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Current-current model for vector-meson photoproduction and related processes

Alfred Actor
Abteilung Physik, Universität Dortmund, Dortmund, West Germany
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We extend the Pomeron-photon analogy (i.e., the idea that the couplings of Pomerons and photons to hadronic matter are similar) to vector-meson photoproduction and related processes, such as ρ electroproduction and Compton scattering. To avoid directly coupling the Pomeron to photons, which is awkward in this model, we follow conventional vector-meson-dominance ideas and approximate the photon's hadronic component by the ρ meson. We then make a current-current ansatz for the ρ elastic scattering amplitude. Spin is fully taken into account. We find that s-channel helicity conservation (SCHC) cannot be realized within this framework if the effective current is conserved; SCHC requires that either this current has a nonconserved component or the forward cross section vanishes. If an effective current with two components is introduced—one a conserved electromagnetic central component, and the other a nonconserved peripheral component—then all the main features of ρ photoproduction can be accounted for. The same model gives a good qualitative description of pp, πp, and Kp elastic scattering, and of diffractive dissociation; in other words, it provides a unified description of all diffractive phenomena. Both components of the effective current have a simple explanation in terms of the quark-parton model.

I. INTRODUCTION

An interesting way to approximate the effects of Pomeron exchange is offered by the current-current model. In this model one makes the following ansatz for the amplitudes describing a diffractive reaction $a+b\rightarrow c+d$ at high energy:

$$f^{(s)}_{cdab}\propto\langle c|V_\mu|a\rangle\langle d|V^\mu|b\rangle. \tag{1.1}$$

Here V_μ is an effective-current operator which, since it is *a priori* undefined, might be either conserved or nonconserved. The amplitude (1.1) has several good features in addition to its simplicity:

(1) Spin 1 is effectively "exchanged" in the t channel because a four-vector V_μ contains spin-1 and spin-0 components (the latter is absent for a conserved current). The spin-1 component corresponds to a fixed pole in the t-channel J plane at $J=1$. Such a fixed pole is, at best, an approximation to a more detailed description which satisfies t-channel unitarity. However, this approximation may be a rather good one for many purposes since it leads to diffractive total and differential cross sections which are energy independent. Experimentally one knows that departures from such constant behavior are relatively weak at high energies.

(2) The amplitude (1.1) is factorized. If one ignores a possible spin-0 component in the current V_μ—this is legitimate at high energy—then it is easy to show that Eq. (1.1) becomes

$$f^{(s)}_{cdab}\propto g_{ca}(t)g_{db}(t). \tag{1.2}$$

The form factors $g_{ca}(t)$ and $g_{db}(t)$ determine the shape of the cross section. Experimentally one knows that factorization in elastic scattering and diffractive excitation is valid, at least for the momentum-transfer region $0<|t|<0.5$ GeV^2 (see Ref. 1). Therefore factorization is a desirable feature in the model.

(3) Inelastic diffractive reactions tend to conserve t-channel helicity.[2] This t-channel helicity conservation (TCHC) property is nearly exact in some cases [e.g., $N^*(1690)$ production] and less so in others [e.g., $N^*(1520)$ production]. The general trend is that TCHC becomes an increasingly good approximation as the mass of the diffractively produced system increases. Even when TCHC is not valid, it seems that t-channel helicity flip is limited to zero or one; larger spin flips appear to be absent. This behavior is an intrinsic feature of the current-current model. In the t channel (i.e., $D+b\rightarrow c+A$ with $D=\bar{d}$, $A=\bar{a}$) the current-current amplitudes are

$$f^{(t)}_{cADb}\propto\langle cA|V_\mu|0\rangle\langle 0|V^\mu|Db\rangle$$
$$\propto g_{cA}(t)g_{Db}(t). \tag{1.3}$$

Because of the spin-1 nature of V_μ, t-channel helicity flip is limited, for any value of t, to $|c-A|\leq 1$, $|D-b|\leq 1$. The various form factors present in the vertices determine the relative strengths of helicity-flip 1 and helicity-flip 0. The important point is that within the limited range of spin structure available in the model one can accommodate all the data on this aspect of diffractive excitation.

(4) Particle spin is fully taken into account in the current-current model. In this respect the model is an improvement on the Wu-Yang model[3] from which it originates. Wu and Yang suggested that elastic pp scattering amplitudes can be approximated by a product of two form factors. Tak-

ing these to be proton electromagnetic form factors, they obtained an acceptable fit to the pp cross section. Of course eikonalization[4] improves the fit to the data because elastic pp scattering is a central process, and rescattering effects are important. For diffractive excitation such rescattering effects (and thus eikonalization) should be less important because these inelastic processes are not central. Therefore the amplitude (1.1) may accurately reproduce the t dependence of inelastic diffractive cross sections, because these probably do not have diffractive minima.

In this article we investigate the suitability of the ansatz (1.1) for vector-meson photoproduction and some related processes (vector-meson electroproduction and Compton scattering). For these reactions one has good reason to think that s-channel helicity is approximately conserved at the vector-meson vertex. s-channel helicity conservation (SCHC) is also known to be a property of πN elastic scattering[5] at high energy, and within the current-current model this behavior is related to the smallness of the nucleon isoscalar anomalous magnetic moment. One lacks the corresponding experimental information about the electromagnetic couplings of vector mesons, of course, and so a quantitative test of the specific model (called the Pomeron-photon analogy) in which the current V_μ is assumed to be the isoscalar electromagnetic current is not possible. Nevertheless, one can inquire into the pros and cons of such a model for ρ photoproduction and the other reactions. In doing so we are led to a very interesting result. Namely, the current V_μ in the ansatz (1.1) must have a nonconserved component if SCHC is to be satisfied. (We do not learn this from πN or NN elastic scattering, where SCHC can be satisfied by a conserved current.) This is an important result, because for inelastic diffractive scattering the current V_μ must also have a nonconserved component. Therefore, our study of vector-meson photoproduction provides the final piece of information necessary for the unification of all diffractive reactions, elastic and inelastic, into a single framework. This framework is a current-current model with a two-component current V_μ; one component is a conserved electromagnetic component (the Pomeron-photon analogy), and the other is a nonconserved component about which we shall say more later.

Previous work on the current-current ansatz has shown that qualitative agreement with the inelastic data can be achieved, if the effective current V_μ is nonconserved.[6,7] A conserved current leads to forward dips in all inelastic diffractive cross sections, and experimentally these dips are not observed. It may be that at asymptotic energy the Pomeron really does decouple at zero angle from all inelastic vertices, and the model with a conserved current is an asymptotic model. Even if this is true, we are forced to work with a more complicated model at present energies, because we want the current V_μ to contain the isoscalar electromagnetic current. (In fact, there is some experimental support for this assumption.) Therefore to construct a realistic model we must add another ingredient. This could be an additional term on the right in Eq. (1.1), or it could be a nonconserved component in the effective current V_μ in addition to the electromagnetic component. However it is introduced, the new ingredient must provide the near-forward cross section for all inelastic diffractive processes. Clearly the most economical thing to do is to assume that this nonelectromagnetic component is only important at small momentum transfer, where it can be fitted to the data, and that for larger $|t|$ the electromagnetic component is dominant. In this way one can devise models which are nearly as unambiguous as the strict Pomeron-photon analogy, and which are consistent with the gross features of the data on all diffractive reactions.

