplitudes describing each of these channels will interfere. The sign of this interference will depend on the total I-spin of the system, as well as on the spatial configuration considered. The intensity of the final states may conveniently be plotted as a function of the c.m. kinetic energies E_+, E_- as shown in figure 10 ; this plot has the property that the number of events occuring

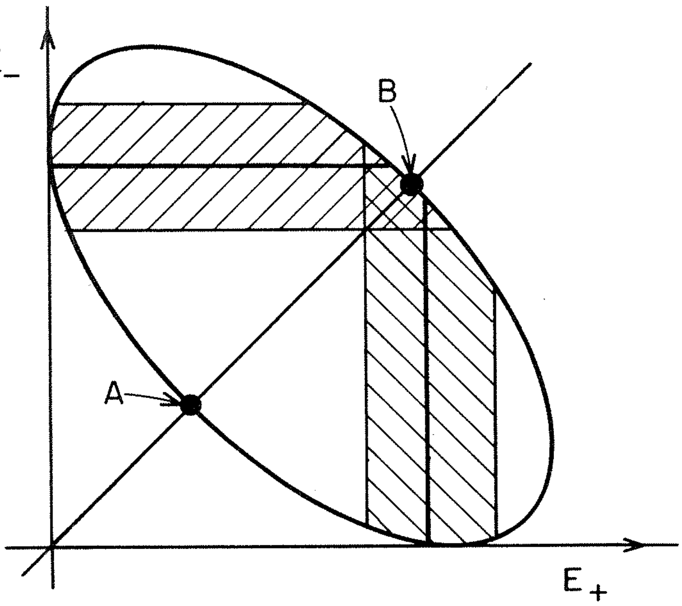

Figure 10 - Sketch illustrating the situation for the reaction $K^- + p \longrightarrow \Lambda + \pi^+ + \pi^-$ at about 850 Mev/c laboratory momentum. The axes specify the barycentric kinetic energies of the π^+ and π^- mesons. The events are restricted to lie in the elliptical region and are observed to be concentred in two overlapping bands corresponding to Λ-π resonant states of total mass 1385 Mev and half-width about 25 Mev.

per unit area is proportional to the square of the matrix-element leading to that configuration. The requirements of momentum conservation limit the events to the enclosed region shown ; the existence of the two channels (3.2) corresponds to the observation of a high intensity of events in the two bands parallel to the axes. There are two spatial configurations of particular interest ; these are displayed in figure 11, and correspond to the positions A and B on figure 10. For configuration B,

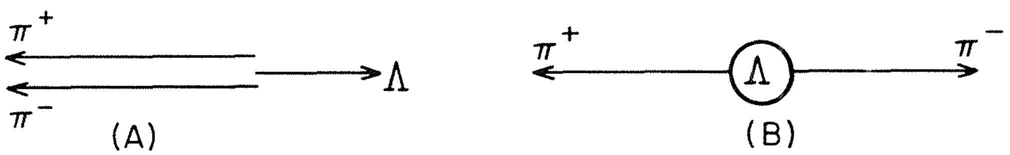

Figure 11 - Two $\Lambda + \pi^+ + \pi^-$ configurations of particular interest (see text).

where the Λ particle is at rest, the amplitude describing the sequence where π^- is the primary pion may be obtained from that where π^+ is the primary pion by an interchange of the pion charge states, followed by a reflection of all axes to bring the π^+ and π^- momentum states back to those of the configuration originally specified. For a state of total isotopic spin I, in which the primary pion carries orbital angular momentum l and the orbital angular momentum of the decay pion in $Y_1^* \longrightarrow \Lambda + \pi$ decay is L, these operations reproduce the original amplitude multiplied by the factor $(-1)^{I+l+L}$; in other words, the interference between the two channels (3.2) will be constructive (as is observed to be the case at 750 and 850 Mev/c) provided that :

$$I + l + L = \text{even}, \tag{3.3}$$

and will be completely destructive at B if $I + l + L$ = odd. For configuration A, the interference is always constructive for I = 0, destructive for I = 1. On the basis of detailed model calculations [71] for various angular momentum configurations (l, L), it has been concluded that the 750 and 850 Mev/c data are not consistent with constructive interference at A. From this it is concluded that I = 1 production is dominant at these momenta and that the configurations of particular interest are $(sP_{3/2})$ and $(pS_{1/2})$, corresponding to the interpretations of Y_1^* as π-Λ isobar and $\bar{K}$-N s-wave bound state, respectively.

At 750 Mev/c, this interpretation is supported by the observations of Prowse et al. [72] on the rate of Y_1^* production in $K^- + n$ collisions, deduced from observations on K^--d interactions. Prowse et al. find the ratio :

$$\sigma\,(K^- + n \longrightarrow Y_1^{*-} + \pi^0)/\sigma\,(K^- + p \longrightarrow Y_1^{*\pm} + \pi^\mp) = 1.4 \pm 0.3, \tag{3.4}$$

to be compared with the value 2 expected if the reaction $\bar{K} + N \longrightarrow Y_1^* + \pi$ proceeded entirely in the I = 1 state. Although the interference at A was not entirely constructive in the 750 Mev/c data, it was apparent that the density of events in the locality of A was greater (by about two standard deviations) than that expected for the I = 1 configurations $(sP_{3/2})$ or $(pS_{1/2})$. Since the (E_+, E_-) distribution is the result of averaging over all orientations of the $\pi^+\pi^-$ plane, the contributions from states of different angular momentum J and parity W do not interfere : consequently I = 0 production in states (J,W) different from the dominant I = 1 production state could be quite appreciable without affecting the degree of symmetry observed between Y_1^{*+} and Y_1^{*-} production on this plot. For example, a 30 % admixture of I = 0 production in the $(pP_{3/2})$ configuration, or the $(sS_{1/2})$ configuration, would still give a good fit to the observed plot, even in the region of A, and would also fit the observed ratio (3.4).

At 850 Mev/c, the data allow less likelihood of such strong I = 0 production, and it will be of interest to see how well the K^-+n reaction rate at this momentum will fit the ratio 2 then expected.

The interference between the two channels (3.2) strongly distorts many of the angular distributions from the forms which would be expected if the Y^* state decayed as an isolated particle (i.e. if the width $\Gamma/2$ were negligible). Figure 12 shows the c.m. angular distribution of the pion resulting from Y_1^* decay (averaged over all production angles), following the production reaction $\bar{K} + N \longrightarrow Y_1^* + \pi$. Since the matrix-elements for the $(sP_{3/2})$ or $(pS_{1/2})$ configurations always give constructive interference for configuration B, it follows naturally that this angular distribution becomes peaked towards forward angles of emission for the decay pions. The distortion from the isotropy which would be expected for the $(pS_{1/2})$ or $(sP_{3/2})$ configurations if $\Gamma/2$ were negligible is

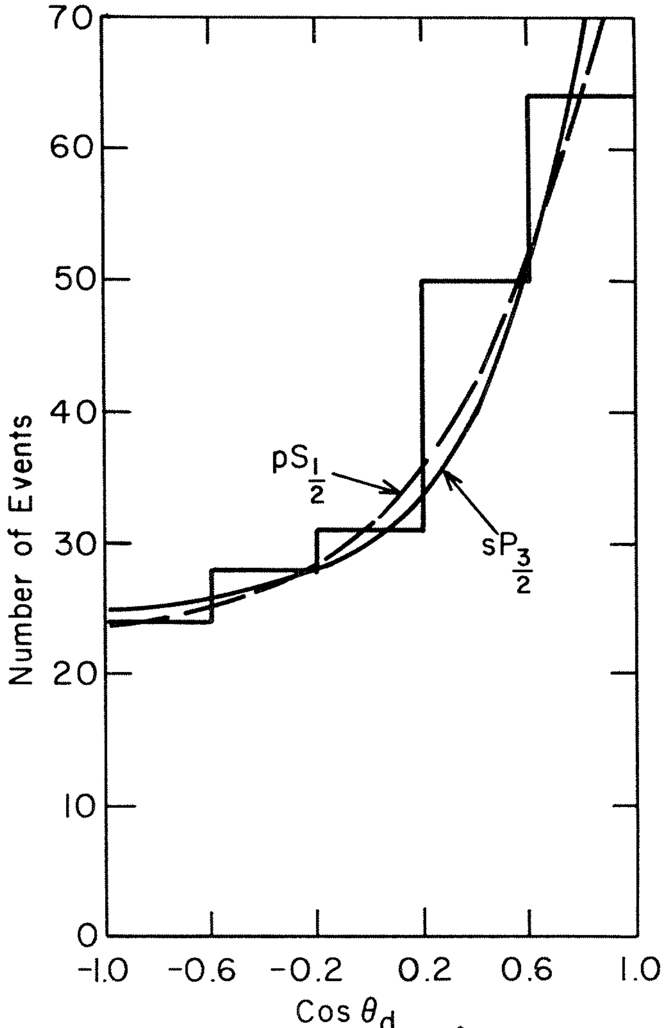

Figure 12 - The c.m. angular distribution of the pion from $Y_1^* \longrightarrow \Lambda + \pi$, relative to the Y_1^* direction of motion, is shown for all events at 850 Mev/c for which the π-Λ total mass lay in the range 1385 ± 25 Mev. Curves calculated for the configurations $sP_{3/2}$ and $pS_{1/2}$ are shown and reproduce the trend quite well.

very strong, and the calculated distributions are actually in good general accord with the observed distributions in this energy range. The same interference effect modifies the Adair distributions, the decay angular distribution for Y_1^* states produced at 0° to the incident K^- direction. Constructive interference for forward decay pions gives rise to a strong backward-forward asymmetry in these decay distributions, an effect which cannot occur with parity conservation for the decay of an isolated state. This distortion makes it very difficult to distinguish between $j = \frac{1}{2}$ and $\frac{3}{2}$ from the observed Adair distributions in this energy region, as shown in figure 13a.

These interference effects are expected to become weaker with increasing $\overline{K}$ momentum, as the two configurations (3.2) have less and less overlap. The Adair distributions are then likely to approximate more closely the idealized distributions for isolated decay. On the other hand, there is also the possibility of a third I = 1 channel contributing to reaction (3.2),

$$K^- + p \longrightarrow \Lambda + \rho^0 \longrightarrow \Lambda + \pi^+ + \pi^-. \qquad (3.5)$$

The threshold for this channel is p_K = 1 190 Mev/c, for m_ρ = 765 Mev. At present there is no direct evidence for this channel.

At 1 150 Mev/c, it is not yet established whether the production is predominantly I = 0 or I = 1 ; in fact, quite appreciable asymmetry is seen here between Y_1^{*+} and Y_1^{*-} states. The data available at this momentum are relatively limited in statistics, and the Adair distributions observed (for $\cos|\vartheta^*| > 0.8$, including both Y_1^* configurations if they both satisfy this criterion) are consistent as shown in figure 13b, but do not exclude the possibility $j = \frac{3}{2}$.

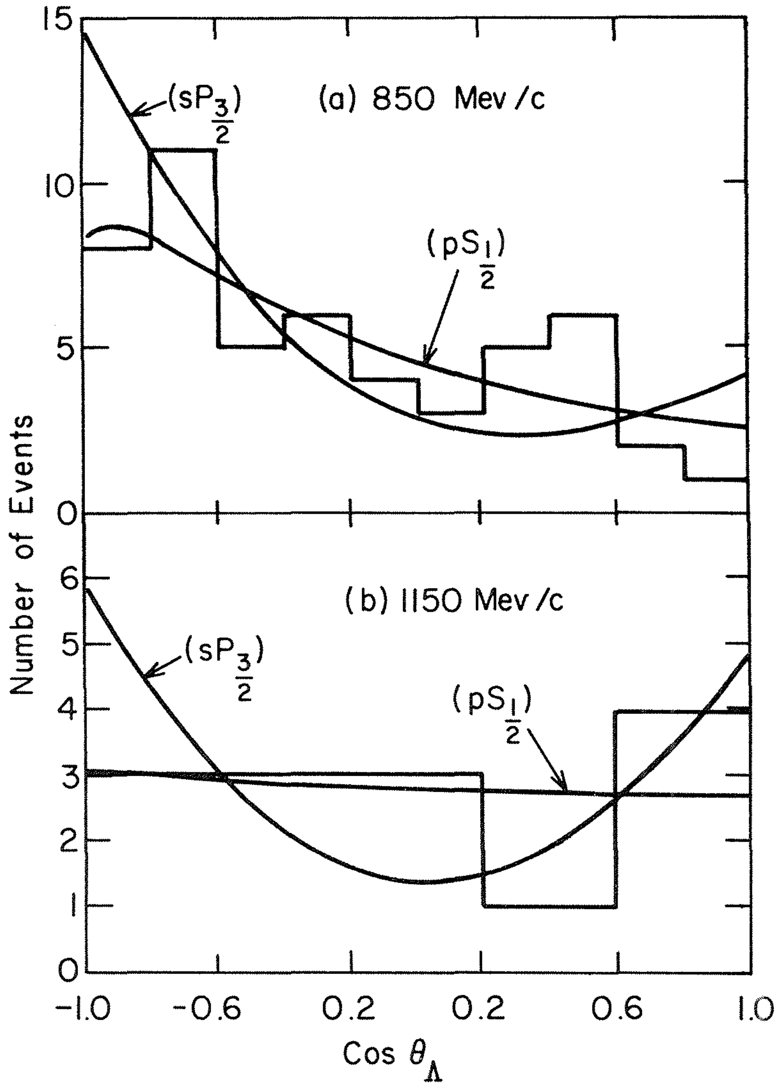

Figure 13 - Adair distributions are given for the Λ hyperon from Y_1^* decay for forward- and backward-produced Y_1^* (integrated over the π-Λ mass range 1385 ± 25 Mev) in K^--p collisions for K^- laboratory momenta 850 and 1150 Mev/c. θ_Λ denotes the angle of emission of the Λ particle in the Y_1^* rest frame, measured from the direction of motion of the Y_1^* in the barycentric frame.

