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Fixed poles of photonic amplitudes involving compositeness*

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The existence and significance of the right-signature $J=0$ fixed pole in vector scattering such as photoproduction and Compton scattering off protons and pions is investigated in the context of gauge-invariant models involving compositeness which have been developed in a previous paper. It is shown that the existence of such a fixed pole depends on the elementarity of the photon and not on the compositeness of the pion or proton target or on the assumption that the charged constituent of the target has no structure. Assuming the ρ meson lies on a Regge trajectory, we show that within the context of these models a $J=0$ fixed pole does not arise in ρ-meson photoproduction.

I. INTRODUCTION

Fixed J-plane poles have attracted renewed interest recently,[1] largely as a result of extensive investigations of photoproduction and electroproduction processes. In strong interactions analytically continued unitarity allows only wrong-signature fixed poles, and only at nonsense values of J.[2] But in electromagnetic and weak interactions of hadrons,[3] unitarity (to first order in the electromagnetic or weak coupling constant) allows current-current amplitudes to have both right- and wrong-signature fixed poles at nonsense values of J. Right-signature fixed poles[4,5] have for a long time been known to be related to the elementarity or compositeness of particles, although the connection has not really been understood. Thus Dashen and Lee[6] comment in their discussion of pion photoproduction that "there is a connection (admittedly fuzzy) between fixed poles in photoproduction and elementarity of the external particles." The relation between fixed poles and elementarity of particles has mostly been investigated in the frame of field-theoretic models[5,7,8] using some kind of ladder approximation. This is natural, because if the model is to exhibit explicitly the distinguishing characteristics between composite particles, i.e., Reggeons, and noncomposite, i.e., elementary particles, then it must be possible within the model to calculate the Regge trajectories on which these bound states lie. It is well known, however, that the generation of Regge trajectories requires at least the summation of a complete sequence of planar[9] ladder diagrams.[5,7]

Recently Drell and Lee[10] established the scaling property in deep-inelastic electron scattering in a model in which the physical nucleon is considered as a bound state of a bare nucleon and a bare meson. They showed that this bound-state formulation provides a fully relativistic generalization of the parton model which is not restricted to the infinite-momentum frame. It is then natural to ask whether this bound-state model for the physical nucleon results in a right-signature fixed J-plane pole in the (real or virtual) Compton scattering amplitude. An investigation of this question[11] has been carried out by S. Y. Lee,[12] who concluded that the existence of the right-signature $J=0$ fixed pole in Compton scattering off protons "seems to be only due to the bound-state nature of the physical nucleon and the existence of local electromagnetic interactions."

In order to understand the origin of the $J=0$ fixed pole, we have investigated a class of gauge-invariant models[13] in the context of the ladder approximation of the Bethe-Salpeter equation, for which the Drell-Lee model[10] is but one example. From this investigation, we conclude here that the existence of the right-signature $J=0$ fixed pole in Compton scattering requires the elementarity of the (real or virtual) photon field and is not due to the compositeness of the target hadron, i.e., nucleon or pion. In effect, we show that such a fixed pole exists only if the vertex function describing the coupling of the photon to a hadron satisfies an inhomogeneous Bethe-Salpeter equation in the ladder approximation, i.e., that the photon is an elementary particle. If a vertex function describing the coupling of a vector meson to hadrons satisfies a homogeneous Bethe-Salpeter equation, i.e., the vector meson is bound and lies on a Regge trajectory, as expected for the ρ meson, then the fixed pole does not exist.

In constructing a bound-state wave function or vertex function for a physical particle such as the nucleon or the ρ meson, there is, of course, a considerable degree of arbitrariness with respect to the nature of its constituents and the binding forces or potentials. Similar to Drell and Lee,[10] we do not explicitly define the binding potential, but assume that it depends only on the square of the 4-momentum transfer, k^2, and falls off asymptotically as k^{-2}. If the constituents are baryons, the potential could very well contain Dirac matrices, e.g., γ_5, for pseudoscalar binding. In such cases, consideration of the potential as a Dirac operator makes the calculation slightly more complicated but does not affect the conclusions. Consequently, unless stated otherwise we will assume the binding potentials to be scalar functions.

Since the existence of a fixed J-plane pole manifests itself by the appropriate amplitude containing a factor ν^J in the Regge limit, i.e., $\nu\to\infty$, an investigation of the relation between fixed poles and compositeness requires an investigation of the high-energy asymptotic (Regge) behavior of the amplitudes. Consequently, the object of this work is to study the asymptotic behavior of the appropriate amplitudes for gauge-invariant models in the context of the ladder approximation of the Bethe-Salpeter equation. In effect this note represents the extension of our previous investigation[5] to more realistic models.

