of π^+ and π^- photoproduction cross sections and have found good agreement with previous data. Also the data presented in this paper agree well with their predictions.

Several authors have fitted the forward peak in π^+ photoproduction using conspiring trajectories, cuts, or a phenomenological background term.[10] Frøyland and Gordon[10] also make predictions for π^- photoproduction. They describe the π^-/π^+ ratio correctly for large momentum transfers; however, they do not find the rise to unity near the forward direction.

We would like to comment on a quark-model prediction by Bialas et al.[11] relating the difference between π^+ and π^- photoproduction to a difference of K^{*0} production cross sections. These authors use a π^-/π^+ ratio of 0.4 and find the quark-model relation to be violated. With our new results which give roughly equal π^+ and π^- cross sections in the forward direction and a smaller difference of the integrated cross sections, the violation is not so evident.

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LORENTZ INVARIANCE IN REGGEIZATION OF PION PHOTOPRODUCTION AMPLITUDES*

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We study the implications of Lorentz and gauge invariance for Reggeization of the pion trajectory in charged pion photoproduction and find that the pion pole need not be introduced as a kinematical singularity. As a consequence of having a dynamical Regge pion pole, we find that effects characteristic of an s-channel nucleon pole appear automatically.

Charged-pion photoproduction $\gamma p \to \pi^+ n$ has recently received considerable attention[1-4] because of the experimental observation[5] of a sharp peak in the forward direction. In the past, forward peaking of many reactions had been explained very nicely in the peripheral model with pion exchange. Whereas a natural pion exchange mechanism exists also in photoproduction, the contribution from this particular single process vanishes in the forward direction; in the usual scheme for Reggeization of helicity amplitudes there is no dynamical pole on the pion trajectory at the position of the pion mass and spin; it is introduced through a kinematical factor.[1,6] This is at least esthetically dissatisfying since this same kinematical factor also occurs in the perturbation amplitude but does not appear in the invariant amplitudes[7] although the pole from the intermediate-state propagator does appear. Furthermore, this same fixed pole would show up at the Regge recurrences on the pion trajectory, and the reasoning that led to the pion kinematical pole implies kinematic ρ and K poles in photoproduction of ρ mesons and K mesons, respectively.

We may proceed as in Ref. 1 by noting that $\gamma N \to \pi N$ is a special case of $\pi N \to VN$ (time reversed). The latter reaction contains a genuine, dynamical pion pole which, we find, can be retained in the limit when the mass of the vector meson m_v tends to zero. In this limit we make the Lorentz invariance requirement that the helicity amplitudes for longitudinally polarized vector mesons should tend smoothly to zero. We find that this "smoothness" restriction is equivalent to demanding gauge invariance,[8] i.e., consistency between masslessness and possible spin states of the photon implies gauge invariance; so it would be superfluous to additionally impose gauge invariance on the S matrix.

The simple example of scalar "pion" photoproduction from a scalar boson K is straightforward and illuminating. The general amplitude for $V(k,\epsilon)+K(p_1)\to\pi(q)+K(p_2)$ can be written

$$T=(q\cdot\epsilon)A(s,t)+(p_1\cdot\epsilon)B(s,t)\equiv T_\mu\epsilon^\mu, \tag{1}$$

where the four-momenta are in parentheses and ϵ is the four-polarization vector of the "photon" with mass m_v and energy $\omega=(m_v^2+\vec{k}^2)^{1/2}$. In the reference frame with the vector meson moving along the z axis, the longitudinal polarization vector has components $\epsilon_0{}^\mu=(\kappa/m_v,0,0,\omega/m_v)$ obtained by a Lorentz transformation from the rest system vector $(0,0,0,1)$, where $\kappa=|\vec{k}|$. The usual steps in Reggeization of helicity amplitudes in the t channel when applied to $\langle 0|T_\mu\epsilon^\mu|\pm 1\rangle$ only will lead to the same kinematic pion pole as discussed in the appendix of Ref. 1. However, in the t channel of the actual photoproduction ($m_v=0$) process under discussion, the invariant amplitude $A(s,t)$ containing the pion pole does not contribute; so it appears that a kinematical pion pole is being forced into the amplitude. But the pion Regge pole does enter the helicity amplitude $\langle 0|T_\mu\epsilon^\mu|0\rangle$ in a term with coefficient varying as m_v^{-1}. If the residue of the pion pole in this longitudinal helicity amplitude remains finite as $m_v\to 0$, which is the case in perturbation theory, then the "smoothness" condition discussed above leads to