For convenience we adopt the device of an effective current V_μ with two components, a conserved electromagnetic component and a nonconserved second component, which is not yet defined. This is the simplest way to preserve factorization. One may think of the nonconserved component as being a parametrization of all effects not encompassed by the first component. One may also seek to give the nonconserved component a definite meaning. We believe that both components have a clear interpretation in terms of the quark-parton model. The conserved electromagnetic component is associated with sea quarks, and it contributes mainly to central reactions. The weaker, nonconserved component is associated with valence quarks, and it contributes mainly to peripheral reactions. We discuss this in more detail in Sec. IV.

An important characteristic of our model (and in models with a two-component Pomeron[8]) is the distinction made between the small-$|t|$ (nonelectromagnetic) region and the larger-$|t|$ (electromagnetic) region. The changeover from small $|t|$ to larger $|t|$ occurs around $|t| \sim 0.1$ to $0.2\ \mathrm{GeV}^2$ (this is easy to estimate), and one can expect different behavior in these two regions. Even for πp and pp elastic scattering, where the conserved component in V_μ does *not* vanish in the forward direction, one may anticipate some difference between the small-$|t|$ and larger-$|t|$ regions because the nonelectromagnetic and electromagnetic components in the current presumably have different properties. The data on pp, πp, and Kp elastic scattering do show such an effect (for a review see

Leith[9,10]). For $|t|<0.15$ GeV^2 the slopes are larger than for $|t|>0.15$ GeV^2. Furthermore, in the forward region there is shrinkage with increasing energy while for $0.15<|t|<0.5$ GeV^2 there is little or no shrinkage. Strictly speaking, we should not talk about shrinkage in the current-current model, which does not allow for any. But we realize the model is only an approximation, whose energy dependence requires corrections. There is no reason to suppose these corrections are the same for both components. Indeed, they seem not to be.

In the following two sections we discuss the elastic reaction

$$\pi+\rho\rightarrow\pi+\rho \tag{1.4}$$

and the related electroproduction process

$$\pi+\gamma^{*}\rightarrow\pi+\rho \tag{1.5}$$

in terms of the current-current model. Our treatment is not quantitative. Our main goal is to determine whether reactions (1.4) and (1.5) fit into the scheme already described, or whether they require some new feature in the model. We are, of course, interested in the transitions $\rho\rightarrow\rho$ and $\gamma^{*}\rightarrow\rho$, and because the model is factorizable, we can just as well use pion targets as nucleon targets. Everything we shall say about $\rho\rightarrow\rho$ and $\gamma^{*}\rightarrow\rho$ is directly applicable to the more complicated reactions $N\rho\rightarrow N\rho$ and $N\gamma^{*}\rightarrow N\rho$, which are of experimental interest.

The main property of reaction (1.5) which we are concerned with is SCHC, which for real photons is known to be nearly exact in the range $0<|t|<0.2$ GeV^2, and accurate to within 10% for $0.2<|t|<1$ GeV^2.[9,11,12] The indications from ρ electroproduction are that SCHC is still valid when the photon becomes virtual.[13] Photoproduction of ω and ϕ mesons also seems to conserve s-channel helicity[9,11,12]; and the data on nucleon Compton scattering[14] are compatible with SCHC. Furthermore, it is reasonable to conjecture that reaction (1.4) also conserves s-channel helicity at high energy. This is a reasonable assumption in view of the fact that SCHC is also a property of πN elastic scattering.

SCHC is, of course, a very strong constraint to impose on any model. In Secs. II and III we shall see what the consequences of this constraint are for the current-current model. It turns out that in the absence of SCHC the model with a conserved current seems to be all right, as it does for πN and NN elastic scattering. But when SCHC is imposed on this model the forward $\pi\rho\rightarrow\pi\rho$ cross section has to vanish. Only if the effective current has a nonconserved component can one construct a realistic model. Therefore the experimental fact that SCHC is observed in ρ photoproduction forces us to introduce a nonconserved component into the effective current.

Two other important features of ρ photoproduction are[9,12]

(a) energy independence of the elastic cross section even at rather low energy, and

(b) nonshrinkage of the forward peak with increasing energy.

Both of these features are nicely compatible with the current-current ansatz, which unavoidably leads to constant $d\sigma/dt$ at sufficiently high energy. Photoproduction of ω and ϕ mesons seems to exhibit the same characteristics (a) and (b), although for these reactions the data are less plentiful. Concerning property (b) it should be mentioned that when the ρ photoproduction data are analyzed in terms of a model with central Pomeron-on- and peripheral f-exchange, the Pomeron contribution to the forward peak turns out to shrink with increasing energy.[15] Therefore one must be cautious when interpreting the data. However, the reaction $N+\gamma\rightarrow N+\phi$, which should be Pomeron-dominated already at rather low energy because of the absence of other t-channel exchanges, does *not* show shrinkage,[9,12] at least in the t range which has been investigated. Therefore the evidence for a constant forward slope in the Pomeron contribution is somewhat contradictory, and we shall have to wait for new experiments, especially on $\gamma\rightarrow\phi$ at small $|t|$, to decide the issue.

Another (preliminary) experimental result is that for quite small $|t|$ (say $|t|<0.2$ GeV^2) the cross section seems to steepen[9,12]. This effect, if it is confirmed by later measurements, would be very similar to the behavior of the pp, πp, and Kp cross sections at small $|t|$ already mentioned. For the latter reactions the cross section is steeper below $|t|\sim 0.15$ GeV^2 than above, and, furthermore, in the small-$|t|$ region there is shrinkage while for larger $|t|$ there is little or none. The same thing may happen in photoproduction. This is an important possibility which will be investigated in future experiments. If this turns out to be correct, then two-component models of the Pomeron will gain in credibility.

This paper is arranged as follows. In Sec. II we perform the mechanical task of constructing the current-current amplitudes for $\pi\rho\rightarrow\pi\rho$ with off-shell ρ mesons. In Sec. III we use the conventional vector-meson-dominance (VMD) approximation for the coupling of photons to hadrons to write the amplitudes for ρ electroproduction in terms of the ones for $\pi\rho\rightarrow\pi\rho$. Here we find that the usual ρ-mass-continuation procedure does not work for the current-current ansatz. One cannot treat the ρ as though it is coupled to a conserved current because this conflicts with SCHC. One cannot satisfy SCHC and current conservation on the ρ leg simultaneously (in the current-current model). This could be a failure of the ansatz, with its fixed $J=1$ exchange structure, but it is also possible that the present continuation methods in VMD are too restrictive. We wish to point out that the reactions which have received the most detailed study, namely

$$N+\gamma^* \to N+\pi,$$

$$N+\rho \to N+\pi,$$

are nondiffractive, and for them Pomeron exchange and SCHC do not play any role. Covariant continuation methods, which seem to be adequate for pion photoproduction, may not be adequate for ρ photoproduction with its remarkable SCHC behavior. We have considered some different continuation procedures, ones which are compatible with the current-current model. One such scheme, which is based on smoothness in the mass dependence of invariant amplitudes, is introduced in Sec. III and is discussed in more detail in the final section of this paper.

In Sec. IV we also compare our model with previous work on the two-component Pomeron. The model is shown to provide a synthesis of many previous ideas and results.

II. ρ ELASTIC SCATTERING

We now turn to a detailed discussion of the quasielastic reaction

$$\pi+\rho \to \pi+\rho, \tag{2.1}$$

where one or both of the ρ mesons can be off-shell. Our primary interest is in the current-current model for this process, but we shall keep the discussion fairly general. First we construct a model with a conserved current and calculate the helicity amplitudes. When SCHC is imposed the model turns out to be too restrictive. Next we try a nonconserved current, and find that in this case SCHC can be satisfied to any desired degree of accuracy. Such a model could therefore provide a fairly realistic description, or parametrization, of $\pi\rho$ elastic scattering. In the following section we discuss the problem of continuing the ρ masses to zero, and extending the model to ρ photoproduction.