In Sec. II we discuss within our models the amplitudes for the scattering of vector mesons off protons and off pions. In Sec. III we first recall briefly the well-known fact that the $J=0$ fixed pole exists for Compton scattering off elementary hadrons, in order to emphasize that the existence of such a fixed pole does not imply that the target hadrons must be bound states. We then go on to consider vector-meson scattering off elementary pions with a nonzero binding potential and demonstrate how the elementarity of both incoming and outgoing vector mesons is necessary for the existence of the fixed pole but independent of charge structure of the pions. We then consider the more complicated case of vector-meson scattering off the charged-meson constituent of a bound-state nucleon and arrive at the same results. This section is concluded with a discussion of how the existence of the fixed pole can be seen to depend on the elementarity of the vector mesons.

We conclude with a discussion of the relevance of our work to the existence of the fixed pole in photoproduction where its existence has been argued by Brandt *et al.*[14] in ρ photoproduction and by groups of authors[15] in π photoproduction. Our conclusion is that the leading fixed pole (i.e., the fixed pole at $J=0$) does not arise in these processes if—as is physically plausible—the produced mesons are composite states which lie on Regge trajectories. An important aspect of our work it to demonstrate explicitly that the infinite sums of planar ladder diagrams necessary for gauge invariance either build up the expected Regge-pole behavior or are of lower order than the Born-like or primary diagrams and thus that the existence of fixed poles may be determined by a study of these primary diagrams alone.

II. KINEMATICS, BETHE-SALPETER EQUATIONS

In ascertaining the existence of a fixed J-plane pole in the scattering of vector particles from hadrons, i.e., $Vh\to V'h$, we are interested in the asymptotic high-energy behavior of the amplitude for fixed values of momentum transfer between the hadrons or vector mesons. In the following it is sufficient to consider zero momentum transfer and thus elastic or quasielastic scattering.

If we designate the 4-momentum of the incident vector particle by q and that of the initial hadron by p, the scattering amplitude in the forward direction satisfying parity and time-reversal invariance can be written as (our metric is $+++-$, so that $p^2=-m_p{}^2$)

$$T_{\mu\nu}(p,q)=g_{\mu\nu}T_1+\frac{p_\mu p_\nu}{m_p{}^2}T_2+q_\mu q_\nu T_3+(p_\mu q_\nu+p_\nu q_\mu)T_4\ . \tag{2.1}$$

In working with photons (real or virtual) gauge invariance requires that

$$T_1+q^2T_3+p\cdot qT_4=0\ ,$$

$$\frac{1}{m_p{}^2}(p\cdot q)T_2+q^2T_4=0\ .$$

For Compton scattering, the forward amplitude can be written as

$$T_{\mu\nu}=\left(g_{\mu\nu}-\frac{q_\mu q_\nu}{q^2}\right)T_1+\frac{1}{m_p{}^2}\left(p_\mu-\frac{p\cdot q\,q_\mu}{q^2}\right)\left(p_\nu-\frac{p\cdot q\,q_\nu}{q^2}\right)T_2\ , \tag{2.2}$$

where T_1 and T_2 are scalar functions of q^2 and $\nu\equiv-(p\cdot q)/m_p$ whose absorptive, i.e., imaginary, parts are proportional to the structure functions W_1 and W_2 that are relevant to deep-inelastic scattering.

In taking the Regge limit, $\nu\to\infty$ with q^2 fixed, it is convenient to work in the rest frame of the target hadron. Consequently, we set

$$p_\mu=(\vec{0},m_p),\quad q_\mu=(0,0,q_3,q_0)\ .$$

In this frame, the invariant amplitudes T_1 and T_2 can be written as

$$T_1 = T_{11}, \quad T_2 = \frac{q^2}{\nu^2 + q^2}\left(T_{11} + \frac{q^2}{\nu^2} T_{33}\right). \tag{2.3}$$

Consequently, in the following we will only be interested in the cases $\nu = \mu = 1$ and $\nu = \mu = 3$.

In the Regge limit of $q_0 = \nu \to \infty$, we have

$$q_3 = q_0 + q^2/2\nu + \cdots = \nu + O(\nu^{-1}) . \tag{2.4}$$

Here we are interested in the forward scattering amplitude, $T_{\mu\nu}$, given by gauge-invariant models within the context of the ladder approximation of the Bethe-Salpeter equation as discussed previously.[13] In order to describe such models, it is necessary to have the Bethe-Salpeter equation appropriate for describing the coupling of a vector particle to either nucleons or mesons, and also the Bethe-Salpeter equation for a nucleon as a bound state of a bare nucleon and a bare meson, or correspondingly, the Bethe-Salpeter equation for a meson as a bound state of a bare nucleon-antinucleon pair.

