SLAC-PUB 355
September 1967

WHAT HAVE WE LEARNED ABOUT ELEMENTARY PARTICLES FROM PHOTON AND ELECTRON INTERACTIONS?*

by

S. D. Drell

Stanford Linear Accelerator Center
Stanford University, Stanford, California

(Paper delivered at the International Symposium on Electron and Photon Interactions at High Energies, SLAC, September 1967)

*Work supported by the U. S. Atomic Energy Commission.

In trying to write this report I find myself in limbo. I am announced as speaking on "What We Have Learned" at the opening of this conference prior to the presentation of many new and more precise experimental results--and in particular before the disclosure of first results from SLAC at significantly higher energies and momentum transfers than heretofore available. My dilemna is that I am aware of these results and therefore know to discard as rubbish much of what I might have said otherwise. On the other hand, I am also unable to present wise or even reasonable interpretations or analyses since no new data has yet been presented. So, I have nothing to say and for the moment I sit in limbo awaiting to be annointed by all the new data during the subsequent sessions of this symposium and awaiting to follow my experimental guides and, to quote from Dante's Divine Comedy[1] [Inferno; Canto 34], to be "issued out, again to see the stars." I would much prefer a divining rod to a divine comedy, however.

I will attempt no completeness of coverage. That comes later in the conference. Several problem areas have been and continue to be of interest to me, and to these I will speak.

Finally, I can humbly advise you to listen closely because the emphasis in my remarks will plainly provide clues as to some of the real live new results you'll learn of starting tomorrow.

Photon and electron interactions have several very charming features that can be exploited: an electromagnetic current enters into a strong interaction blob very discretely and gently via a known conserved current j_μ. We can thus probe for structure, correlations, resonances, and selection rules by introducing a known perturbation and looking for and analyzing the spectrum of responses. Moreover, the small strength of the electromagnetic coupling $\alpha = \frac{1}{137}$ invites application of perturbation calculations and generally it suffices to work to lowest order in the electromagnetic interaction, except for radiative corrections as Dr. Yennie will discuss in the following report.

Studies of electron scattering from nuclei are a beautiful example of this very point, and we will learn much more about this from Dr. Walecka later on. The form factor is measured by elastic scattering

$$F(q) = \int d^3r\, \rho(r)\, e^{i\vec{q}\cdot\vec{r}} \tag{1}$$

and from its q dependence and diffraction minima we learn about nuclear radii and surfaces and first clues as to the forces. Furthermore, high resolution studies of electric and magnetic transitions to individual excited states have taught us individual transition matrix elements, and since we know the operator j_μ we learn about the states instead of vice versa. Then, too, sum rules can be constructed on the basis of very general relations and quantum mechanical principles and depending only on general properties of the interacting systems. Their primary value lies precisely in their freedom from details so that they are valuable first steps into the study of new fields. A familiar one for inelastic electron scattering expresses the electron scattering cross section for electrons of all energies to emerge after a fixed momentum transfer q in a coulomb field of strength Ze.

In the limit of large q^2 it reads

$$\lim_{q^2 \to \infty} \frac{d\sigma}{dq^2} = \frac{4\pi\alpha^2}{q^4} Z \; . \tag{2}$$

Turning to scattering from the proton what have we learned about so-called elementary particle physics from the analogous studies?

The pioneering measurements of Hofstadter[2] in the very beginning showed that the proton was a rather fat, diffuse charge and current distribution with a root mean square radius of ~ 0.8f. Nambu[3] first recognized the need for the existence of an isoscalar vector meson resonance, the ω^0, and Frazer and Fulco[4] subsequently showed in detail the case for the isovector ρ^0 in order to provide a theoretical basis for the form factor behavior--indeed the early form factor work led to predictions that there existed vector mesons of sub-nucleonic mass.

But how far have we come since then?

It is clear from the extensive very beautiful data[5] accumulated over the past decade starting with the original Hofstadter measurements that the form factors of the nucleon fall off rapidly at large q^2 and, at least until tomorrow, we know of two popular and adequate fits to $G_M^p(q^2)$:

$$\left[G_M^p(q^2)\right]^2 \sim \begin{cases} \left[\dfrac{1}{\left\{1 + q^2/.71\ \mathrm{GeV}^2\right\}^2}\right]^2 \\[2ex] \left[\sqrt{4.6q_{\mathrm{GeV}}}\ e^{-4.6q_{\mathrm{GeV}}} + 8\times 10^{-3}\sqrt{2q_{\mathrm{GeV}}}\ e^{-2q_{\mathrm{GeV}}}\right] \\ \qquad\qquad\qquad\qquad\qquad\qquad q \gg 1 \end{cases} \tag{3}$$

Both the dipole fit[6] and the refinement by Schopper[7] of the Wu-Yang[8] exponential decay do well out to $q^2 \sim 10\ \mathrm{GeV}^2$ but we will learn more tomorrow[9] to considerably larger values. I have heard no violent rumors of dips or bumps in G at the larger q^2 values to which it has now been measured--nor of any emergence of qualitatively different new features such as a hard core, so perhaps it makes some sense for us to ask how has theory fared with this type of rapidly and smoothly falling behavior.

The canonical very simple starting point has been for the past decade a dispersion approach[10] with the form factor represented as a sum of Yukawa-like terms [$q^2 > 0$ for scattering measurements]

$$G(q^2) = \int_{4m_\pi^2}^{\infty} \frac{\rho(\sigma^2)d\sigma^2}{\sigma^2 + q^2} \tag{4}$$

with the spectral amplitude $\rho(\sigma^2)$ describing the exchange of one or of a few of the neutral vector mesons or resonant enhancements from the electromagnetic current to the proton line in Fig. 1. Each resonance contributes a bump to $\rho(\sigma^2)$ at its mass $\sigma^2 = M_r^2$, and we write in the approximation of narrow resonances

$$G(q^2) \simeq \sum_{\text{resonance}} \frac{\rho_r(M_r^2)}{q^2 + M_r^2}$$

Evidently a cancellation must be contrived in order to give

$$\int d\sigma^2 \rho(\sigma^2) \approx \sum_r \rho_r(M_r^2) = 0 \tag{5}$$

and lead to an asymptotic behavior decreasing as $1/q^4$ or faster. No one has succeeded, however, using only the observed, known vector masses to fit the large q^2 behavior, the radius

$$< R^2 > \equiv 6 \int \frac{\rho(\sigma^2) d\sigma^2}{\sigma^4} = (.813f)^2 \quad , \tag{6}$$

and the scaling law extended by very recent DESY data[11] out to $q^2 = 3$ GeV^2:

$$\frac{G_{EP}}{\frac{1}{\mu_p} G_{MP}} = .90 \pm .24 \ . \tag{7}$$

The most elementary way to achieve a $1/q^4$ fall off is to provide a theoretical argument which orders you to multiply a Yukawa form for the vector meson propagator by a similar Yukawa form for the form factor describing the vector coupling to the nucleon line (see Fig. 2) perhaps via quark wave functions or by constructing a Lagrangian formalism coupling the electromagnetic current directly to the amplitude of the vector meson potential. An elaborate analysis of the latter type has been given by Kroll, Lee, and Zumino[12] who provide an explicit statement of what vector dominance means in a Lagrangian field theory. They assume that the entire hadronic electromagnetic current operator is identical to the vector meson field amplitude and derive thereby

$$G^{\gamma}(q^2) = \frac{m_v^{\ 2}}{m_v^{\ 2} + q^2} \, G^{V}(q^2) \tag{8}$$

where G^V is the vector meson-nucleon form factor and when subjected to a dispersion analysis would be treated as in Eq. (4). $G^V(q^2)$ contains contributions from all but the vector meson pole itself at $-q^2 = m_V^2$ and a priori there is no reason for its value at large q^2 to decrease as $1/q^2$. In fact, the general dispersion approach has a very severe limitation when applied to a study of the behavior at large q^2. This is because the dispersion integral converges only very slowly with all contributions being essentially equally weighted in Eq. (4) up to large masses $\sigma^2 \sim q^2$. In contrast, the mean square radius calculation of Eq. (6) has a $1/\sigma^4$ convergence factor to enhance the low lying resonance terms. It may be reasonable to assume that only a few low lying resonance contributions dominate in Eq. (6) for $< R^2 >$ but it is certainly an extravagant optimism to extend that same assumption to calculating Eq. (4) for large values of $q^2 \gtrsim 10$ BeV^2.

Experience has proved that theorists often find a very useful guiding light when most desperately needed by returning to what Gell-Mann has referred to as their theoretical laboratory of the Schrodinger equation with an interaction described by a superposition of Yukawa potentials. With this assumption it has proved possible to derive double dispersion relations[13]--or the Mandelstam representation--as well as Regge behavior[14] in potential theory to gain better insight into the relativistic problem. It may not be so bad an idea to turn here also in search of more light on the form factor problem. We usually shy away because the proton is not a loosely bound system as is the deuteron or a heavy nucleus. In the latter cases our intuition has a comfortable graphical representation and dispersion theoretic relativistic sheen in the notion of the anomalous threshold and it is respectable to view, and calculate, the deuteron as a loosely bound system of neutron plus proton as illustrated in the reduced graph of Fig. 3. Not so for the proton--is it a heavy core plus a meson, or several bound quarks?