Notation used throughout the paper is summarized in Figs. 1 and 2. The s channel is called $a+b\to c+d$; thus $s=(p_a+p_b)^2$ and

$$t=k^2, \tag{2.2}$$

with

$$k=p_d-p_b=p_a-p_c. \tag{2.3}$$

The masses

$$m_c=m_a=m_\pi \tag{2.4}$$

are equal. The masses m_d, m_b of the two ρ mesons in Eq. (2.1) may be equal (elastic $\pi\rho$ scattering) or unequal (e.g., ρ photoproduction).

The current-current amplitude for reaction (2.1) is

$$f^{(s)}_{0d0b}=\lambda\langle c|V_\mu|a\rangle\langle d|V^\mu|b\rangle, \tag{2.5}$$

where the current matrix elements are

$$\langle c|V_\mu|a\rangle=(p_c+p_a)_\mu f_\pi(t), \tag{2.6}$$

$$\langle d|V_\mu|b\rangle=\epsilon^*_{d_\alpha}(p_d)\epsilon_{b_\beta}(p_b)M_\mu^{\alpha\beta}, \tag{2.7}$$

with

$$\begin{aligned} M_{\mu\alpha\beta}=&(k^2p_{b_\mu}-k_\mu k\cdot p_b)[g_{\alpha\beta}h_1(t)+p_{b\alpha}p_{d\beta}h_2(t)]\\ &+(k\cdot p_d g_{\mu\alpha}-p_{d\mu}k_\alpha)p_{d\beta}h_3(t)\\ &+p_{b\alpha}(k\cdot p_b g_{\mu\beta}-p_{b\mu}k_\beta)h_4(t), \end{aligned} \tag{2.8}$$

and λ is a constant. The current (2.6) is conserved because $m_c=m_a=m_\pi$. The vector-meson current (2.7) is conserved because

$$k^\mu M_{\mu\alpha\beta}=0 \tag{2.9}$$

by construction. At the pomeron-ρ vertex there are four form factors $h_i(t)$. This is because the

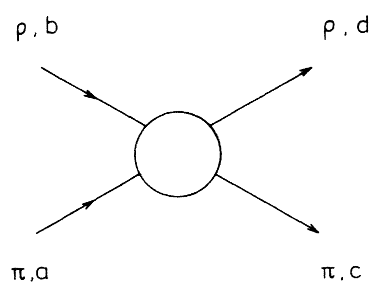

FIG. 1. The s channel $\pi+\rho\to\pi+\rho$.

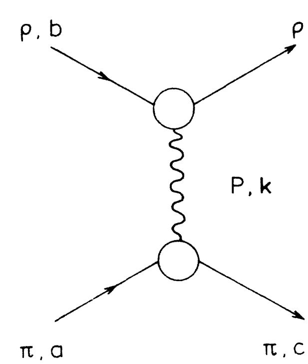

FIG. 2. One-Pomeron-exchange amplitude for $\pi+\rho\to\pi+\rho$.

Pomeron has spin 1 in the model with a conserved effective current. As we shall discuss later, there are two more form factors if the current is not conserved.

From Eqs. (2.5)–(2.8) the helicity amplitudes are found to be

$$f^{(s)}_{0d0b}=\epsilon^*_{d\alpha}(p_d)\epsilon_{b\beta}(p_b)M^{\alpha\beta}\,, \tag{2.10}$$

with

$$\begin{aligned}M_{\alpha\beta}&=\lambda f_\pi(t)(p_c+p_a)^\mu M_{\mu\alpha\beta}\\&=g_{\alpha\beta}A_1+p_{b\alpha}p_{d\beta}A_2+p_{c\alpha}p_{d\beta}A_3+p_{b\alpha}p_{a\beta}A_4\\&\quad+p_{c\alpha}p_{a\beta}A_5+p_{b\alpha}p_{b\beta}A_6+p_{c\alpha}p_{b\beta}A_7\,,\end{aligned} \tag{2.11}$$

and the invariant amplitudes are

$$\begin{aligned}A_1&=p_b\cdot(p_c+p_a)\lambda k^2 f_\pi h_1\\&\approx st\lambda f_\pi h_1\,,\end{aligned} \tag{2.12}$$

$$\begin{aligned}A_2&=\lambda f_\pi[p_b\cdot(p_c+p_a)(k^2h_2+h_3-h_4)\\&\qquad-k\cdot p_d h_3-k\cdot p_b h_4]\\&\approx s\lambda f_\pi(th_2+h_3-h_4)\,,\end{aligned} \tag{2.13}$$

$$\begin{aligned}A_3&=2k\cdot p_d\lambda f_\pi h_3\\&=(t+m_d^{\,2}-m_b^{\,2})\lambda f_\pi h_3\,,\end{aligned} \tag{2.14}$$

$$\begin{aligned}A_4&=2k\cdot p_b\lambda f_\pi h_4\\&=-(t-m_d^{\,2}+m_b^{\,2})\lambda f_\pi h_4\,,\end{aligned} \tag{2.15}$$

$$A_5=0\,. \tag{2.16}$$

A_6 and A_7 do not contribute directly to the helicity amplitudes because

$$p_b\cdot\epsilon_n(p_b)=0\;; \tag{2.17}$$

only if the M function (2.11) satisfies the condition

$$p_b^\beta M_{\alpha\beta}=0 \tag{2.18}$$

are the invariant amplitudes A_6, A_7 related (by constraints) to physical amplitudes. Equation (2.18) corresponds to an incoming ρ meson which is coupled to a conserved current. If one tries to construct a model along these lines then terms proportional to $p_{b\beta}$ have to be included in the vertex (2.8). We shall not assume that the condition (2.18) is satisfied, and therefore we can simply ignore A_6, A_7.

The amplitude A_5 is proportional to the t-channel double-flip amplitude (see the Appendix). Therefore it must be zero in the current-current model, which contains no admixture of spin-2 exchange in the t channel.

Calculating the helicity amplitudes (2.10) we find

$$\begin{aligned}4m_am_bf^{(s)}_{0d0b}&=4m_dm_b\epsilon^*_d(p_d)\cdot\epsilon_b(p_b)A_1\\&\quad-(-1)^{d+b}T_{db}^{\;2}d_{0d}^{\;1}(\chi_d)d^1_{0b}(\chi_b)A_2\\&\quad-\delta_{d0}(-1)^bS_{cd}T_{db}d^1_{0b}(\chi_b)A_3\\&\quad+\delta_{b0}(-1)^dS_{ab}T_{db}d^1_{0d}(\chi_d)A_4\\&\quad+\delta_{d0}\delta_{b0}S_{cd}S_{ab}A_5\,.\end{aligned} \tag{2.19}$$