As shown in Fig. 1 the vector-meson vertex function Γ is assumed to satisfy either an inhomogeneous ($Z^V \neq 0$) or a homogeneous ($Z^V = 0$) Bethe-Salpeter equation of the form

$$\Gamma_\mu^{\pi V \pi}(p, p+q, q) \equiv \Gamma_\mu^{\pi V \pi}(p) = Z^V (2p+q)_\mu + \int d^4x\, W^\pi(x)\, \Pi_R(p+q+x)\, \Gamma_\mu^{\pi V \pi}(p+x)\, \Pi_R(p+x) \tag{2.5}$$

for the coupling of a vector particle of 4-momentum q to two π mesons, and

$$\Gamma_\mu^{NVN}(p, p+q, q) \equiv \Gamma_\mu^{NVN}(p) = Z^V \gamma_\mu + \int d^4x\, W^N(x)\, P_R(p+q+x)\, \Gamma_\mu^{NVN}(p+x)\, P_R(p+x) \tag{2.6}$$

for the similar coupling of a vector particle to a nucleon-antinucleon pair. In these equations $W(x)$ is the binding potential, which we assume falls off asymptotically like $O(x^{-2})$, and the subscript R means dressed or renormalized in accordance with the ladder approximation contained in these equations. The bare meson and nucleon propagators are defined, respectively, as

$$\Pi^{-1}(k) = k^2 + \mu^2 - i\epsilon,$$

$$P^{-1}(k) = -(ik\cdot\gamma + m - i\epsilon) .$$

Gauge invariance for a photon field requires that these vertex functions satisfy the following generalized Ward identities:

$$\begin{aligned} q_\mu \Gamma_\mu^{\pi\gamma\pi}(p) &= q\cdot(2p+q) = \Pi^{-1}(p+q) - \Pi^{-1}(p) , \\ q_\mu \Gamma_\mu^{N\gamma N}(p) &= q\cdot\gamma = i[P^{-1}(p+q) - P^{-1}(p)] , \end{aligned} \tag{2.7a}$$

if the photon vertex is bare. For structured photon vertices [such as those defined by (2.5), (2.6) for $W\neq 0$] the Ward identities must contain the appropriately renormalized pion or nucleon propagators and so read

$$\begin{aligned} q_\mu \Gamma_\mu^{\pi\gamma\pi}(p) &= \Pi_R^{-1}(p+q) - \Pi_R^{-1}(p) , \\ q_\mu \Gamma_\mu^{N\gamma N}(p) &= i[P_R^{-1}(p+q) - P_R^{-1}(p)] . \end{aligned} \tag{2.7b}$$

(Note that the models discussed in Ref. 13 require renormalized charged particle propagators only in the case of Compton scattering with structure in the photon vertices.)

The wave function for a bound-state nucleon built from a bare nucleon and a bare pion is assumed to satisfy the homogeneous ($Z^N = 0$) Bethe-Salpeter equation

$$\phi^N(k, k+p, p) \equiv \phi^N(k) = \int d^4x\, W^{\pi N}(x)\, P_R(p+k+x) \times \phi^N(k+x)\, \Pi_R(k+x) \tag{2.8}$$

and that for a bound-state pion built from a bare nucleon-antinucleon pair

$$\Gamma^{N\pi N}(k, k+p, p) \equiv \Gamma^\pi(k) = \int d^4x\, W^{NN}(x)\, P_R(p+k+x) \times \Gamma^\pi(k+x)\, P_R(k+x) . \tag{2.9}$$

In writing these equations, it was assumed that the bound particle is incoming. The equations for outgoing bound particles are similar in nature and need not be written out explicitly.[13]

III. COMPTON SCATTERING IN VARIOUS MODELS

We now calculate the invariant amplitudes T_1 and T_2 for various gauge-invariant models to investigate the origin of the $J=0$ right-signature fixed pole. The asymptotic behavior of the am-

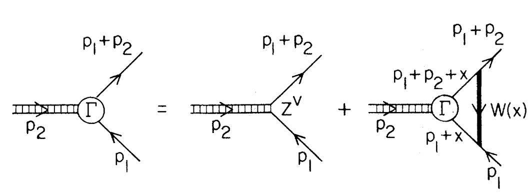

FIG. 1. Bethe-Salpeter equation describing the coupling of a vector particle with 4-momentum p_2 to two particles of 4-momentum p_1 and $p_1 + p_2$.

plitudes, in general, will be of the form (as will be shown)

$$T_1 \sim (1+e^{-i\pi\alpha})\,\nu^{\alpha} + R_1(q^2)\ ,$$
$$\nu^2 T_2 \sim (1+e^{-i\pi\alpha})\,\nu^{\alpha} + R_2(q^2)\ , \tag{3.1}$$

where $R_1(q^2)$ and $R_2(q^2)$ are the fixed-pole residue functions for the invariant amplitudes T_1 and $\nu^2 T_2$, respectively.