Nevertheless what can we learn in our private Galilean lab. For a spinless system in an s state, Eq. (1) can be reduced to a radial integral and with the good smooth properties of the superposition of Yukawa potentials

$$V(r) = \int_{a_{min}}^{\infty} \sigma(a) \frac{e^{-ar}}{r} \, da \tag{9}$$

we can invoke the Riemann-Lebesgue lemma to find asymptotically[15]

$$\begin{aligned} F(q) &= \frac{4\pi}{q} \int_{0}^{\infty} \sin qr \, (r\rho(r))dr \\ &\underset{q \to \infty}{\sim} \frac{g(0)}{q^2} - \frac{g''(0)}{q^4} + \frac{g^{iv}(0)}{q^6} + \ldots ; \end{aligned} \tag{10}$$

where $g(r) \equiv r\rho(r)$; i.e., the wave function behavior at the origin $r \to 0$ determines the limiting form factor behavior at $q \to \infty$. We all recall that, for the class of potential represented by Eq. (9) with $\int \sigma(a)da$ finite, $\psi(r) \sim$ const as $r \to 0$. Therefore automatically $g(0) = 0$ and the form factor falls off as rapidly as $1/q^4$ if $\int \sigma(a)da$ exists, i.e., providing there is no stronger than a $1/r$ singularity in the growth of the potential at the origin. With this behavior we have what is now called a superconvergence relation for the electromagnetic form factor, Eq. 5.

In fact, the coefficient of $1/q^4$ is proportional to $\int \sigma(a)da$ and if this were to vanish we might find a super-superconvergent behavior. The point I want to make is simply that: from the point of view of the Schrodinger equation, a fall off with $1/q^4$ or faster[16] seems most natural, and I wish I understood where it gets lost on the way to the dispersion relation which makes a fall off faster than $1/q^2$ look like an accidental cancellation. Clearly we have to develop more insight in this problem!! Why does the natural and--till tomorrow--observed $1/q^4$ or more rapid fall off lose out in a relativistic dispersion approach?

There is another ashamedly primitive and successful use of the Schrodinger equation and that is to the mean square radius calculation which probes the outer edges, not the inner reaches of the wave function. In application to the proton we are without the anomolous threshold to dignify the effort; nevertheless note in analogy with the deuteron, for which

$$< R^2 >_D \; = \; \frac{1}{2M_{red}\,\epsilon_{Bind}} \; \approx \; (3f)^2$$

where $M_{red} = M_p = \frac{1}{2} M_D$ is the reduced mass and $\epsilon_{Bind} \sim 2.2$ MeV is the deuteron binding energy, we have for the proton

$$< R^2 >_p \; \approx \; \frac{1}{2\mu\;(200\text{ MeV})} \; \approx \; (.8f)^2 \qquad (11)$$

where we have arbitrarily set the binding energy of the proton to 200 MeV, a one significant figure for the mean photodisintegration energy between the 140 MeV threshold and the 300 MeV resonance. This is so reasonable a result as to make you wonder if there isn't a message in it somewhere--some call back to simpler, low-energy, non-relativistic ideas!

Encouraged by this, we look for a suitably dignified relativistic dispersion approach that does no violence to basic field theoretic notions and does, however, emphasize the intuitive and low energy features of the interaction. As we remarked earlier, the traditional dispersion approach of continuing in the photon mass as in Fig. 4 converges poorly for large values of q^2. Thus approximation of the absorptive part by a few low-lying resonances is an unreasonable one. However, there is another route which offers more promise and that is the one developed first by Bincer[17] who formulated the dispersion relation of the electromagnetic vertex as a function of the proton mass.

In this approach the appropriate form factor is expressed as

$$G(q^2) = \frac{1}{\pi} \int \frac{dW^2}{W^2 - M^2} \ \mathrm{Im}\ G(W^2, q^2) \tag{12}$$

where Im G is the amplitude for a virtual photon of mass q^2 to be absorbed by a nucleon and form a real intermediate state of total mass W which can then couple to an off-shell proton of the same mass as in Fig. 4. To the extent that this absorptive amplitude is dominated by its low mass contributions,[18] $W \sim M$, we can approximate it by the threshold photopion production amplitude times the nucleon pion coupling strength. For real photons the exact low energy behavior of the photopion production amplitude is known and is given by the Kroll-Ruderman theorem.[19]

Applying this idea earlier to the calculation of the nucleon g-2 value, Pagels[20] and I found that both the isovector character of the nucleon moment and its approximate numerical value were reproduced fairly well when we retained only the contribution to the absorptive amplitude between $M \lesssim W \lesssim 1.5M$ and used the threshold theorems. The usual grief which befalls the perturbation calculations were found to be in the high mass contributions $1.5M < W < \infty$ which the perturbation approximation severely distort. This threshold dominance view also reproduces second and fourth order electron g-2 values rather well and has made a definite prediction on the α^3 contribution, since confirmed and refined by more detailed studies of Parsons.[21] In this application, the low energy theorem for Compton scattering--i.e., the Thomson limit plus the Gell-Mann, Goldberger, and Low[22] result for the term linear in the frequency--replaces the Kroll-Ruderman theorem.

Since the photon mass is only a parameter in this calculation and the low energy weighting represented by the deonominator in Eq. (12) is still present we may extend these studies to finite q^2 values. Dennis Silverman[23] and I have looked at this in recent weeks with encouraging initial results. One which I can report is the following: We apply Eq. (12) to calculate the magnetic moment form factors for the proton and neutron. The charge form factors are not calculated but assumed to be given by subtracted dispersion relations. We have then two relations between four unknown functions. With assumption that the neutron's electric form factor vanishes, $G_E^N(q^2) \equiv 0$ as is consistent over the measured range, and in the limit $q^2 \gg 4M^2$, we have found a scaling law for the Sachs[24] form factors

$$G_M^P(q^2) \sim -3G_M^N(q^2) \sim G_E^P(q^2)$$

Since the experimental information[25] supports a not too different scaling law up to $q^2 \leq 3M^2$

$$G_M^{\,P}(q^2) \sim -\tfrac{3}{2}G_M^{\,N}(q^2) \sim 3G_E^{\,P}(q^2)$$

we can be, and are, encouraged to work harder and press on. There is an experimental implication here that will be of great interest to learn from future data that we hope may be forthcoming before long: do the threshold electropion production amplitude and the nucleon form factors show the same behavior as a function of photon mass for large q^2 as predicted in this calculation?

Independent of any calculation, the behavior of the electropion production amplitude--and indeed of all inelastic electro-production amplitudes--is of fundamental interest. In particular we have a lien on the results of the measurement of the pion's electromagnetic form factor. This is accomplished by working in a particular kinematic region that emphasizes the contribution of the pion current contribution as done by Akerlof et al.[26] Dispersion theorists have "explained" the large nucleon radius of ~ 0.8f by making the pion "fat" via the ρ-resonance and it had better be so found! The results so far that quote a radius of 0.8 ± 0.1f are in this context very satisfying.

For a discussion of possible sum rules analogous to Eq. (2) as predicted by current algebra or quark models, we await tomorrow's report by Dr. Bjorken and experiments that remain to be reported at a future conference.

A totally different theoretical basis for the form factor discussions is the quark model for the Wu-Yang[8] relation

$$\left(\frac{d\sigma}{dt}\right)_{pp} = \left[\left(\frac{d\sigma}{dt}\right)_{q^2=0}\right]_{pp} \left[G(q)\right]^4 \tag{13}$$

where $\left(\frac{d\sigma}{dt}\right)_{pp}$ is the differential cross section for proton plus proton scattering at any high energy, and $G(q^2)$ is the electromagnetic form factor. Wu and Yang proposed this relation on the basis of the following physical idea: Form factors tell you how likely or probable it is for a proton to stay together when hit. If we assume that the way the proton shatters is independent of what hits it, then so is its probability to remain intact for a given momentum transfer q^2. This is given by $G^2(q^2)$ in electron proton scattering and by $[G(q^2)]^4$ in proton-proton elastic scattering since both protons must avoid being shattered. Eq. (13) is then a no parameter relation between experiments.

This view has been extended to the quark model by Kokedee and Van Hove[27] who have argued for small quarks forming a large hadron one can neglect rescattering corrections to the quark-quark interactions. They thereby derive the successful additivity rules for hadron processes as well as relation, Eq. (13). How well this model fares at the large q^2 values to which electron proton scattering measurements have now been extended, we will learn from Dr. Taylor's[9] report.

Let us turn next to two body reactions to review what we have learned from recent experiments on $\gamma + N \rightarrow \begin{cases} \pi + N \\ \rho + N \\ \gamma + N \end{cases}$. What clues do we see which may be joined with the full body of ideas and concepts that have proved of value in analyzing hadron reactions at high energies in order to provide some basis for an understanding of these reactions?