Here we have used the formulas

$$2m_dp_b\cdot\epsilon_d(p_d)=(-1)^dT_{db}d^1_{0d}(\chi_d)\,, \tag{2.20}$$

$$2m_dp_c\cdot\epsilon_d(p_d)=\delta_{0d}S_{cd}\,, \tag{2.21}$$

$$2m_bp_d\cdot\epsilon_b(p_b)=-(-1)^bT_{db}d^1_{0b}(\chi_b)\,, \tag{2.22}$$

$$2m_bp_a\cdot\epsilon_b(p_b)=\delta_{0b}S_{ab}\,, \tag{2.23}$$

where the crossing angles[16] χ_d, χ_b are defined by

$$S_{cd}T_{db}\sin\chi_d=2m_d\sqrt{\phi}\,, \tag{2.24}$$

$$\begin{aligned}S_{cd}T_{db}\cos\chi_d&=-(s-m_c^{\,2}+m_d^{\,2})(t+m_d^{\,2}-m_b^{\,2})\\&\quad-2m_d^{\,2}(m_c^{\,2}-m_a^{\,2}-m_d^{\,2}+m_b^{\,2})\,,\end{aligned} \tag{2.25}$$

$$S_{ab}T_{db}\sin\chi_b=2m_b\sqrt{\phi}\,, \tag{2.26}$$

$$\begin{aligned}S_{ab}T_{db}\cos\chi_b&=(s-m_a^{\,2}+m_b^{\,2})(t-m_d^{\,2}+m_b^{\,2})\\&\quad-2m_b^{\,2}(m_c^{\,2}-m_a^{\,2}-m_d^{\,2}+m_b^{\,2})\,,\end{aligned} \tag{2.27}$$

and S_{ij} and T_{ij} are given by

$$S_{ij}^{\;2}=s^2-2s(m_i^{\,2}+m_j^{\,2})+(m_i^{\,2}-m_j^{\,2})^2\,, \tag{2.28}$$

$$T_{db}^{\;2}=t^2-2t(m_d^{\,2}+m_b^{\,2})+(m_d^{\,2}-m_b^{\,2})^2\,. \tag{2.29}$$

ϕ is the usual physical boundary function.[17]

For large s the amplitudes (2.19) are

$$2f^{(s)}_{11}\approx-(1+\cos\theta)A_1-(\phi/s^2)A_2\,, \tag{2.30}$$

$$2f^{(s)}_{1,-1}\approx(1-\cos\theta)A_1+(\phi/s^2)A_2\,, \tag{2.31}$$

$$\begin{aligned}2\sqrt{2}m_bf^{(s)}_{10}&\approx-\sqrt{s}\sin\theta A_1\\&\quad+(\phi/s^2)^{1/2}(t-m_d^{\,2}+m_b^{\,2})A_2\\&\quad-\sqrt{\phi}A_4\,,\end{aligned} \tag{2.32}$$

$$\begin{aligned}2\sqrt{2}m_df^{(s)}_{01}&\approx\sqrt{s}\sin\theta A_1\\&\quad-(\phi/s^2)^{1/2}(t+m_d^{\,2}-m_b^{\,2})A_2\\&\quad+\sqrt{\phi}A_3\,,\end{aligned} \tag{2.33}$$

$$\begin{aligned}4m_dm_bf^{(s)}_{00}&\approx2[\tfrac{1}{2}s(1-\cos\theta)-m_d^{\,2}-m_b^{\,2}]A_1\\&\quad+[t^2-(m_d^{\,2}-m_b^{\,2})^2]A_2\\&\quad-s(t-m_d^{\,2}+m_b^{\,2})A_3\\&\quad-s(t+m_d^{\,2}-m_b^{\,2})A_4+s^2A_5\,.\end{aligned} \tag{2.34}$$

Here we have used the approximations

$$T_{db}\cos\chi_d \approx -(t+m_d^{\,2}-m_b^{\,2}) ,$$
$$T_{db}\cos\chi_b \approx (t-m_d^{\,2}+m_b^{\,2}) . \tag{2.35}$$

Furthermore, when $|t| \ll s$ the following approximations are acceptable:

$$\sin\theta \approx 2(|t|/s)^{1/2} , \tag{2.36}$$

$$\cos\theta \approx 1-2|t|/s , \tag{2.37}$$

$$\phi \approx s^2|t| . \tag{2.38}$$

Let us now consider the restrictions implied by SCHC, beginning with the double-flip amplitude (2.31). Because $1-\cos\theta \approx 2|t|/s$, we see that A_1 does not contribute to leading order in this amplitude. (To leading order all the helicity amplitudes are proportional to s.) Therefore we only have to consider the A_2 term, and for the double-flip amplitude to be negligible it is necessary that

$$A_2 \approx 0 . \tag{2.39}$$

Next consider the single-flip amplitudes (2.32) and (2.33). For these to be negligible the conditions

$$2A_1 \approx -sA_3 \approx -sA_4 \tag{2.40}$$

have to be satisfied. The nonflip amplitudes are therefore

$$2f_{11}^{(s)} \approx (1-|t|/s)sA_3 , \tag{2.41}$$

$$4m_d m_b f_{00}^{(s)} \approx (|t|+m_d^{\,2}+m_b^{\,2})sA_3 + s^2A_5 . \tag{2.42}$$

Note that there is no SCHC constraint on the invariant amplitude A_5, which contributes only to $f_{00}^{(s)}$. The fact that $A_5=0$ in the current-current model means that both $f_{11}^{(s)}$ and $f_{00}^{(s)}$ are determined by

$$A_3 = (m_d^{\,2}-m_b^{\,2}-|t|)\lambda f_\pi h_3 ,$$

which has a zero at $|t| = m_d^{\,2}-m_b^{\,2}$. In the case of elastic scattering the zero occurs at $t=0$ and both nonflip amplitudes vanish in the forward direction. This means the total cross section is zero. We see that SCHC cannot be satisfied by a conserved current. One or the other of these conditions must be relaxed.

If the effective current is not conserved then the nonflip amplitudes do not vanish at $|t| = m_d^{\,2}-m_b^{\,2}$, even with exact SCHC. To show this we go back to Eq. (2.8), and replace this tensor by one appropriate to a nonconserved current V_μ. A suitable choice is

$$M'_{\mu\alpha\beta} = p_{b\mu}(g_{\alpha\beta}h'_1 + p_{b\alpha}p_{d\beta}h'_2) + g_{\mu\alpha}p_{d\beta}h'_3 + g_{\mu\beta}p_{b\alpha}h'_4 + k_\mu(g_{\alpha\beta}h'_5 + p_{b\alpha}p_{d\beta}h'_6) . \tag{2.43}$$

The pion matrix element is

$$\langle c|V_\mu|a\rangle = (p_c+p_a)_\mu f_\pi + k_\mu f_s . \tag{2.44}$$

There are now six form factors at the vector-meson vertex. However, h'_5 and h'_6 correspond to spin-0 exchange in the t channel and so these two form factors are unimportant at high energy. For the same reason, $f_s(t)$ in Eq. (2.44) is unimportant. The M function still has the form (2.11), but now the invariant amplitudes are

$$A_1 = p_b\cdot(p_c+p_a)\lambda f_\pi h'_1 + \lambda f_s(k\cdot p_b h'_1 + k^2h'_5) \approx s\lambda f_\pi h'_1 , \tag{2.45}$$

$$A_2 = \lambda f_\pi[p_b\cdot(p_c+p_a)h'_2 - h'_3 - h'_4] + \lambda f_s(k\cdot p_b h'_2 - h'_3 + h'_4 + k^2h'_6) \approx s\lambda f_\pi h'_2 , \tag{2.46}$$

$$A_3 = 2\lambda f_\pi h'_3 , \tag{2.47}$$

$$A_4 = 2\lambda f_\pi h'_4 , \tag{2.48}$$

$$A_5 = 0 . \tag{2.49}$$

Imposing the SCHC constraints (2.39) and (2.40) we again find Eqs. (2.41) and (2.42) for the nonflip amplitudes. Both are proportional to $A_3 = 2\lambda f_\pi h'_3$, which has no zero at $t=0$. Therefore the model with a nonconserved current is qualitatively acceptable.

