Vol. B7 (1976) ACTA PHYSICA POLONICA No 12

INTUITIVE DISCUSSION OF HIGH ENERGY PHOTON INTERACTION*

By D. R. Yennie

Laboratory of Nuclear Studies, Cornell University**

(Presented at the XVI Cracow School of Theoretical Physics, Zakopane, May 27 — June 7, 1976)

High energy photon interactions are discussed in terms of the hadronic structure of the photon. It is shown how certain qualitative features of the data may be understood from this point of view, and some indications of its limitations are given.

1. Introduction

There are two rather different intuitive views of the nature of photon interactions with hadrons. In the first, it is noted that since the photon coupling strength to the hadrons is small, higher order electromagnetic effects may be neglected; and the electromagnetic interaction is regarded as taking place between a rather simple object (the real or virtual photon) and a very complex, strongly interacting system. In this view, all the complicated things that are observed are due to the properties of the hadronic sytem. In the other view, the photon is viewed as being something like a hadron itself a small fraction ($\sim 1/137$) of the time and its interactions are very similar to other hadronic interactions, but are much reduced in magnitude. Then the interactions are viewed as a mutual property of two objects which are themselves very complicated in structure.

To illustrate this point, we anticipate the fact that there is an interaction which permits a photon to transform to a vector meson (such as the ϱ^0, ω or ϕ)

$$\gamma \leftrightarrow V \tag{1.1}$$

and that this vector meson somehow couples to other hadrons in turn. Such an interaction could affect the charge distribution as seen in electron scattering experiments, as illustrated in Fig. 1a. In this case, it has been usual to regard the modification due to (1.1) as being associated entirely with the target: V causes a change in the distribution of charge which is measured by the electron through a virtual photon. In this case it is worth noting

* Supported in part by the National Science Foundation.

** Address: Laboratory of Nuclear Studies, Cornell University, Ithaca, NY 14853, USA.

that the photon is usually *highly* virtual (and spacelike: its energy $\ll$ its momentum). Another process is illustrated in Fig. 1b. Here a real, high-energy photon transforms into a V before interacting with a target to produce a multiparticle final state. In this case it is more natural, as we shall argue in this paper, to regard the V as part of a "hadronic structure" of the photon since in particular the conversion to V often takes place well outside

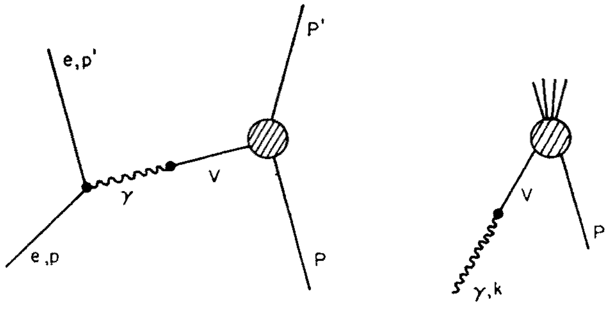

Fig. 1. Two situations where a photon interaction proceeds through an intermediate vector meson. (a) In elastic electron scattering, the vector meson is usually regarded as providing a modification of the nucleon's charge and magnetic moment distribution. (b) In a high energy photon interaction, the conversion takes place outside the target and the vector meson is regarded as part of the photon's structure

the normal structure of the target. We have described two extreme situations where the physical interpretation seems rather plausible, but there are also intermediate cases where the interaction and the structures must be considered as a unity. While it is strictly speaking meaningless to discuss the photon's structure independently of its subsequent interaction, we shall later analyze the structure of a high-energy free photon, having in mind a target at rest with which it will ultimately interact.

In a complete theory of high-energy electromagnetic interactions, the intuitive picture of how a collision takes place plays an interesting but not a crucial role. Lacking such a theory, the intuitive picture acquires a more important significance. It substitutes for a theory and aids us in developing models for interpreting and correlating experimental facts. One line of models (vector-meson-dominance or generalized- (extended-) vector--dominance) places the emphasis on the hadronic structure of the photon, while another (parton) concentrates on the structure of the target in an infinite momentum frame. The hadronic structure models are more natural and successful in describing the facts in one kinematic domain, and the parton models in another; they are not in conflict, but each seems "unnatural" when employed beyond its appropriate kinematic region.

It is the aim of these lectures to provide an introduction to the main ideas of the hadronic structure of the photon, starting with a very brief review of the experimental evidence which suggests such structure. Much of the material incorporated here has been drawn freely from a comprehensive review by Ted Bauer, Frank Pipkin, Robin Spital and myself [1]. Because this review should soon become available, I have not made a great effort here to be historically accurate or complete. My views have been influenced by my three collabo-

rators and by many other colleagues at Cornell and SLAC with whom I have discussed these questions in recent years. Among these, I would like particularly to mention Ashok suri, Stanley Brodsky, Kurt Gottfried, Nikola Jurisic, Winfried Schmidt, and Tung-Mow Yan.

The similarities of photon and hadron interactions at high energies

We want to mention some of the qualitative similarities of high energy photon interactions to those of hadrons. It will be seen that, except for the magnitudes of the cross sections, these interactions are very similar. In fact, there is an old rule-of-thumb that one can estimate photon cross sections crudely by multiplying corresponding pion cross sections by a factor of order α/π. In summary, these similarities are [2]:

(*i*) Total cross sections on nucleons: Both show spectacular resonances at low energies, and above about 3 GeV they level out and become structureless, apparently tending to a constant at high energies (a logarithmic dependence at high energies is still possible). The cross section on the neutron is nearly the same as that on the proton, and the difference seems to vanish as $E \to \infty$; thus the photon cannot interact primarily with the charge of the target. In magnitude, the total photon cross section is asymptotically about 1/220 times the average of the pion cross sections, i.e. it is smaller by approximately the fine structure constant in order of magnitude.

(*ii*) Diffractive processes: Diffractive photoproduction is the analogue of diffractive hadronic scattering. It is about the same fraction of the total cross section ($\sim 20\%$) and has similar angular dependence (roughly exponential, with similar slopes). Elastic scattering of photons also appears to be primarily diffractive; in particular,the amplitude has a small real part. This is significant because, while hadronic amplitudes would be expected to be nearly imaginary due to the almost complete absorption at small impact parameters, there is no obvious reason why the photon amplitude could not have a significant refractive part. The absence of such a real part is of course related to the rapid approach of the photon cross section to a constant at high energies.

(*iii*) Two-body reactions: These have similar features such as peaks and dips which seem to be governed by common rules.

(*iv*) Inclusive cross sections: Both display sharply falling P_T^2 distributions and comparable longitudinal momentum distribution properties.

In both the two-body reactions and the inclusive cross sections, the similarity is particularly striking if the cross sections are divided by the corresponding total cross sections.

(*v*) Absorption by nuclei: While the total photon cross section on nuclei is of course much smaller than that of hadrons, the high-energy ($\gtrsim 10$ GeV) A dependence is similar, corresponding to a strong shadowing effect.

All these features suggest a picture in which the photon acts like a hadron a small fraction ($\sim\alpha$) of the time. The simplest framework for describing these features, and one which works remarkably well overall, is vector meson dominance (VMD). In this model, the photon is assumed to be a well-defined linear superposition of the ϱ^0, ω, and ϕ mesons before interaction. With the present best knowledge for these mesons to be in the photon, it turns out that they account for a large part of the total photon cross section; in fact, the ϱ^0 alone accounts for about 2/3 of this total. Specific processes which have a small cross section which decreases rapidly with energy, such as $\gamma N \to \pi N'$, do not agree in all details with predictions from this model, but at least qualitatively the vector mesons seem to play an important role in them. The shadowing effect of the total cross section in nuclei [3, 4] is smaller than expected from VMD. This could indicate that a small fraction of the photon's interaction (say 20%) is not due to a hadronic component, or is due to hadronic components of such high masses that they have not yet saturated at current energies.

2. *A perturbation model of the dressed photon*

As we have just seen, from a purely phenomenological viewpoint, photon interactions with hadrons bear many remarkable similarities to purely hadronic interactions. At a very crude level, this could be understood if the physical photon $|\gamma\rangle$ were a superposition of two types of states: a bare photon $|\overset{0}{\gamma}\rangle$, which at high energies accounts for a small, or perhaps negligible, portion of the interaction; and a small — of order $\sqrt{\alpha}$ — hadronic component $\sqrt{\alpha}|h\rangle$ which undergoes conventional hadronic interactions. That is, we expect the important part of the physical photon state to be expressible as

$$|\gamma\rangle \cong \sqrt{Z_3}\,|\overset{0}{\gamma}\rangle + \sqrt{\alpha}\,|h\rangle, \tag{2.1}$$

where Z_3 is introduced to assure the proper normalization of $|\gamma\rangle$; all states in (2.1) have the same 3-momentum $\boldsymbol{k}$. Invariance considerations dictate that $|h\rangle$ should have the same symmetry quantum numbers as the photon, i.e. $J^{PC} = 1^{--}$, $Q = B = S = 0$. The copious photoproduction of the vector mesons ϱ^0, ω and ϕ suggests that they provide very important contributions to $|h\rangle$. The restrictive assertion that these three mesons are the sole hadronic constituents of the photon, and that the bare component $|\overset{0}{\gamma}\rangle$ cannot interact with hadrons, is the hypothesis of vector meson dominance (VMD) in its most naive and clear-cut form. The less restrictive assumption that all interactions result from $|h\rangle$ which has more constituents than ϱ^0, ω, and ϕ, is referred to as generalized vector dominance (GVD).