The most simple process at first glance is the photoproduction of ρ^{o}, ω, and φ mesons since zero strangeness neutral vector mesons have quantum numbers in common with the photon. We may, therefore, think of them as being photoproduced with a forward diffraction peak--i.e., the photon materializes as one of its heavy vector brothers with the same internal quantum numbers. The physical idea is simple. The most likely event to occur when a photon impinges on a nuclear target is inelastic meson production; diffraction scattering is just the shadow of these inelastic channels and the forward peak is the statement of coherence among them for each inelastic channel amplitude f_{an} to return to the same initial state a:

$$\mathrm{Im}\, f_{aa} = \sum_n f_{an} f^*_{na} \delta^4(p_n - p_a) = \sum_n \left| f_{an} \right|^2 \delta^4(p_n - p_a)$$

Since the photon and the ρ^{o} have quantum numbers in common (for the isovector part of the photon) and since the ρ mass leads to a negligibly small longitudinal momentum transfer of $\frac{m_\rho^2}{2k} \sim$ 50 MeV corresponding to a coherence length of 4f at k = 5 BeV, we may expect the classical diffraction character to be evident, and indeed it is up to 6 GeV, the present limit of measurements[28]

$$\left(\frac{d\sigma_d}{d\Omega}\right)_{0^\circ} \propto k_\gamma^{\,2} \; ;$$

$$\frac{d\sigma_d}{dt} \propto e^{+at} \quad ; \quad (t \equiv -q^2 < 0 \text{ for scattering}) \tag{14}$$

$$\text{and } \sigma_d = \text{const} \approx 15\text{-}20 \ \mu b \ .$$

These results are in accord with theoretical anticipations and with simple models.[29] A general and very useful way to describe them is by a vector dominance model that couples photons directly with the neutral vector mesons[30] of zero strangeness as in Fig. 5. Photoproduction of neutral vector mesons is then directly proportional to hadron diffraction amplitudes. With these models we can relate different amplitudes to each other in terms of a single amplitude coupling photons to the vector mesons and thereby construct direct relations between observable processes that can be checked against experiment.

Thus the pion electromagnetic form factor is computed from the graph of Fig. 6. If we denote the coupling constant of the photon to the ρ^0 meson by $\frac{e}{f_\rho} m_\rho^{\,2}$ and the coupling strength of the ρ to the pion current by $g_{\rho\pi\pi}$ vector dominance tells us $f_\rho \cong g_{\rho\pi\pi}$ or

$$\frac{f_\rho^{\,2}}{4\pi} \cong \frac{g_{\rho\pi\pi}^2}{4\pi} = \frac{12\Gamma(\rho\to\pi\pi)}{m_\rho[1-4\mu^2/m_\rho^{\,2}]^{\frac{3}{2}}} \approx 2.5 \tag{15}$$

where $\Gamma(\rho\to\pi\pi) \sim 130$ MeV is the ρ width. The same idea applied to the electron (or muon) decay branching ratio of the ρ as in Fig. 7 gives[31]

$$\Gamma(\rho\to e\bar{e}) \cong \frac{1}{3}\,\frac{\alpha^2}{[f_\rho{}^2/4\pi]}\,m_\rho$$
$$\simeq \left(\frac{2.5}{[f_\rho{}^2/4\pi]}\right)\quad 4\times 10^{-5}\,\Gamma(\rho\to\pi\pi) \tag{16}$$

This is in agreement with the reported branching ratios to within a factor of 2 if we use the value given by Eq. (15) for the coupling strength and supports the extrapolation of vector dominance between the photon to the ρ pole. Using the optical theorem we can relate the forward diffraction cross section to the total ρ-nucleon cross section, $\sigma_{\rho N}$.

$$\left(\frac{d\sigma_{\gamma\rho}}{dt}\right)_{0^\circ} = \frac{\alpha}{[f_\rho{}^2/4\pi]}\;\frac{1}{16\pi}\,[\sigma_{\rho N}]^2 \tag{17}$$

With $\left(\frac{d\sigma_{\gamma\rho}}{dt}\right)_{0^\circ} \simeq 150\ \mu b/\mathrm{GeV}^2$ and $f_\rho{}^2/4\pi \simeq 2$, we derive from Eq. (17) $\sigma_{\rho N} \simeq 30$ mb in reasonable accord with the quark model or SU_6 prediction that $\sigma_{\rho N} \simeq \frac{2}{3}\,\sigma_{NN} \simeq \sigma_{\pi N} \sim 25$ mb. These and other relations will become more precise with further data at higher energies, and I look forward to more accurate comparisons later in this conference.[32]

For further consistency of the diffraction or vector dominance model, we turn to the A dependence of forward ρ^{o} photoproduction in complex nuclei. The forward amplitude for production on the individual nucleons is attenuated as the wave propagates toward the edge of the nuclear matter sphere. From the A dependence of the forward production cross sections in Be, C, $A\ell$, Cu, Ag, and Pb, the nuclear mean free path can be determined.[33] The resulting value[34] of $\sigma_{\rho N} = 31 \pm 3$ MeV is in comfortable agreement with the prediction quoted above, as is the measured nuclear density parameter; you will hear about these results in more detail from Dr. Pipkin's and Dr. Ting's reports.

What we have learned then is that diffraction production has met with good qualitative successes thus far in application to the ρ^{o}. Less well established is the evidence for ω^{o} diffraction production which is predicted to be smaller by a factor of 1/9 on the basis of the usual assignment of the photon as an SU_3 octet and a U spin singlet together with the usual ω-φ mixing. Furthermore, the ω cross section is more difficult to measure due to the three body decay. Higher energies will help here since with its forward angle dependence on k_{γ}^{2} as in Eq. (14) the diffraction process should grow to a dominant role relative to the π^{o} exchange contribution.

We look forward eagerly to data on φ^{o} production at high energies because suppression of this reaction to a level of a few tenths of a microbarn as currently measured is an outstanding problem. Factors of 10-20 reduction from simple quark models are required at present, and more generally Harari[35] has shown one must strain seriously (if not violate) all theoretical models because it is so small a cross section.

Aside from this φ problem, let us take from this discussion the vector dominance idea and the applicability of the diffraction model to photon processes. In modern Regge language, we refer to the Pomeranchuk trajectory with internal quantum numbers of the vacuum to formalize the notions of classical diffraction scattering. Pictorially we think of what is going on as follows [Fig. 8]. The amplitude near the forward direction is $a(t) \sim is^{\alpha(t)} \sim is$ for $t \to 0$. This has several implications to look for in future experiments:

1) Near $t = 0$, does the ratio of the helicity flip to the helicity non-flip amplitude, corresponding to the ratio of longitudinally to transversely polarized ρ^0's, decrease as $1/s$ for large and increasing s?

2) Does the diffraction peak shrink--i.e., is $\frac{\alpha(t < 0)}{\alpha(0)} < 1$ or is there a fixed pole at $J = 1$ as if the Pomeron has a flat trajectory? The evidence from πp and pp scattering is not crystal clear here, and the unitarity arguments which provide theoretical arguments for a finite slope in hadron processes are absent in the electrodynamic ones since we are working only to lowest order in $\alpha = \frac{1}{137}$.

3) Does the behavior $\frac{d\sigma}{d\Omega} \sim k_\gamma{}^2$ of Eq. (13) remain valid at higher energies and become clearly in evidence also for ω and φ mesons? No other vectors decaying to a $\pi^+\pi^-$ ($G = +1$; $C = -1$) have been observed[36] to be diffraction produced in the $(\text{mass})^2$ range of 0.35-1.2 GeV^2. Perhaps "daughters" will make their debut at still higher masses.

4) The cross section for f^0 photoproduction is small at DESY energies[37] but if diffraction production is possible it might grow appreciably at the highest SLAC energies. This is a new twist on C violation in electromagnetic processes of strongly interacting particles since such a C violating component is needed if this is to occur via the diffraction channel with exchange of the vacuum quantum numbers.[38] The only other mechanism exhibiting a growth in the forward differential cross section with k_γ^2 would be the so far unobserved exchange of elementary vector mesons. The Primakoff effect is of much too small a magnitude.

We march next into π production. This amplitude has been studied with synchrotrons since the early 1950's and has provided crucial evidence for determining pion quantum members as well as establishing conclusively the existence of the 3-3 resonance. Since then, much more understanding of fundamental importance with respect to quantum numbers has been gained from studies in the resonance region below 2 GeV. More recent interest has also focused on the higher energy regions as we search for evidence that the Regge pole, or moving pole, hypothesis for the scattering amplitude as a function of angular momentum j is applicable and the domain of success for these notions can be extended to embrace photon processes.

For inelastic processes such as photopion production, the relevant trajectories are those lying highest and with quantum numbers different from the vacuum; hence we are here not talking about the Pomeron. For π^+ photoproduction, trajectories on which lie the π itself, as well as the charged ρ, A_1, A_2 and the B [if it has the quantum numbers $J^P = 1^+$; $G = +1$; $C = -1$] can be exchanged; and for π^0 photoproduction only the odd charge conjugation neutral ones, the B, φ, ω, and ρ contribute as shown in Fig. 9.

Can the high energy behaviors be explained in terms of these t channel exchanges or is a vector dominance model, as discussed earlier and represented by Fig. 10, more useful so that we can relate the data directly to π production of transversely polarized ρ's, ω's, and φ's? Or is the moving pole hypothesis either in conflict or just incomplete?

Some Regge ideas applied to the non-vacuum channel have met with significant success in photon amplitudes; viz.

1) Harari[39] has used Regge pole asymptotics for t channel exchanges in virtual forward Compton scattering and standard assumptions as to the position of the intercepts of the trajectories on which they lie to give a natural and simple explanation of the failure of calculations of $\Delta I = 1$ mass splittings such as the neutron-proton mass difference and of the success of the $\Delta I = 2$ ones such as the $\pi^+ - \pi^0$ mass splitting.