$$(k\cdot q)A(s,t)=-(k\cdot p_1)B(s,t). \tag{2}$$

This relation can be interpreted somewhat like a conspiracy (perhaps electromagnetic conspiracy) condition. In the t channel the helicity amplitudes $\langle 0|T|\pm 1\rangle$ would normally contain only the B part, but because of Eq. (2) the dynamical pion pole in $A(s,t)$ can be introduced into these amplitudes. All of the undesirable features of fixed kinematical poles thus are simply avoided. The Lorentz invariance condition (2) suggests that the pion Regge pole in photoproduction exists in a conspiratorial arrangement with the Regge contributors to $B(s,t)$, the latter being somewhat daughterlike since they are down by s^{-1} relative to A because of the $(k\cdot p_1)\sim s$ factor. In conventional notation,[1] the helicity amplitudes are written

$$\langle 0|T_\mu\epsilon^\mu|+1\rangle=\sin\theta\bar{f}_{01}=(p/\sqrt{2})\sin\theta B$$

$$\langle 0|T_\mu\epsilon^\mu|0\rangle=m_v\bar{f}_{00}=m_v^{-1}[\kappa(q_0+k_0)A$$
$$+(E\kappa-pk_0\cos\theta)B]. \tag{3}$$

Using Eq. (2), we may express the invariant amplitude B in terms of either $\bar{f}_{01}$ or $\bar{f}_{00}$. The result is

$$B=2\sqrt{2}(t-4M_K^2)^{-1/2}\bar{f}_{01}, \tag{4}$$

$$B=-\bar{f}_{00}2(t-\mu^2)^2/[(s-u)(t+\mu^2)+(t-\mu^2)]. \tag{4$'$}$$

We next remove, according to the usual procedure,[1,9] the kinematical factors to get the kinematical-singularity-free helicity amplitudes

$$\bar{f}_{01}=(t-4M_K^2)^{1/2}\tilde{f}_{01};\ \bar{f}_{00}=(t-\mu^2)^{-1}\tilde{f}_{00}, \tag{5}$$

where $\tilde{f}_{00}$ has the dynamical pion pole and $\tilde{f}_{01}$ is finite at $t=\mu^2$. Explicity displaying the pion Regge-pole contribution, we have from Eq. (4$'$)

$$B_\pi=-\beta_\pi\frac{1+e^{-i\pi\alpha}}{\sin\pi\alpha}$$
$$\times\frac{2(t-\mu^2)}{(s-u)(t+\mu^2)+(t-\mu^2)^2}P_\alpha(\cos\theta) \tag{6}$$

and from Eqs. (2) and (6)

$$A_\pi = \beta_\pi \frac{1+e^{-i\pi\alpha}}{\sin\pi\alpha} \frac{2(s-M_K^2)}{(s-u)(t+\mu^2)+(t-\mu^2)^2} P_\alpha(\cos\theta), \tag{7}$$

where β_π is the usual residue function and α the pion trajectory function. Therefore, we see that the invariant amplitude $A(s,t)$ contains the dynamical pion pole times a kinematical factor, the latter becoming a constant as $t\to\mu^2$; $B(s,t)$, however, although it contains the dynamical pion factor, becomes a direct-channel K pole as $t\to\mu^2$. Thus, we see that consistent Reggeization of the pion-exchange amplitude implies automatic inclusion of a direct-channel pole, and correspondence between low-energy theorems and Regge theory is obtained. That effects of direct-channel poles are included when a process is described by Regge exchange amplitudes is nicely illustrated in a recent analysis by Schmid[10] of πN charge exchange.