At 750 and 850 Mev/c (and also at 1 150 Mev/c), the Y_1^* production angular distribution is essentially isotropic. This is compatible with the $(sP_{3/2})$ configuration, but requires $J = \frac{1}{2}$ to be dominant for the $(pS_{1/2})_J$ configuration. On the other hand, the geometrical limit for the reaction cross section for a state of angular momentum J is $2.85\left(J + \frac{1}{2}\right)$ mb. at 850 Mev/c, whereas the observed cross section is 3.2 ± 0.3 mb. The largeness of the observed cross section does argue against the dominance of a $J = \frac{1}{2}$ production state, even though a substantial fraction of this cross section might well come from states of other (J,W). This means that, with $j = \frac{1}{2}$ for Y_1^*, the interpretation of the production must be quite complicated ; with $j = \frac{3}{2}$, the most natural interpretation is I = 1 $(sP_{3/2})$ production.

The Y_1^* state has also been observed to play a role in the π^- + p reactions [34, 35],

$$\pi^- + p \longrightarrow \begin{cases} \Lambda + \pi^- + K^+, & (3.6a) \\ \Lambda + \pi^\circ + K^\circ, & (3.6b) \end{cases}$$

in the K^--d capture reaction [56],

$$K^- + d \longrightarrow p + Y_1^{*-} \longrightarrow p + \Lambda + \pi^- , \tag{3.7}$$

and in the K^--He^4 capture reaction [72],

$$K^- + He^4 \longrightarrow He^3 + Y_1^{*-} \longrightarrow He^3 + \Lambda + \pi^- . \tag{3.8}$$

The statistics on reactions (3.6) are relatively limited and the Y_1^* resonance does not show up prominently in the data. The Adair distributions observed are compatible with isotropy. In reaction (3.7) the effects of Y_1^* production are clearly seen but the Y_1^* events cannot be individually separated from $p + \Lambda + \pi^-$ configurations which result from the final state reaction $\Sigma + N \longrightarrow \Lambda + N$. The Y_1^* produced in reaction (3.8) travels a mean distance of order 10^{-13} cm before decay, whereas the r.m.s. radius of He has been established by Hofstadter and Collard [73] to be about 1.7×10^{-13} cm. In this situation, final state interactions are likely to be of considerable importance in distorting the Y decay angular distribution. For low He^3 momenta, the distribution is observed to be asymmetric, the Λ hyperon preferring to follow the He nucleus, presumably a result of the attractive forces in the Λ-He^3 systems. For large He^3 momenta, the Λ-π^- angular distribution is found to be essentially isotropic, although statistics are then rather limited. If the K^--He^4 capture were from s-orbitals, isotropic decay would require $j = \frac{1}{2}$ for the Y_1^* state. However the estimates of Day [2] would allow about 10 % of K^--He^4 capture to occur from p-orbitals, and it is possible that the fraction might be larger than this ; with p-orbital capture, isotropic Y_1^* decay would also be possible with $j = \frac{3}{2}$. To sum up, we must conclude that it is difficult to obtain any clear-cut conclusion on the Y_1^* spin from the existing data.

The strongest evidence for the Y_o^* resonant state is that presented by Alston et al. [59] from the study of the reactions,

$$K^- + p \longrightarrow \begin{cases} \Sigma^{\pm} + \pi^{\mp} + \pi^+ + \pi^- , & (3.9a) \\ \Sigma^o + \pi^o + \pi^+ + \pi^- , & (3.9b) \end{cases}$$

for K^- momentum 1 150 Mev/c. Little evidence of the Y_o^* state has been found in the study of other reactions. The reaction $Y_o^* + \pi$ can take place only from the I = 1 component of the $\overline{K} + N$ interaction : no indication of this state is observed in the K_2^o-p interactions at 975 Mev/c [74], which are entirely I = 1, nor in the K^--p interactions in the range 750-850 Mev/c, which are believed to occur dominantly with I = 1. If Y_o^* and Y_1^* have $j = \frac{1}{2}$ and $j = \frac{3}{2}$, respectively, and if these reactions are dominated by the I = 1 state $(J, w) = \frac{3}{2} -$ —— perhaps as the result of a resonant affect —— then Y_o^* production would occur with an s-wave pion, whereas Y_o^* production would require emission of a p-wave pion : the latter might be sufficiently less probable for Y_o^* production to be weak relative to Y_1^* production in this energy range. At 1 150 Mev/c, the absence of the $Y_o^* + \pi^o$ reaction might be understood as an indication that the reactions at this momentum are predominantly in the I = 0 channel, which may be made plausible by pointing out that the momentum 1 150 Mev/c is well up on the high-energy wing of the marked I = 0 K^--p resonance observed by Kerth [42] at 1 815 Mev. In the K^--d and K^--He^4 capture reactions, the identification of the Y_o^* state is made difficult by the fact that the Y_o^* mass is almost coincident with the peak in the distribution of π-Σ c.m. energies which would be predicted from phase space and the nucleon momentum distribution in these nuclei, and no conclusion has yet been reached concerning the presence or absence of Y_o^* in these reactions. In view of the paucity of evidence for the Y_o^* state at 1 405 Mev, it is not surprising that there are no indications on the Y_o^* spin and parity at present.

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EXPERIMENTAL STATUS OF STRANGE ISOBARS

F. T. SOLMITZ

I - INTRODUCTION -

This has been an extremely fruitful year for the study of resonances, particularly resonances involving strange particles. The first experimental evidence for such a strange resonance was presented at the Rochester Conference a year ago. In the course of the year this resonance has been confirmed in half a dozen other experiments ; at the same time, evidence for 4 or 5 additional resonances has accumulated.

Let us start with some quite simple general considerations : the same phenomenon is referred to either as isobar or as resonance. The concept of resonance is most immediately applicable to the case of two interacting particles. If there exists an attractive force between two particles, the phase shift of the system may pass through 90° for some energy ; the scattering cross-section will then pass through a maximum at (or near) that energy.

The well known pion-nucleon resonance in the (3/2, 3/2) state is an illustration of such a simple resonance. In the case of the $\bar{K}$-nucleon system we may have hyperon production in addition to elastic scattering. Schematically :

$$\bar{K} + N \begin{cases} \rightleftarrows \pi + \Lambda \\ \qquad \downarrow\uparrow \\ \rightleftarrows \pi + \Sigma \end{cases}$$

A bump in the cross-section for any one of the possible scattering or reaction process may now be attributed to the influence of either $\bar{K}$-N forces or π-Λ forces or π-Σ forces. If one could turn off the coupling between the various channels, a resonance peak in the $\bar{K}$-N channel, for instance, would be attributable to $\bar{K}$-N forces ; but in the presence of coupling such a resonance has repercussions in all the other channels that are open. It is then more natural to speak of the existence of an isobar, that is, a comparatively long-lived state with certain quantum numbers and certain partial decay rates into the various open channels. It is the task of the theoretician to "explain" the presence of the isobar and its properties in terms of forces between pairs of particles, or if he wishes, in terms of more general models. All the interesting experimental information relating to an isobar is contained in the specification of :

a) its mass M,

b) width Γ, related to its lifetime τ by $\Gamma = \hbar/\tau$,

c) strangeness,

d) isotopic spin I,

e) spin S,

f) parity P,

g) branching fractions into various channels.

Table I summarizes all the experimental information on strange isobars that has been either published, or presented at this conference.

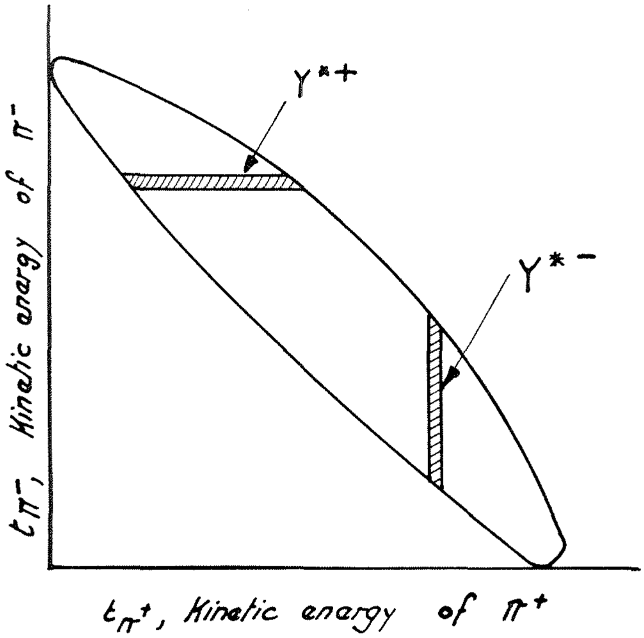

Figure 1 - Schematic of an "ideal" Dalitz plot for the state $\Lambda\pi^+\pi^-$.

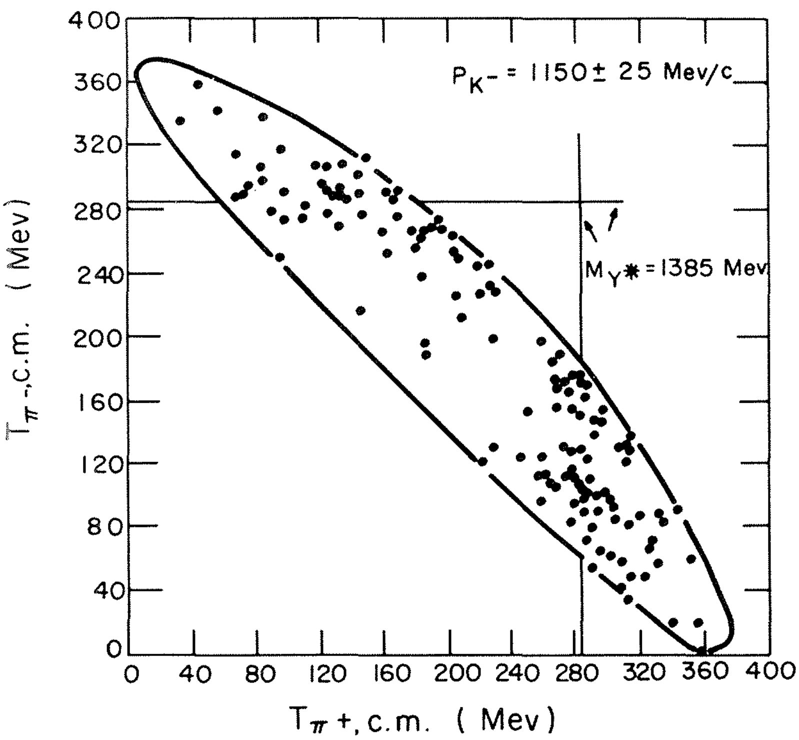

Figure 2 - Dalitz plot for the reaction $K^- + p \longrightarrow \Lambda + \pi^+ + \pi^-$ for K^- of 1150 Mev/c [1].

Table I

Summary of available data on strange isobars

Mass (Mev)	Half-width (Mev)	I-spin	Spin	Parity	Decay products	Branching fraction	References
Baryon isobars:							
1385	± 25	1	?	?	$\Lambda + \pi$ $\Sigma + \pi$	96 % 4 %	1 - 9
1405	± 10	0	?	?	$\Sigma + \pi$ $\Lambda + 2\pi$	?	10 - 12
1525	± 20	0	$\geqslant 3/2$	?	$\Sigma + \pi$ $\Lambda + 2\pi$ $K^- + p$ $\bar{K}^\circ + n$	?	13
1815	± 60	0	?	?	many	?	14
Meson isobar:							
885	± 8	1/2	?	?	$K + \pi$	100 %	15 and 8

Let us begin with a discussion of the baryon isobars, i.e. the Y^*'s. It is clear that for such isobars with a mass less than $(m_{K^-} + m_p) = 1432$, the $\bar{K}$ - N channel is not open. Since one cannot in practice perform pion-hyperon scattering experiments, one must look to more complicated processes for evidence of such isobars. In a reaction in which the final state consists of 3 or more particles, one can study the rate as a function of the relative energy of some pair of particles. An isobar may then manifest itself as a maximum in the rate for a particular energy of such a pair. If we want to study pion-hyperon resonances we can look for instance at reactions of the type :

$$\bar{K} + N \longrightarrow (\pi + Y) + \pi' \qquad \text{(i)}$$

or :

$$\bar{K} + (\text{Nucleus}) \longrightarrow (\pi + Y) + (\text{Nucleus})' \qquad \text{(ii)}$$

or :

$$\pi + N \longrightarrow (\pi + Y) + K \qquad \text{(iii)}$$

For the moment we shall consider reactions involving only 3 particles in the final state. Due to the equations of energy and momentum conservation, such a 3-body configuration at a given total energy can be described in terms of just two independant parameters (we are ignoring the orientation of the 3-body system with respect to the direction of the incident particle, and possible polarization effects). It turns out to be convenient to pick the kinetic energies of two of the 3 particles as the two parameters. Thus, for instance, in the case of the $(\Lambda\,\pi^+\pi^-)$ system we may simply specify t_{π^+} and t_{π^-} ; a collection of configurations of the $(\Lambda\,\pi^+\pi^-)$ system is then represented as a series of points in the t_{π^+} - t_{π^-} plane. These points are kinematically confined to the interior of a roughly elliptical region (see figure 1). It can be shown that such a "Dalitz-plot" should be uniformely populated if the transition matrix element is independant of the energies of the final particles [16]. There is of course no reason to suppose that this be the case, but in general one would expect a more or less smooth distribution in the absence of strong final-state interactions. If on the other hand two of the final particles, say the Λ and π^-, are the decay products of a comparatively long-lived isobar, Y^{*-}, then the π^+ produced along with the Y^{*-} will have a fixed energy, and the events will be confined to a vertical line on the Dalitz-plot ; for a finite life-time of the Y^{*-}, the line will of course have a certain width. Similarly, a Y^{*-} isobar would populate a horizontal band and a π-π resonance might give rise to an increased population along a line sloping at - 45°, corresponding to a fixed Λ kinetic energy.