2) A dip in the π^0 photoproduction angular distribution at a momentum transfer of $t = -0.6$ GeV^2 is observed[40] to persist for photon energies ranging from 2 to 5 GeV, and this suggests an origin associated with a non-sense point in the presumably dominant ω-exchange trajectory. Thus in the crossed, or t, channel the $\gamma + \pi^0$ form a state of unit helicity since the photon is transverse. However, at a value of t at which the ω trajectory crosses $j = 0$ it acts like a spin 0 particle under the 3 dimensional (Euclidean) rotations and cannot support a unit of helicity--thus the non-sense zero. For arbitrary fractional j we may think of a superposition of all integer j being present--this is just a completeness statement; but the contribution vanishes at $\alpha_\omega(t) = 0$ which by linear trajectory extrapolations occurs near $t = -0.6$ GeV^2. By the way, inelastic μ or e scattering can always be analyzed

as Hand[41] has spelled out in detail to provide a sizable and known proportion of longitudinal photons for which this non-sense zero does not occur. An experiment is hereby advertized! An appreciable fraction of longitudinal photons can be achieved with large energy losses and the filling in of this non-sense zero studied.

So we might optimistically expect from these fragmentary clues that Regge will fill the bill of fare for photoproduction. To anticipate, we may look at π^+ photoproduction which has the four invariant amplitudes of CGLN[42] corresponding to incident photon and proton helicities parallel or anti-parallel and the final neutron with or without helicity flip. It was pointed out some time ago[43] that this spin structure of the amplitude suggests that things might be especially interesting at 0°.

Real photons, being transversely polarized, introduce into the production amplitude a unit of (spin) angular momentum along their direction of motion. This unit of spin cannot be carried off by a zero spin pion produced at precisely the forward angle $\theta = 0^\circ$ since its orbital angular momentum is normal to its direction of motion. A unit of spin must therefore be transmitted to the target. This requirement suppresses the contribution at $\theta = 0^\circ$ of the t channel exchanges that are normally assumed to dominate the high energy, low momentum transfer behavior of this process. Of the four invariant CGLN amplitudes, only the one flipping the proton helicity and taking up the photon spin can contribute. To be more precise, we are interested in production angles $\theta < \theta_p = \mu/\omega$ which is the characteristic angle for the peripheral photoproduction events leading to pions of mass μ and energy ω. To explain in some detail what might be taking place here,[44] we follow the Van Hove[45] approach for simplicity and work directly in the s-channel, or energy channel, of the photoproduction amplitude, building up a Regge exchange of arbitrary fractional angular momentum from a sum of all integral spin exchanges. The results are identical with the usual approach of making helicity decompositions in the t channel and has the virtues that it is closely related to single particle exchange diagrams, that the complicated helicity crossing matrices are avoided along with their treacherous behaviors for mass zero particles, and that gauge invariance is manifest.

The crucial point in this analysis is that the photon wave function is $F_{\mu\nu} \sim \epsilon_\mu k_\nu - \epsilon_\nu k_\mu$ for the gauge invariant field amplitude. Since $F_{\mu\nu}$ is antisymmetric under the interchange $\mu \longleftrightarrow \nu$ both of its indices cannot be contracted with those of the symmetric wave function $\psi^{\mu_1 \cdots \mu_J}$ representing the (arbitrary) spin J particle being exchanged. Therefore, independent of what happens at the nucleon vertex of Fig. 9, one of the indices of $F_{\mu\nu}$ must be contracted with the momentum vector of the pion, q_α, in forming a scalar (or pseudoscalar) for the upper vertex. This gives

$$(\epsilon_\mu k_\nu - \epsilon_\nu k_\mu) q_\mu \quad \sim \quad \begin{cases} k\epsilon \cdot q \approx \omega^2\theta \\ \\ k \cdot q = \frac{1}{2}(\mu^2 + \omega^2\theta^2) \end{cases}$$

The first term vanishes at $\theta = 0$ since it does not conserve angular momentum in the forward direction; the second term does conserve J_z

but is reduced by two powers of the energy and thus we find a dip from such an interaction--i.e., the amplitudes at $\theta = 0^{\circ}$ and at $\theta = \theta_p$ have the ratio

$$\frac{A(\omega, \theta \sim 0^{\circ})}{A(\omega, \theta \sim \theta_p = \mu/\omega)} \sim \mu/\omega \ll 1 \ . \tag{18}$$

Reggeizing these amplitudes in no way alters the dip--it only tempers the energy variation to the familiar s^{α}.

To avoid this dip we must bring in the vector potential of the photon $A_\mu \sim \epsilon_\mu$ and this leads us to consider the pion convection current and the t channel exchange of the pion itself which requires special treatment.

The photon coupling to the pion current via the vector potential $A_\mu \sim \epsilon_\mu$ is not gauge invariant by itself. To maintain gauge invariance, one must also include the s-channel contribution of the nucleon current as illustrated in Fig. 11. In perturbation theory for the process $\gamma + p \rightarrow \pi^+ + n$ one has the amplitude

$$A \sim e g_{\pi NN} \, \bar{u}(p') \gamma_5 \left\{ \frac{\epsilon \cdot q}{k \cdot q} - \frac{\epsilon \cdot p}{k \cdot p} - \frac{i\sigma_{\mu\nu}}{2k \cdot p} \epsilon^{\mu} k^{\nu} \right\} u(p) \ . \tag{19}$$

The first term comes from pion exchange, the second and third terms come from the s channel nucleon Born diagram. The gauge invariant convection current part of Eq. (19), $\left(\frac{\epsilon \cdot q}{k \cdot q} - \frac{\epsilon \cdot p}{k \cdot p}\right)$, vanishes in the forward direction. The remaining part of Eq. (19) has no forward dip and gives in the non-relativistic limit the Kroll-Ruderman[19] theorem for the threshold s wave photoproduction of π^+'s. However, such a $\gamma_5 \sigma_{\mu\nu} \epsilon^\mu k^\nu$ contact term describes the emission of pions in low (s and p) waves only. In the conventional Regge-peripheral view, one considers that high energy reactions proceed via the cumulative effect of many high partial waves (large impact parameters) and that production in low partial waves plays a relatively minor role, being severely suppressed by absorption mechanisms. Hence, we might not feel that such contact terms will be important at high energies and in particular they of course do not appear in neutral π^0 photoproduction. Analogous contact terms can be present in pure strong interaction processes; experience seems to indicate that they are not important.

Presumably at high energies and at finite momentum transfers $-t \sim 0.1 - 0.3$ GeV^2 a bump begins to grow with increasing s according to Eq. (18). In advance of seeing any very high energy data, this is our main expectation and expresses the fact that the contact term leads to $\frac{d\sigma}{dt} \sim 1/s^2$ whereas the exchange of a Regge trajectory with $\alpha(t) > 0$ would lead to a contribution $\frac{d\sigma}{dt} \sim s^{2\alpha(t)-2} > s^{-2}$.

In the spirit of this philosophy we are tempted to conclude as a general result that a forward dip must occur[44] in the high energy production of single pions. What could dull this temptation:

First of all, data. In π^0 photoproduction at 3 GeV from DESY [and continued to 5.8 GeV in data reported to this conference] a dip is indeed seen at very forward angles before tying in to a Primakoff peak for Coulomb production.[46] However, in contrast the π^+ data at 2.1 GeV fails to show a drop[47] and in fact indicates a narrow peak rising at 0°.

Since there is some weak evidence of a dip beginning to form at 4 GeV and at the forward direction[48] we can only say (in public at the time of presenting this report) that, without a divining rod revealing what will be reported by Dr. Richter later this week, the high energy behavior is not clear. At the very high energy of SLAC secondary beam measurements[49] the angular distribution of 7 GeV π^+'s from a Be target produced by an 18 GeV bremmstrahlung beam displays a forward peak, but this is attributable to vector or tensor meson production followed by a decay to a π^+--viz. A_1 or $A_2 \to \rho + \pi^+$.

However, we shall stay with the π^+ cross section behavior at 0° because of its great interest to a theoretical understanding of what is going on. The questions to focus on are: What is happening at 0°? What is the energy variation at finite momentum transfers ($\sim 0.2\ \mathrm{GeV}^2$) once the special angular momentum constraints at 0° are lifted? And how can it be correlated with known Regge trajectories? Can we theoretically avoid a dip at zero degrees--independent of whether or not data orders us to?

The phenomenon of conspiracy of Regge trajectories offers one possible means of avoiding a dip at forward angles. Since it will come up often during the following discussions, I will spend a few mintues trying to expose the mysterious conspirators. In practice, conspiracy[50] means the following as we illustrate by a concrete example.

Consider the exchange of trajectories of various quantum numbers and parities. We break down the contributions in terms of the four CGLN amplitudes.[42] For the exchange of a trajectory with the normal spin-parity relation $P = (-)^J$, we calculate two amplitudes depending on whether we choose scalar or vector type coupling to the nucleon line.

$$\begin{aligned} &\text{Scalar type coupling} \quad A \sim a(t)s^{\alpha-1}[-tM_A + M_B - 2M_D] \\ &\text{Vector type} \qquad\qquad\quad \sim b(t)s^{\alpha-1}\, M_D \end{aligned} \tag{20}$$

$a(t)$ and $b(t)$ are arbitrary residue functions of the momentum transfer t and α denotes the "angular momentum" of the trajectory or trajectories. We are working only to leading order in s for high energies. For the exchange of an abnormal trajectory with $P = (-)^{J+1}$ and with γ_5 coupling, we calculate similarly to leading order in s

$$A \sim -c(t)s^{\alpha-1}\, M_B \tag{21}$$

Only the amplitude M_A is finite in the forward direction, which coincides with $t = 0$ in the high energy limit to which these remarks are confined, since it is the contribution from the helicity states that conserve angular momentum at $t = 0$. We normally would wish to rule out the possibility of poles in the residue functions at $t = 0$ in order to avoid finite contributions to A (or even singularities) from amplitudes M_B and M_D that do not even conserve angular momentum at $t = 0$. The amplitudes in Eqs. (20) and (21) vanish in this case for $t = 0$. However, we can "conspire" to introduce such residue poles if the α's in Eqs. (20) and (21) are all equal at $t = 0$ and if the residues are related by

$$\begin{aligned} 2a(t) &= b(t) = 2c(t) \\ &\sim \frac{1}{t} \quad \text{for} \quad t \to 0 \end{aligned} \tag{22}$$

The total amplitude is then finite in the forward direction

$$A_{total} \sim s^{\alpha-1} M_A$$

and has no angular momentum violating parts. If this be the case, there will be no forward dip in π^+ photoproduction.