We consider first briefly the usual Born amplitudes, which correspond to the hadrons being elementary and the binding forces zero, in order to demonstrate that the existence of the fixed pole does not depend on the hadrons being bound states.

For Compton scattering off an elementary proton, the forward scattering amplitude is (in standard notation[13])

$$T_{\mu\nu} = \bar{u}(p)\,[\gamma_\nu P(p+q)\,\gamma_\mu + \gamma_\mu P(p-q)\,\gamma_\nu]\,u(p)\ . \tag{3.2}$$

Thus, using $\bar{u}(p)\,(ip\cdot\gamma + m) = 0$, $(ip\cdot\gamma+m)\,u(p)=0$, and the fact that $p\cdot\gamma\gamma_1 = -\gamma_1 p\cdot\gamma$ for $\vec{p}=0$, we have

$$T_1 = \frac{\nu^2}{q^2}\,T_2 = -\frac{(2p\cdot q)^2}{(2p\cdot q)^2 - q^4}\,(\bar{u}u/m_p)\ ,$$

and the residue functions for the fixed pole are

$$R_2 = q^2 R_1 \equiv q^2 C\ , \tag{3.3}$$

where C is a real constant independent of q^2, i.e., $C = -\bar{u}u/m_p = -2$. This last relationship ensures that the longitudinal amplitude T_L defined as

$$T_L = \left(\frac{\nu^2}{q^2} + m_p{}^2\right) T_2 - T_1$$

does not have a $J=0$ fixed pole.

For Compton scattering off a charged elementary meson with no binding forces (i.e., $W^{\pi}=0$) the forward scattering amplitude is

$$T_{\mu\nu} = (2p+q)_\mu\,(2p+q)_\nu\,\Pi(p+q) + (2p-q)_\mu(2p-q)_\nu\,\Pi(p-q) - 2g_{\mu\nu}\ , \tag{3.4}$$

the last term being the seagull contribution which is necessary for gauge invariance. Consequently,

$$T_{11} = -2,\quad T_{33} = -2 + \frac{2(q^2+\nu^2)q^2}{q^4 - (2p\cdot q)^2}\ , \tag{3.5}$$

and the residue functions for the fixed pole are

$$R_2 = q^2 R_1 = q^2 C\ ,$$

where C is a real constant independent of q^2, i.e., $C=-2$. For T_L the same conclusion is seen to hold as for Compton scattering off nucleons.

Thus, the fact that the $J=0$ right-signature fixed-pole residue function for Compton scattering off elementary hadrons, i.e., protons and pions, as considered above, is nonzero—as was also discussed by S. Y. Lee[12] and by Brodsky *et al.*[8] in the context of a composite nonperturbative parton model—shows clearly that the existence of the fixed pole is independent of whether the target hadron is elementary or composite.

We will now consider Compton scattering off elementary pions with a nonvanishing binding potential W (equivalent to gluon exchange[16]), to illustrate our point that the charge form factor of the elementary target need not be 1, but can be described by a structured vertex function, and to indicate how the existence of the fixed pole is dependent on the elementarity of the photon field.

The gauge-invariant model[13] for this process consists of the diagrams shown in Fig. 2 where the charged particle propagators are assumed to be renormalized in accordance with (2.7b). It is important to realize that the binding potential W acting between the incoming and outgoing mesons is the same as that which acts between a leg and the intermediate meson and gives the photon vertex its structure; similarly, W gives the charged-particle propagator its structure. If we were to assume that the photon is not elementary but a bound state of two mesons, i.e., $Z^{\gamma}=0$ in Fig. 1, then it is this potential that binds the two mesons to form the bound state.

We define the primary diagrams $D^P_{\mu\nu}$, $C^P_{\mu\nu}$, and $S^P_{\mu\nu}$ as the first diagrams in the expressions for these quantities, i.e., the diagrams for which there is no potential W acting between the ingoing and outgoing meson legs. Then in the forward direction the ladder sums for $D_{\mu\nu}$, $C_{\mu\nu}$, and $S_{\mu\nu}$ can be written