Adair [17] has argued that a clumping of events in the general region of the two shaded bands of figure 1 might be due not to a resonance, but to centrifugal barrier and angular momentum effects. It would seem a good idea to subject any evidence for a new isobar to the most careful scrutiny to make sure that one is not dealing with another type ot phenomenon. However, the 1385 Mev Λ-π resonance first reported at Rochester has survived Adair's challenge : firstly, the Dalitz plots do not really very closely resemble the distribution predicted by Adair's model ; secondly, the influence of the Λ-π resonance has been found to play an important role in all the reactions of the type (i), (ii) and (iii) which have been studied to date.

II - Y_1^*, 1385 Mev ISOBAR -

Let us briefly review the Y_1^* evidence : figure 2 shows the Dalitz-plot of the data for the reaction :

$$K^- + p \longrightarrow \Lambda + \pi^+ + \pi^-$$

for an incident K^- momentum of 1150 Mev/c (this is the data first reported at Rochester). We see fairly clearly the clustering along the indicated horizontal and vertical line. Figure 3 shows the Dalitz-plot for the same reaction but at lower K momentum, 850 Mev/c. Here the situation is somewhat confused due to the fact that the two lines come very close to each other. This confusion is more striking at a K momentum of 760 Mev/c, were the two lines actually cross inside the "ellipse" (see figure 4). Nevertheless, it should be noted that most of the events fall within the indicated bands, although the latter cover only a comparatively small fraction of the area of the "ellipse".

Another reaction of the type (i), namely :

$$\bar{K}^\circ + p \longrightarrow \Lambda + \pi^+ + \pi^\circ,$$

has veen studied by Martin et al. [3] ; here again the influence of the Λ-π resonance is quite striking.

Next we turn to reactions of the type (ii), i.e. interactions of $\bar{K}$'s with nuclei. Such reactions have been studied with K incident on d, He, and C, and in each case the Λ-π resonance was found to have an important effect (2, 5, 6). As an illustration, let us consider the reaction :

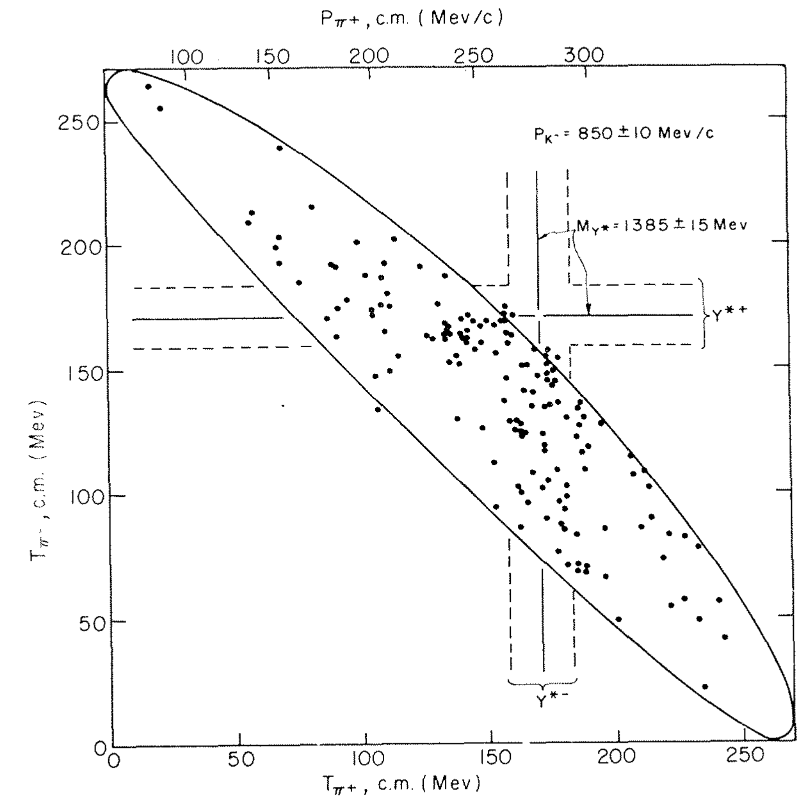

Figure 3 - Dalitz plot for the reaction $K^- + p \longrightarrow \Lambda + \pi^+ + \pi^-$ for K^- of 850 Mev/c [4].

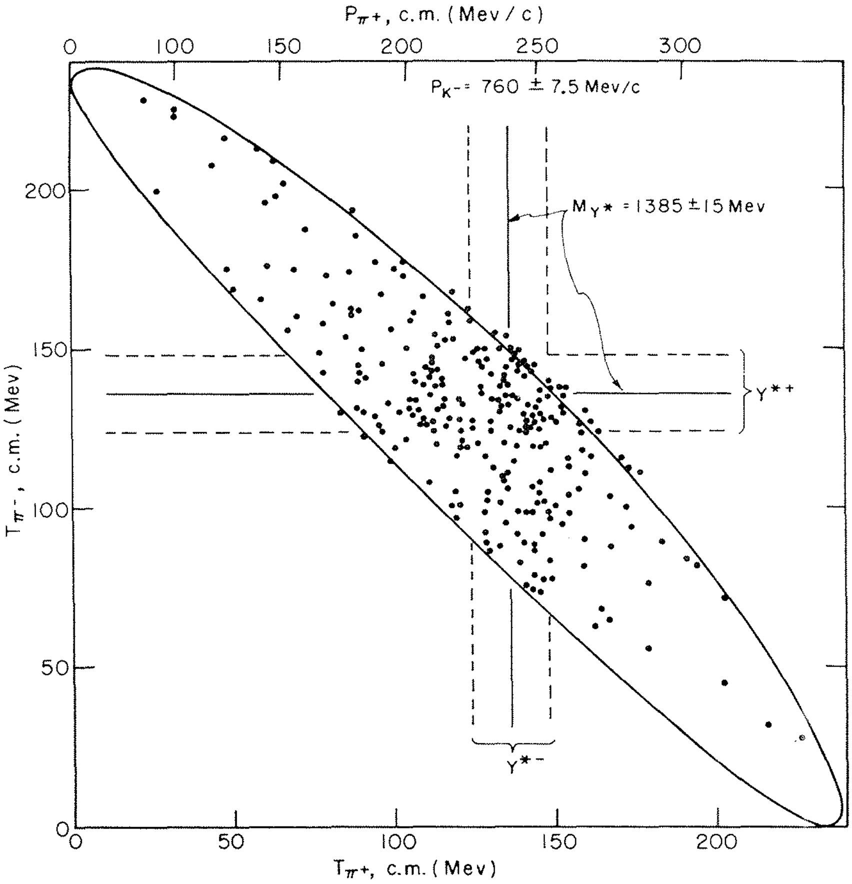

Figure 4 - Dalitz plot for the reaction $K^- + p \longrightarrow \Lambda + \pi^+ + \pi^-$ for K^- of 760 Mev/c [4].

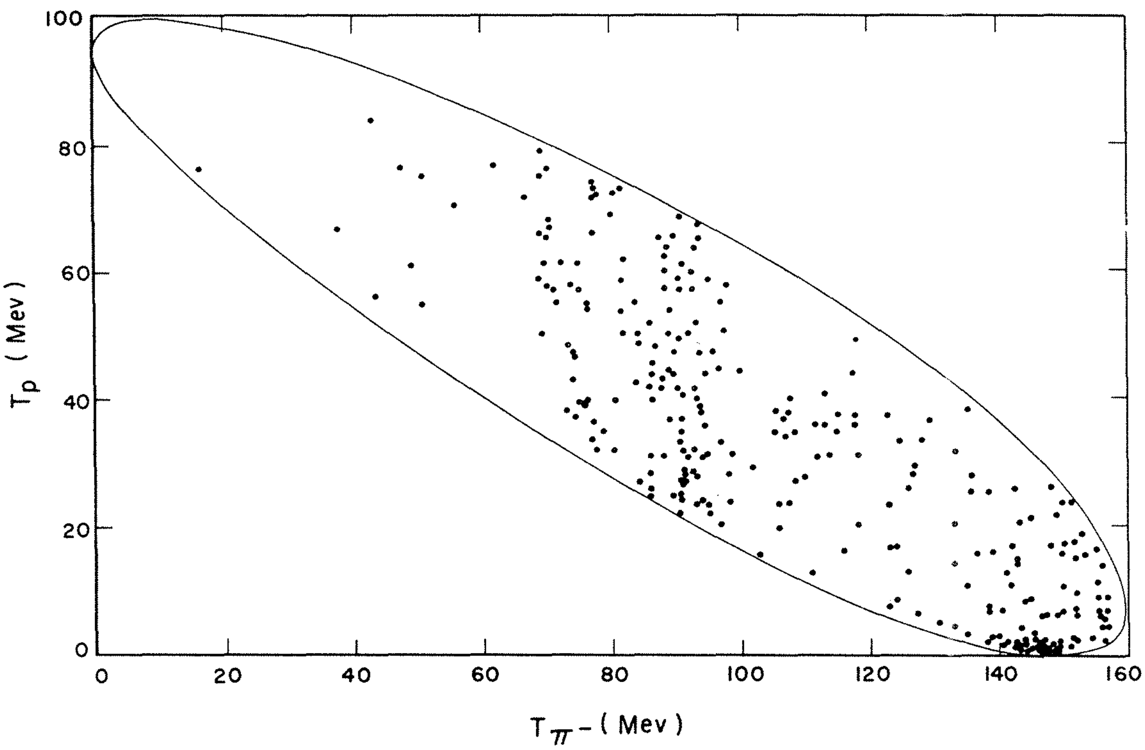

Figure 5 - Dalitz plot for the reaction $K^- + p \longrightarrow \Lambda + \pi^- + p$ for K^- mesons captured at rest [2].