The theoretical case for the possible occurrence of such a conspiracy has a group theoretical basis and scattering models exhibiting it can be constructed with the Bethe-Salpeter equation.[50] The photoproduction data will go a long way to deciding whether Regge conspiracy is of relevance here.

Independent of the particular dip question, we also are intensely curious to learn how the differential cross section varies with energy at finite t values. Is there any evidence of Regge exchanges with positive intercepts $\alpha(t) > 0$? How relevant are the simple Born terms of Eq. (19) for the small t high energy data. Is the dip filling Kroll-Ruderman term important[51] indicating that fixed poles are prominant? Perhaps the vector dominance ideas discussed earlier are applicable here--in which case one should be able to correlate the small angle pion photoproduction cross sections with the analogous ones of pion production of transversely polarized neutral vector mesons[52] (ρ^0, ω, φ) through the connection illustrated in Fig. 10. Clearly photoproduction--also of K mesons and at backward angles corresponding to baryon trajectory exchanges--will command attention on center stage at the next photon conference because of its fundamental importance.

Once the notion of a fixed pole or singular residue in conjunction with conspiracy is introduced, we might turn more cautiously and critically elsewhere for evidence that pure simple Regge pole hypothesis fails to fill the bill for photon processes.

One simple example first discussed by Mur[53] is the elastic Compton scattering from a proton. In the forward direction we expect, in complete analogy with πp and pp elastic scattering and as also invoked for ρ^o photoproduction earlier, that we will find a forward diffraction peak corresponding to Pomeron exchange. However, a Pomeron leading to a constant total cross section, σ_t, at high energies must have a Regge trajectory intersecting at $\alpha_p(0) = 1$ and thus behaving under three dimensional rotations as a vector. It can then not couple to 2γ's any more than a vector π^o could have decayed to 2γ's. More precisely in the forward direction the photon cannot flip helicity and an incident right circularly polarized γ (rhγ) must emerge as a rhγ simply by angular momentum conservation. Upon crossing to the t channel and the process $\gamma + \gamma \rightarrow p + \bar{p}$, the emerging rh$\gamma$ crosses to a ℓhγ incident and the two incoming γ's form a system with two units of helicity. This cannot however be deposited upon a Pomeron of unit spin if $\alpha_p(0) = 1$. If the Pomeron does not couple or if we must contrive to make $\alpha_p(0) < 1$, we do not predict a constant σ_t at high energies and we lose in an instant, the motivating charm of the Pomeranchuk trajectory in Reggeism. Originally it was designed to reproduce in hadron physics the classical diffraction picture in the classical problem of light scattering.

We also run into the following fundamental contradiction.[53] The Pomeron in ρ^o photoproduction leads to an inelastic cross section $\sigma_{inel} \sim \frac{1}{\ln s}$ if the diffraction peak shrinks. However, if it is absent from forward elastic Compton scattering, then by the optical theorem $\sigma_t \sim s^{\alpha_x(0)-1}$ where $\alpha_x(0) < 1$ is the intercept of the next highest lying trajectory. We are then led to a contradiction since $\sigma_t < \sigma_{inel}$ and simple pure Regge behavior is once more on the ropes for photon problems. Either we must remove the Pomeron altogether from all inelastic channels, or we must give it a singular residue to cancel its non-sense zero at $t = 0$ for forward Compton scattering, or we must abandon the classical diffraction analogy of $\alpha_p(0) = 1$. Could it be that there is a fixed pole at $J = 1$? Other arguments for a fixed pole at $J = 1$ have been presented to make current algebra predictions compatible with Regge asymptotic behavior.[54]

Once this Pandora's box is opened, we have a new ball game and several experiments acquire enhanced interest. First of all, the energy dependence of the total absorption cross section of virtual high energy γ's from inelastic electron or muon scattering[41] will show up any differences between the forward diffraction amplitude for transverse quanta and for longitudinal ones that are free of the Pomeron's non-sense zero. Secondly, a study of forward Compton scattering at low energies ($< \mu = 140$ MeV) can reveal whether we are led to the requirement of a subtraction constant at infinite energy for the real part of the forward non-spin flip Compton amplitude from a proton. To amplify this observation we write the amplitude for forward Compton scattering from a proton as given by Gell-Mann, Goldberger, and Thirring[55]

$$f(\nu) = f_1(\nu)\, \underline{e}'^{*} \cdot \underline{\epsilon} + i\underline{\sigma} \cdot \underline{e}'^{*} \times \underline{e}\, \nu\, f_2(\nu) \tag{23}$$

where ν is the photon energy and e and e' are the transverse polarization vectors of the incident and outgoing photon. The dispersion relation for $f_1(\nu)$ usually appears as

$$f_1(\nu) = -\alpha/M + \frac{\nu^2}{2\pi^2} \int_\mu^\infty \frac{d\nu' \sigma_t(\nu')}{\nu'^2 - \nu^2} \tag{24}$$

where the exact classical Thomson limit is introduced as a subtraction constant at zero energy and $\sigma_t(\nu')$ is the total photoabsorption cross section by the proton. We are working to lowest order in $\alpha = \frac{1}{137}$ but to all orders in the strong interactions and the threshold in the dispersion integral is $\mu \simeq 140$ MeV, the threshold for photopion production. In the present context, total cross sections means the photoabsorption to form hadron final states, and the large but well understood Bethe-Heitler processes are excluded. Whether or not the spin dependent amplitude $f_2(\nu)$ requires a subtraction in its dispersion relation, its zero energy limit is exactly known in terms of the proton charge, mass, and anomalous moment k_p to be[22]

$$f_2(0) = -\frac{\alpha}{M^2} k_p^{\,2} \; ; \quad k_p = 1.79 \tag{25}$$

Combining equations (23), (24), and (25), we have an exact result for the forward angle differential elastic Compton cross section

$$\begin{aligned} \lim_{\nu^2 \to 0} \left(\frac{d\sigma}{d\Omega}\right)_{0^\circ} &= \left|f_1(\nu)\right|^2 + \nu^2 \left|f_2(\nu)\right|^2 \\ &= \frac{\alpha^2}{M^2} \left[1 - \left(\frac{\nu}{\mu}\right)^2 \left\{ \frac{\mu^2 M}{\alpha \pi^2} \int_\mu^\infty \frac{d\nu' \sigma_t(\nu')}{\nu'^2} - \frac{\mu^2}{M^2} k_p^{\,4} \right\} + 0\left[\left(\frac{\nu}{\mu}\right)^4\right]\right] \end{aligned} \tag{26}$$

The coefficient of the low energy slope is already known very accurately from measured photoabsorption cross sections[56] up to 6 GeV since the integral converges rapidly and to one significant figure

$$\frac{\mu^2 M}{\alpha\pi^2}\int_{\mu}^{6\ \mathrm{GeV}}\frac{d\nu'\sigma_t(\nu')}{\nu'^2}-\frac{\mu^2}{M^2}k_p^{\ 4}=+0.7\ . \tag{27}$$

Further refinement in this number will result from measurements at higher energies but in any case the changes will be small.[57] Evidently there is a sizable and measurable slope with $(\mathrm{energy})^2$ to be measured and checked against the very general assumptions that are the input into the forward dispersion relations for scattering of light (relativity, macroscopic causality, and unitarity).

The only possible source of disagreement between the predictions of Eqs. (26) and (27) and experiment, short of a theoretical catastrophe of the highest order, could come about as follows: Due to the contribution from a t channel exchange of an "elementary particle" of fixed spin 2 contributing to the real part of the forward spin independent amplitude, we must add a term $\lambda\nu^2$ on the right hand side of Eq. (24)--or more generally a real polynomial in ν^2 without disturbing the low energy Thomson limit. We may not welcome such a contribution, and we may not understand whence it originates, but evidently it would not be the first appearance of corrections to simple Reggeism in processes with photons.[58] On general principles it cannot be ruled out--in particular, we cannot fall back on the usual unitarity arguments that are invoked at this point in hadron amplitudes since we are working only to lowest order in α. I view an experimental confrontation of Eqs. (26) and (27) as a problem of very high urgency in "medium energy" photon physics.

This brings me finally to the end of this report which I close by noting that our faith in the electrodynamic current j_μ has been unquestioned, and we have learned from many beautiful and heroic experts that the Dirac-Maxwell QED is a singularly lovely theory--even to distances of $\lesssim$ a nucleon Compton wavelength. This is an extrapolation down by a factor $\sim 10^6$ in the scale of sizes from the domain of its origin, and in this new realm no firm evidence of a granularity in the space-time structure or of the vacuum itself or of any other breakdown of QED has appeared. The few potential shadows on the horizon of this lovely picture I leave for Dr. Yennie's talk.