The above discussion for the rather special example can be translated directly to the more complicated and physically interesting problem of pion photoproduction on nucleons. For $V(k,\epsilon)+P(p_2)\to\pi^+(q)+N(p_1)$ the amplitude is

$$T=\overline{N}(p_2)T_\mu\epsilon^\mu N(p_1)=\overline{N}(p_2)\sum_{i=1}^{6}U_i(s,t,u)u_iN(p_1), \tag{8}$$

where $\overline{N}(p_2)$ and $N(p_1)$ are the nucleon spinors and the u_i's are defined in the s channel as

$$u_1=\gamma_5(\gamma\cdot k)(\gamma\cdot\epsilon),\qquad u_2=\gamma_5(P'\cdot\epsilon),\qquad u_3=-\gamma_5(q\cdot\epsilon),$$

$$u_4=\gamma_5(\gamma\cdot\epsilon),\qquad u_5=-\gamma_5(\gamma\cdot k)(q\cdot\epsilon),\qquad u_6=\gamma_5(\gamma\cdot k)(P'\cdot\epsilon), \tag{9}$$

with $P'=p_1+p_2$. Just as above, we begin by considering the longitudinal helicity amplitudes which contain the pion pole term in $U_3(s,t,u)$. These amplitudes involving longitudinally polarized "photons" with $m_v\neq 0$ in terms of the invariant amplitudes are

$$\langle 0|T|\tfrac12\tfrac12\rangle=\tfrac12[-2m_vp\cos\theta U_1+m_v^{-1}\{-4Ep\omega\cos\theta U_2-4E^2\kappa U_3$$
$$-2M\kappa U_4-4ME\omega\kappa U_5-4\omega^2p\cos\theta MU_6\}], \tag{10}$$

$$\langle 0|T|-\tfrac12\tfrac12\rangle=\tfrac12 m_v^{-1}\sin\theta[2p\omega U_4+4Ep\kappa^2U_5+4\omega p^2\kappa\cos\theta U_6], \tag{11}$$

which we define as Z_1 and $Z_2\sin\theta$, respectively, where M, E, and p are the nucleon mass, energy, and the magnitude of the momentum, respectively, and θ is the t-channel scattering angle.

Each of the invariant U_i can be expressed in terms of s, t, and u and kinematical-singularity-free, parity-conserving (KSFPC) helicity amplitudes. Those involving transversely polarized photons ($\tilde{X}_i$, $\tilde{Y}_i$) are essentially the $\tilde{f}_{cd,ab}{}^t(s,t)$ of Ref. 1:

$$\tilde{X}_{1,2}=[2\sqrt2\sin\theta K_{1,2}]^{-1}(\langle+1|T|+\tfrac12+\tfrac12\rangle\pm\langle+1|T|-\tfrac12-\tfrac12\rangle)$$

$$\tilde{Y}_{1,2}=[2\sqrt2K_{3,4}]^{-1}[\langle 1|T|-\tfrac12\tfrac12\rangle/(1-\cos\theta)\pm\langle 1|T|+\tfrac12-\tfrac12\rangle/(1+\cos\theta)], \tag{12}$$

where the K_i, with $\lambda\equiv[t-(m_v+\mu)^2]^{1/2}[t-(m_v-\mu)^2]^{1/2}$, are

$$K_1=\lambda;\quad K_2=[(t-4M^2)/t]^{1/2};\quad K_3=\lambda/\sqrt{t};\quad K_4=(t-4M^2)^{1/2}. \tag{12'}$$