These ideas may be formulated more precisely by considering the structure of the state $|\gamma\rangle$ in conventional time-independent perturbation theory [5]. To lowest order in the electromagnetic interaction H' between photons and hadrons

$$\sqrt{\alpha}\,|h\rangle = \mathop{\mathrm{S}}_{n} \frac{|n+\rangle\,\langle n+|H'|\overset{0}{\gamma}\rangle}{\nu - E_n}\,, \tag{2.2}$$

where $|n+\rangle$ is a completely interacting hadronic state and $|\overset{0}{\gamma}\rangle$ includes the exact hadronic vacuum, together with a bare photon of momentum $\boldsymbol{k}$. The energy of the photon is ν ($\equiv |\boldsymbol{k}|$ for real photons and $\sqrt{\boldsymbol{k}^2 - Q^2}$ for virtual ones, and the energy of the hadronic state E_n may be written $\sqrt{\mathcal{M}^2+\boldsymbol{k}^2}$, where $\mathcal{M}$ is the mass of the total hadronic state $|n+\rangle$. $\underset{n}{\mathsf{S}}$ stands for sums and integrals over all appropriate labels in n. The "unpertubed" Hamiltonian is composed of the complete hadronic Hamiltonian plus the Hamiltonian of the free e.m. field. The perturbation is

$$H' = e \int J_\mu(x) A^\mu(x) d^3x, \tag{2.3}$$

where J_μ is the hadronic electromagnetic current.

As mentioned earlier, the invariance properties of H' assure that the states $|n+\rangle$ have the photon's symmetry quantum numbers. Possible examples of such states are $\pi^+\pi^-$, $\pi^+\pi^-\pi^0$, $K\bar{K}$, $N\bar{N}$, etc. Since the vector mesons ϱ^0, ω, and ϕ are not stable under strong interactions, they actually do not appear as separate terms in the sum in (2.2). However, to the extent that they may be regarded as narrow resonances in the $\pi^+\pi^-$, $\pi^+\pi^-\pi^0$, and $K\bar{K}$ modes respectively, their contributions may be separated off and replaced by equivalent elementary vector meson states, which may then be regarded to interact more or less like distinct physical particles (see below). We then recover the contributions expected from VMD or GVD.

Although expression (2.2) in general depends on ν, at high energies it may be simplified and given a more familiar appearance. When $\nu \gg \mathcal{M}$ (or $\nu \gg \sqrt{\mathcal{M}^2+Q^2}$ in the case of virtual photons), the energy denominator of (2.2), may be approximated

$$\begin{aligned} \nu - E_n &= -\mathcal{M}^2/2\nu \qquad \text{(real photons, } \nu \gg \mathcal{M}) \\ &= -(\mathcal{M}^2+Q^2)/2\nu \ \text{(virtual photons, } \nu^2 \gg \mathcal{M}^2+Q^2). \end{aligned} \tag{2.4}$$

Naively, one might expect that this small energy denominator causes the hadronic term to increase with ν. Instead, when the center-of-mass motion is separated out and the dependence of the phase space on ν is taken into account, a constituent of fixed mass $\mathcal{M}$ is found to have a strength which is independent of ν [5]. We may then say that each constituent is "frozen in", and rewrite the hadronic component as

$$\sqrt{\alpha}\,|h_S\rangle = +(2\pi)^{3/2} e \int_{\mathcal{M}_1{}^2}^{\sim \nu^2} \frac{d\mathcal{M}^2}{\mathcal{M}^2+Q^2} \underset{n_i}{\mathsf{S}} |\boldsymbol{k}, \mathcal{M}, n_i\rangle \langle \boldsymbol{k}, \mathcal{M}, n_i | \varepsilon \cdot \boldsymbol{J}(0) | \text{vac}\rangle, \tag{2.5}$$

where the sum now extends over all internal quantum numbers holding $\boldsymbol{k}$ and $E_n = \sqrt{\boldsymbol{k}^2+\mathcal{M}^2}$ fixed. The label h_s means that (2.5) is valid only for the low mass components of the photon. The factor $1/(\mathcal{M}^2+Q^2)$ resembles a meson propagator.

Eq. (2.5) is the desired form of the hadronic component, and it will be used for all our discussions of hadron-mediated photon interactions. For transversely polarized photons, the matrix element occuring in (2.5) is independent of $\boldsymbol{k}$. This is the "freezing in" just emphasized. For longitudinally polarized photons (possible only for $Q^2 \neq 0$), the matrix element depends on $\boldsymbol{k}$ in a well-defined way. Taking this together with the $\boldsymbol{k}$-dependence of the polarization vector of the photon, one again finds that the longitudinal structure is frozen in. (Gauge invariance requirements make the treatment of the interaction of longitudinally polarized photons somewhat more subtle however. The bare photon interaction is essential for gauge invariance, but it is possible for cross sections to be dominated by the hadron-mediated contribution, see Ref. [5], Sec. 4.) We note also that the hadronic states occurring in (2.5) are precisely the ones seen in $e^+e^- \to$ hadrons and the probability per unit $\mathcal{M}^2$, $P(\mathcal{M}^2, Q^2)$, of their occurrence is given by the cross section of the latter reaction [6] namely (for transverse photons)

$$P(\mathcal{M}^2, Q^2) = \frac{(2\pi)^3 e^2}{(\mathcal{M}^2+Q^2)^2} \mathop{\mathrm{S}}_{n_i} |\langle \mathbf{0}, \mathcal{M}, n_i | \varepsilon \cdot J(0)\, |\mathrm{vac}\rangle|^2$$

$$= \frac{\mathcal{M}^4}{(\mathcal{M}^2+Q^2)^2} \, \frac{\sigma_{\mathrm{tot}}(e^+e^- \to X(\mathcal{M}))}{4\pi^2\alpha}\,. \tag{2.6}$$

Briefly, the following is known about the e^+e^- annihilation experiments. For masses below about 1.5 GeV, the resonances associated with the familiar vector mesons ϱ^0, ω, and ϕ completely dominate the final states, confirming that they are the most important low mass members of the photon's hadronic constituent. Additional very broad resonances have been observed (possibly excited states of these familiar vector mesons) and very recently some very narrow resonances and threshold behavior at high energies have been discovered (new particles (J/ψ), and probably new physics). At higher masses, most theoretical expectations had been that $P(\mathcal{M}^2, 0)$ would decrease, probably as const/$\mathcal{M}^2$ (see, for example, Ref. [7]). This behavior known as "scaling" means the dependence of $\sigma_{\mathrm{tot}}(e^+e^- \to X(\mathcal{M}))$ can be simply predicted from dimensional analysis, ignoring the importance of any mass except $\mathcal{M}$, the center-of-mass energy. At the time of writing, it is not yet clear whether the expected scaling behavior will be attained, but it seems plausible that it will. In any case, it is clear that higher masses give a large contribution to the photon's structure.

Let us consider the low mass resonances in a little more detail. One might suspect that such states dominate (2.5) at low Q^2 because of the factor $1/(\mathcal{M}^2+Q^2)$. As indicated previously, the hypothesis that such resonances do dominate (2.5) is called VMD. If a resonance is very narrow, we may replace the $\mathcal{M}^2$ in the denominator of (2.5) by its value at the position of the resonance and take it outside the integral. The remaining integral, including contributions from all decay channels (n_i), is then *defined* to be the appropriate vector meson state. Thus, we find

$$\sqrt{\alpha}\, |h_S\rangle_{\mathrm{res}} = \sum_V \frac{e}{f_V} \frac{m_V^2}{m_V^2+Q^2} |V\rangle, \tag{2.7}$$

where $|V\rangle$ is an appropriately normalized state and the normalization constant is chosen by convention to be e/f_V for real photons. As will be discussed in Sec. 3, when the ϱ^0 resonance in the $\pi^+\pi^-$ channel is studied in detail [5], it is found that $|\varrho^0\rangle$ defined in this way corresponds to a spatially localized state. That is, it is very similar to any other physical particle. It is expected that the other vector meson states have the same property.

This naive VMD hypothesis (2.7) has great predictive power because it connects scattering amplitudes involving high energy photons, $\gamma N \to X$, to analogous amplitudes involving vector mesons, $V\mathrm{N} \to X$. We find, assuming the bare photon does not interact,

$$\langle X|S|\gamma \mathrm{N}\rangle = \sum_V \frac{e}{\bar{f}_V} \frac{m_V^2}{Q^2+m_V^2} \langle X|S|V\mathrm{N}\rangle, \tag{2.8}$$

where "S" is the S matrix operator of formal scattering theory. When restricted to the ϱ^0 mode, (2.8) expresses the "ϱ^0-photon" analogy. Eq. (2.8) is illustrated by Fig. 1b.

Returning now to some general consequences of VMD, by measuring the reactions $e^+e^- \to \pi^+\pi^-$, $\pi^-\pi^0\pi^+$, K^+K^-, and $K\bar{K}$ one determines the important parameters $\bar{f}_\varrho$, $\bar{f}_\omega$ and $\bar{f}_\phi$, and also makes internal tests of VMD. For example, studying the "bump" in the $\pi^+\pi^-$ mode attributed to the ϱ^0, it is readily seen that its width is given by the $\varrho - \pi^+\pi^-$ coupling strength $f_{\varrho\pi\pi}$, and that its overall normalization is determined by the $\gamma - \varrho^0$ coupling strength $\left(\frac{e}{\bar{f}_\varrho} m_\varrho^2 \text{ in VMD}\right)$. These two couplings are related by the pion form factor at $\mathscr{M}^2 = 0$. The data seem to be quite consistent with this normalization/width relationship predicted by VMD [8], when account is taken of the finite width effects [9]. Analyses for $\bar{f}_\omega$ and $\bar{f}_\phi$ in the other modes are similar but complicated somewhat by a variety of branching ratios.