FOOTNOTES AND REFERENCES

1. Dante Alighieri, c.f. John Bartlett, Familiar Quotations, 13th edition (Little Brown and Company, Boston), 1955, p. 76.

2. R. Hofstadter, Annual Review of Nuclear Science 7, 231 (1957); R. Hofstadter, F. Bumiller, and M. Yearian, Reviews of Modern Physics 30, 482 (1958).

3. Y. Nambu, Phys. Rev. 106, 1366(L), (1957).

4. W. Frazer and J. Fulco, Phys. Rev. Letters 2, 365 (1959); Phys. Rev. 117, 1609 (1960).

5. c. see reports of R. Hofstadter and Richard Wilson to the Second International Symposium on Electron and Photon Interactions at High Energies, Hamburg, June 1965 [Vol. 39, Springer Tracts in Modern Physics (Springer-Verlag), 1965].

6. W. Albrecht et. al., Phys. Rev. Letters 18, 1014 (1967); M. Goitein et. al., Phys. Rev. Letters 18, 1016 (1967).

7. H. Schopper, "Some Remarks Concerning Proton Structure," CERN 67-3 (February 16, 1967); to be published.

8. T. T. Wu and C. N. Yang, Phys. Rev. 137, B708 (1965).

9. For a continuation to $q^2 = 25\ \mathrm{GeV}^2$, see the report of R. Taylor to this conference.

10. G. Chew, S. Gasiorowicz, R. Karplus, and F. Zachariasen, Phys. Rev. 110, 265 (1958); P. Federbush, M. Goldberger, and S. Treiman, Phys. Rev. 112, 642 (1958).

11. W. Bartel, B. Dudelzak, H. Krehbiel, J. M. McElroy, U. Meyer-Berkhout, R. J. Morrison, H. Nguyen-Ngoc, W. Schmidt, and G. Weber, DESY 67-18 (May 1967); to be published.

12. N. M. Kroll, T. D. Lee, and B. Zumino, Phys. Rev. 157, 1376 (1967). For earlier work, see especially J. J. Sakurai, Ann. Phys. (N.Y.) 11, 1 (1960); M. Gell-Mann and F. Zachariasen, Phys. Rev. 124, 953 (1961).

13. J. Bowcock and A. Martin, Nuovo Cimento 14, 516 (1959); R. Blankenbecler, M. Goldberger, N. Khuri, and S. Treiman, Ann. Phys. (N.Y.) 10, 62 (1960).

14. T. Regge, Nuovo Cimento 14, 951 (1959); 18, 947 (1960).

15. M. J. Lighthill, Fourier Analysis and Generalized Functions (Cambridge University Press, Cambridge), 1958, p. 56.

16. For a three quark system, $g(r) \equiv r\rho(r)$ in Eq. 10 is replaced by $\tilde{g}(r) \equiv r \int d^3s |\psi(\underline{r},\underline{s})|^2$. Again $\tilde{g}(0) = 0$ for a wide class of potentials and more generally if $\psi(\underline{r},\underline{s})$ is finite for $\underline{r}$ and $\underline{s} \rightarrow 0$ the form factor falls off more rapidly than $1/q^5$. S. D. Drell, A. C. Finn, and M. H. Goldhaber, to be published; see also, Phys. Rev. 157, 1402 (1967).

17. A. M. Bincer, Phys. Rev. 118, 855 (1960).

18. There is no unphysical region. For simplicity we neglect the pion mass, $\mu/M \rightarrow 0$.

19. N. M. Kroll and M. A. Ruderman, Phys. Rev. 93, 233 (1954).

20. S. D. Drell and H. R. Pagels, Phys. Rev. 140, B397 (1965).

21. R. Parsons, to be published.

22. F. Low, Phys. Rev. 96, 1428 (1954); M. Gell-Mann and M. L. Goldberger, Phys. Rev. 96, 1433 (1954).

23. To be published.

24. F. Ernst, R. Sachs, and K. Wali, Phys. Rev. 119, 1105 (1960).

25. C.f., Ref. 11; and R. Budnitz, J. Appel, L. Carroll, J. Chen, J. R. Dunning, Jr., M. Goitein, K. Hanson, D. Imrie, C. Mistretta, J. K. Walker, and Richard Wilson, Phys. Rev. Letters 19, 809 (1967).

26. C. W. Akerlof, W. W. Ash, K. Berkelman, C. A. Lichtenstein,

A. Ramanauskas, and R. H. Siemann, to be published; c.f., Phys. Rev. Letters 16, 147 (1966).

27. J. J. J. Kokkedee and L. Van Hove, Nuovo Cimento 42, 711 (1966); L. Van Hove, Lectures at 1966 Scottish Universities Summer School, edited by T. Preist and L. Vick (Penum Press, New York), 1967.

28. L. J. Lanzerotti et. al., Phys. Rev. Letters (to be published); Cambridge Bubble Chamber Group, Phys. Rev. 146, 994 (1966); J. G. Asbury et. al., Phys. Rev. Letters (to be published).

29. S. M. Berman and S. D. Drell, Phys. Rev. 133, B791 (1964); and S. D. Drell, Proceedings of the Second International Symposium on Electron and Photon Interactions at High Energies, loc. cit., p. 71.

30. M. Ross and L. Stodolsky, Phys. Rev. 149, 1172 (1966); and c.f. Ref. 12.

31. Y. Nambu and J. J. Sakurai, Phys. Rev. Letters 8, 79 (1962); M. Gell-Mann, D. Sharp, and W. Wagner, Phys. Rev. Letters 8, 261 (1962).

32. See the reports of Drs. Harari, Pipkin, Ting, and Weinstein for more up to date data and detailed analyses.

33. M. Ross and L. Stodolsky, op. cit.; S. D. Drell and J. S. Trefil, Phys. Rev. Letters 16, 552, 832E (1966).

34. J. Asbury, op. cit. See also the similar prediction from the strong absorption model; Y. Eisenberg et. al., Phys. Letters 22, 217 (1966).

35. H. Harari, Phys. Rev. 155, 1565 (1967). Among recent attempts to reconcile the quark model with the small φ cross section, we cite K. Kajantie and J. S. Trefil, Phys. Letters 24B, 106 (1967).

36. J. Asbury et. al., to be published.

37. R. Erbe et. al., Nuovo Cimento 48, 262 (1967) have measured an f^{o}

photoproduction cross section of 2.5 ± 1.3 μb for γ energies in the interval 2.5 GeV < k_γ < 3.5 GeV, and 0.6 ± 0.4 μb for 3.5 GeV < k_γ < 5.8 GeV.

38. For more on C violation in electromagnetic interactions of the hadrons, see the report of T. D. Lee. This specific proposal and its qualitative behavior and estimates grew out of a discussion by F. E. Low and the author in a SLAC scheduling committee meeting (unpublished).

39. H. Harari, Phys. Rev. Letters 17, 1303 (1966).

40. M. Braunschweig et. al., Phys. Letters 22, 705 (1966); G. Bolon et. al., Phys. Rev. Letters 18, 926 (1967).

41. L. Hand, Phys. Rev. 129, 1834 (1963); SLAC Report No. 25 (Summer Study Report), 1963.

42. G. Chew, M. Goldberger, F. Low, and Y. Nambu, Phys. Rev. 106, 1345 (1957).

43. A. Berkov, E. Zhizhin, V. Mur, and Yu. Nikitin, Soviet Physics JETP 18, 1091 (1964); in Russian JETP 45, 1585 (1963).

44. This discussion follows closely S. D. Drell and J. D. Sullivan, Phys. Rev. Letters 19, 268 (1967).

45. L. Van Hove, Phys. Letters 24B, 183 (1967).

46. N. Braunschweig, W. Braunschweig, D. Husmann, K. Lükelsmever, and D. Schmitz, to be published.

47. G. Buschhorn et. al., Phys. Rev. Letters 18, 571 (1967).

48. J. Dowd, D. Caldwell, K. Heinloth, and T. Sherwood, Phys. Rev. Letters 18, 414 (1967).

49. A. Boyarski et. al., Phys. Rev. Letters 18, 363, 1967.

50. For general discussion of conspiracy, see contributions in "Comments in Nuclear and Particle Physics," March 1967 by G. Chew and by M. Goldberger. For particular application to photoproduction, see R. Sawyer, Phys. Rev. Letters 18, 1212 (1967); P. K. Mitter, to be published; B. Diu and M. LeBellac, Orsay preprint Th/198, to be published; S. Frautschi and Lorella Jones, Phys. Rev. (to be published); J. Ader, M. Capdeville, and Ph. Salin, CERN Preprint Th. 803, to be published; J. Frøyland and D. Gordon, Phys. Rev. Letters (to be published).

51. M. B. Halpern, Phys. Rev. 160, 1441 (1967) has found the contribution of the Kroll-Ruderman term to be an order of magnitude larger than that of the lowest lying six resonances to the real part of the forward photoproduction amplitude.

52. D. Beder, Phys. Rev. 149, 1203 (1966); B. Diu and M. Le Bellac, Phys. Letters 24B, 416 (1967).

53. V. D. Mur, Soviet Physics 17, 1458 (1963) [in Russian, JETP 44, 2173 (1963)]; 18, 727 (1964) [in Russian, JETP 45, 1051 (1963)]; H. D. I. Abarbanel and S. Nussinov, Phys. Rev. 158, 1462 (1967); H. K. Shepard, Phys. Rev. 159, 1331 (1967)

54. J. Bronzan, I. Gerstein, B. Lee, and F. Low, Phys. Rev. Letters 18, 32 (1967); V. Singh, Phys. Rev. Letters 18, 36 (1967).