The SCHC conditions (2.39) and (2.40) imply that only one form factor at the ρ vertex is independent:

$$-h'_1 \approx h'_3 \approx h'_4 \approx h'(t), \quad h'_2 \approx 0 . \tag{2.50}$$

Of course, these are not exact conditions, and they refer only to the dominant s-channel-helicity-conserving contribution. If we think of the current V_μ as being the sum of an electromagnetic component and a nonconserved component then the form factor $h'(t)$ in Eq. (2.50) can be written

$$h'(t) = \beta(t) + (m_d^{\,2}-m_b^{\,2}-|t|)h_{\rm em}(t) , \tag{2.51}$$

where $\beta(t)$ corresponds to the nonelectromagnetic component.

Finally, we mention a very characteristic feature of the current-current model. For $\pi\rho$ elastic scattering at relatively small $|t|$, Eqs. (2.41) and (2.42) become

$$f_{11}^{(s)} \approx f_{00}^{(s)} \approx sA_3/2 . \tag{2.52}$$

Therefore if SCHC is valid then the ρ helicity states ±1 and 0 are equally populated. This is because $A_5=0$ in the model. By factorization the same is true for ρN elastic scattering.

III. ρ PHOTOPRODUCTION

The reactions of experimental interest are ρ photoproduction and electroproduction,

$$\pi + \gamma^* \rightarrow \pi + \rho\,, \tag{3.1}$$

where the photon γ^* is real or virtual (spacelike). Following the conventional idea that vector mesons dominate the interactions between photons and hadrons (see Fig. 3), we assume that the γ^* turns into an off-shell ρ meson which then quasielastically scatters from the pion. This means the amplitudes for reaction (2.1) determine the amplitudes for reaction (3.1), up to some phenomenological factor. However, the $\pi\rho \rightarrow \pi\rho$ amplitudes have to be continued off-shell, and this is not a unique procedure. Several continuation schemes have been proposed in the extensive literature on vector-meson dominance, and some of these shall be described below. As we shall see, the current-current model is only compatible with a continuation in which the ρ meson is *not* coupled to a conserved current, even when it has zero mass. This excludes all of the standard covariant treatments, and so we must look for a new continuation procedure. We decide to continue the *on-shell invariant amplitudes* for elastic scattering to nonphysical values of the ρ mass, leaving the *helicity coefficients of these invarient amplitudes unchanged.* Then a massless ρ can have helicity zero as well as helicity ± 1, but of course these ρ amplitudes with helicity zero are decoupled from photon amplitudes.

Reaction (3.1) is best studied within the context of the electroproduction reaction

$$e + \pi \rightarrow e + \pi + \rho\,, \tag{3.2}$$

which is depicted in Fig. 4. The photoproduction amplitudes are obtained by taking the $q^2 \rightarrow 0$ limit of the electroproduction amplitudes and simultaneously discarding the electron vertex. The amplitude for reaction (3.2) is (see Fig. 4)

$$T = \frac{e^2}{q^2}\,\overline{u}(q_1)\gamma_\mu u(q_2)\langle cd|J^\mu|a\rangle\,, \tag{3.3}$$

where $q = q_2 - q_1$, and J_μ is the electromagnetic current. Current conservation implies that

$$q^\mu\langle cd|J_\mu|a\rangle = 0\,, \tag{3.4}$$

and of course one trivially verifies that

$$\begin{aligned} q^\mu\overline{u}(q_1)\gamma_\mu u(q_2) &= (m_2 - m_1)\overline{u}(q_1)u(q_2) \\ &= 0\,, \end{aligned} \tag{3.5}$$

because the electron masses $m_1 = m_2$ are equal. Using the identity

$$g_{\mu\nu} = \sum_n (-1)^n \epsilon^*_{n\mu}(q)\epsilon_{n\nu}(q) - \frac{q_\mu q_\nu}{q^2} \tag{3.6}$$

for spacelike q_μ we can rewrite Eq. (3.3) in the form

$$T = \frac{e^2}{q^2}\sum_n (-1)^n \Gamma_n T_{dn}\,, \tag{3.7}$$

where

$$\Gamma_n \equiv \epsilon^*_{n\mu}(q)\overline{u}(q_1)\gamma^\mu u(q_2)\,, \tag{3.8}$$

$$T_{dn} \equiv \epsilon_{n\mu}(q)\langle cd|J^\mu|a\rangle\,. \tag{3.9}$$

The amplitudes T_{dn} are helicity amplitudes for the virtual reaction (3.1). In the limit $q^2 \rightarrow 0$ the longitudinal helicity state $n = 0$ decouples from the transverse-polarization states $n = \pm 1$. To see this we recall the formula, valid for spacelike q_μ,

$$(-q^2)^{1/2}\epsilon_{0\mu}(q) = (|\vec{q}|,\, -q_0\hat{q})\,. \tag{3.10}$$

When $q^2 \sim 0$ one can verify that

$$(-q^2)^{1/2}\epsilon_{0\mu}(q) \approx q_\mu - \frac{1}{2|\vec{q}|q_0}\,q^2(|\vec{q}|, q_0\hat{q})\,. \tag{3.11}$$

Then, because of Eq. (3.5),

$$\Gamma_0 \approx (-q^2)^{1/2}\overline{u}_1\gamma_0 u_2/|\vec{q}|\,, \tag{3.12}$$

and Γ_0 is $O(\sqrt{q^2})$. The longitudinal-photon amplitudes T_{d0} also vanish like $O(\sqrt{q^2})$ at $q^2 \sim 0$ because of gauge invariance. Therefore the $n = 0$ term in Eq. (3.7) is $O(q^2)$ relative to the $n = \pm 1$ terms.

In the VMD approach one assumes that the amplitudes T_{dn} for $\pi\gamma^* \rightarrow \pi\rho$ and $f^{(s)}_{db}$ for $\pi\rho \rightarrow \pi\rho$ are re-

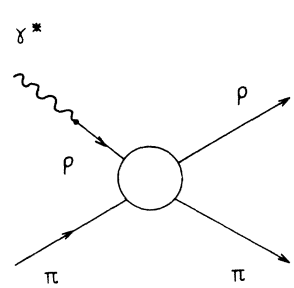

FIG. 3. Vector-meson-dominance diagram for ρ photoproduction and electroproduction.

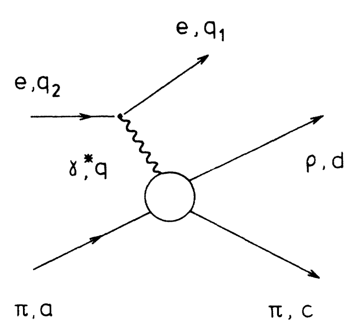

FIG. 4. One-photon-exchange diagram for ρ electroproduction.

lated by[18]

$$T_{d,\pm1}=\frac{e}{2\gamma_\rho}F(q^2)f^{(s)}_{d,\pm1}\,,$$
$$T_{d0}=\frac{e}{2\gamma_\rho}F(q^2)(-q^2)^{1/2}\frac{1}{m_\rho}f^{(s)}_{d0}\,, \tag{3.13}$$

where γ_ρ is a phenomenological constant, and $F(q^2)$ is a form factor normalized to unity, $F(0)=1$. Conventionally[18]

$$F(q^2)=m_\rho^2/(m_\rho^2-q^2),$$

a result following directly from the current-field identity

$$J^{\text{em}}_\mu=\sum_V \frac{e}{2\gamma_V}m_V^2\phi^V_\mu\,, \tag{3.14}$$

where ϕ^V_μ is the vector-meson field and $V=\rho,\omega,\phi$. [Equations (3.13) should of course have additional terms corresponding to ω and ϕ. But we are not interested here in quantitative work with the VMD model, and these terms shall be ignored.] As mentioned earlier, we do not assume that the ρ is coupled to a conserved current, and therefore $\partial^\mu\phi^\rho_\mu \neq 0$. This means a term with the form $-\partial_\mu\Phi$, where Φ is a scalar operator, has to be added to the right-hand side of Eq. (3.14) to ensure that $\partial^\mu J^{\text{em}}_\mu=0$. This extra term $-\partial_\mu\Phi$ does not contribute to any physical amplitude, and therefore Eqs. (3.13) still hold.