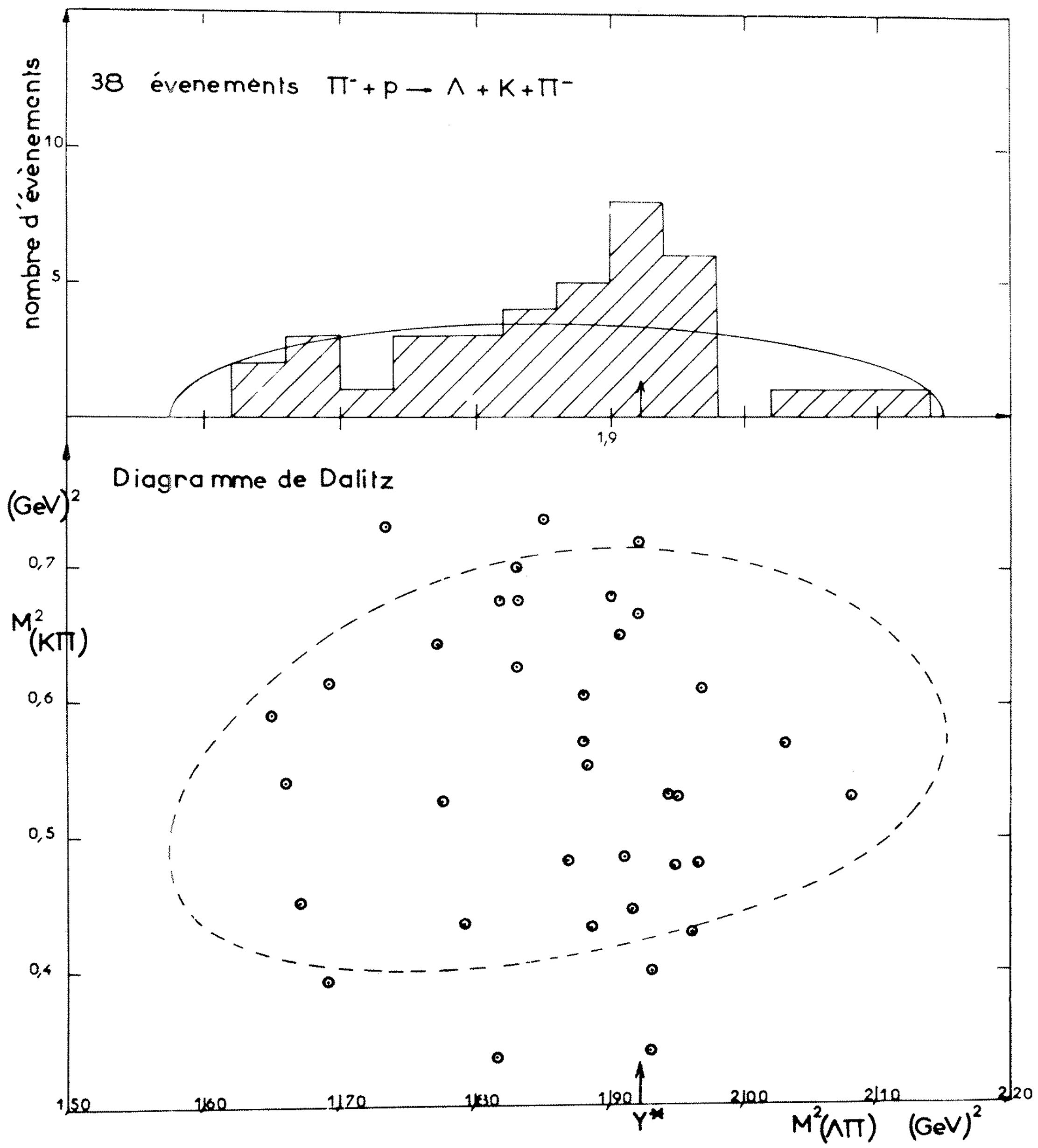

Figure 6 - Dalitz plot for the reaction $\pi^- + p \longrightarrow \Lambda + \pi + K^+$ for π^- of 1.6 Gev/c [7].

$$K^- + d \longrightarrow \Lambda + \pi^- + p.$$

The data of Dahl et al. [2] are shown in figure 5. The interpretation is somewhat complicated : firstly, the K^- can interact with the neutron to give $\Lambda + \pi^-$; the original proton will show up in the final state as a very low energy "spectator" proton ; figure 5 clearly shows a cluster of events in which the proton has a very low energy ; secondly, other events must be attributed to the now well known conversion processes, e.g. :

$$K^- + (p, n) \longrightarrow \Sigma^+ + \pi^- + n, \qquad \Sigma^+ + n \longrightarrow \Lambda + p.$$

In these events the π^- has an energy of about 90 Mev ; they form a vertical band on the Dalitz plot. Most of the remaining events on the Dalitz plot, about one third of the total, have protons of about 30 Mev kinetic energy, corresponding to a mass of the $\Lambda\pi$ system of about 1385 Mev.

No evidence for the Y_1^* in reactions of type (iii), i.e. pion-nucleon interactions, has yet appeared in print, but three reports contaning such evidence were presented at this conference.

Rogozinski [7] presented the preliminary results of a collaboration experiment involving groups at Saclay, Orsay, Bologna and Bari, - a study of the interactions of 1.6 Gev/c π^- in hydrogen. Figure 6 gives the Dalitz plot for the reaction :

$$\pi^- + p \longrightarrow \Lambda + \pi^- + K^+ ;$$

note that there the scales are labelled with squares of masses of two-body systems rather than kinetic

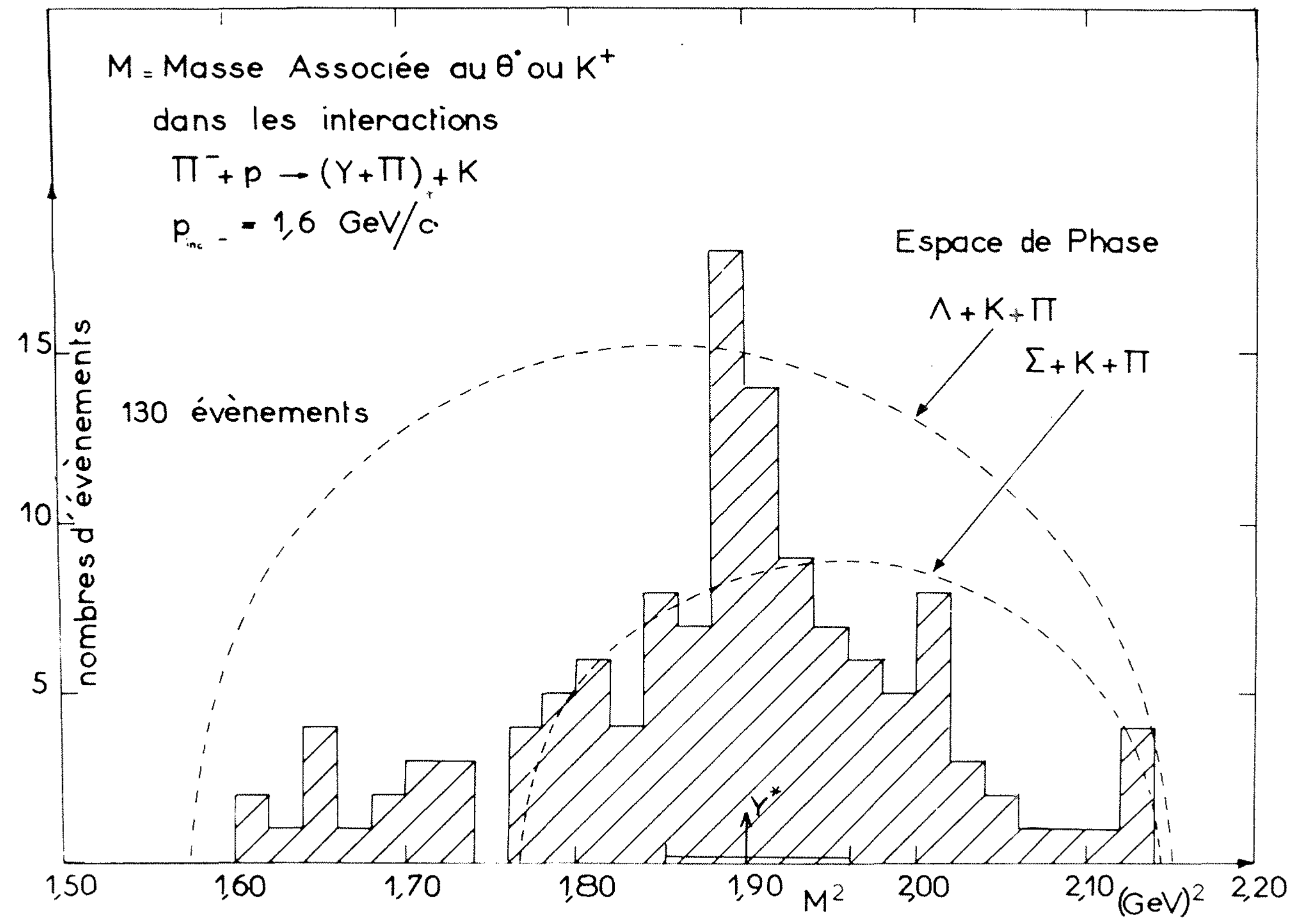

Figure 7 - Y^* mass plot for all reactions of the type $\pi^- + p \longrightarrow Y^{*\pm}_{0} + K^{+}_{0}$ for π^- of 1.6 Gev/c [7].

energies ; this turns out to be equivalent since there is a linear relation between the square of the mass of a two-body system and the kinetic energy of the third particle. One finds, for instance :

$$t_K = \frac{(E - M_K)^2}{2E} - \frac{1}{2E} M^2_{(\pi, \Lambda)} ;$$

here E is the total c.m. energy available to the three particles. Figure 7 gives a mass plot for the pion-hyperon system including data for all the possible charge states. One sees a distinct bump at a point corresponding to the 1385 Mev resonance.

Walker [8] presented preliminary results of an experiment in which 1.9 Gev/c π^- were incident on hydrogen. Figure 8 shows the Λ - π^- mass plot for the $\Lambda \pi^- K^+$ final state. Again we see evidence of a bump around 1385 Mev.

Finally Stroot [9] reported on a π^--p counter experiment performed at CERN (all the data previously discussed come from bubble chamber experiments). In this experiment, the rate of K^+ produced at a given laboratory angle is measured as a function of incident π^- momentum. It turns out that the mass of the particle or particle system produced along with the K^+ is a monotonic

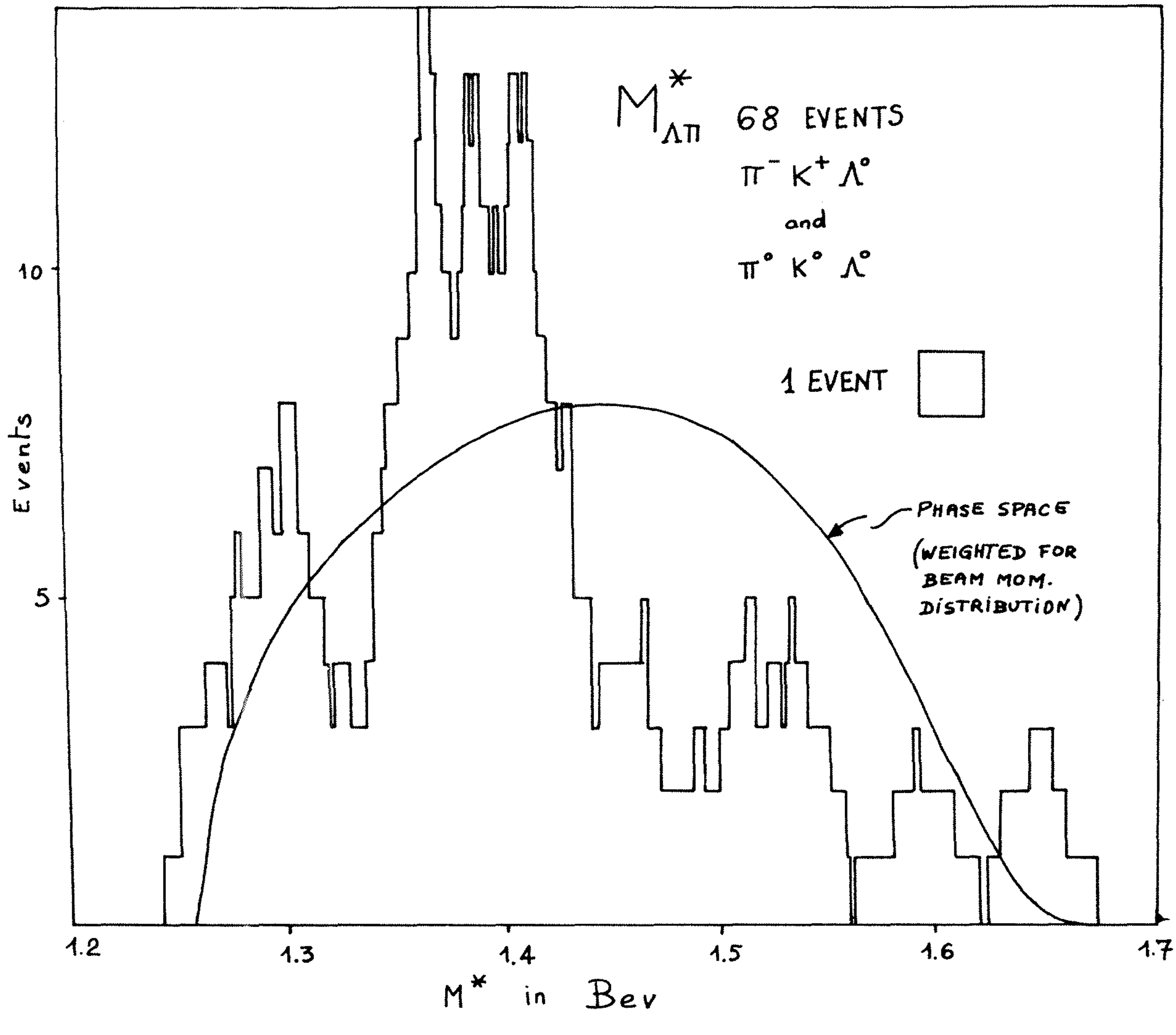

Figure 8 - Λ - π mass plot for the reaction $\pi^- + p \longrightarrow \Lambda + \pi^- + K^+$ for π^- of 1.9 Gev/c [8].

function of the π^- momentum. This particle system has strangeness - 1 ; its isotopic spin I must be 1 or greater, since its charge, and hence I_z, is equal to - 1. The experimental data (figure 9) show a bump at the position of the Σ^- mass, and at about 1385 Mev ; in addition there is a third bump around 1580 Mev ; further work will be needed to establish whether or not this third bump represents a new resonance.