A process of great interest to us in this paper, which is very easy to analyze from the viewpoint of the photon's hadronic component, is the photoproduction of the vector mesons themselves. The target modifies the hadronic constituents (essentially by absorption

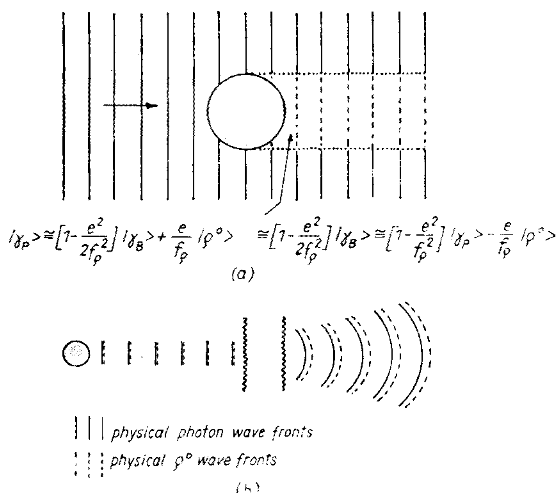

Fig. 2. Compton scattering and photoproduction of ϱ^0 as a version of shadow scattering. (a) In the shadow region, the ϱ^0 component is absorbed out, leaving a bare photon. The state may be re-expressed in terms of physical particle states. (b) The modified portion of the wave evolves and gives real Compton scattering and ϱ^0 production. These two amplitudes are related to the ϱ^0 shadow scattering amplitude through VMD coefficients. Of course, other photon constituents also contribute to Compton scattering

at small impact parameters) so that behind the target the state corresponds to a slightly attenuated incident photon state plus physical hadrons. The mechanism is illustrated in Fig. 2a for the ϱ^0 part of the photon. If we now make the further assumption that constituents are not mixed by the scattering, i.e. the "diagonal approximation", we can view ϱ^0 production by an incident real photon as shown in Fig. 2b. Under these circumstances then, we clearly have the result

$$f_{\gamma V} = \frac{e}{\bar{f}_V} f_{VV}, \tag{2.9}$$

for any vector meson V in the photon's hadronic component. In practice we sometimes define a phenomenological photoproduction coupling by $\hat{f}_V$ in place of $\bar{f}_V$ in (2.9). The difference between these allows for extrapolation between the vector meson and photon masses, as well as for non-VMD type processes. Comparison of experimental values of $\bar{f}_V^2/4\pi$ and $\hat{f}_V^2/4\pi$ then provides a test of VMD. In the case of the ϱ^0, the VMD relation is quite good; for the ω and ϕ, the situation is still unclear and is complicated by $\omega-\phi$ mixing.

3. *The dipion constituent of the photon*

There are two reasons why it is important to study the dipion constituent of the photon. The first is that it dominates the real photon cross section. Since it includes the ϱ^0 contribution, we may make the simple VMD estimate

$$\sigma_\gamma^{(2\pi)} \cong \frac{e^2}{\bar{f}_\varrho^2}\, \sigma_\varrho, \tag{3.1}$$

where $\sigma_\varrho \cong 27$ mb and $\bar{f}_\varrho^2/4\pi \cong 2.5$, leading to about 80 μb, out of a total cross section of ~ 120 μb for $E_\gamma = 6$ GeV. The second reason is that the dipion constituent is more amenable to analysis than some of the other constituents (although the same analysis would be directly valid for K pairs). In particular, it will be possible to find out something of the internal spatial structure of the dipion system and show explicitly the "photon--shrinking" effect with increasing Q^2 [10]. This analysis will provide some intuitive guidance about the possible behavior of other constituents.

We assume a specific model for the dipion matrix element in (2.6), namely that it is given by VMD: the photon couples to the ϱ^0 which propagates and finally decays to a pion pair. The ϱ^0 propagator is modified to take into account the vacuum polarization bubbles due to the dissociation of the ϱ^0 into pion pairs. Thus we have

$$\langle \boldsymbol{q}_+ \boldsymbol{q}_- - | \varepsilon \cdot \boldsymbol{J}(0) | \text{vac} \rangle = \boldsymbol{\varepsilon} \cdot (\boldsymbol{q}_+ - \boldsymbol{q}_-) F_\pi(\mathcal{M}^2)/(2\pi)^3,$$

where

$$F_\pi(\mathcal{M}^2) = \frac{-(m_\varrho^2 - \Pi(0))}{\mathcal{M}^2 - m_\varrho^2 + \Pi(\mathcal{M}^2)}. \tag{3.2}$$

The coupling of the ϱ^0 to the pion pairs is called $f_{\varrho\pi\pi}$, while the coupling of the photon to the ϱ^0 is called $-em_\varrho^2/\bar{f}_\varrho$. Their product is adjusted so that $f_{\varrho\pi\pi}/\bar{f}_\varrho = [1-\Pi(0)/m_\varrho^2]$, making the pion form factor 1 at zero momentum transfer. The vacuum polarization effects are incorporated in $\Pi(\mathcal{M}^2)$, whose imaginary part gives the width, according to

$$\operatorname{Im} \Pi(\mathcal{M}^2) = m_\varrho \Gamma_\varrho(\mathcal{M}^2) = \tfrac{1}{3} \frac{f_{\varrho\pi\pi}^2}{4\pi} \frac{q_\pi^3}{\omega_\pi}, \tag{3.3}$$

where $\omega_\pi = \frac{1}{2}\mathcal{M}$ and $q_\pi = (\omega_\pi^2 - m_\pi^2)^{1/2}$. The real part of Π is related to the imaginary part by a twice subtracted dispersion relation chosen so that the propagator has a simple normalization at the mass of the ϱ^0:

$$\operatorname{Re} \Pi(m_\varrho^2) = \operatorname{Re} \Pi'(m_\varrho^2) = 0. \tag{3.4}$$

Using the form (3.3), it is possible to evaluate Π explicitly using dispersion relations (or equivalent techniques). This has been done by Gounaris and Sakurai and Vaughn and Wali [9] (with somewhat different normalization conventions), and it turns out that the additional factor $(1-\Pi(0)/m_\varrho^2)$ relative to VMD is precisely the normalization correction they obtained in their analysis of $e^+e^- \to \pi^+\pi^-$.

We summarize the results of this model [5]:

(1) The total dipion probability, which is given by

$$P_{2\pi} = \frac{e^2}{\pi \bar{f}_\varrho^2} \int_{\mathcal{M}_{\text{th}}^2}^{\infty} \frac{d\mathcal{M}^2}{(\mathcal{M}^2+Q^2)^2} \frac{m_\varrho^4 \operatorname{Im} \Pi(\mathcal{M}^2)}{|\mathcal{M}^2 - m_\varrho^2 + \Pi(\mathcal{M}^2)|^2}$$

is easily evaluated by contour integration with the result

$$P_{2\pi}(Q^2) = \frac{e^2}{\bar{f}_\varrho^2} \frac{[1+\Pi'(-Q^2)]\,[F_\pi(-Q^2)]^2}{[1-\Pi(0)/m_\varrho^2]^2}. \tag{3.5}$$

Π and Π' are readily evaluated, and for $Q^2 = 0$ one finds approximately

$$P_{2\pi}(0) = \frac{e^2}{\bar{f}_\varrho^2}(1-0.39\Gamma_\varrho/m_\varrho), \tag{3.6}$$

or nearly the VMD result. However, as will be seen later, there is no unique separation of this probability into a pure ϱ^0 term plus a non-resonant contribution. Aside from finite width corrections, the Q^2 dependent probability

$$P_{2\pi}(Q^2) \sim \frac{e^2}{\bar{f}_\varrho^2}\left(\frac{m_\varrho^2}{m_\varrho^2+Q^2}\right)^2, \tag{3.7}$$

which is the VMD expression.

It should be remarked that had we ignored the existence of the ϱ^0 resonance, $P_{2\pi}$ would have turned out to be infinite (logarithmically divergent). Thus the resonance causes the dipion contribution to saturate at relatively low energies. However, it will become increasingly clear that it is not correct to assume that the non-resonant dipion continuum has been completely replaced by the ϱ^0. The new picture is that the dipion contribution should be thought of mainly as a ϱ^0 meson with a little bit of non-resonant two-pion state attached.

(2) To bring out this physical picture more clearly, we next study the internal spatial structure of the dipion state. There is of course no way to do this at extremely small distances, where the strong interactions will completely obscure the picture. Instead, we make the assumption that outside some radius the pions are sufficiently well separated that it makes sense to use a two pion scattering wave function. This might become reasonable when the distance between pions is greater than about 1 F. In the center-of-mass frame for the two pion system, the P-state wave function takes the form

$$|\boldsymbol{q}_+\boldsymbol{q}_-\pm\rangle \sim 2\omega_\pi \hat{q}_\pi \cdot \hat{r} e^{\pm i\delta_1} \frac{\partial}{\partial(q_\pi r)} \frac{\sin(q_\pi r+\delta_1)}{q_\pi r},$$

where the factor $2\omega_\pi$ has been introduced because of the relativistic normalization used, and r is the spatial distance between the pions. With this ansatz and certain further approximations, it is possible to evaluate this part of the wave function in the rest frame of the dipion system. The remarkable result is

$$f(\boldsymbol{r}) = F_\pi(-Q^2)\boldsymbol{\varepsilon}\cdot\nabla\left\{\frac{1}{r}\exp\left[-(m_\pi^2+\tfrac{1}{4}Q^2)^{1/2}r\right]\right\}. \tag{3.8}$$

To see the significance of this result, we first set $Q^2 = 0$ and note that the resulting wave function is completely independent of the ϱ^0 meson! Physically, this means that the probability associated with the ϱ^0 is contained entirely inside the interaction region. This result could have been anticipated from the theory of decaying states, in which a superposition of energy eigenfunctions from the vicinity of a resonance leads to a localized state. If we had neglected the strong interactions of the pions, the result (3.8) would have been valid at all radii (except that some assumptions would have been less justified at smaller radii where higher masses would be much more important). The divergence of $P_{2\pi}$ in that case is undoubtedly associated with the strong singularity in the wave function. As a result of the interactions, this singular part of the wave function is somehow eliminated in favor of the ϱ^0 structure. This reinforces the approximate picture that the dipion constituent of the photon may be regarded as a ϱ^0 meson core surrounded by a two pion cloud. However, the dipion constituent is a complete unit and there is no way to make a unique physical separation into these two parts.