55. M. Gell-Mann, M. L. Goldberger, and W. Thirring, Phys. Rev. 95, 1612 (1954).

56. C.f. reports to this conference of E. Lohrmann and F. Pipkin.

57. The ν^4 terms can be estimated using the f_1 and f_2 dispersion relations and are $< 10\%$ at $\nu \sim 100$ MeV.

58. A. H. Mueller and T. L. Trueman, Phys. Rev. 160, 1296, 1306 (1967); H. Abarbanel, F. Low, I. Muzinich, S. Nussinov, and J. Schwarz, Phys. Rev. 160, 1329 (1967).

FIGURE CAPTIONS

Fig. 1 - Dispersion graph for the absorptive amplitude, $\rho(\sigma^2)$. V^o denotes the neutral zero-strangeness vector resonances.

Fig. 2 - Diagram for calculation of the electromagnetic form factor as a product of the vector meson propagator multiplied by the vector meson form factor.

Fig. 3 - Reduced graph for the deuteron electromagnetic vertex. The mass inequality $M_D^2 > M_p^2 + M_n^2$ allows such two dimensional reduced graphs to be drawn. [c.f. J. D. Bjorken and S. D. Drell, RELATIVISTIC QUANTUM FIELDS (McGraw-Hill Book Company, New York, 1965) p. 235].

Fig. 4 - Dispersion graph for the absorptive part, Im $G(W^2,q^2)$ in Eq. (12).

Fig. 5 - Graph for diffraction amplitude to photoproduce neutral vector mesons of zero strangeness by the vector dominance model.

Fig. 6 - Graph for ρ^o exchange contribution to the pion electromagnetic form factor.

Fig. 7 - Graph for lepton decay of the ρ^o.

Fig. 8 - Diffraction photoproduction of the ρ^o via exchange of the Pomeranchuk trajectory.

Fig. 9 - Relevant trajectory exchanges for photoproduction amplitudes.

Fig. 10- Vector dominance model relating photopion production to pion production of transversely polarized neutral vector mesons.

Fig. 11- Pion current and nucleon pole contributions to π^+ photoproduction.