Concerning Eqs. (3.13) we note the following important points:

(a) T_{dn} is defined in terms of a gauge-invariant M function $\tilde M_{\alpha\mu}$ [see Eq. (3.9)]

$$T_{dn}=\epsilon^*_{d\alpha}(p_d)\epsilon_{n\mu}(q)\tilde M^{\alpha\mu}\,, \tag{3.15}$$

where

$$q^\mu\tilde M_{\alpha\mu}=0. \tag{3.16}$$

When $q^2=0$ the longitudinal-photon amplitudes T_{d0} vanish like $O(\sqrt{q^2})$ because of the gauge-invariance condition (3.16). This is entirely necessary because a massless photon is a physical particle, which, on grounds of Lorentz invariance alone, cannot have a zero-helicity state. The transverse-photon amplitudes $T_{d,\pm1}$ survive the limit $q^2\to0$ of course.

(b) $f^{(s)}_{db}$ is defined in terms of the M function $M_{\alpha\beta}$ in Eq. (2.11) which is *not* gauge invariant. Therefore the amplitudes $f^{(s)}_{d0}$ need not vanish at zero mass. A massless ρ meson is not a physical particle, and there is no reason why it should not have zero helicity. Only physical particle states are classified as representations of the Lorentz group. When a particle state is continued off-shell, everything depends on the continuation, which is not unique.

Our continuation scheme is very simple. We begin with the on-shell $\pi\rho\to\pi\rho$ scattering amplitudes

$$2f^{(s)}_{1,\pm1}\approx\mp(1\pm\cos\theta)A_1\mp(\phi/s^2)A_2, \tag{3.17}$$

$$2\sqrt2\,m_\rho f^{(s)}_{10}\approx-\sqrt s\,\sin\theta A_1+t(\phi/s^2)^{1/2}A_2-\sqrt\phi\,A_4, \tag{3.18}$$

$$2\sqrt2\,m_\rho f^{(s)}_{01}\approx\sqrt s\,\sin\theta A_1-t(\phi/s^2)^{1/2}A_2+\sqrt\phi\,A_3, \tag{3.19}$$

$$4m_\rho^2f^{(s)}_{00}\approx[s(1-\cos\theta)-m_\rho^2]A_1+t^2A_2-st(A_3+A_4)+s^2A_5. \tag{3.20}$$

Then we continue the invariant amplitudes $A_1,\ldots,A_5$ to $p_b^2=m_b^2=q^2$ (whatever this entails). In other words, we assume that $A_1,\ldots,A_5$ are smooth functions of m_b^2. The mass $m_b=m_\rho$ in all of the invariant amplitude coefficients in Eqs. (3.17)–(3.20) remains unchanged. Only the invariant amplitudes are continued. Therefore the longitudinal amplitudes $f^{(s)}_{d0}$ do not vanish at $m_b=0$, nor are they infinite there; they are just two of the functions one has obtained from the on-shell scattering amplitudes by mass continuation of the invariant amplitudes.

Evidently the Lorentz-transformation property of the on-shell amplitudes $f^{(s)}_{db}$ is not affected by the continuation to $m_b^2=q^2$. Nevertheless, we are assuming in Eqs. (3.13) that these amplitudes are proportional to electroproduction amplitudes T_{db} which have a different Lorentz-transformation property. To avoid a conflict with Lorentz invariance we must admit that our procedure is not frame-independent, and we must choose a particular frame in which it is supposed to hold. The natural choice is the s-channel c.m. frame. Equations (3.13) are therefore defined in this frame. The functions $f^{(s)}_{db}$ are obtained from (3.17)–(3.20) as described, and the electroproduction amplitudes T_{db} are then determined in this frame. While one may regret that covariance has been lost, this is not an unusual development in treatments of the hadronic interactions of photons (see e.g. Yennie[19]). We shall see later that by sacrificing covariance we have made a substantial gain in simplifying the coupling of Pomerons with photons.

Having given a definite meaning to Eqs. (3.13), our construction of a current-current model for reaction (3.1) is formally complete. The electroproduction amplitude (3.7) is determined via the VMD relations (3.13) in terms of $\pi\rho\to\pi\rho$ amplitudes. The latter are given in Eqs. (3.17)–(3.20) (or by exact formulas in Sec. II) in terms of the invariant amplitudes A_i, and these are given as functions of the form factors $h_i'(t)$ associated with a nonconserved current V_μ by Eqs. (2.45)–(2.49).

There is no difficulty with SCHC, which implies the $h_i'(t)$ approximately satisfy the conditions (2.50). To develop a rough model for ρ photoproduction one could first disregard SCHC breaking, so the only important amplitude is

$$f_{11}^{(s)} \approx -A_1 \approx -s\lambda f_\pi(t) h'(t),$$

and the cross section has the form

$$\frac{d\sigma}{dt} \sim |\lambda f_\pi h'|^2.$$

Choosing $f_\pi(t)$ to be the pion electromagnetic form factor, or the form factor determined from πN elastic scattering by means of the Chou-Yang model,[20] one could then obtain the form factor $h'(t)$ from $d\sigma/dt$.

We remind the reader that $h'(t)$ corresponds to a current with two components,

$$h'(t) = t h_{\rm em}(t) + \beta(t),$$

where $h_{\rm em}(t)$ is electromagnetic and $\beta(t)$ is nonelectromagnetic. The near-forward cross section is proportional to $|\beta(t)|^2$, while for larger $|t|$ (say $|t| > 0.1$ or $0.2\ \mathrm{GeV}^2$) the electromagnetic component is dominant. To estimate the changeover value $t = t_c$, one can fit $\beta(t)$ to the small-$|t|$ data and $h_{\rm em}$ to the data at larger $|t|$. An easy way to obtain an upper bound on t_c is to approximate $h_{\rm em}$ by an exponential and calculate the position of the maximum of the function $|t h_{\rm em}|^2$, which must be larger than t_c. The data[9,12] on vector-meson photoproduction do suggest there is some sort of changeover in the region $|t| \sim 0.1$ to $0.2\ \mathrm{GeV}^2$, with a larger slope at the smaller $|t|$ values, as in elastic pp scattering. If this is true, then one may also expect more energy dependence in the small-$|t|$ region than for larger $|t|$, as is observed in the pp case. This would mean the function $\beta(t)$ is not energy independent, but rather it decreases weakly with increasing energy. Such behavior is acceptable because we have given the nonconserved component of the current V_μ no interpretation up till now. Indeed, we have only included it in the current-current ansatz for convenience because of factorization. Therefore we need not worry about some residual energy dependence in the function $\beta(t)$. In contrast, the electromagnetic form factor $h_{\rm em}(t)$ must be strictly energy independent. This means (for sufficiently high energy) the cross section for $|t| > |t_c|$ must not shrink.

IV. DISCUSSION

(1) The most attractive feature of the mass-continuation scheme we employ is that it simplifies the Pomeron-photon coupling problem. Well-known complications[21] attend this coupling when the photon is treated as an external leg on the strong-interaction amplitude. In the VMD approach one encounters essentially the same complications when one assumes that the ρ couples to a conserved current. We do not make this assumption; therefore our strong-interaction amplitude does not satisfy current conservation on the ρ leg, and the longitudinal ρ amplitudes do not vanish when the ρ mass is continued to zero. Only the invariant amplitudes are continued in the ρ mass, while all helicity coefficients multiplying the invariant amplitudes are left unchanged. The Lorentz-transformation properties of the helicity amplitudes are not affected by the mass continuation. Whether the photoproduction amplitudes obtained in this way are on the whole more useful, or less useful, than ones obtained by other mass-continuation procedures is a question we cannot answer here. However, we shall see shortly that the usual continuation based on the current-field identity does not work for our model.