The next point of interest is the photon shrinking which is evident in (3.8) [10]. As Q^2 increases, the two pion tail of the photon is rapidly drawn in. The range is again independent of the ϱ^0, but the strength depends on the ϱ^0 through the pion form factor. While it cannot be proved by this type of argument, it seems plausible that the whole structure may continue to shrink as Q^2 increases. This means that although one may choose to regard the ϱ^0 as a well-defined object of definite internal size, it is still possible to superpose states, including the ϱ^0, which have a spatially smaller structure. *If this view is correct, then the* *hadronic component of the photon cannot be regarded a sum of separately interacting constituents. Rather, all the constituents would have to act in unison, corresponding to an object whose structure, and hence interactions depends on Q^2 in a non-trivial way.* This would be in distinction to the diagonal assumption of the generalized VMD hypothesis, which has eachcomponent interacting independently in the total cross section. In a sense, it is possible that as far as interactions are concerned *the whole is less than the sum of its parts.*

(3) What are the experimental consequences of the dipion part of the photon's structure? It has already been pointed out that it gives approximately two-thirds of the total photon cross section. The dipion component also shows up as diffractive photoproduction of $\pi^+\pi^-$ pairs. This is a very important process experimentally as it accounts for about 15% of the total photon cross section, or about 20% of the isovector part of the total

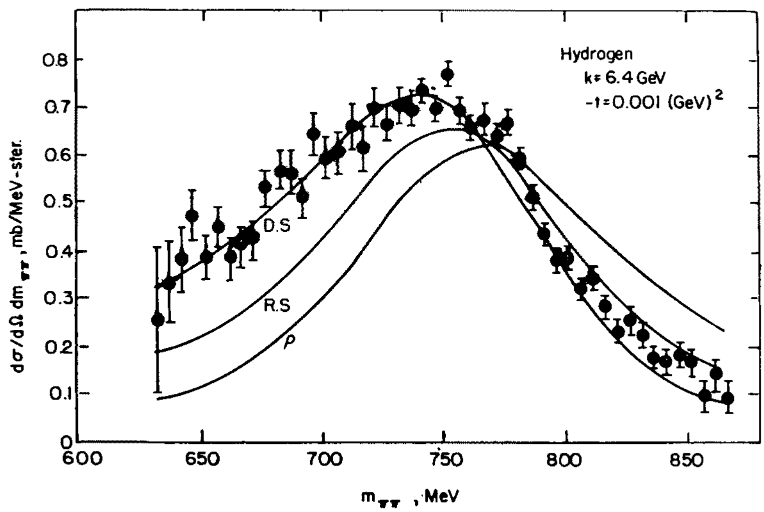

Fig. 3. The mass distribution of pion pairs in photoproduction is compared with various models arising from the hadronic structure of the photon

photon cross section. This corresponds reasonably well to the usual ratio of hadronic total elastic cross section to total cross section. In more detail, the observed mass distribution of pion pairs is qualitatively similar to that in the original structure of the photon. An example is shown in Fig. 3; the data are from the DESY-MIT group [11]. The curve labeled ϱ^0 represents the pure ϱ^0 part of the data as fitted by Spital and Yennie [12]. It is obvious that the resonance peak is strongly skewed toward lower masses by an interfering background. The simplest model for such a mass distribution would be that each mass component in the photon is individually absorbed by the target (i.e. the diagonal assumption). If all components experienced the same absorption, the resulting amplitude for pair production would be

$$\text{M.E.} \propto \frac{1}{\mathcal{M}^2}\,\frac{-m_\varrho^2}{\mathcal{M}^2-m_\varrho^2+\Pi(\mathcal{M}^2)}\,. \tag{3.9}$$

The square of this leads to the phenomenological Ross–Stodolsky formula for the mass distribution [13]. This is shown in Fig. 3 as the curve labeled R–S.

Evidently, while the Ross–Stodolsky formula does have the right qualitative behavior, it is not quantitatively very good. We next describe the connection of (3.9) with another popular treatment of the mass distribution, namely the Söding model [14]. A simple rearrangement of (3.9) gives

$$\text{M.E.} \propto \frac{-1}{\mathcal{M}^2 - m_\varrho^2 + \Pi(\mathcal{M}^2)} + \frac{1}{\mathcal{M}^2}\, \frac{\mathcal{M}^2 - m_\varrho^2}{\mathcal{M}^2 - m_\varrho^2 + \Pi(\mathcal{M}^2)}\,. \tag{3.10}$$

The first term of (3.10) may be interpreted as the pure ϱ^0 term. For example, it can be shown that it does not contribute to the two-pion tail in configuration space. In configuration space, the entire contribution from the first term of (3.10) is contained inside the interaction region.

The second term of (3.10) closely resembles the usual Drell amplitude [15] for pair production as modified by the double counting correction of Bauer and Pumplin [16]. Because of their different spatial structure, we would expect these two contributions to be differently absorbed. The Söding model accounts for this by assigning the two amplitudes independent strengths, interpreted as being proportional to σ_ϱ and $\sigma_{\pi^+} + \sigma_{\pi^-}$, respectively; in fact, the second term does turn out experimentally to be enhanced relative to the first. The curves, labeled D–S in Fig. 3 correspond to this model, with $\sigma_{\pi^+} + \sigma_{\pi^-}$ taken to be $2\sigma_\varrho$ for simplicity; it fits the data quite well.

It is clear that this general picture gives a good account of the main feature of pion pair photoproduction from nucleons and nuclei. However, it is equally clear that there is very little hope of constructing a theory which will predict the mass spectrum perfectly as a function of energy, momentum transfer, and nucleus. Such a theory would treat the dipion component as a unit rather than make an artificial distinction between a pure ϱ^0 and nonresonant pions. Therefore, it has been necessary to resort to a phenomenological description for the purpose of interpreting experiments (Spital and Yennie [12]). This description concentrates on determining the ϱ^0 cross section and gives up any attempt to extract useful quantitative information from the background region.

(4) The other piece of experimental evidence for the dipion structure is given by the inclusive pion spectrum. The picture is that one pion hits the target and interacts in some way which is not detected while the other flies free and is detected. This is reminiscent of the original Drell process [15]. It is easy to see that Drell's result corresponds to precisely this picture. His original formula was

$$d\sigma_{\gamma\to\pi^-} = \frac{\alpha}{2\pi}\, \frac{\sin^2\theta}{(1-\beta\cos\theta)^2}\, \frac{d\Omega}{4\pi}\, \frac{\omega(k-\omega)}{k^3}\, d\omega\, \sigma_{\pi^+}^{\text{tot}}. \tag{3.11}$$

Without dificulty, this may be re-expressed in term of the intermediate two-pion total mass with the result

$$d\sigma_{\gamma\to\pi^-} \cong \frac{\alpha}{2\pi}\, \frac{d\mathcal{M}^2}{\mathcal{M}^4}\, q_\perp^2\, \frac{dq_\parallel}{k}\, \sigma_{\pi^+}^{\text{tot}} \tag{3.12}$$

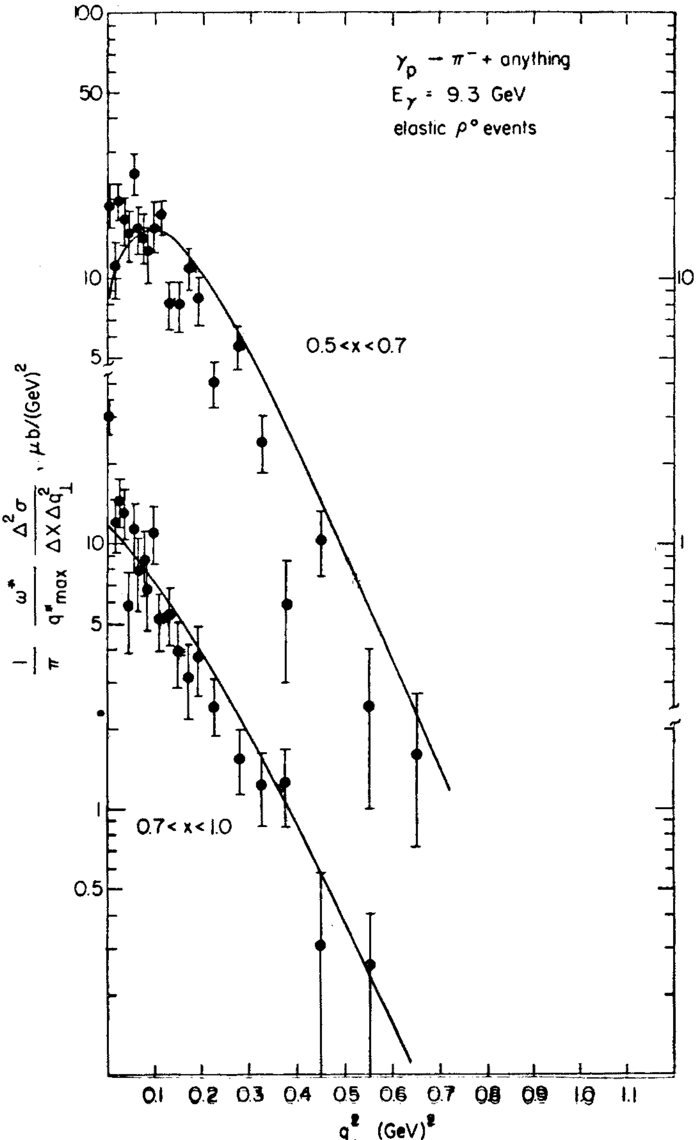

Fig. 4. Inclusive π^- photoproduction, ϱ^0 events ($\gamma p \to \pi^+\pi^- p$, $m_{\pi^+\pi^-} < 1$ GeV) only

which is precisely what we would have anticipated from (2.5), omitting the effect of the ϱ^0. It seems very strange that the ϱ^0 should not appear in this expression, but a more complete analysis [5] taking the ϱ^0 into account yields the same result. The result of such an exercise, hovever is that it makes one suspicious about using Drell's expression for large $\mathscr{M}^2$ ($\lesssim 2$ GeV^2, say) without providing any guidance as to the mechanism which cuts it off.