$$\begin{aligned}
D_{\mu\nu}(p,q) &= D^P_{\mu\nu}(p,q) + \int d^4x\,\Pi_R{}^2(p+x)\,R((p+x)^2, x^2) \times D^P_{\mu\nu}(p+x,q)\ ,\\
C_{\mu\nu}(p,q) &= C^P_{\mu\nu}(p,q) + \int d^4x\,\Pi_R{}^2(p+x)\,R((p+x)^2, x^2) \times C^P_{\mu\nu}(p,x,q)\ ,\\
S_{\mu\nu}(p,q) &= S^P_{\mu\nu}(p,q) + \int d^4x\,\Pi_R{}^2(p+x)\,R((p+x)^2, x^2)\,S^P_{\mu\nu}(p+x,q)\\
&= S^P_{\mu\nu}(p,q)\left[1+\int d^4x\,\Pi_R{}^2(p+x)\,R((p+x)^2,x^2)\right],
\end{aligned} \tag{3.6}$$

where

$$\begin{aligned}
D^P_{\mu\nu}(p,q) &= \Gamma^{\pi\gamma\pi}_\mu(p,p+q,q)\,\Gamma^{\pi\gamma\pi}_\nu(p,p+q,q)\,\Pi_R(p+q),\\
C^P_{\mu\nu}(p,q) &= \Gamma^{\pi\gamma\pi}_\mu(p,p-q,q)\,\Gamma^{\pi\gamma\pi}_\nu(p,p-q,q)\,\Pi_R(p-q),\\
S^P_{\mu\nu} &= -2(Z^{\gamma})^2 g_{\mu\nu}\ ,
\end{aligned} \tag{3.7}$$

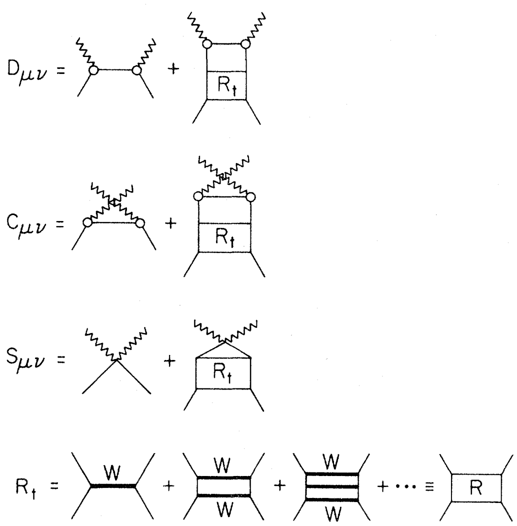

FIG. 2. The amplitudes $D_{\mu\nu}$, $C_{\mu\nu}$, and $S_{\mu\nu}$ required for a gauge-invariant description of Compton scattering off an elementary pion with a neutral binding potential W. The photon vertex function is assumed to satisfy an inhomogeneous Bethe-Salpeter equation as shown in Fig. 1 with the same binding potential $W(x)$. In Compton scattering the charged-particle propagators are assumed to be appropriately renormalized.

and the two-meson scattering amplitude R is shown in Fig. 2.

Since $S_{\mu\nu}^{P}$ is independent of q, $S_{\mu\nu}$ is also independent of q and will just be a constant times $(Z^{\gamma})^2$. Consequently, we need only to consider the expressions for $D_{\mu\nu}$ and $C_{\mu\nu}$.

We first consider the photon to be elementary, i.e., $Z^{\gamma}=1$. In the limit of large ν, the leading contributions can be seen to be those due to diagrams where each $\Gamma_{\mu}^{\pi\gamma\pi}$ is replaced by the inhomogeneous term of its Bethe-Salpeter equation, as we shall demonstrate.

In order to understand why this is possible, consider iterating Eq. (2.5) for $\Gamma_{\mu}^{\pi\gamma\pi}$. The mth term in the resulting series will involve $m-1$ integrations and will contain the product of $m-1$ propagators containing q. For the smooth Bethe-Salpeter kernels we are considering here the renormalized propagators can be calculated in certain approximations.[17] Their asymptotic behavior depends on the bare mass and the renormalization constant and is a standard result in renormalization theory[17]:

$$\Pi_R^{-1}(k)=Z\Pi^{-1}(k), \quad |k^2|\to\infty .$$

Here the propagators involving q play a crucial role in the asymptotic ν region. In effect, each propagator containing q will introduce a factor (ν^{-1}) into the asymptotic behavior of the contribution being considered. Consequently, since the region of integration is restricted by the potential and the other propagator, each successive term in the iteration expansion for $\Gamma_{\mu}^{\pi\gamma\pi}$ will be an order of ν^{-1} less important than the preceding term in the large ν region. In the limit of large ν, $\Gamma_{\mu}^{\pi\gamma\pi}$ for $Z^{\gamma}\neq 0$ can thus be replaced by the first term in the iteration which is just the inhomogeneous term of its Bethe-Salpeter equation. An argument equivalent to that presented here appears in the work of Biswas *et al.*[18]

Replacing $\Gamma_{\mu}^{\pi\gamma\pi}$ by $(2p+q)_{\mu}$ reduces the primary graphs to the Born graphs considered in the previous model. Consequently, forgetting the contributions due to the integrals in $D_{\mu\nu}$ and $C_{\mu\nu}$, there will be a fixed pole coming again from $S_{\mu\nu}$.