It is evidently a good idea to perform mass continuations in terms of invariant amplitudes, because these amplitudes are free from kinematic singularities and constraints (i.e., singularities and constraints which depend explicitly on the external masses). In the current-current model one expresses the invariant amplitudes as products of covariant form factors. The form factors at the vector-meson vertex are therefore the functions one continues in the ρ mass.

This version of the VMD approximation for the strong interactions of photons can be used for Compton scattering and photon-photon scattering as well as for electroproduction. For Compton scattering, the strong-interaction amplitude is the one for $\rho N \to \rho N$. For $\gamma\gamma \to \gamma\gamma$ and $\gamma\gamma \to$ hadrons the strong-interaction amplitudes are $\rho\rho \to \rho\rho$ and $\rho\rho \to$ hadrons, respectively. Only the on-shell invariant amplitudes are continued in each case, and Lorentz transformation properties are unaffected. There is no need for a fixed pole in any amplitude.

Factorization of photon-hadron cross sections is obvious in our approach to the Pomeron-photon coupling problem. For example, the prediction $\sigma_{\gamma\gamma}\sigma_{NN} \sim \sigma_{\gamma N}{}^2$ follows from the corresponding prediction $\sigma_{\rho\rho}\sigma_{NN} \sim \sigma_{\rho N}{}^2$. When one works directly with current amplitudes (i.e., ones with external photon legs), this factorization property is obscured by the presence of fixed poles.[22]

Of course, a price must be paid for all this convenience. Our model is frame-dependent, and it rests heavily on the VMD phenomenology. However, the latter has been qualitatively rather successful, and for our purpose here (namely to investigate the Pomeron-photon analogy) it should be reliable enough.

(2) Let us take pion Compton scattering as an example. The basic strong-interaction amplitudes are the elastic $\pi\rho\to\pi\rho$ amplitudes (3.17)–(3.20). Starting with on-shell amplitudes, we continue the invariant amplitudes A_i to $m_a = m_b = 0$, leaving the invariant amplitude coefficients unchanged. Therefore the spin structure of the helicity amplitudes is unchanged. The relevant $\rho\to\rho$ amplitudes for on-shell photons are

$$2f^{(s)}_{1,\pm 1} \approx \mp(1\pm\cos\theta)A_1 \mp(\phi/s^2)A_2. \tag{4.1}$$

The other three $\rho\to\rho$ amplitudes $f^{(s)}_{10}$, $f^{(s)}_{01}$, and $f^{(s)}_{00}$, which are decoupled from the Compton scattering amplitudes, remain nonzero. In the current-current model the invariant amplitudes are given by Eqs. (2.45)–(2.49). Only the form factors $f_\pi(t)$ and $h'_{1,2}(t)$ are important for the Compton scattering amplitudes

$$T_{1,\pm 1} = \left(\frac{e}{f_\rho}\right)^2 f^{(s)}_{1,\pm 1}. \tag{4.2}$$

To build SCHC into the model we only have to assume that $h'_2 \sim 0$ so that $A_2 \sim 0$. When the photons are continued off-shell then Eq. (4.2) must be generalized to include longitudinal-photon states, according to the usual VMD formulas (3.13). This is trivial to do.

We reemphasize the point that the spin structure of the on-shell scattering amplitudes is not affected by the mass continuation. These amplitudes thereform transform like on-shell amplitudes for any value of the ρ mass. The connection between s- and t-channel $\pi\rho\to\pi\rho$ helicity amplitudes[16]

$$f^{(s)}_{db} = \sum_{D',b'} d^1_{D'd}(\chi_d) d^1_{b'b}(\chi_b) f^{(t)}_{D'b'}, \tag{4.3}$$

is unaffected by the mass continuation. The crossing angles χ_d, χ_b are, of course, independent of any numerical change in the invariant amplitudes.

Things are very different when the photons are treated as external particles in the strong-interaction amplitude. Then one writes a crossing equation such as Eq. (4.3) for the Compton amplitudes, but with $m_\rho = 0$ so that $\chi_d = \pi$ and $\chi_b = 0$, and

$$T^{(s)}_{1,\pm 1} = T^{(t)}_{1,\mp 1}. \tag{4.4}$$

The forward Compton amplitude $T^{(s)}_{1,1}$ is therefore equal to the t-channel double-flip amplitude $T^{(t)}_{1,-1}$. In the Regge model this function has a nonsense zero at $t = 0$[21]. Therefore a fixed pole at $t = 0$ must be present in $T^{(t)}_{1,-1}$ because σ_{tot} for photoproduction does not vanish at high energy. All the unphysical longitudinal-photon amplitudes in both channels are zero.

(3) The conventional approach to the ρ-mass continuation is to assume that the ρ couples to a conserved current.[18,23,24] Then the M function (2.11) satisfies the constraint

$$p^\beta_b M_{\alpha\beta} = 0, \tag{4.5}$$

in which case the invariant amplitudes A_6, A_7 are related to physical amplitudes by the relations

$$(m_a^2 + m_b^2 - t)A_3 + (s - m_a^2 - m_b^2)A_5 + 2m_b^2 A_7 = 0, \tag{4.6}$$

$$2A_1 + (m_a^2 + m_b^2 - t)A_2 + (s - m_a^2 - m_b^2)A_4 + 2m_b^2 A_6 = 0. \tag{4.7}$$

These conditions guarantee that when $m_b \to 0$ the amplitudes $f^{(s)}_{d0}$ vanish like $O(m_b)$. The incoming ρ is being treated like a physical particle even when it is off-shell, and its helicity can only be ± 1 when its mass is zero.

When $m_b = 0$, as for photoproduction, the constraint (4.6) becomes

$$(t - m_a^2)A_3 = (s - m_a^2)A_5. \tag{4.8}$$

In the current-current model $A_5 = 0$, and therefore $A_3 = 0$. Here the model runs into conflict with SCHC, which requires that $2A_1 + A_4 \approx 0$ and $A_3 \approx A_4$. In fact, there is no way to satisfy SCHC when the initial ρ is coupled to a conserved current. This mass continuation cannot be used together with the current-current ansatz.

The conflict with SCHC can be made even more apparent if we adopt the continuation scheme of Cho and Sakarai[24] by assuming the invariant amplitudes $A_1, \ldots, A_7$ are independent of m_b. Then the two constraints (4.6) and (4.7) are really four,

$$(m_a^2 - t)A_3 + (s - m_a^2)A_5 = 0,$$
$$A_3 - A_5 + 2A_7 = 0,$$
$$2A_1 + (m_a^2 - t)A_2 + (s - m_a^2)A_4 = 0,$$
$$A_2 - A_4 + 2A_6 = 0,$$

and $A_5 = 0$ leads immediately to $A_3 = A_7 = 0$, which is incompatible with SCHC.

Models with nonzero A_5 do not have this problem. Consider the conventional Regge model with a moving Pomeron pole $\alpha_P(t) = 1 + \alpha'_P t$. The amplitude A_5 is proportional to the t-channel double-flip amplitude (see Appendix) and therefore it contains the nonsense factor

$$\alpha_P(\alpha_P - 1) = (1 + \alpha'_P t)(\alpha'_P t), \tag{4.9}$$

which vanishes at $t = 0$. Unless the residue contains a fixed singularity at $t = 0$ the amplitude A_5 must vanish there. Compatibility with SCHC implies that the fixed singularity is indeed present. Note that the current-current model corresponds to $\alpha'_P \equiv 0$, and in this case one cannot save the situation by introducing a fixed pole.