Next, we turn to the question of confirming the presence of this process in the data. So far this has been done only in a very rough way, but it should be possible to improve

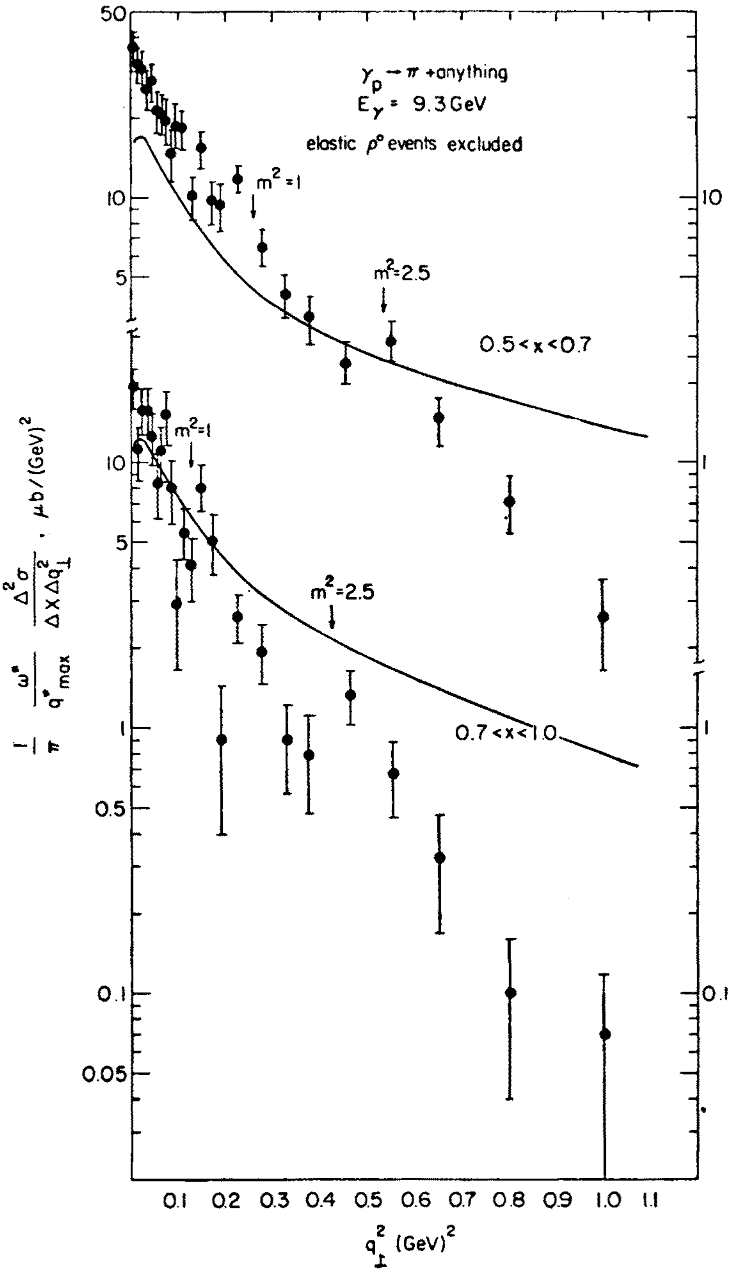

Fig. 5. Inclusive π^- photoproduction with ϱ^0 events excluded. The theoretical curves are for the Drell process and a few typical values of $\mathscr{M}^2$ for the dipion are indicated

on the treatment which will be described here. Intuitively it seems clear that the experimental region where this mechanism should dominate is the one where the π^- carries off most of the energy of the photon. In any case, if the process does not show up in an important way in this region, there is certainly no hope that it could be important for lower energy pions. There are simply too many mechanisms for producing π^- mesons with a small fraction of the photon's momentum, whether or not we believe in complete dominance by the hadronic component of the photon.

In order to compare theory and data (9.3 GeV data from Moffeit et al. [17]), it is also necessary to take into account the special contribution from diffractive production of pion pairs, particularly through the ϱ^0. The process

$$\gamma + P \to \varrho^0 + P \to \pi^+ + \pi^-$$

actually contributes about half the pions seen near the upper end of the spectrum. This part of the data has been separated out experimentally [17] and is shown in Fig. 4, along with a crude theoretical estimate. A similar estimate has been made of the "non-ϱ^0" events in Fig. 5. While this cannot be regarded as a striking fit to the data, it certainly indicates the presence of the dipion component of the photon, particularly in the high-x region where it is expected to dominate.

4. *Space-time view of photon interactions*

Our discussion so far has been phrased in the language of stationary state perturbation theory. Our intuitive discussion of photon interactions assumed tacitly a *separability* of the problem into two stages: The internal structure of the photon is first worked out independently of the target, and then the various components are permitted to interact with the target. For this picture to be valid, the virtual hadrons must be present for a time sufficiently long so as to undergo collisions as if they were ordinary, "real" hadrons. They must be able to propagate like physical particles into the target as if they had become detached from the photon. They would then reflect their photonic origin only in their symmetry quantum numbers and the specifics of the superposition.

In free space, the virtual hadrons never could become detached. The physical photon is always making transitions back and forth between a bare photon and hadronic states in such a way as to produce the steady state (2.1). Such transitions are called vacuum polarization fluctuations, and we may regard interactions with the target as arising from the interception of a photon in its temporary hadronic state. For these interactions to satisfy the separability conditions it is necessary that the fluctuation lasts a typical time (called the formation time) which is relatively long compared to the interaction time with the target. The formation time may be estimated from a venerable uncertainty principle argument to be

$$t_f \sim \left| \frac{1}{\nu - E_n} \right| \sim \frac{2\nu}{Q^2 + \mathscr{M}^2}. \tag{4.1}$$

For fixed Q^2 and $\mathscr{M}^2$, this is directly proportional to the photon's energy. For example, the ϱ^0 component inside a 10 GeV photon has a formation distance of order 7 F.

In this example, the distance over which the photon can travel in its hadronic phase is much larger than the typical dimensions of hadrons. Since the interaction time is likely to be characterized by the nucleon radius, the separability condition should be well satisfied for contributions arising from the ϱ^0 constituent. When the formation time becomes shorter (large $\mathcal{M}^2$ and/or large Q^2), the hadron-mediated interactions may become practically indistinguishable from interactions due to the bare photon term.

Spital and Yennie [18] have studied the separation of photon interactions into direct and hadron-mediated terms using time-dependent perturbation theory. They show that the hadron-mediated term is a superposition over hadronic states of the amplitude for the

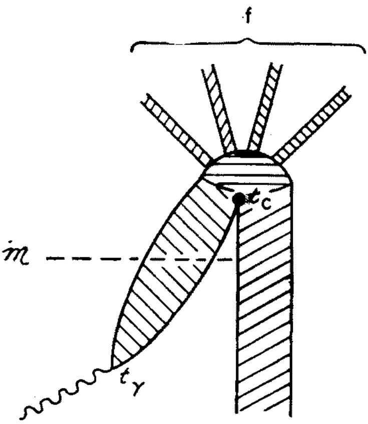

Fig. 6. Space-time illustration of the process $\gamma p \to f$ as mediated by the hadronic structure of the photon. The photon converts to hadrons at time t_γ and those hadrons start to interact with the target at time t_c

photon to contain a hadron times the amplitude for that hadron to interact. The space--time picture of such interactions is illustrated in Fig. 6. A photon converts at time t_γ to hadrons of mass m which subsequently interact with the target at time t_c. A prior integration over space assures that the hadron momentum is the same as the photon's momentum $\boldsymbol{k}$. The main contribution to this integral comes from the region neat the light cone along the direction of the photon. The time-dependence of this amplitude is given simply by

$$e^{i(t_c - t_\gamma)(\nu - \sqrt{k^2 + \mathcal{M}^2})} \sim e^{-i(t_c - t_\gamma)(Q^2 + \mathcal{M}^2)/2\nu} \equiv e^{-i(t_c - t_\gamma)/t_f}. \tag{4.2}$$

Thus we see that t_f gives the rate of oscillation of this time dependence. Integration of t_γ from $-\infty$ to t_c (with suitable damping) yields

$$\frac{t_f}{i} = \frac{1}{i} \frac{2\nu}{(Q^2 + \mathcal{M}^2)},$$

which, aside from a factor i^{-1}, is just the energy denominator appearing in (2.2). We see from this that t_f plays the role of a *coherence time* — although the integration extends to $-\infty$, only hadrons produced within an interval t_f can act together in an interaction.

The discussion so far is somewhat artificial in that constituents of different masses are likely to interfere in producing the final state $|f\rangle$; i.e., we must integrate over the masses appearing in (4.2). To do this in general, we need some information about the mass-dependence of the two factors associated with Fig. 6. The first factor, which is the amplitude for the photon to convert to a state of mass $\mathcal{M}$, is simply $\langle n^+|H'|\overset{0}{\gamma}\rangle$ as given in (2.2) and information about it is available from $e^+e^- \to \mathcal{M}$, as was discussed in Sec. 2. In the case of the dipion, we even have a fairly good idea of the general structure of the state which is present just before the target. The second factor is more problematic. To the extent that the mediating hadronic states are dominated by simple resonances (ϱ^0, ω, ϕ) and the final state is simple, it may be possible to obtain information about it by studying inverse reactions in which a resonance is produced. For the moment, we simply note that both factors, as well as the phase space of the internal integration, can produce an $\mathcal{M}^2$--dependence. Thus we have argued the plausibility for an expression of the form

$$\int_{-\infty}^{t_c} dt_\gamma \int_{\mathcal{M}_t^2}^{\text{,,}\infty\text{``}} \frac{d\mathcal{M}^2}{2\sqrt{k^2+\mathcal{M}^2}}\, e^{-i(t_c-t_\gamma)(\mathcal{M}^2+Q^2)/2\nu} G_f(\mathcal{M}^2). \tag{4.3}$$

If the t_γ-integration is carried out first, we obtain

$$-i \int_{\mathcal{M}_t^2}^{\text{,,}\infty\text{``}} \frac{d\mathcal{M}^2}{\mathcal{M}^2+Q^2}\, G_f(\mathcal{M}^2) \tag{4.4}$$

which has an obvious plausible interpretation. In particular, if G_f peaks strongly at various resonances, we recover the VMD expression (2.8).