The kinematical factors K_i are not unique; Eq. (12′) (giving the factors of Ref. 1 when $m_v \to 0$), or those of Ref. 3, or another set can be deduced from the U_i's depending on when in the calculation (relative to taking the limit $m_v \to 0$) the K_i are extracted from the helicity amplitudes. The factors similar to Eq. (12′) also can be taken from Z_1 and Z_2 of Eqs. (10) and (11):

$$\tilde{Z}_1 = m_v^{-1}\lambda\sqrt{t}Z_1, \quad \tilde{Z}_2 = m_v^{-1} \times [t/(t-4M^2)]^{1/2} Z_2, \qquad (13)$$

so that Eq. (10) may be rewritten in the form

$$\xi(s-u)U_2 + \lambda^2 U_3 = -4m_v^2[2(t-4M^2)^{-1} \times (s-u)(t\tilde{X}_1 + 2M\tilde{Y}_1) - 4M\tilde{Y}_2 + \tilde{Z}_1/t + 2M\xi\tilde{Z}_2/t], \qquad (14)$$

where $\xi \equiv t + m_v^2 - \mu^2$. Evaluating at $t = \mu^2 - m_v^2$ ($\xi = 0$), we obtain an expression for U_3 in terms of the helicity amplitudes which is exactly that following from Eqs. (8) and (9) providing Z_1 and Z_2 contain the "kinematical" factor m_v as in Eq. (13). At $t = (m_v \pm \mu)^2$ or $\lambda = 0$, Eq. (14) yields an expression for U_2; however, the result for U_2 from Eqs. (8) and (9) in terms of the KSFPC helicity amplitudes of Eqs. (12) and (12′) is not consistent with the former and can only be consistent if the usual kinematical factor K_2 multiplying $\tilde{X}_2$ (F_2 of Ref. 3) in U_2 (essentially A_2 of Ref. 3) is wrong. The additional factor must lead to another m_v multiplying $\tilde{X}_2$; otherwise, the longitudinal helicity amplitudes will not vanish as $m_v \to 0$ [a possibly more correct kinematic factor is ξK_2; cf. Eq. (15)].

The amplitudes U_2 and U_3 are not "dynamically" independent; as $m_v \to 0$, the contribution of the longitudinal helicity amplitudes to observables must vanish, and at $m_v = 0$ Eqs. (10) and (11) lead to the usual gauge-invariance relations between U_2 and U_3 and between U_4, U_5, and U_6, respectively. The present approach should, therefore, allow for an essentially simultaneous treatment of pion photoproduction and electroproduction, and for some insight into the relation of Reggeization and the rho-photon analogy. This subject will be treated elsewhere.

Given (the kinematically independent) U_2 and U_3 in terms of helicity amplitudes following from Eqs. (8) and (9), one can see immediately by inspection that a relation similar to Eqs. (4) and (4′) will result from Eq. (14) relating the dynamical pion pole in $\tilde{Z}_1$ to the transverse helicity amplitude $\tilde{X}_2$. Above, the gauge-invariance relation Eq. (2) could have been written in a more general form with an analytic function φ times m_v^2 on the right. We generalize Eq. (2) for the pion-photoproduction problem to

$$U_3 = -[p\kappa\cos\theta/E\omega]U_2 + m_v^2\varphi, \qquad (15)$$

which for $m_v = 0$ is simply $(P'\cdot k)U_2 + (k\cdot q)U_3 = 0$. Substituting (15) into Eq. (14) we obtain the sum rule [analogous to equating (4) and (4′)]

$$p\kappa\cos\theta X_2 = \omega[\kappa p m_v^{-1}Z_1 + M\omega m_v^{-1}Z_2 + 2\kappa^2E^2p\varphi - MY_2]. \qquad (16)$$