The residue functions R_1 and R_2 will again obey $R_2=q^2R_1=q^2C$, but the constant C which is real and independent of q^2 will be modified by the potential.

It is interesting to see what role the integrals in $D_{\mu\nu}$ and $C_{\mu\nu}$ play in determining the asymptotic behavior of T_1 and T_2. For $D_{\mu\nu}$ the integral of interest is

$$I_{\mu\nu}^{D}=\int d^4y\,\Pi_R{}^2(y)\,R(y^2,(y-p)^2)\times(2y+q)_{\mu}(2y+q)_{\nu}\Pi_R(y+q)\ , \tag{3.8}$$

where the variable of integration has been changed from x to $y=p+x$. The quantity $-y^2$ is, in effect, the mass of the meson whose scattering from a meson of mass μ^2 is described by R.

The asymptotic behavior of this integral can be obtained by finding where in the y space the integrand is maximal. For simplicity we shall here consider the integral for unrenormalized propagators. For renormalized propagators—such as those calculated by Böhm[17] in the zero-width approximation (in which the cuts in the exact propagator are approximated by poles)—the argument is similar, though more complicated because of the larger number of poles, the coupling also becoming renormalized. Thus, considering unrenormalized propagators, the conditions that the integrand be maximal are

$$y^2=O(\mu^2),\quad (y+q)^2=O(\mu^2)\ . \tag{3.9}$$

Subtracting these two constraints and using $q_0=q_3+O(\mu^2/\nu)=\nu\to\infty$ gives

$$q^2+2\nu(y_3-y_0)=O(\mu^2)\ . \tag{3.10}$$

Thus

$$y_3 = y_0 + O\left(\frac{\mu^2}{\nu}\right) \quad . \tag{3.11}$$

The first of the two constraints (3.9) then can be written

$$O(\mu^2) = y^2 = y_3^{\,2} - y_0^{\,2} + y_\perp^{\,2} = 2y_0\, O\left(\frac{\mu^2}{\nu}\right) + y_\perp^{\,2} \tag{3.12}$$

Consequently, the region of interest in the y_3, y_0 plane consists of a region of width $O(\mu^2/\nu)$ about $y_3 = y_0$ which extends out to $y_0 = O(\nu)$.

The integral thus takes the form

$$I^D_{\mu\mu} \propto \frac{1}{\nu}\int^{\nu} d\nu'\, R(\mu^2, 2\mu\nu')\left[(2\nu' + \nu)\,\delta_{\mu 3} + O(\mu^2)\,\delta_{\mu 1}\right]^2 \,, \tag{3.13}$$

where the only values of the indices of interest are $\mu = \nu = 1$ and $\mu = \nu = 3$.

From the great amount of work done on Regge poles and ladder diagrams for meson-meson scattering,[2, 19] the function R can be assumed to have an asymptotic behavior $(2\mu\nu')^{\alpha(t=0)}$. The various possible values of α can be determined by solving the bound state Bethe-Salpeter equation with pion constituents, i.e., Eq. (2.5) with $Z^V = 0$.[5, 19] Consequently,

$$I^D_{\mu\mu} \sim \nu^{\alpha}\left[\nu^2 \delta_{\mu 3} + O(\mu^2)\delta_{\mu 1}\right] \,. \tag{3.14}$$

Since the expression for $C_{\mu\nu}$ can be obtained by replacing q by $-q$ in that for $D_{\mu\nu}$, we have

$$T_{\mu\nu} = D^P_{\mu\nu} + C^P_{\mu\nu} + S_{\mu\nu} + R_{\mu\nu} \,, \tag{3.15}$$

where

$$R_{11} \sim \frac{1}{\nu^2} R_{33} \sim \nu^{\alpha} + (-\nu)^{\alpha} = (1 + e^{-i\pi\alpha})\,\nu^{\alpha} \,.$$

Thus for an elementary photon, i.e., $Z^{\gamma} \neq 0$, with interactions due to the potential W, the invariant amplitudes T_1 and T_2 have the asymptotic behavior

$$\begin{aligned} T_1 &\sim (1 + e^{-i\pi\alpha})\,\nu^{\alpha} + C \,, \\ \nu^{\alpha} T_2 &\sim (1 + e^{-i\pi\alpha})\,\nu^{\alpha} + q^2 C \,, \end{aligned} \tag{3.16}$$

where C is a real constant independent of q^2 given by [see the remarks following (3.5) and also (3.6)]

$$C = -2\left(1 + \int R(y^2, (y-p)^2)\, \Pi_R^{\,2}(y)\, d^4y\right) .$$

These results are in agreement with those of Brodsky *et al.*,[8] although, of course, the motivation of their investigation was different.