It is interesting to study the effect of the mass integration on the time dependence. In general, this depends on the form of G_f, which would be expected to vary with the process under consideration. There is one situation in which we feel we have an understanding of this mass dependence, namely when the mediating system is in resonance. In this case, we expect

$$G_f \approx \frac{\frac{1}{\pi}\, h_f(k) m_V \Gamma_V}{(\mathcal{M}^2-m_V^2)^2+m_V^2\Gamma_V^2}. \tag{4.5}$$

(Briefly: the matrix element for the photon to convert provides a factor $D = (m^2-m_V^2 - im_V\Gamma_V)^{-1}$ and the particles propagating into the target provide a factor D^*. The factor $m_V\Gamma_V$ comes from the internal phase space integration. h_f represents the interaction of the resonance with the target to produce the final state f.) When this expression is inserted into (4.3), the $\mathcal{M}^2$ integration may be *estimated* to be

$$\int_{-\infty}^{t_c} \frac{dt_\gamma}{2\sqrt{k^2+m_V^2}}\, e^{-i(t_c-t_\gamma)\left(\frac{Q^2+m_V^2}{2\nu} - i\frac{m_V\Gamma_V}{2\nu}\right)}. \tag{4.6}$$

We see that the effect of the mass integration has led to an exponential damping of the integral in the past. (Note: As in the theory of decaying states, there are also small terms which decrease as inverse powers of $(t_c - t_y)$ when a more precise calculation is made; these terms are associated with threshold regions in $\mathcal{M}^2$). As long as $m_V\Gamma_V \ll Q^2+m_V^2$, the most relevant time is still the formation time. In fact, if the time integration is carried out now, the result is

$$-i\,\frac{h_f}{Q^2+m_V^2-im_V\Gamma_V} \tag{4.7}$$

which is essentially the VMD expression (the $-im_V\Gamma_V$ is spurious; it results from some of our approximations).

More generally, we expect each process and each mediating channel to be associated with some build-up time ($2\nu/m_V\Gamma_V$ in the resonance case). Roughly speaking, if a range of masses $\Delta\mathcal{M}^2$ contributes together, this buildup time should be given by a kind of uncertainty relation

$$t_b \sim \frac{2\nu}{\Delta\mathcal{M}^2}\,. \tag{4.8}$$

As long as this time is much larger than the formation time (i.e. $\Delta\mathcal{M}^2 \ll \mathcal{M}^2+Q^2$) the picture of hadron-mediated interactions should be valid: The photon converts in a leisurely way to hadrons over a long time, t_b, and those hadrons appearing within a time t_f act together on the target to produce the final state f. On the other hand, if in a given process the build-up time were to become comparable or smaller than the formation time ($\Delta\mathcal{M}^2 \sim \sim \mathcal{M}^2+Q^2$), this picture breaks down and our usual intuition fails. This could well be the situation with the constituents associated with the scaling region in $e^+e^- \to X(\mathcal{M}^2)$, where not a great deal of structure is apparent (aside from the new phenomena). The mediating hadrons in this region may bear no resemblance to the usual ones, and the superposition may make them appear as some more elemental matter, such as quark pairs.

We note that each process serves as a kind of filter, selecting out of the photon the particular constituents it needs. Only the resonance contributions would appear to have fairly well-defined build-up and formation times. Thus there is no unique time associated with the mediating hadronic matter. The effects of superposition will require the development of new intuition. This has been discussed already in terms of the shrinking photon. It could be that as Q^2 increases, the whole dipion structure shrinks — that is, the structure as a whole might become smaller than the ϱ^0 alone, due to cancellations. This should have experimental consequences: (1) In electroproduction of ϱ^0's on hydrogen, the slope parameter B (in e^{Bt}) should decrease with Q^2. The present published data are inconclusive, but you should watch for some data from SLAC which will appear soon. (2) The total cross section for ϱ^0 might decrease more rapidly with Q^2 than as predicted by VMD. There are indications of this, but the situation is confused because of $t_{\min}$ effects and uncertainties in how to handle certain kinematic factors. (3) Electroproduction of ϱ^0's on nuclei (A-dependence) might reveal that σ_ϱ depends on Q^2. The point is that the matter produced in the nucleus might propagate out before separating into the ϱ^0 and other constituents; it might have a smaller absorption cross section as a unit. Incidentally, this experiment would be interesting as well since the cross sections for longitudinal ϱ^0's apparently is smaller than that for transverse ϱ^0's, as seen from production on hydrogen. This could be related to the internal geometrical structure of the ϱ^0, which might retain some features of the two-pion state suggested by (3.8).

There are two ways of making manifest this long-range behaviour of hadron-mediated interactions. One is through photon shadowing in nuclei. The experimental situation is summarized by Professor Friedman in his lectures. While the effect appears to be present, especially for real photons, the experimental situation is still somewhat confused. The other is through its consequences for the total photon cross sections and the deep--inelastic structure functions, which we turn to now.

The argument begins with the fact that the propagation of light through matter is completely determined by the commutation properties of the electromagnetic current density J_μ. This permits us to relate the imaginary part of the forward Compton scattering amplitude (or, equivalently, through the optical theorem, the photon-nucleon cross section) to the space-time behavior of the absorption and emission of photons. The derivation is quite well-known, and we shall not reproduce it here. Rather we shall describe briefly the arguments by which the long-range behavior is inferred from the data.

The result of the analysis is that the forward Compton scattering amplitude may be expressed as a Fourier transform of a matrix element of the commutator of two current operators. Then, using various symmetry properties, it is possible to express the inelastic structure functions W_1 and W_2 as space-time integrals. For example,

$$\nu W_2(k^2, \nu) = 2Q^2 \nu M \int d^4y \sin(k \cdot y)\theta(y \cdot P)\theta(y^2)f_2(y^2, y \cdot P). \tag{4.9}$$

y is the 4-vector space-time separation between the points of absorption and re-emission of the virtual photon, and $y \cdot P\,(=y_0 M$ in the laboratory frame) is called the "range".

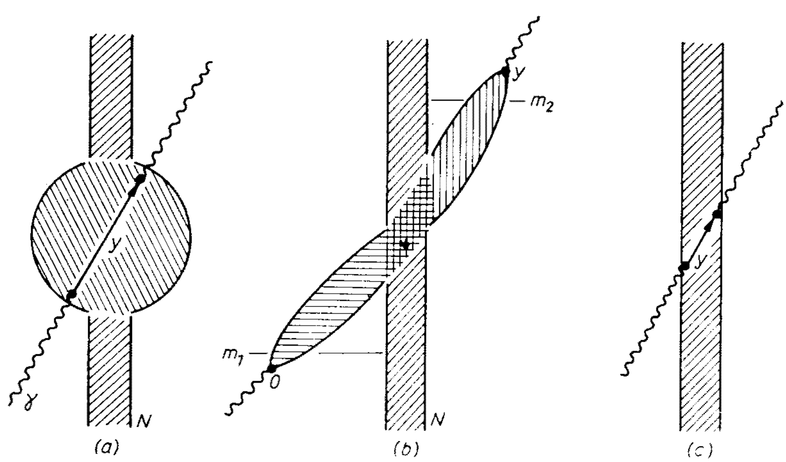

Fig. 7. The space-time dependence of forward Compton scattering. (a) The complete amplitude. (b) The hadron-mediated contribution. (c) The short-range contribution

This is illustrated in Fig. 7. Note that y is *not* the difference between the coordinate of the photon's disappearance and the nucleon position.

Simplest analyses of (4.9) assume $f_2(y^2, y \cdot P)$ is regular and non-vanishing near the "light cone" ($y^2 = 0$). The singularity structure of the current commutator that leads to this behavior is "canonical". (If one evaluates the commutator with free field operators, the singularities obtained are called canonical.) These assumptions with $f_2(0, y \cdot P)$ a non-

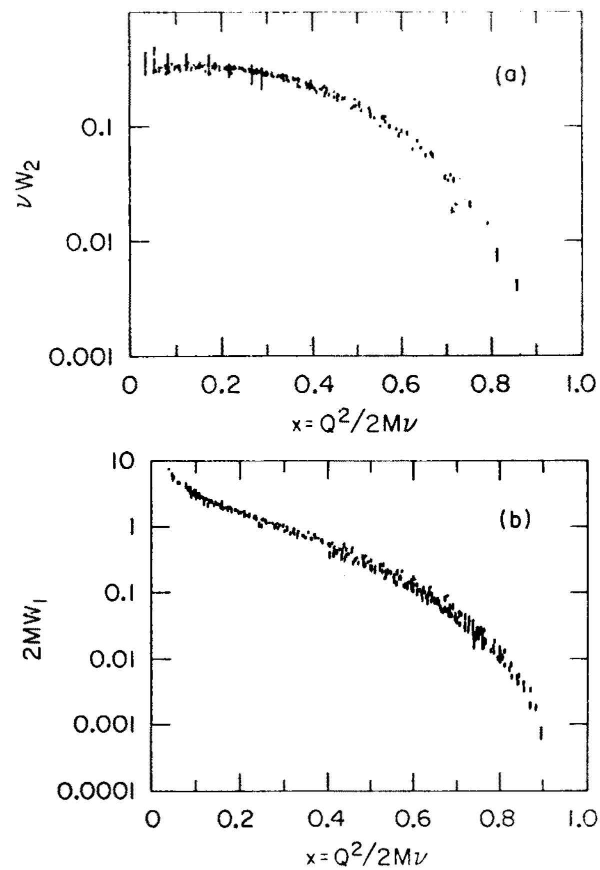

Fig. 8. Example of the scaling of $2MW_1$ and νW_2 [19]

-trivial function of $y \cdot P$ lead to Bjorken scaling [19]. This means that as $Q^2, \nu \to \infty$ with $x = Q^2/2M\nu$ fixed the structure functions MW_1 and νW_2 become functions $F_1(x)$ and $F_2(x)$ and do not depend on Q^2 and ν separately. The connection between canonical dimensions and scaling has been shown by Brown, Jackiw, van Royen, and West, Brandt, and Leutwyler and Stern [20]. Fig. 8 illustrates that the data appear to have this behavior, at least over the SLAC kinematical range.