This equation cannot be reduced to a trivial identity by choosing an appropriate φ as long as the generalized gauge condition Eq. (15) is regarded as nontrivial. The requirement that Eq. (11) for $Z_2 \to 0$ as $m_v \to 0$ leads to a relation between $\tilde{Y}_2$ and $\tilde{Z}_2$, allowing (16) to be written as an equation between $\tilde{X}_2$ and $\tilde{Z}_1$ which contain the analytic functions $P_\alpha'(\cos\theta)$ and $P_\alpha(\cos\theta)$, respectively. A connection between two such amplitudes can only be realized through relations like $zP_\alpha'(z) = \alpha P_\alpha(z) + P_{\alpha-1}'(z)$, $z = \cos\theta$. Thus, the residue function in either $\tilde{X}_2$ or $\tilde{Z}_1$ must contain α, and to retain the pion pole at $t = \mu^2$ we select the solution with $\tilde{X}_2$ containing α^{-1}. A more careful analysis of Eqs. (4) and (4′) shows that such an α dependence is really present in the simple example and Eqs. (6) and (7) hold only to the leading power of z.

In summary, the dynamical pion pole appearing in "massive" photon reactions can be retained as $m_v \to 0$. This is desirable as Reggeization of the pion through a kinematical singularity obscures the meaning of the pion electromagnetic form factor measured by isolating the pion pole in a process such as $eN \to eN\pi$. Also, correspondence with low-energy theorems appears automatically; e.g., Eq. (15) shows that at $t = \mu^2$, U_2 has a nucleon pole $(s-M^2)^{-1}$. This relation with perturbation theory along with electroproduction will be discussed elsewhere.

We close with the following remarks: (i) When the KSFPC helicity amplitudes are expressed in terms of the invariant U_i's the conspiracy relation at $t = 0$ between X_2 and Y_1 (F_2 and F_3 of Ref. 3) amounts to the trivial relation $A_1 = A_1$ (notation of Ref. 3). (ii) A similar comment applies to the conspiracy relation between $\tilde{Z}_1$ and $\tilde{Z}_2$ in vector-meson production by pions [indicated by

t^{-1} factors in Eq. (14)].

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EXPERIMENTAL EVIDENCE AGAINST THE EXISTENCE OF STRANGE LEPTONS*

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Arguments based on experimental evidence are presented against the existence of strange leptons, as recently suggested by Weiner.

In a recent paper, Weiner has made the interesting suggestion that strange leptons exist and that strangeness is thereby conserved in semileptonic decays of strange particles.[1] The selection rules in semileptonic decays are deduced from this hypothesis without the necessity of invoking the "absence of neutral leptonic currents" and the "$\Delta S = \Delta Q$ rule." Weiner suggests that in addition to the usual leptons there exist strange leptons according to three possibilities:

(i) There exist neutral strange leptons, ν^S and $\bar{\nu}^S$, with strangeness -1 and $+1$, respectively.

(ii) There exist charged strange leptons μ^{+S}, e^{+S} and μ^{-S}, e^{-S} with strangeness $+1$ and -1, respectively.

(iii) There exist both the charged and neutral strange leptons listed above.

We present here direct experimental evidence againt possibility (ii) by showing that μ^-'s from the decay $K_L^0 \to \pi^+ + \mu^- + \bar{\nu}$ undergo nuclear capture with the same rate as μ^-'s from π^- decays. We also discuss evidence against the other two possibilities.

For purposes of reference we shall designate muons from $K_{\mu 3}^0$ decay as $\mu_K^\pm$ and those from $\pi^\pm$ decay as $\mu_\pi^\pm$. If Weiner's second possibility is correct, strangeness is conserved in $K_{\mu 3}^0$ decay with

$$K_L^0 \to \pi^\pm + \mu_K^{\mp S} + \bar{\nu}(\nu).$$

The mean life for μ_π^- in carbon is shorter than that for μ_π^+ because some of the μ_π^- undergo nuclear capture in the reaction

$$\mu_\pi^- + p \to n + \nu.$$

If strange neutrinos do not exist this process