The situation for a bound-state vector particle is easily obtained by considering the effect of setting Z^{γ} equal to zero in the previous derivation. First and most important, there is no seagull contribution since $S^P_{\mu\nu}$ is proportional to $(Z^V)^2$, and thus there is no fixed pole unless it comes from $D_{\mu\nu}$ and $C_{\mu\nu}$. But the effect of setting Z^V equal to zero means that $\Gamma^{\pi V\pi}_{\mu}(p, q)$ will fall off[20] at least as fast as ν^{-1} and thus the primary graphs in $C_{\mu\nu}$ and $D_{\mu\nu}$ will, like the corresponding Born graphs, give at most contributions of $O(\nu^{-2})$ to T_1 and $\nu^2 T_2$. However, the above method of obtaining the asymptotic behavior of the integrals in $D_{\mu\nu}$ and $C_{\mu\nu}$ involved setting the constituent legs of the vertex functions close to their mass shell in the dominant region of integration. Consequently, in the integral the (almost) on-shell vertex functions will contribute in the same way as those for an elementary photon and the Regge behavior will again be obtained.

Thus, in conclusion the results can be written as

$$T_1 \sim (1 + e^{-i\pi\alpha})\,\nu^{\alpha} + C \,,$$

$$\nu^2 T_2 \sim (1 + e^{-i\pi\alpha})\,\nu^{\alpha} + q^2 C \,,$$

where

$$C = -2(Z^V)^2\left(1 + \int R(y^2, (y-p)^2)\, \Pi_R^{\,2}(y)\, d^4y\right) .$$

This demonstrates how the existence of the fixed pole in Compton scattering depends on the photon being elementary ($Z^{\gamma} = 1$) and not a bound state ($Z^{\gamma} = 0$). It also shows that the elementary proton or pion may have a charge form factor which is different from 1.

The existence of the $J = 0$ fixed pole for Compton scattering off a bound nucleon with local electromagnetic interactions but with vertices without charge structure has been demonstrated by S. Y. Lee[12] in his consideration of the Drell-Lee composite nucleon model.[10] To give structure to the photon coupling in the alternative version of the Drell-Lee model, in which the charged bare particle is a meson, the possibility of interaction, i.e., a potential or gluon exchange between the internal meson lines, must be considered. (We consider only this case here, because it is the natural extension of the above considerations.) Such a model, to be gauge invariant, must include an infinite number of diagrams where the potential acts between the internal meson legs. This is, of course, the same mechanism as that considered in the Compton scattering off an elementary meson with a nonzero potential. Designating the sums over exchanges in the s and t channels by R_s and R_t respectively, the gauge-invariant model[13] consists of the diagrams shown in Fig. 3.

This model is clearly an order of magnitude more complicated than the Drell-Lee model and the Compton scattering off an elementary pion as just considered. Consequently, we will only sketch how the model gives T_1 and T_2 the expected Regge-

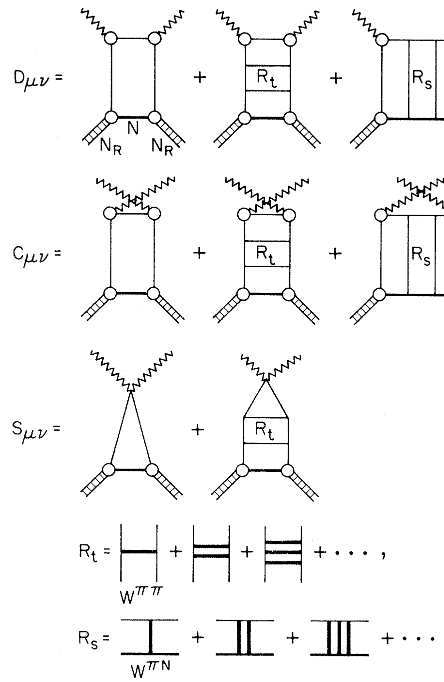

FIG. 3. The various diagrams whose sum is gauge-invariant in the case that interactions are permitted between internal mesons and nucleons (responsible for the physical nucleon being bound) and between the internal mesons themselves (responsible for the structure of the photon vertex). Note that gauge invariance requires the charged-particle propagators to be appropriately renormalized in Compton scattering.