We are particularly interested in the long range (large $y \cdot P$) behavior of f_2, and its possible connection with the hadronic structure of the photon. Ioffe [21] first showed that the experimental data on inelastic electron scattering required a long range. The general idea of his argument is as follows. Consider the exponent of (4.9) and rewrite $k_3 \approx \nu + Q^2/2\nu$ with $|\boldsymbol{k}| = k_3$. Then

$$k \cdot y \cong \nu(y_0 - y_3) - \frac{Q^2}{2\nu}\, y_3. \tag{4.10}$$

For large ν, rapid oscillations in the $y_0 - y_3$ dependence force the major contribution to come from the vicinity of the light cone. If, in addition, f_2 were short-range in $y \cdot P$ (hence in y_3), the result of the integration would depend on Q^2 only in the combination $Q^2/2\nu$. Ioffe showed that this is incompatible with certain features of the data and hence concluded that f_2 must be long-ranged. Pestieau, Roy, and Terazawa [22] then made a more quantitative

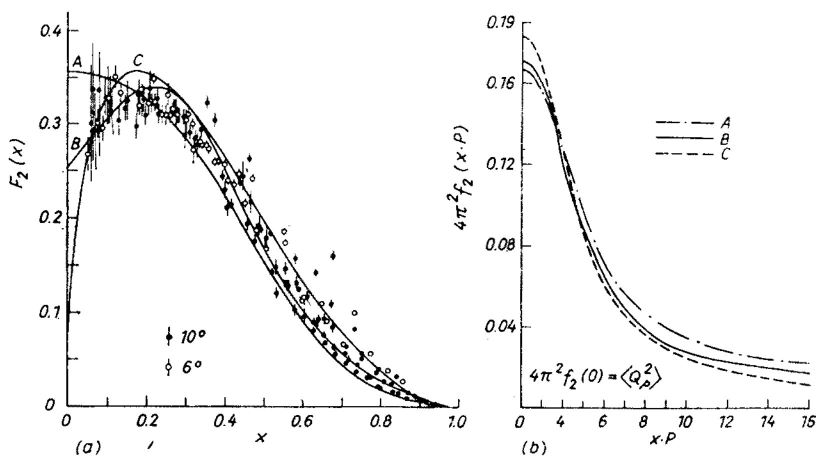

Fig. 9. From the space-time analysis by Pestieau, Roy, and Terazawa [22]. (a) Possible scaling fits to νW_2. (b) The dependence of $f_2(0, y \cdot p)$ for each of the fits

analysis using a formula of Jackiw, van Royen and West [20] for inverting the scaling limit $F_2(x)$ of νW_2 to find $f_2(0, y \cdot P)$. Their results depend, of course, on the assumed form of F_2 since that quantity could only be guessed from finite energy data. The long range behavior of f_2 turned out to be extremely sensitive to the small x behavior of $F_2(x)$, as might have been expected from Ioffe's argument. Their results are shown in Fig. 9. If $\lim_{x \to 0} F_2(x) \neq 0$, the result is that $f_2 \propto y_0^{-1}$ for large y_0. Such a behavior is necessary to overcome the factor of Q^2 appearing explicitly in (4.9). For smaller y_0, the function changes behavior at $y_0 \sim R_p$ and remains finite as $y_0 \to 0$.

The study of the long-range behavior was elaborated further by suri and Yennie [23]. They draw the following conclusions:

(1) If f_2 is short-ranged, F_2 approaches 0 at least as fast as x^2 for small x. Experimental data plotted vs x is shown in Fig. 8. While it suggests validity of scaling, it appears more nearly to approach a constant for small x, and is certainly not proportional to x^2.

(2) If f_2 is short-ranged, the real photon cross section should drop as $1/\nu^2$ for high energies. Present data indicate that $\sigma_\gamma(\nu)$ approaches a positive constant ($\sim 100\ \mu b$) as $\nu \to \infty$.

(3) Since (1) and (2) appear to be in sharp disagreement with present data, long--range terms are necessary; and an asymptotic term like $h^{(1)}(y^2)/y_0$ (large y_0) will yield both a constant real cross section and an F_2 which is non-vanishing for small x.

$$f_2 \propto \frac{h(y^2)}{y_0}, \quad \sigma_\gamma = \text{const. and } \lim_{x\to 0} F_2(x) \neq 0. \tag{4.11}$$

(4) Long-range terms with the structure $h^{(a)}(y^2)/y_0^a(1 < a < 3)$ yield Regge type energy behavior of the total cross section: $1/\nu^{a-1}$.

To reinforce our conclusions that the long-range behavior expected from VMD or GVD is demanded by the data, let us look at the diffractive contributions in more detail. In

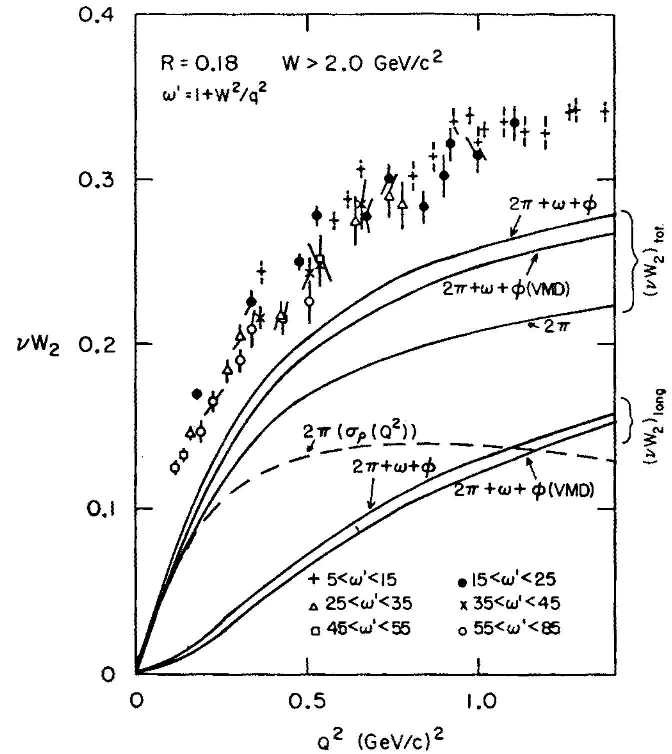

Fig. 10. νW_2 vs Q^2 in the small x region. The data is from Friedman and Kendall [24], and the various VMD contributions (now somewhat obsolete) are from Yennie [5]

the space-time analysis, these correspond to long-range terms with $a = 1$. We would expect that if x is not too large, it should be possible to ignore the short distance smoothing of f_2. The usual guess is that $(R_p Q^2/2\nu) \ll 1$ where R_p is the radius of the proton ($\sim 1/2m_\pi$).

Thus we require $x \ll \frac{2m_\pi}{M} = 0.3$. In practice, $x < 0.1$ seems adequate and we expect the pure diffractive term to dominate and yield

$$\nu W_2^{(\mathrm{diff})} = Q^2 G^{(1)}(Q^2). \tag{4.12}$$

A compilation of data plotted as function of Q^2 from this region is shown in Fig. 10 and they seem to fall along a universal curve independent of ν as expected. Of course, other Regge contributions, as well as short range smoothing, could be causing small systematic deviations from the ultimate curve (for $x \to 0$). The reader should note that Fig. 10 is not incompatible with scaling since it contains only small Q^2 data. Fig. 11 extends the data to higher Q^2 using the data of Watanabe et al. [24] (this figure was kindly supplied by S. Herb from his Cornell dissertation.) It appears that the high ω muon scattering data are

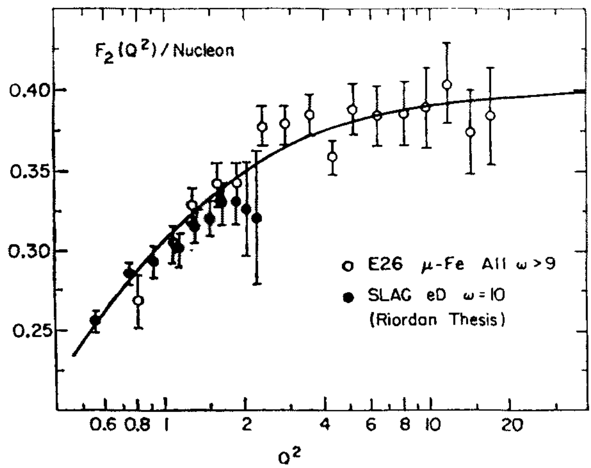

Fig. 11. νW_2 vs Q^2 in the small x region as extended to higher Q^2 by Herb

compatible with this diffractive picture and join smoothly with the older SLAC data. A more complete analysis along this line is being prepared by L. Hand for presentation at Trieste in July.

It seems plausible to interpret these long range ($a = 1$) terms with hadron-mediated interactions since the photon is never inside the nucleon as seen in the laboratory frame. Can one also verify that diffractive hadron-mediated terms have this long-range behavior? The answer is of course yes. This has been done by Pestieau and Urias [25] and by Spital and Yennie [18]. They find that such terms are of the expected form $h(y^2)/y_0$.

If we do wish to "find VMD" in the data, we must first of all, in line with the above discussion, restrict ourselves to small x. Moreover, we should also restrict ourselves to small Q^2, which (see Eq. (2.5)) places greatest emphasis on the photon's low mass components. Interestingly enough, if one looks at the small x region of the present SLAC data presented by J. Friedman and H. Kendall (1972), where machine energy limitations link small x with small Q^2, it appears that VMD contributions can account for a large fraction of the data [5]. This is shown in Fig. 10 (these curves should be redone with new parameters, but the general impression is correct.)

As we have seen, there is ample evidence that the small x region of inelastic electron scattering is associated primarily with the hadronic structure of the photon. Within this framework (usually called GVD), it is possible to fit the data in many possible ways, in all of which at small Q^2 the ϱ^0, ω, and ϕ would make the dominant contributions. Unlike the original difficulty that VMD appeared to predict too low a cross section, the problem with GVD is how to suppress the effects of the very abundant higher mass constituents that we have up to this point ignored. To see the problem most clearly, consider the total probability of the hadronic component, Eq. (2.6), using the data from e^+e^- interactions up to the currently available masses (up to around 5 GeV). This turns out to be of order 1.3α, with the ϱ_0 contributing 0.4α [26]. Unless the trend of this data changes drastically, the total hadronic component could be several times α. Clearly, it is not possible that this whole component be absorbed with a cross section of order 25 mb, which would lead to a total photon cross section of ~ 250 μb, as well as a large overestimate of νW_2. The cross sections of some of the higher mass constituents must be considerably smaller.