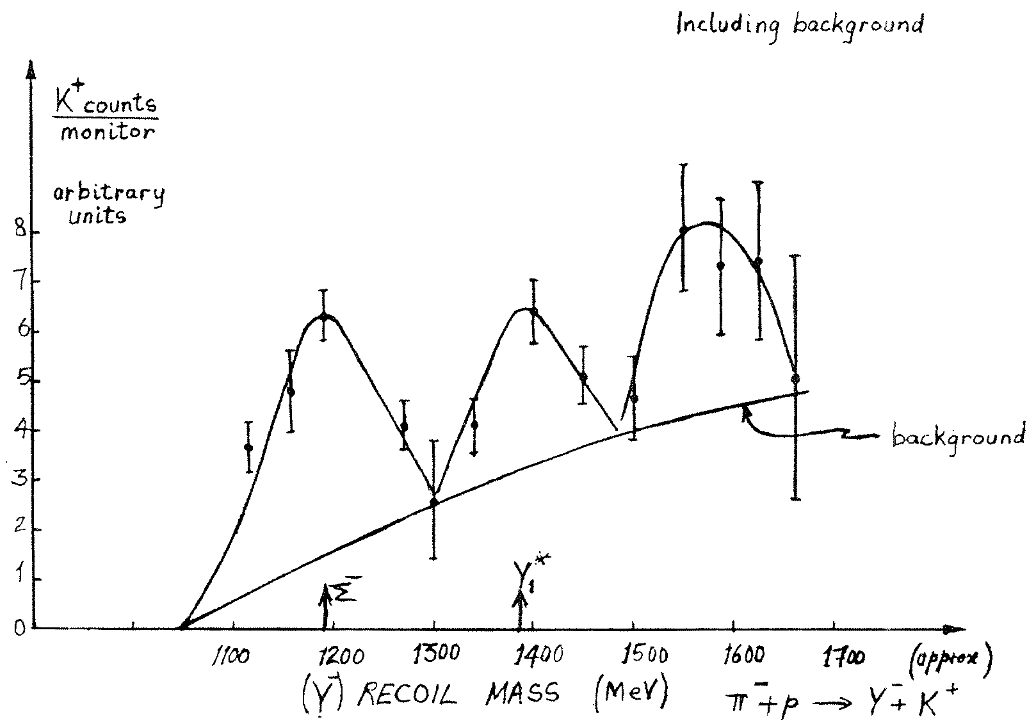

Figure 9 - Rate of K^+ produced as a function of Y^- mass in the reaction $\pi^- + p \rightarrow Y^- + K^+$; here Y^- designates the particle system produced along with the K^+ [9].

The evidence for the existence of the Y_1^* isobar is apparently overwhelming. In proving its existence, we obtain at once its mass, 1385 Mev, its strangeness, - 1, and its isotopic spin, 1 ; the latter two properties follow from the fact that isobar decays into $\Lambda + \pi$. The width is also easily obtainable, provided the experimental resolution is adequate. A half-width of about 25 Mev seems to be compatible with most of the experimental observations [18].

Since the sum of the Σ and π masses is less than 1385 Mev, one would expect the Y_1^* to decay into $\Sigma + \pi$ some fraction of the time. One can look for such an effect in, for instance, the $\Sigma\ 2\ \pi$ final states of K^- p interactions ; such an examination [10, 11] shows no clear-cut effect with presently available statistics, and leads to an upper limit of the Σ / Λ branching fraction of a few per cent.

One might hope to obtain the spin and parity of the isobar from a study of the angular distribution and polarization of its decay products ; this is hard even for such comparitively stable particles as Λ's and Σ's. In the case of the Y_1^* there are additional difficulties ; it happens not to be a good approximation to treat the Y_1^* as real particle (at least when it is produced in K^- p reactions at moderately low energy). The angular distribution of the Λ's with respect to the line of flight of the Y^* (figure 10) gives an indication of this. For the decay of a particle via strong (hence, presumably, parity conserving) interactions, this distribution should be symmetric about 90° ; the data show a marked asymmetry. Dalitz and Miller [18] have shown that this effect can be explained in terms of interference arising from the requirements of Bose statics for the two pions present in the final state. However, the problem becomes so complicated that one is unable to drawn any definite conclusions from the presently available data [19].

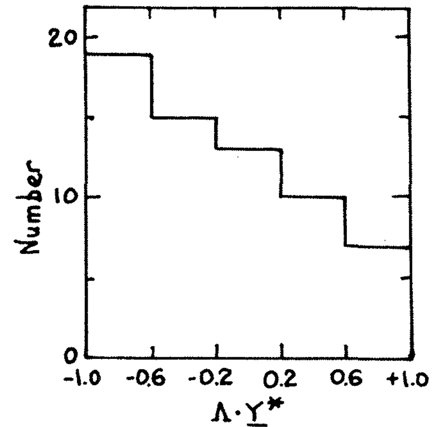

Figure 10 - Angular distribution of the decay Λ with respect to the line of flight of the Y^* [4].

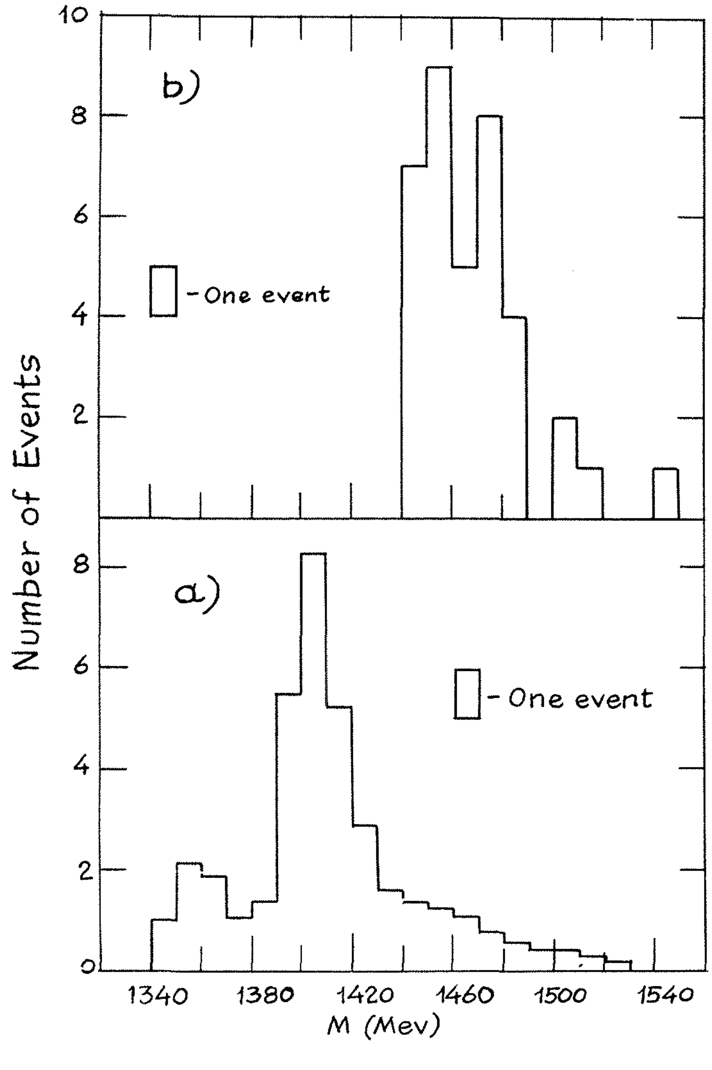

Figure 11 - Y_0^* effects in K^- absorption at flight (37 events) [12].
a) observed $Y_0^* \longrightarrow \Sigma^\pm + \pi^\mp$ invariant mass distribution.
b) expected M-distribution for K^- absorption free protons (no Y^* production).

III - 1405 Mev I = 0 ISOBAR -

A strangeness - 1 isobar with isotopic spin zero would be expected to manifest itself in neutral $\Sigma\pi$ system, i.e. $\Sigma^+\pi^-$, $\Sigma^-\pi^+$, or $\Sigma^o\pi^o$. Alston et al. [10] and Bastien et al. [11] have examined K^- p reactions leading to a Σ and several π's in search for evidence of such an isobar. The events with 2 π's do not show convincing evidence for the existence of an isobar. The analysis of the events with 3 π's is evidently somewhat more complicated ; if we consider, for instance, the state $\Sigma^+\pi^-\pi^+\pi^-$, we can imagine the Σ^+ to be associated with either one of the π's. Alston et al. [10] find that in almost all of 32 events examined there is a neutral $\Sigma\pi$ combination with a mass in a narrow band around 1405 Mev (half width 10 Mev). In addition, about half of 16 events which are most probably examples of the reaction :

$$K^- + p \longrightarrow \Sigma^o + \pi^o + \pi^+ + \pi^-$$

have a $\Sigma^o\pi^o$ mass not far form 1405 Mev (actually, the peak appears to be about 1390). The authors conclude from the relative number of $\Sigma^+\pi^-$, $\Sigma^-\pi^+$ and $\Sigma^o\pi^o$ events, that the indicated isobar must have isotopic spin zero.

Eisenberg et al. [12] have presented independent evidence for a $\Sigma\pi$ resonance at a mass of about 1405 Mev. This evidence is obtained from a study of reactions of the type :

$$K^- + (\text{emulsion nucleus}) \longrightarrow \Sigma^{\pm} + \pi^{\mp} + \text{other particles.}$$

The $\Sigma\pi$ mass spectrum for 37 K^- absorptions in flight is shown in figure 11a. Figure 11b shows the distribution of c.m. energies of the K^- and an assumed free, stationary proton for the same 37 events. If one were indeed dealing with reactions of the type $K^- + p \longrightarrow \Sigma^{\pm} + \pi^{\mp}$, one would expect the distribution of figure 11a to be similar to that of figure 11b, with some widening and some lowering of the average energy, due to Fermi motion of the proton and excitation of the residual nucleus. Instead one observes a curve with a considerably sharper peak and a much reduced average energy. The authors conclude that this effect must be attributed to a $\Sigma\pi$ resonance at about 1405 Mev.

IV - 1525 Mev I = 0 ISOBAR -

Evidence for another I = 0 isobar with a mass of 1525 Mev has been found by Ferro-Luzzi et al. [13]. The K^- p absorption cross-section in the I = 0 state goes through a maximum for incident K^- momentum of about 400 Mev/c (see figure 12). At this momentum the angular distribution of the reaction products, as well as the angular distributions for elastic and charge exchange scattering, are strongly anisotropic, indicating that the apparent isobar has a spin greater than 1/2.

V - 1815 Mev I = 0 ISOBAR -

Kerth [14] has presented evidence for a third I = 0 isobar. This evidence consists of a bump in the total K^- p cross-section (the I = 0 assignment is based on the absence of a bump in the K^- n cross-section at the same energy).

VI - 885 Mev K RESONANCE -

The only K-π resonance that has been found so far is the one reported by Alston et al. [15] and based on a study of the reactions :

$$K^- + p \longrightarrow \bar{K}^o + \pi^- + p,$$
$$\longrightarrow K^- + \pi^o + p,$$
$$\longrightarrow K^- + \pi^+ + n$$

The analysis is similar to that of the $\Lambda\pi$ resonance. The ratio of the rates for the first two reactions leads to an assignment of I = 1/2. Walker [8] presented additional evidence for this resonance, based on the study of π-p reactions at 1.9 Gev/c.

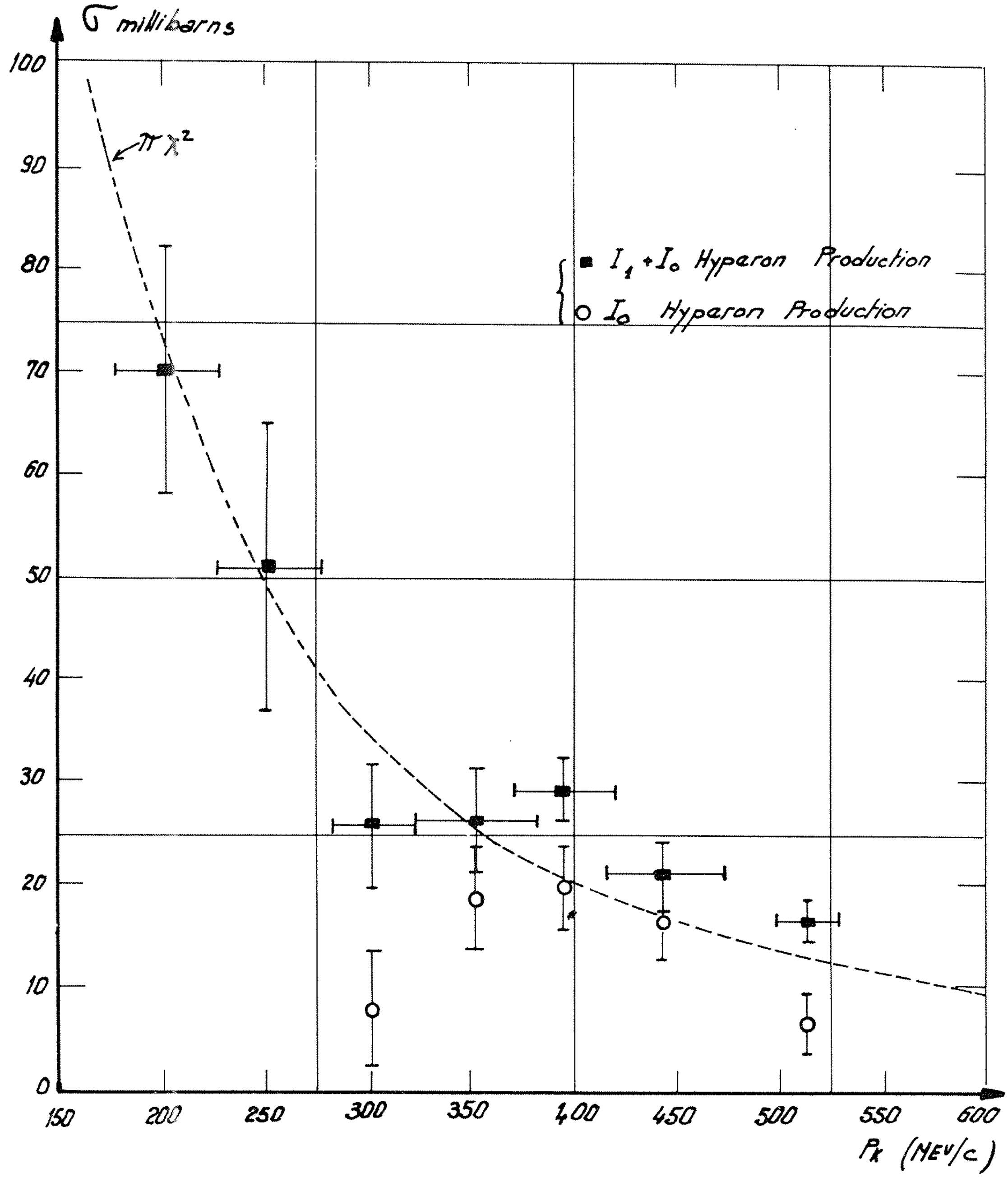

Figure 12 - Hyperon production cross-section in low energy K^- p interactions [13].