pole contribution, but gives a fixed $J=0$ pole (originating from $S_{\mu\nu}$ in Fig. 3) only in the case that $Z^{\gamma}\neq 0$, i.e., only if the photon is elementary and not a bound state. The last point is trivial since the fixed pole in such models comes from the seagull-type diagrams which are proportional to $(Z^{\gamma})^2$. Similar to the model discussed previously the primary diagrams for $D_{\mu\nu}$ and $C_{\mu\nu}$ and those involving R_s (see Ref. 12) give no asymptotic contribution to T_1 and T_2. As in the previous discussion the diagrams involving R_t give the expected Regge contribution irrespective of whether the photon vertex function is assumed to satisfy a homogeneous or an inhomogeneous Bethe-Salpeter equation, i.e., irrespective of whether the vector particle is a bound state or an elementary particle (the reason being, as we have seen, that the integrand of the integral for the relevant diagrams is maximal when the internal mesons are close to their mass shell values).

In models where the charged constituent is an elementary nucleon and nonzero potentials are considered, the Dirac matrices are an additional complication. In such models, it is important to observe for $q_0\to\infty$ that although $P(q+x)$ is of order ν^0, a product $P(q+x)\,P(q+x')$ is of order ν^{-1} for finite x and x'. It is with this mechanism that higher order R_s type diagrams can be seen to be negligible with respect to the primary diagrams as is the case for the corresponding Drell-Lee model.[10,12]

In investigations of more complicated gauge invariant models, the same conclusion is found as illustrated here. That is, the existence of a $J=0$ fixed pole in Compton scattering is independent of the nature of the target hadron and the charge form factor of the elementary constituent, but depends solely on the photon being an elementary particle and not a bound state.[21]

This work strongly casts doubt on the arguments of Brandt *et al.*[14] and others[15] that there could be a $J=0$ fixed pole in ρ or even π photoproduction. Since models of the nature considered here can give a gauge-invariant description for ρ or π photoproduction,[13] and since in the Regge limit, i.e., $\nu\to\infty$, the squares of the 4-momenta (i.e., the external masses) are unimportant and the kinematics is essentially the same as considered here, there will not be a fixed pole unless the ρ or π meson is an elementary particle, i.e., one which does not lie on a Regge trajectory. This, however, would be hard to accept.

The absence of fixed poles in hadronic and photoproduction processes has also been obtained by Blankenbecler *et al.*[22] in the context of models utilizing the infinite-momentum frame. In the context of the Cambridge nonperturbative parton model fixed J-plane poles have been investigated by Landshoff and Polkinghorne[23] and Hughes and Osborn.[24] These groups of authors have also taken into account classes of gluon-exchange diagrams.

We conclude with some remarks on the vector-meson-dominance model. The question is whether fixed J-plane poles are in conflict with this model. We will recall that the basic hypothesis of this model is the current-field identity which connects the hadronic electromagnetic current with the fields of the vector mesons V. This identity allows us to re-express the t-channel current-current helicity amplitude $F_{NN,\gamma\gamma}$, or the amplitude $\tilde{F}_{NN,\gamma\gamma}$ which is free of kinematic singularities in s, in terms of the corresponding photoproduction amplitudes $F_{NN,\gamma V}$ or $\tilde{F}_{NN,\gamma V}$. Thus

$$\tilde{F}_{NN,\gamma\gamma}=\sum_{V}\tilde{F}_{NN,\gamma V}C_{V\gamma}\,, \tag{3.17}$$

where $C_{V\gamma}$ are coefficients. The statement that there is no fixed pole at $J=J_0$ in the photoproduction of Reggeized vector mesons V means[4,5] that

$$0=\int_{z_0}^{\infty} dz\,[\mathrm{Im}\tilde{F}_{NN,\,\gamma V}(z+i\epsilon,t) \mp(-1)^n\,\mathrm{Im}\tilde{F}_{NN,\,\gamma V}(z-i\epsilon,t)]\,P_{n-J_0-1}(z)\ , \tag{3.18}$$

where n is the larger of the moduli of the helicity differences of the ingoing and outgoing t-channel states and P is a Legendre polynomial. The relation (3.17) then implies

$$0=\int_{z_0}^{\infty} dz\,[\mathrm{Im}\tilde{F}_{NN,\,\gamma\gamma}(z+i\epsilon,t) \mp(-1)^n\,\mathrm{Im}\tilde{F}_{NN,\,\gamma\gamma}(z-i\epsilon,t)]\,P_{n-J_0-1}(z)\ .$$

Thus, unless the bare photon contributions are also introduced, the vector-meson-dominance model predicts a vanishing fixed-pole residue.[25]

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