At the present time, there is no generally accepted explanation for this apparently small interaction of the higher mass components. Perhaps it indicates that the matter which is initially created in e^+e^- annihilation is not strongly interacting until some time has passed. That is, the higher mass part of (2.2) represents a localized state which is completely unlike our usual picture of hadronic matter. Such a state may not have sufficient time to evolve and interact while passing through a nucleon.

One possibility is that only a small portion of the total hadronic state actually interacts with the target. Since the constituent state is virtual, the intuitive idea is that time dilation prevents high laboratory momentum components from ever evolving into fully dressed particles with full laboratory cross sections. Hence, we concentrate on constituents with small laboratory momentum, In naive parton models, where the virtual photon's consitituents are seen as parton-antiparton pairs, limiting target interactions to appropriate slow or "wee" partons [27] reduces the constituent's cross section by a factor of $1/\mathscr{M}^2$ [28]. This is known as the "aligned jet" model. (In a more abstract context the physics was essentially described in the covariant parton model by Landshoff, Polkinghorne and Short [29]. Such models are consistent with scaling in $\sigma_{e^+e^-}$ and are quite compatible with a reasonable photon cross section and scaling in deep-inelastic electron scattering. (For spin 1/2 partons these models are also consistent with a small value of R.)

Another possible view is that the entire structure of the photon becomes very point-like in this high mass region and it could pass through the nucleon with very little chance of interaction unless it has a very central collision. Since higher mass contributions are relatively emphasized at higher Q^2 (as is obvious from (2.5)), this view would lead to a relatively greater fraction of events with large transverse momentum.

Still a third possibility exists, possibly in conjunction with the previous two, that matter is being produced in $e^+e^- \to X$ with different quantum numbers and therefore intrinsically small interactions with nucleons (e. g. "charmonium" or heavy leptons). While future clarification of this third possibility will come chiefly from storage ring experiments, it is also possible that experiments in nuclei could shed light on these mechanisms.

I enjoyed the friendly and stimulating atmosphere of the School and wish to thank W. Czyż and A. Białas for inviting me and L. Maximon for arranging the financial support.

REFERENCES

[1] T. Bauer, F. Pipkin, R. Spital, D. R. Yennie, article in preparation for *Reviews o Modern Physics*.

[2] More complete discussions may be found in the 1971 and 1973 Erice Lectures by Sakurai: J. J. Sakurai, in *Properties of the Fundamental Interactions*, A. Zichichi, ed. Editrice Compositori, Bologna 1973, Part A, p. 243; J. J. Sakurai, in *Laws of Hadronic Matter*, A. Zichichi, ed., Academic Press, New York 1975.

[3] Early theoretical papers on photon shadowing: L. Stodolsky, *Phys. Rev. Lett.* **18**, 135 (1967); S. J. Brodsky, J. Pumplin, *Phys. Rev.* **182**, 1794 (1969); K. Gottfried, D. R. Yennie, *Phys. Rev.* **182**, 1595 (1969); B. Margolis, C. L. Tang, *Nucl. Phys.* **B10**, 329 (1969); M. Nauenberg, *Phys. Rev. Lett.* **22**, 556 (1969).

[4] Experimental papers on photon shadowing: H. Meyer, B. Naroska, J. M. Weber, M. Wong, V. Heynen, E. Mandelkow, D. Notz, *Phys. Lett.* **33B**, 189 (1970); V. Heynen, H. Meyer, B. Naroska, D. Notz, *Phys. Lett.* **34B**, 651 (1971); D. O. Caldwell, V. B. Elings, W. P. Hesse, G. E. Jahn, R. J. Morrison, F. V. Murphy, D. E. Yount, *Phys. Rev. Lett.* **23**, 1256 (1969); **24**, 796 (E); D. O. Caldwell, V. B. Elings, W. P. Hesse, R. J. Morrison, F. V. Murphy, B. W. Worster, D. E. Yount, *Phys. Rev. Lett.* **25**, 609, 902(E) (1970); D. O. Caldwell, V. B. Elings, W. P. Hesse, R. J. Morrison, F. V. Murphy, D. E. Yount, *Phys. Rev.* **D7**, 1362 (1973); T. A. Armstrong, W. R. Hogg, G. M. Lewis, A. W. Robertson, G. R. Brookes, A. S. Clough, J. M. Freeland, W. Galbraith, A. F. King, W. R. Rawlinson, N. R. S. Tait, J. C. Thompson, D. W. L. Tolfree, *Phys. Lett.* **34B**, 535 (1971); *Phys. Rev.* **D5**, 1640 (1972); *Nucl. Phys.* **B41**, 445 (1972); G. R. Brookes, A. S. Clough, H. H. Freeland, W. Galbraith, A. F. King, T. A. Armstrong, W. R. Hogg, G. M. Lewis, A. W. Robertson, W. R. Rawlinson, N. R. S. Tait, J. C. Thompson, D. W. L. Tolfree, *Phys. Rev.* **D8**, 2826 (1973).

[5] D. R. Yennie, *Rev. Mod. Phys.* **47**, 311 (1975).

[6] V. N. Gribov, *Soviet Phys. Zh. Eksp. Teor. Fiz.* **57**, 1306 (1969) (*JETP* **30**, 709 (1970)); S. J. Brodsky, J. Pumplin, *Phys. Rev.* **182**, 1794 (1969); D. M. Ritson, *Phys. Rev.* **D3**, 1267 (1971); K. Fujikawa, *Phys. Rev.* **D4**, 2794 (1971).

[7] J. D. Bjorken, *SLAC Report* No. 167, Vol. I, 1 (1974).

[8] D. Benaksas, G. Cosme, B. Jean-Marie, S. Jullian, F. Laplanche, J. LeFrancois, A. D. Liberman, G. Parrour, J. P. Repellin, G. Sauvage, *Phys. Lett.* **39B**, 289 (1972); **42B**, 507 (1972).

[9] G. J. Gounaris, J. J. Sakurai, *Phys. Rev. Lett.* **21**, 244 (1968); M. T. Vaughn, K. C. Wali, *Phys. Rev. Lett.* **21**, 938 (1968).

[10] H. Cheng, T. T. Wu, *Phys. Rev.* **183**, 1324 (1969); J. D. Bjorken, J. B. Kogut, D. E. Soper, *Phys. Rev.* **D3**, 1382 (1971); J. B. Kogut, *Phys. Rev.* **D5**, 1152 (1972).

[11] This data was supplied privately by Profs U. Becker and S. Ting.

[12] R. Spital, D. R. Yennie, *Phys. Rev.* **D9**, 126, 138 (1974).

[13] M. Ross, L. Stodolsky, *Phys. Rev.* **149**. 1172 (1966); *Phys. Rev. Lett.* **17**, 563 (1966).

[14] P. Söding, *Phys. Lett.* **19**, 702 (1966).

[15] S. D. Drell, *Phys. Rev. Lett.* **5**, 278 (1960).

[16] T. H. Bauer, *Phys. Rev.* **D3**, 2671 (1971); J. Pumplin, *Phys. Rev.* **D2**, 1859 (1970).

[17] K. C. Moffeit, J. Ballam, G. B. Chadwick, M. Della-Negra, R. Gearhart, J. J. Murray, P. Seyboth, C. K. Sinclair, I. O. Skillicorn, H. Spitzer, G. Wolf, H. H. Bingham, W. B.

Fretter, W. J. Podolsky, M. S. Rabin, A. H. Rosenfeld, B. Windmolders, G. P. Yost, R. H. Milburn, *Phys. Rev.* **D5**, 1603 (1972); the data separated into ϱ^0 and non-ϱ^0 events was supplied privately by Dr Moffeit.

[18] R. D. Spital, D. R. Yennie, *Nucl. Phys.* **106**, 269 (1976).

[19] J. D. Bjorken, *Phys. Rev.* **179**, 1547 (1969); the example in Fig. 19 is from R. E. Taylor, in *1975 Lepton-Photon Symposium*, W. T. Kirk, ed. Stanford Linear Accelerator Center, Stanford, California, p. 679.

[20] L. S. Brown, in *1969 Boulder Lectures in Theoretical Physics*; R. Jackiw, R. van Royen, G. B. West, *Phys. Rev.* **D2**, 2473 (1970); R. A. Brandt, *Phys. Rev. Lett.* **23**, 1260 (1969); R. A. Brandt, *Phys. Rev.* **D1**, 2808 (1970); H. Leutwyler, J. Stern, *Nucl. Phys.* **B20**, 77 (1970).

[21] B. L. Ioffe, *Phys. Lett.* **30B**, 123 (1969).

[22] J. Pestieau, P. Roy, H. Terazawa, *Phys. Rev. Lett.* **25**, 402 (1970).

[23] A. suri, D. R. Yennie, *Ann. Phys.* (N. Y.) **72**, 243 (1972).

[24] The compilation of data shown in Fig. 10 is from J. Friedman, H. Kendall, *Ann. Rev. Nucl. Sci.* **22**, 203 (1972); the data in Fig. 11 is from Y. Watanabe, L. N. Hand, S. Herb. A. Russell, C. Chang, K. W. Chen, D. J. Fox, A. Kotlewski, P. F. Kunz, S. C. Loken, M. Strovink, W. Vernon, *Phys. Rev. Lett.* **35**, 898 (1975).

[25] J. Pestieau, J. Urias, *Phys. Rev.* **D8**, 1552 (1973).

[26] I wish to thank Mr. David Kaplan for this evaluation.

[27] R. P. Feynman, *Phys. Rev. Lett.* **23**, 1415 (1969).

[28] J. D. Bjorken, *SLAC Report* No. 167, Vol. I, 1 (1974); N. N. Nicolaev, V. I. Zakharov, *Phys. Lett.* **55B**, 397 (1975).

[29] P. V. Landshoff, J. C. Polkinghorne, R. D. Short, *Nucl. Phys.* **B28**, 225 (1971).