VII - ACKNOWLEDGMENTS -

The author is indebted to Y. Eisenberg, M. Ferro-Luzzi, A. Rogozinski, J.P. Stroot, and W. Walker for informing him of unpublished results.

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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 :

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

Here :
$$Z_1 = \begin{pmatrix} \Sigma^+ \\ \dfrac{\Lambda - \Sigma^\circ}{\sqrt{2}} \end{pmatrix} \qquad Z_2 = \begin{pmatrix} \dfrac{\Lambda^\circ + \Sigma^\circ}{\sqrt{2}} \\ \Sigma^- \end{pmatrix}$$

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 :

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

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/ $\Sigma^- + p \longrightarrow \dfrac{\Lambda^\circ + n}{\Sigma^\circ + n}$ 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 $\Sigma^- + p \longrightarrow \dfrac{\Lambda^\circ + n}{\Sigma^\circ + n}$ 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 ± .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 $K^- + p \longrightarrow \dfrac{\pi^+ + \Sigma^-}{\pi^- + \Sigma^+}$ ratio between threshold and 200 $\dfrac{\text{Mev}}{c}$. 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 $\dfrac{\text{Mev}}{c}$. (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)	$\Gamma/2$	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.

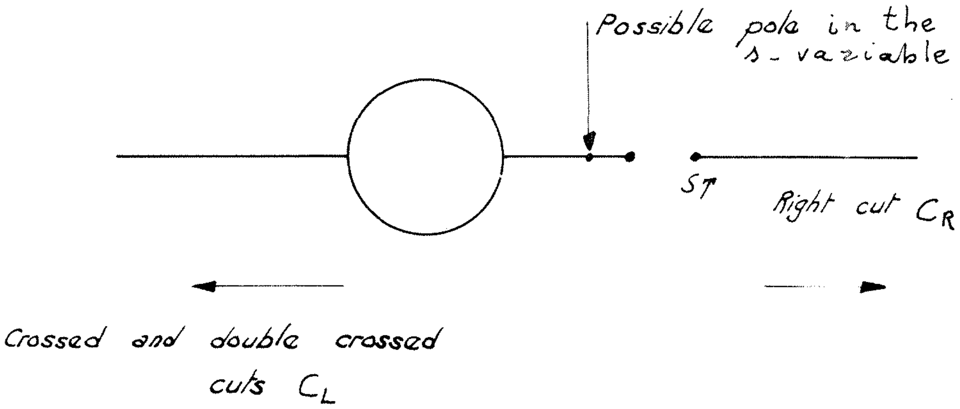
The object of this section is not to present these calculations but simply to review qualitatively the essential resonance formulation covering both πY and $\bar{K}N$ interactions. The pattern in all these cases is that of the first Chew-Low calculation of the 3,3 resonance rewritten in terms of the 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 $\underline{13}$, 1143, (1959) ; Phys. Rev. Letters $\underline{5}$, 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 :

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

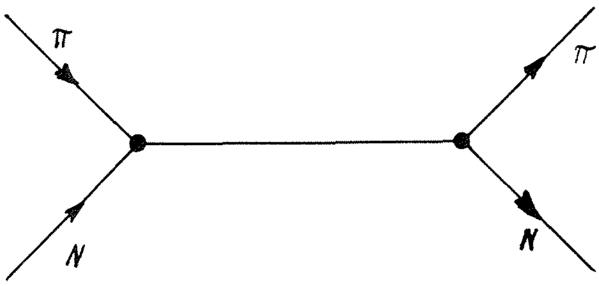
The condition for a resonance is that δ_l increases through $\left(n + \frac{1}{2}\right)\pi$ at $s = s_r$. Thus at the resonance energy :

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

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)	$g_{\pi\Lambda\Sigma} = g_{\pi NN}$ $g_{\pi\Lambda\Sigma} \neq g_{\pi\Sigma\Sigma} \neq 0$ $\delta = \frac{g_{\Sigma\Sigma}^2 - g_{\Lambda\Sigma}}{g_{\Sigma\Sigma}^2 + g_{\Lambda\Sigma}}$	$g_{\pi\Sigma\Sigma} = 0$
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

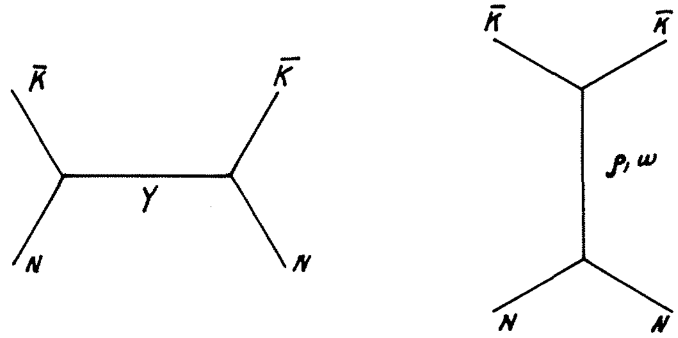
$$\begin{pmatrix} T_{\bar{K}N\to\bar{K}N} & T_{\pi\Sigma\to\bar{K}N} \\ T_{\bar{K}N\to\pi\Sigma} & T_{\pi\Sigma\to\pi\Sigma} \end{pmatrix}$$

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

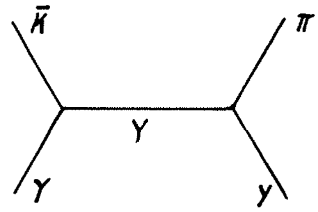
In a static approximation note that the direct Y pole for $\pi Y \longrightarrow \pi Y$, $\bar{K}N \longrightarrow \bar{K}N$ and $\bar{K}N \longrightarrow \pi Y$ appears in $J = \frac{3}{2}$ state if $p(\Lambda\Sigma) = +1$ (assuming with Dalitz that $p(KN\Lambda) = -1$) and in $J = \frac{1}{2}$ state 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 $\frac{R}{S_T - S_{\rho,\omega}}$.

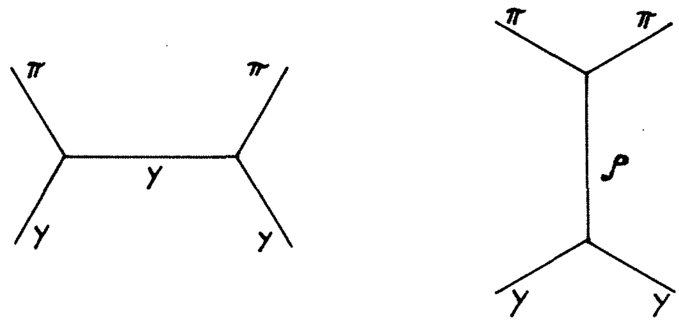
$\bar{K}N \rightarrow \bar{K}N$



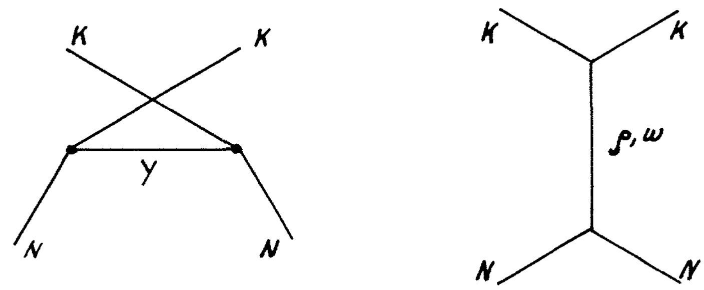
$\bar{K}N \rightarrow \pi Y$



$\pi Y \rightarrow \pi Y$



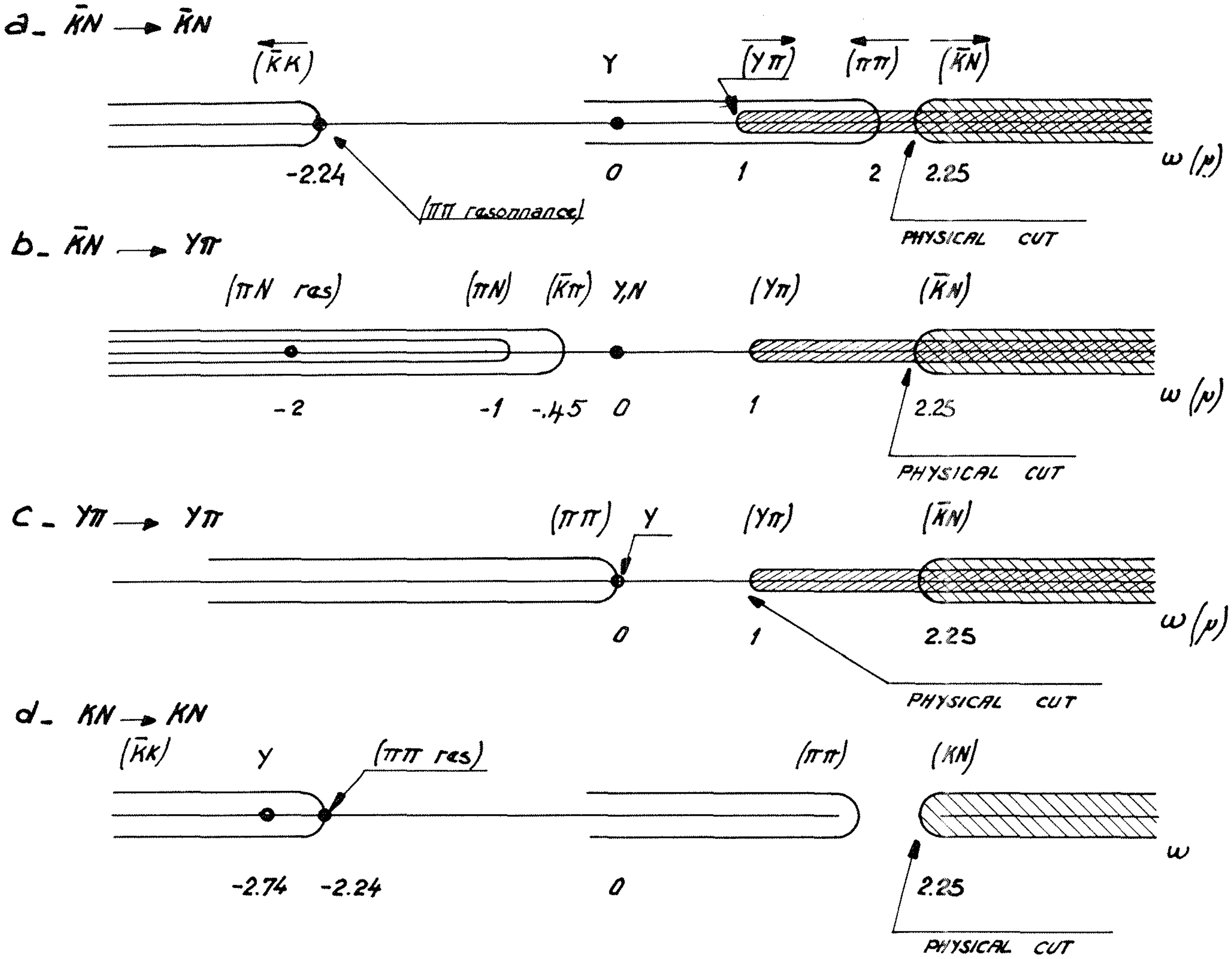
$KN \rightarrow KN$



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 :

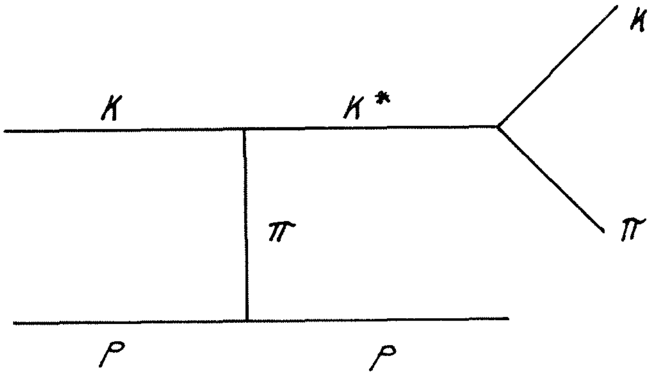
$$K^* \longrightarrow K + \gamma$$

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 :

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

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 :

$$\begin{aligned} \sigma &= 1.32 \text{ mb for } J = 1 \\ &= .105 \text{ mb for } J = 0 \end{aligned}$$

assuming $\Gamma/2$ = 8 Mev.