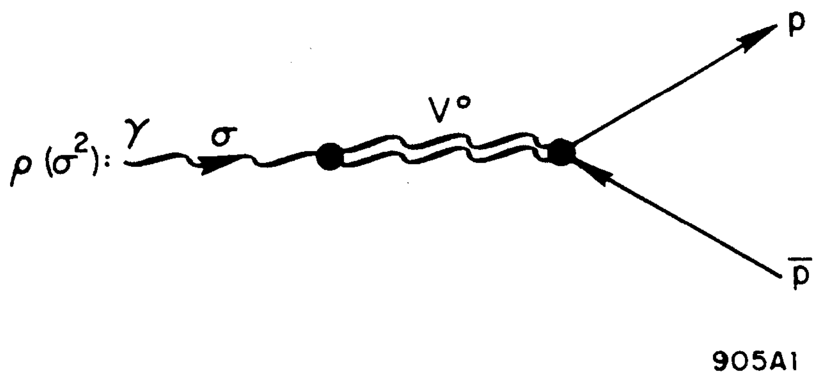


Fig. 1

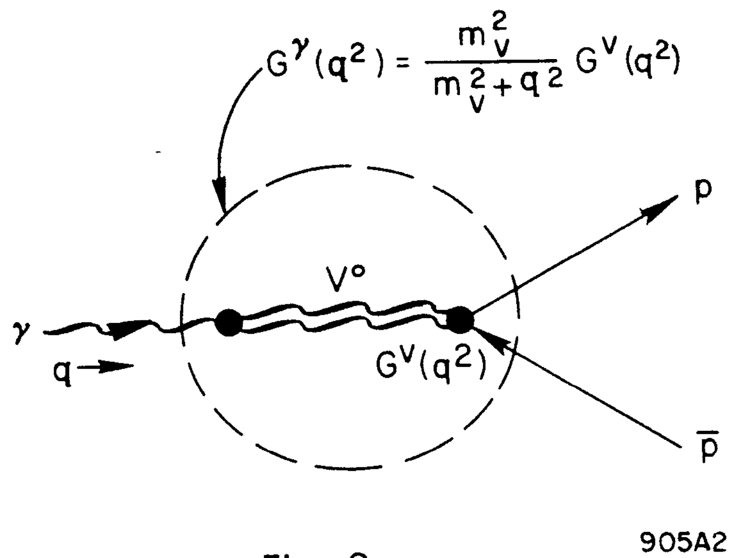


Fig. 2

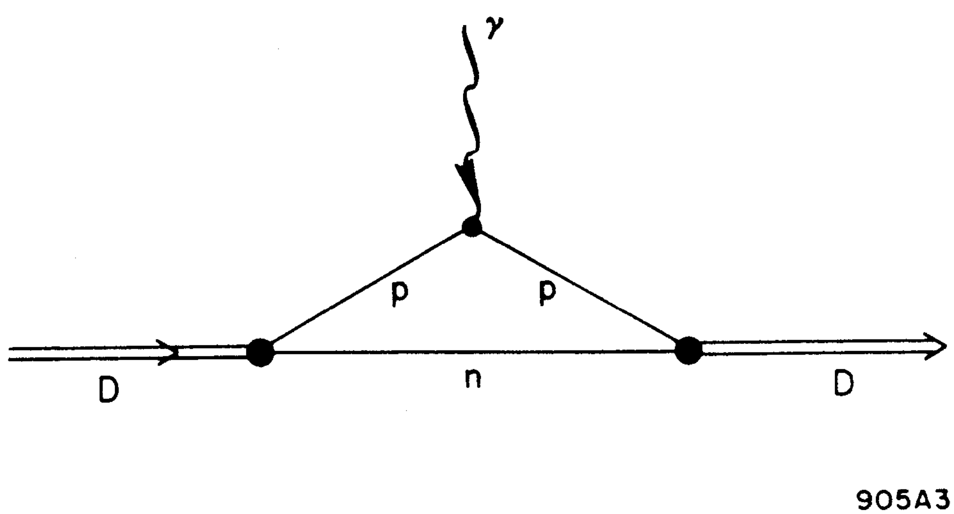


Fig. 3

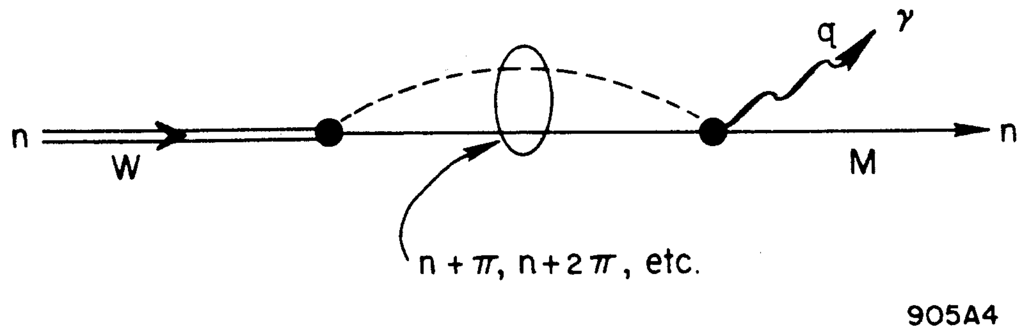


Fig. 4

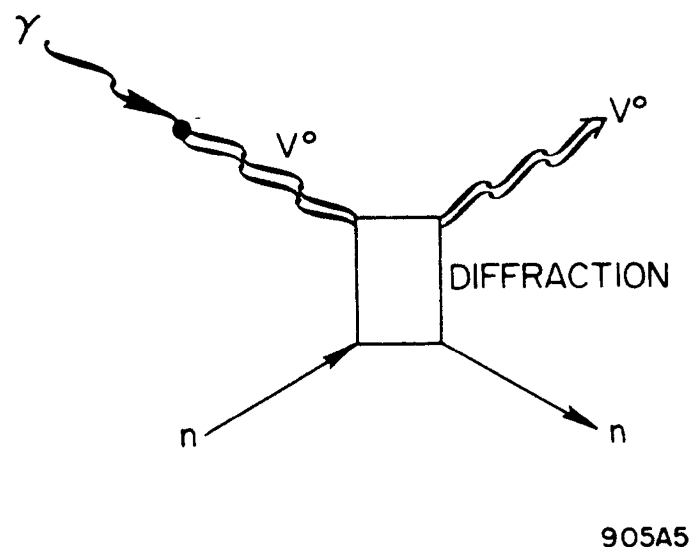


Fig. 5

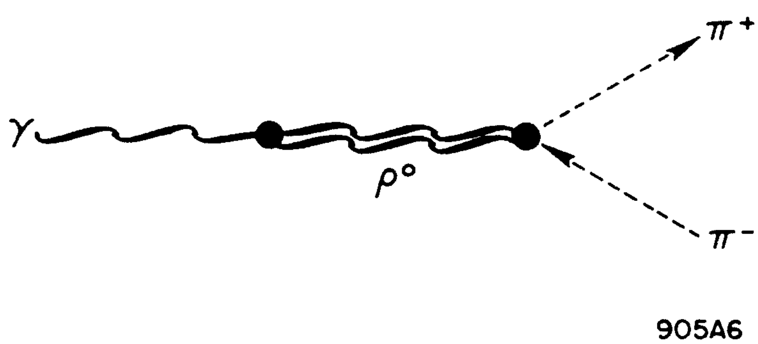


Fig. 6

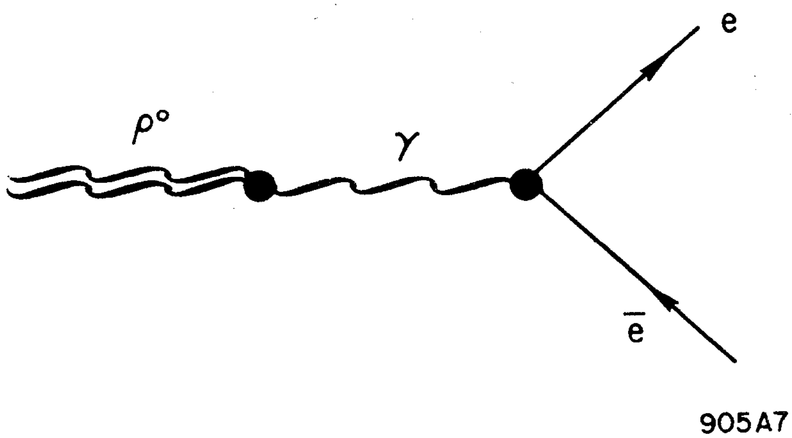


Fig. 7

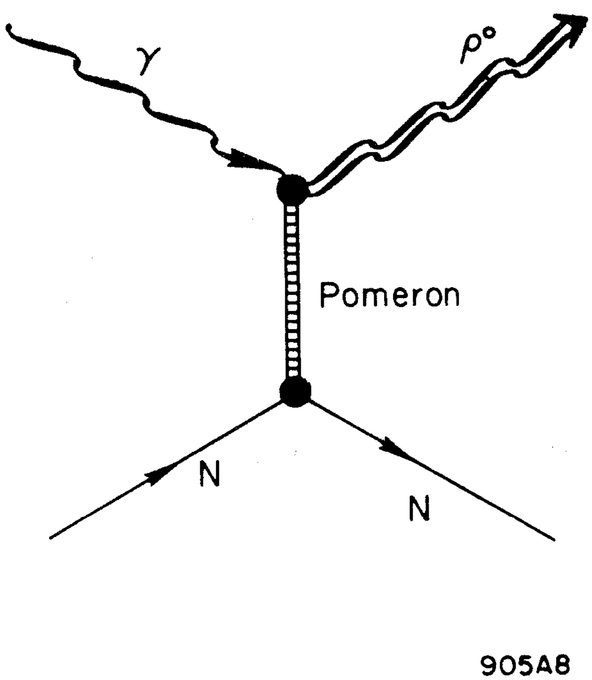
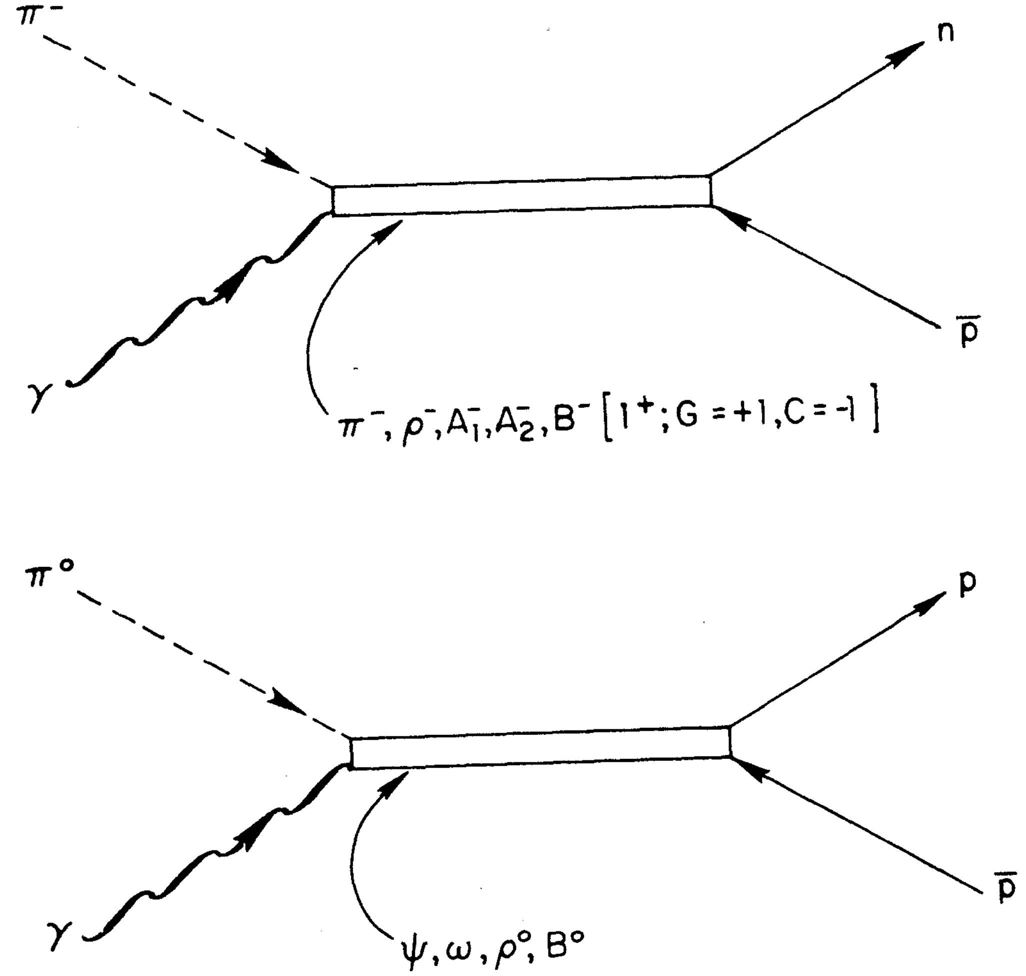


Fig. 8


π^-
n
$\bar{p}$
γ
$\pi^-, \rho^-, A_1^-, A_2^-, B^- \left[1^+; G = +1, C = -1\right]$
π^0
p
$\bar{p}$
γ
$\psi, \omega, \rho^0, B^0$

905A9

Fig. 9

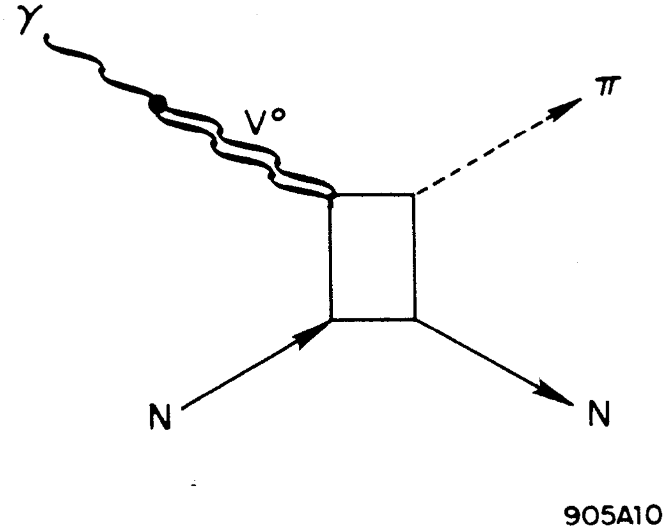


Fig. 10

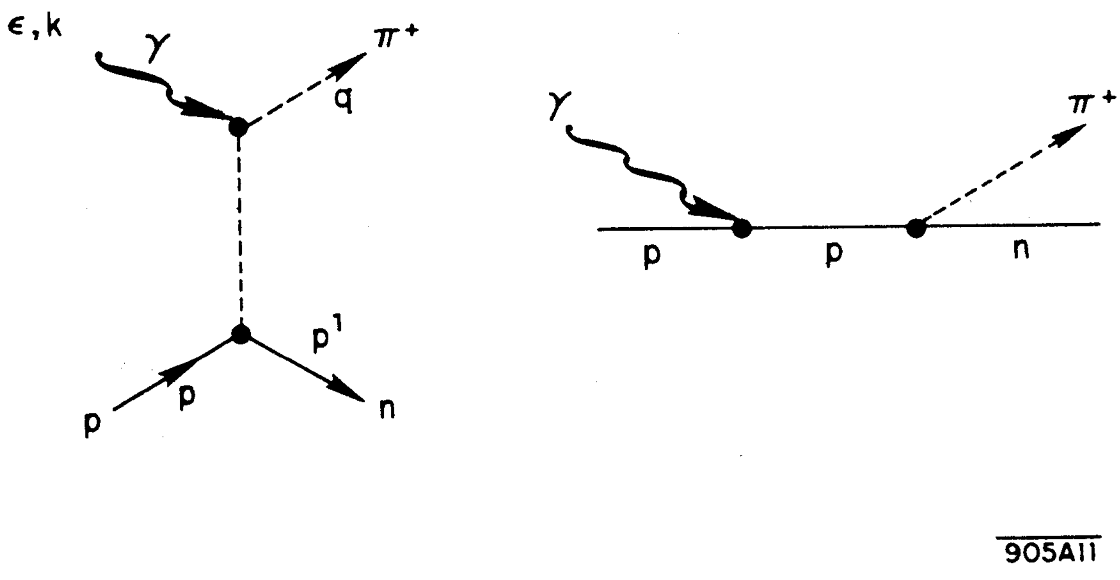


Fig. 11