The experimental value is $\sim$ 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 :

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

J = 0 K^* would imply the outgoing mesons must both be in K_1^o or K_2^o 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 :

$$\begin{array}{lll} N^*(I = 3/2); & N^*\left(I = \frac{1}{2}\right); \; \Xi^*\left(I = \frac{1}{2}\right); \; \Xi^*(I = 3/2), \\ Y_1^*(I = 1) & ; \; Y_o^*(I = 0) \; ; \; Y_2^*(I = 2); \\ X^*(I = 1) & ; \; Z^*(S = 3, \; I = 1) \\ \quad (S = +1) \; ; & \end{array}$$

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

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

5 - BALL-FRAZER MECHANISM -

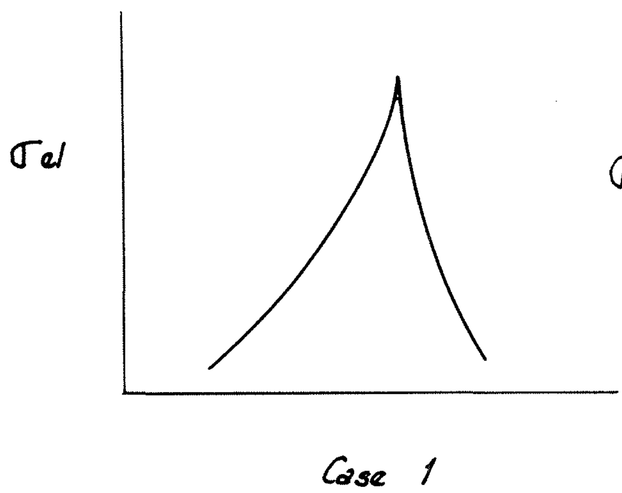
The discussion of K^* leads us naturally to the consideration of the highest known hyperon resonance at 1815 Mev. The most striking feature of this resonance is that it occurs at $K^* + N$ threshold in the $K^- + p$ channel.

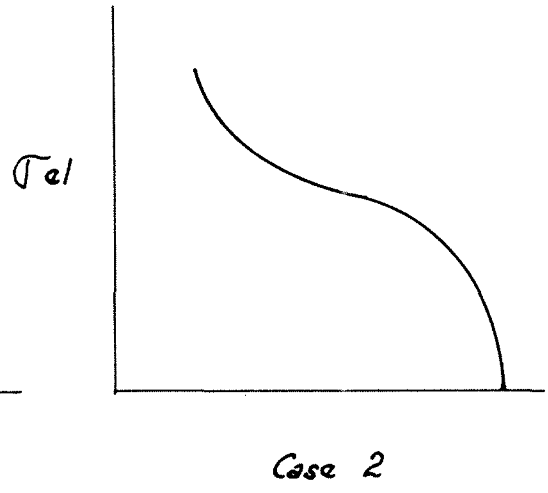
The existence of a resonance near an inelastic threshold has been noted before. The third pion-nucleon resonance $N^{(3)}$ occurs close to $N + \rho$, $N + \omega$ as well as $\Sigma + K$ thresholds (Baz, Kiev Conference, 1959), while $N^{(4)}$ the $I = 3/2$ resonance at 1900 Mev falls at $Y_1^* + K$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

CONCLUSIONS OF THE CONFERENCE

R. P. FEYNMAN

What can I say ?

I did not have any time during the conference to discover anything new on my own. Most of the talks have been summaries themselves. Almost everybody heard them and there is no point in going over it again. I also did not think that it was a good idea for me to tackle those particular points which I did not understand and make comments about them because then everybody would want the same opportunity to put so many questions at one time. I therefore dont have much to say. But I will talk a long time anyway.

I did not conceive my talk as a summary of the conference so much as a discussion of conclusions of the conference in some wider sense. What I want to do is to talk more about the flavour of the meeting. I want to ask what is most characteristic of the meeting - What new position are we in at the present time - What kind of things do we expect in the future ?

At each meeting it always seems to me that very little progress is made. Nevertheless, if you look ever any reasonable length of time, a few years say, you find a fantastic progress and it is hard to understand how that can happen at the same time that nothing is happening in anyone moment (zeno's paradox).

I think that it is something like the way clouds change in the sky - They gradually fade out here and build up there and if you look later it is different. What happens in a meeting is that certain things which were brought up in the last meeting as suggestions come into focus as realities. They drag along with them other things about which a great deal is discussed and which will become realities in focus at the next meeting.

The thing most characteristic of this meeting is the bringing into focus of the reality of a few resonances and the beginning a philosophy of such resonances. Of course every bump or wiggle in every curve now becomes an other resonance. There will then be the accomplishement at the next meeting, of having gotten rid of some of them, of having created a few more ones and, of course, the discussion of other new bumps which will go on to the meeting after that.

I will try to give a general picture of physics as it seems to me at present and will then discuss strong interaction in more detail.

It is clear to everyone today that physics is almost entirely in the hands of the experimenters. I think, nevertheless, that we should appreciate that theory is supposed to have a predictive value. What do we mean by prediction ? One wonders for example, how Herodotus could believe in the oracle of Delphi, in his time, as he was an intelligent man. What really happens is that each of the predictions of the oracle are in vague language and they become particularly clear when the event occurs afterwards, so you can see how it works. The high priests of Babylon used to predict things by looking at the liver of a sheep. And why ? Because in the complexity of the arrangement of the veins, interpreted correctly, they could tell what the future would be. It is that complexity, and the possibility of reinterpreting later the arrangement of the veins, that permits the power of the priests to be maintained. The diagrams which go with my name appear to me like the veins of a liver of a sheep. It is always possible to follow the right lines after the events.

If you try to test modern theory on its predictive value, you find that it is very weak. Of course, there are many sucesses. There is the success of the Chew-Low theory. There is the fact that a resonance in one reaction will be a resonance somehow in another reaction as illustrated at this meeting by the Electromagnetic form factor analysis. This is a sort of small amount of prediction. There are also the rules of strangeness and isotopic spin, and the coupling of the weak interactions as long as strangeness does not change but, otherwise, it is not too good. One might say that there is sometimes a succesful prediction of a certain new particle ; but like the Greeks who said everything, whatever we discover, the Greeks have discoevered it first, and whatever particle one would discover, someone, somewhere, has already predicted it.

When a new particle or a new fact is discovered, I notice that all the theorists do one of two things : they either form a group, or disperse.

I am familiar with a number of experimental physicists and they are sort of men of the earth. Therefore, I have always suspected that, one day, working far way from theorists, close to their big machines, they will get the idea of a new experiment ; an experiment which will test the oracle. They would like to see what would happen, just for the fun of it, if they falsly report that there exists a certain bump, or an oscillation in a certain curve, and see how the theorists predict it. I know these men so well that the moment I thought of that possibility I have honestly always been concerned that some day they will do just that. Then you can imagine how absurd the theoretical physicists would sound, making all these complicated calculations to demonstrate the existence of such a bump, while these fellows are laughing up their sleeves. For this reason, I have found myself almost incapable of making calculations of the type that most other people make. I am afraid they will catch me out !

With the existence of the resonances, which way will things go in future theories ? The first interesting or surprising thing, is that the resonances, the new states or particles, (which views I think, are more or less equivalent) whatever they are appear to be fairly narrow. I would think that, because of the strong interactions, all the resonances are wide.

In nuclei, we also have strong interactions and very narrow resonances. The existence of these resonances depends however, on the fact that the nucleus is complicated, and contains a large number of particles. In order to get an α particle out, you have to wait until all the neutrons and protons locally can form an α particle, which they only to once in a great while. The resonance is then narrow for any escape channel, and this comes from the complexity of the nuclei, in other words the large number of possible states.

It is then also possible that the narrowness of these new resonances or particles, is the result of a very large number of possible states, which are similar aside from the energy value. The large number of possible conditions (like K^-N, or $\Sigma^-\pi^\circ$, or Σ° to the ones found in Nuclei $\bar{K}\,K^+\,\pi^-$ etc.) in which such a given particle can be, makes it hard for it to find the particular combination into which it is supposed to disintegrate. I believe that this idea, in this particular application, is originally due to Heisenberg.

If we now have a resonance which is fairly narrow in width, it may be that it is a fundamental particle, and the Σ for instance, which is stable, may be a derived particle. The difference between one and the other is obviously only the energy value and, from a theoretical standpoint, it is not supposed to be deep in the physics that the lowest numerical value of the energy is The Thing. Another picture, then, might then be that these resonances are the fundamental particles and that the Λ, nucleons, Σ's, are derived particles, which happen to have a lower energy, so that they are stable. From the number of new particles being found here, the rash of new theories that we can expect, is very great !

We already see however a new kind of physics which may be in fact successful in a way. Because of the very large number of states involved, the amplitude to feed any particular term may be relatively small, so that, there may be a resurgence, with mild success, of a kind of perturbation theory : One could pick on a certain process, so that for a peripheral collision say, among the various diagrams, the transfer from one of the systems to the other has to be via a state with a given set of quantum numbers. But the amplitude to generate this particular state may be small for the same reason that resonances are narrow. If it is small enough for the particular channel, a kind of perturbation theory may be applicable. I suspect that the future theories of interactions will revert gradually back, until the dispersion theory has converted itself into essentially the perturbation theory, with intermediate states being resonances, and with perhaps a kind of Heitler damping modification (I have been helped by some remarks of Gell-Mann). I do not know however what will happen. I am just trying to guess.

Now, an interesting question comes up : when is a particle a fundamental one, and when is it a composite one ? For instance, every particle may be made of proton, neutron and Λ. The fundamental particles can as well be however the neutron, Σ' and Ξ and so on. One can start from nearly any set of particles, at first sight, and make others. The first time this question arose, so far as I know historically, was with the theory of the π-meson of Fermi-Yang. They assumed that a proton and an anti-neutron interact very strongly from an unknown force. (An obvious suggestion is a vector meson like the photon, but coupled to the nucleonic charge instead of the electric charge with a strong coupling, and a large rest mass). The binding energy is so great compared to the nucleon mass, that the mass of the pion is less than 2 M by a terrific amount. Fermi is the first one who asked the question : how can I distinguish experimentally between the possibility that the π is a fundamental particle or a derived particle in this sense ?

First, what does it mean ? In the old fashioneal spirit, the problem is to write the Lagrangian of the world and to deduce its consequences. But what field variable are to be put in it ? Do we have to start with operators for the nucleon, the proton and the π, or do we start with operators for the proton, the neutron, and an interaction theory, and get the π out ? Fermi tried to answer this question but I don't think it has ever been answered. I should remark that the kind of argument that has been made is something like this : suppose, for example, that something is held in a potential well and has got one bound state. If we go to very high energy scattering states, there is practically no effect of the potential. Now let us take the potential off and count how many states there are from zero energy up to a very high energy, in a box of very large size. Let us now turn the potential on gradually. The upper state will not move because it is not affected by the potential. All the lower states will have their phases shifted in some way, so that they fit into the box with a slightly different wave length. The total number of states will not be changed. This particular object, then, if it is one of these state, is really a bound object and not a new object. To find this, we would compare the number of states in the absence of potential, with the number of states for the excited particle not bound.

With such a theory is the Deuteron a bound object or a new particle ? Everybody would guess that the deuteron raises no question but with the π meson it is not so easy. Looking at proton neutron scattering would settle the deuteron question. The π problem might then also be solved by looking at proton antineutron scattering.

I would like however to propose a principle which, I suppose, will be found correct in the future : it will never be possible to tell which are the fundamental particles, and which are the derived ones. That is a fundamentally unanswerable question. If the theory ever does involve something like a Lagrangian (which I do not think it may necessarily do), with certain field operators in it which correspond to certain fundamental particles, someone else will be able to write as well an other Lagrangian, which will have different fields in it which correspond to other basic particles, but, nevertheless all the physical conclusions derived from either Lagrangian will be the same. Whitout the Lagrangian picture, the idea is : no matter how the physics is written, ultimately, it will be of such a nature that it will never be possible to tell which are the fundamental and which are the derived particles among the strongly interacting ones. That is the proposition.

The only time I ever found this idea of any use was in the theory of β decay. When I was trying to figure out the law of β decay I supposed that this interaction would be most simply represented using two-component wave functions even for the proton and neutron. The two-component equation is simple enough for a particle, with an electrodynamic vector coupling. But the π meson nucleon interaction (a γ_5 or $\gamma_5 \not{q}$ interaction) is quite difficult to express in a two-component wave equation. Therefore, temporarely I considered the π to be a Fermi-Yang derived particle, held together by some neutral vector field. It tu ns out then that with such a theory the β decay rate of the π^+ into the π° can be exactly calculated from the Nucleon β decay rate. If the π is a fundamental particle in itself it is, a priori, unlikely that one would find that particular rate. But suppose we take the following point of view : Nature is so constituted that it will be impossible to tell which of the two theories is right. If the $\pi^+ \longrightarrow \pi^\circ$ rate did not agree with the rate calculated from the Yang-Fermi model, we could conclude this model is wrong. Since I assume such a decision is impossible I predict the β decay rate of the π^+ into the π° will be found to be the rate predicted by the Yang-Fermi model, that of the conserved Vector Current formula.

To maintain this view one would have to consider that any argument such as the phase-shift argument mentioned earlier is wrong in some way. This argument has been made in term of a potential. This is too simple and, because of inelastic scattering the phase might well be never definite ; or they might change logarithmicaly as we go up in energy. I do not know how to develop the argument in a relativistic theory but if I knew I would make the hypothesis that Nature is constituted so that one can't make that test.

I would also mention the point of view of Chew which he expressed at La Jolla. I cannot explain it very well but it is consonant with the idea that you will never be able to tell which particle are the "true" ones. He sets up a set of interconnected dispersion theory equations. No one particle is better than the other. It is then necessary to add something to this hierarchy of equations in order to limit them in some way. Chew proposed that it is built in such a way that somehow it is "as close to trouble" as it could be. I do not know how much success he will have.

Other forms of theoretical attacks have been presented here by Pr. Salam. These theories start out with a symmetry among all the particles or among a set of particles which are taken to be fundamental. The postulated symmetries are however not true in fact. Let us consider first the strongly interacting particles. Along these lines one would say that the neutron, the proton, the Λ, the Σ and the Ξ are all equal ! At this stage everybody is in agreement that it is really a beautiful theory. The only trouble is that it is not true. Nature is really unsymmetrical. There is a problem how to put the asymmetry in, (the difference in masses for instance). The theory then always get dirty.

We use to laugh at the Greeks who claimed that the planets had to go in circles because it was a perfect figure. If they were talking in the modern times they would use group theoretic arguments and would imply that from the point of view of the planet the sun looks always the same, or that we have invariance under a combined time displacement and rotation. But the planets do not go in circles ! Nature is not "symmetrical" and the question is why not ? Let us consider the following possibility. Let us suppose that, with regards to these particles, Nature is really unsymmetrical in the beginning and that any near symmetry that we see is due to the complexity of all the interactions. To go to the extreme let us also consider the weak interactions and take the extreme view that in the beginning there is not any, even approximate, parity conservation in the fundamental law but that, in the complicated interactions of everything, it all averages out somehow that the parity conservation law is almost perfect. Going back to this planet picture we would say that the tidal forces make the orbit look more and more like a circle though it really isn't a circle.

This is my point of view. There is an other one, mainly due to Heisenberg which is far deeper and more likely. According to it, things are symmetrical in a certain sense and get unsymmetrical in a very interesting way. To give a hint how this works, let me consider the simple example of a ferromagnet. If you write the Hamiltonian of the whole ferromagnetic system, disregarding the lattice structure, it does not make any difference in which direction you quantize the spins. The final state however has the lowest energy when all the spins are quantized parallel. An excited state, obtained when you turn over the spin of one of the particles, say, is then unsymmetrical. The energy of the excited states depends on the orientation of this spin. Although the original writing of the interactions of all the particles has no axis of symmetry in it, the lowest state has lost this symmetry. Such a picture is then applied to the vacuum, supposed to be the lowest state of the world. The total spin of the ferromagnet analog can be in any one of a million directions. We then get a degeneracy. It points however in one direction because the whole world is polarized and we understand this way that things which, at first sight, should have the same energy may not, in fact, have the same energy.

To give an example, I would like to remind you of the following argument against the theories that state that for every conserved quantity there is an associate vector meson with a conservation law analogous to electrodynamics. The conservation law is related with gauge invariance and these new mesons should then also be gauge invariant. Since gauge invariance is usually believed to imply that the mass is zero, the first prediction of theses theories would be that all these mesons have zero mass, a point which is disregarded. It is however possible that the lowest state of the system considered is of such a polarized nature, that, when one excites an extra meson it may have a mass. If one considers a photon and renormalizes its propagator for the electromagnetic interactions, there is a renormalization of charge but, according to gauge invariance, no renormalization of the mass. Schwinger pointed out to me however, that one can use gauge invariance to prove that the mass of the real photon is equal to zero, only if one assumes that in the complete dressed photon, there is a finite amplitude to find the undressed one. If the dressed meson wave function is now orthogonal to the undressed one it may well be that such a vector meson has a non zero mass. It is interesting to see if that can ever be done. It would make Heisenberg idea work in a specific way in a particular case. This would explain the mass of the vector mesons.

There is now one point that should be made with regards to the theory of global symmetry. It is unsymmetrical as far as the masses of the strongly interacting particles are concerned. (The symmetry of isotopic spin is however almost correct and it is interesting to wander where such a symmetry comes from. It is important because that may be a real symmetry). A theory like global symmetry goes on to say that even including strangeness, things are nearly symmetrical. People who start with a symmetrical theory, with a general rule, are always beset by the fact that the lack of symmetry of Nature is not as small as they would like. Since one cannot calculate anything with strong coupling there is no way to honestly compare with experiment. One notices of course that the symmetry is in fact somewhat broken but only qualitative arguments can be developed. All these theories have stayed with us such a long time only because of the difficulty of deriving predictions. There is however a very close symmetry about which we have no idea except perhaps the view of Heisenberg of the polarized vacuum. This is the extremely close symmetry between neutron and proton which is from the energetic stand point so accurate. From all the strong interactions one hardly notices the deviation and yet, they are completely different ! One is charged, the other is neutral. They are as different as night and day, they are as different as a neutron and a proton ! It is very interesting to have nature have this very close symmetry and then this lopsided thing put on top of it. It remains an absolute mystery.

Let us now consider electrodynamics. It seems to be working fine and the reports in this conference have checked it again. Various tests have shown that it is correct to almost one GEV of reciprocal momentum. We must however remember that there are several things that we still do not understand in this field. Why for instance is it tied on in an asymmetrical way ? Why is the coupling constant 1/137 and why are all the charges the same ? These are things which we do not understand. The rule connecting charge, isotopic spin and strangeness is also strange !

I would also like to talk about another principle which has never been stated in a complete way and which, I think, is originally due to Gell-Mann. This is the principle of minimal electromagnetic coupling. It is usually possible to tell intituively what is meant by saying a given coupling is or is not in accordance with this principle even though this principle has not been formulated precisely so far. Let us consider an example. The proton is observed to have an anomalous magnetic moment but according to this principle we would assume that this is entirely due to the meson cloud. We would not suppose that part of the anomalous magnetic moment of the proton is due to a real Pauli term. In other words there is no anomalous couplings that are not due to complexity in the interactions. In the fundamental Lagrangian of the world, to take again this picture, the coupling is as simple as it could be. Such a principle "works" in the sense that electrons and muons are both coupled minimaly. There is no other electromagnetic coupling than the one we get when replacing ∂_μ by $\partial_\mu - A_\mu$ wherever it appears in the Lagrangian, written without any electromagnetic interaction.

Such a principle can however be presented only vaguely so far, since this fundamental Lagrangian picture is not reliable. To put it in a nutshell we would like to say that there is no electromagnetic coupling except through the intermediary of some other particles themselves being coupled in a simple manner.

Let us now consider the weak interactions. We then meet the leptons. From the point of view of strong interaction the mere existence of the leptons is a miracle ! Why did nature bother to make electrons and neutrinos and, furthermore, to make the μ. What is the μ ? We can summarized all its properties up to now in saying that it satisfies the Dirac equation with its own mass. Why has the μ got this mass ? How would you predict it à la oracle of Delphi once you know it ? Of course, since it is known to 5 decimal places you note a great quiet-ness on this point.

The weak interactions have been discussed here elegantly by Lee. I then lean on this discussion. The interactions considered fits nicely and people are tempted not to worry about a number of points. Everybody is happy about parity violation though parity is actually violated in a very queer way. Whereas we would have expected that, because parity is almost conserved, a small violation of parity conservation would mean what there is a small amplitude to go from a state of given parity to a state of opposite parity, we find that whenever there is a weak interaction it links a given state to two states of opposite parities with equal amplitude. One could built such a theory with two component wave functions, somebody else could use Chirality ! I know that it makes no difference, that it just means to multiply the wave function by $1 + \gamma_5$, but I do not understand it.

The fact that the μ can be replaced everywhere in the interactions by an electron make these particles equivalent. The same value found for the coupling of the muon neutrino, electron neutrino and proton neutron pairs is also usually "explained" by the hypothesis of universal interaction but all this does not settle the question. If it really did one would be clearer about how the strange particles are coupled.

Another interesting problem arises from the fact that the nucleon wave function has to be multiplied by $1 + \gamma_5$; it is not the antinucleon wave function. In other words why is the proton a particle and not an anti-particle ? The rule could as well have been, instead our actual point of view, that the $1 + \gamma_5$ term has to be associated with the negative particle in the neutral-charged pair. All these points have been sources of worries when the theory was born, but people do not seem to worry about it any more now that the theory fits.

When the strangeness does not change, everything fits extremely well, including the disintegration of the π, but I must say we just don't understand the weak interactions at all.

With respect to the π decay I would like to make a few comments. If one assumes that the π meson is coupled through a pseudo vector coupling instead of the pseudo scalar coupling usually considered, then one finds (adding, however, another assumption) the same result for the decay rate as derived previously by Goldberger and Treiman for pseudo-scalar coupling but incorrectly because of all the terms neglected. With a pseudo-vector coupling such terms are not neglected and the answer still agrees with experiment. For the π Nucleon effective interaction Hamiltonian we would then write $\gamma_5 \not{\phi}$ instead of γ_5. This has quite different results in perturbation theory. A γ_5 alone gives a large S wave scattering. It is actually very small and one has to claim, though it has never been proved, that the S wave scattering somehow damps itself out. With a $\gamma_5 \not{\phi}$ coupling on the contrary there is very little S wave scattering in lowest order. Beside this there are some predictions to make. Such a coupling implies that the interaction goes to zero when the frequency and the momentum of the π go to zero, which is not necessarily the case in the usual γ_5 theory. There is then an asymptotic place where you can test the theory. I found that it does not quite work. The S wave is very small but not small enough. As mentioned by Pr Matthews some residual S-wave scattering is however supposed to come from one of the virtual states involving one of the new resonances. Such virtual states were of course not included. I then still think that it is consistent to assume that the actual coupling is equivalent to $\gamma_5 \not{\phi}$. I know it is not renormalizable but there is no theorem that says that all theories must be renormalizable. There is only a theorem that says that all the theories are wrong at high energy !

Let us now turn back to weak interactions and consider the strangeness changing processes. When one has such success with the strangeness not changing processes, it is tantalizing to try to generalize it. Of course both types of processes have to do with each other : parity is violated, the rates have the right order of magnitude. It is the same phenomena. One then assumes that the coupling are the same ; the proton lambda pair enters the interaction as the electron neutrino one. I know that such an assumption incidently desagree with experiment by a factor 5 but I am not too worried about that at the moment. There is also a possible coupling with the proton -Σ° pair and from a certain point of view, take global symmetry say, the $\frac{\Lambda^\circ + \Sigma^\circ}{\sqrt{2}}$ combination is one of the particles rather than either the Λ° or Σ°. If this particular "particle" then enters the coupling alone we would then get an extra factor 0.5 for the rate. If we take an other kind of symmetry you would get other $\sqrt{2}$ factors. In addition these amplitudes are renormalized by strong interactions. In other words, we cannot make any prediction with regards to universality unless we can identify the particles somehow. This we do not know how to do. What is done is then to take some kind of symmetry and then see what the predictions are ; whether it looks equally beautiful for the weak decay coupling. In every case it fails in one way or an other quantitatively and I have never been able to get the right rules. One of the very difficult points is the slow rate of K^+ into $\pi^\circ e^+$ and ν. At present, experiments are telling us the amplitudes of particular reactions and we know that there is some kind of complexity. We think that it is likely to be complexity among the strongly interacting particles and simplicity in β decay but we cannot unravel it at the moment.

Of course getting a symmetry helps a lot, but we now have some evidence from Fry and his collaborators, that, when strangeness increases, charge does not necessarily increase. We will see if that is substantiated with more statistics. There are other rules observed among weak decays. For instance, when the leptons are not involved, the isotopic change by 1/2 is, by all odds, the dominant process. This is rather mysterious so far. Finding the correct law of the weak interactions is a fascinating problem, because it works already so well for the strangeness not changing processes. I hope that future experiments will continue to give us information and that tell us which of our "simplest ideas" are wrong.

This is a kind of general summary of the present position and how I think things will go, more or less, in the future.

IMP. LOUIS-JEAN — GAP

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