(4) The current-current model for diffractive

reactions discussed in this paper is based on the experimental facts of factorization, nearly-constant cross sections, centrality and approximate SCHC for elastic scattering, and peripherality and a strong trend toward TCHC in diffractive excitation. All of these characteristics are realized in a natural way within the model by giving the effective current two components:

(a) a central, electromagnetic component which is proportional to the isoscalar electromagnetic current (this is the Pomeron-photon analogy), and

(b) a peripheral, nonconserved component which is not identified with a known current. (One possibility would be the $I=0$ part of the neutral, vector weak current, but too little is known about this current for this identification to be verified. Nevertheless, this is an attractive possibility in view of the fact that unified theories of weak and electromagnetic interactions seem to be physically valid and useful. An analogy involving photons should thereefore, perhaps, involve Z bosons as well. In a separate article we shall discuss some implications of this extended Pomeron-photon analogy.)

Independent of the nature of the effective current, the model provides factorization and constant cross sections automatically. For pp, πp, and Kp elastic scattering the central electromagnetic component is dominant, but the peripheral component is also present and accounts for some of the finer features observed. SCHC is presumably due to the electromagnetic component (i.e., to the small value of the $I=0$ anomalous magnetic moment of the nucleon). For diffractive excitation, the electromagnetic component is suppressed in the forward direction by current conservation, and therefore the peripheral component is dominant. This component accounts for nearly all of the small-$|t|$ cross section, and therefore it is responsible for TCHC, and the slope-mass effect, and other characteristics of diffractive excitation.

Two-component models of the Pomeron have been introduced previously (see e.g. Hartley and Kane[8]). The current-current ansatz differs from these models mainly in that it is an operator formulation of the Pomeron which is applicable to any diffractive reaction (elastic, inelastic, or inclusive). The ansatz automatically provides each type of diffractive reaction with the correct qualitative features. Previous formulations[8] are not models which are applicable to all cases. Either the discussion is limited to elastic scattering, or, when more generality is sought, a prescription is given for each type of diffractive reaction separately. By introducing a two-component effective current we have gained something and lost something. What we have gained is the unification already described. But we have sacrificed Regge behavior and other pieces of the accumulated hadron phenomenology which enable accurate data fits to be made. This phenomenology can be put back in by hand. Whether this should be done, or some other course followed, is an open question.

Mention must be made of the model constructed by Abarbanel, Drell, and Gilman,[25] which is also a two-component model of the Pomeron. These authors achieved excellent fits to the pp elastic scattering data using an electromagnetic component of the Wu-Yang type plus a Reggeized Pomeron component which shrinks with increasing energy. At very high energies this Reggeized Pomeron only contributes at small $|t|$, where shrinkage still occurs, and for larger $|t|$ the electromagnetic component, which does not shrink, is dominant.

(5) The presence of both central and peripheral components in the effective current can be explained by a quark-parton model. The central component is most important when the two hadrons collide at small impact parameter rather than edge against edge. We know such collisions are often elastic, with no disturbance of the "valence" quarks. Presumably these tend to be found near the surface of the hadron. The "sea" quarks, i.e., those with a very small fraction of the hadron's momentum, are distributed in an *a priori* unknown fashion within the hadron. As viewed from one hadron, the other hadron is a disk with valence quarks mostly on the edge and sea quarks toward the center. In a collision at small impact parameter the two quark seas pass through each other. We summarize the effects of the sea-sea interaction by the central, electromagnetic component in the current. At the quark level, a pointlike sea quark in one hadron exchanges an $I=0$ vector Pomeron with a sea quark in the other hadron. The sea-sea interaction is then a product of quark distributions which can be written in current-current form. These quark distributions are essentially electromagnetic form factors. Rescattering is of course important and therefore eikonalization or some equivalent procedure should be applied to the basic current-current interaction.

A central collision may, of course, be highly inelastic. Indeed it is known that multiplicity tends to increase with decreasing impact parameter.[26] This is another indication that sea quarks rather than valence quarks, whose number is fixed, play the dominant role at small impact parameter.

(6) The peripheral component in the effective current is associated with valence quarks. These congregate near the surface of hadrons. One indication that this is the case comes from nondiffractive reactions, such as πN charge exchange, in which valence quarks participate actively. These

reactions show strong peripheral behavior. Roughly speaking, it seems that a hadron is unlikely to survive being struck near the edge, while it may well survive a blow near the center. By survive, we mean to retain all its valence quarks. These can easily be lost, as when charge is exchanged. They can also be lost in diffractive excitation, which is an interaction between valence quarks in one hadron and sea quarks in the other. (See Lubatti and Moriyasu[27] for a discussion of diffractive excitation at the quark level.) The small multiplicities observed in diffractive excitation and the prevalence of two-body threshold effects[27,28] strongly support this valence-sea interpretation of diffractive excitation. The overlap of valence and sea quarks is maximum for intermediate values of impact parameter, and therefore diffractive excitation is moderately peripheral.

In our model one of the cross terms, namely the product of the nonconserved current (at the inelastic vertex) times the conserved current (at the elastic vertex), is responsible for diffractive excitation. Since it is a product of two impact-parameter distributions, one of them peripheral and the other one broad and mainly central, this amplitude is moderately peripheral, as it should be.

(7) Double diffraction dissociation in our model is provided by the product of the nonconserved current with itself. The amplitudes will clearly be quite strongly peripheral, since they are the product of two peripheral b distributions. Experimental results on double diffraction dissociation have recently become available.[29] It is still too early to say if this type of reaction is more peripheral than single diffraction dissociation (in our model it is); however, factorization seems to hold, which is at least encouraging.

(8) Finally, we mention some other points concerning the peripheral component in our model.

(a) Since the electromagnetic component is energy-independent, we must attribute rising total cross sections to the peripheral component, which does contribute to elastic scattering in the near-forward direction. According to the data this contribution shrinks (i.e., must be Reggeized) and therefore the interpretation seems to be consistent. The peripheral component gives the entire near-forward cross section in diffractive excitation, and one knows these cross sections also rise,[30] at about the same rate as the total cross sections. Again the model is consistent.

(b) The peripheral component provides most, or all, of the forward cross section in the reactions $\gamma N \to \rho N$, $\gamma N \to \gamma N$, and $\gamma\gamma \to \gamma\gamma$ with real or virtual photons. By analogy with elastic pp scattering, one expects the near-forward cross sections for these reactions to have different characteristics from the cross sections for larger $|t|$. There seems to be some evidence for this in $\gamma p \to \rho p$.[9,12] Just as for elastic scattering and diffractive excitation, one expects that cross sections for the reactions above should rise, and at a similar rate.

(c) The cross section for inclusive electron-proton scattering

$$e + p \to e + \text{anything}$$

is proportional to the (imaginary part of) the forward Compton amplitude. For limited $|q^2|$ we have seen that, if SCHC holds, this amplitude is dominated by the peripheral component and not by the central component. Let us now consider the Bjorken limit $s \to \infty$, $|q^2| \to \infty$ with q^2/s finite. It is clear that a current-current amplitude with its fixed-energy behavior and factorized structure is very poorly-suited to provide scaling in the variable q^2/s. Therefore the current-current model cannot be continued to the scaling region from the region of limited $|q^2|$. This does not come as a surprise, for one cannot describe a highly-spaclike strongly-interacting photon as an assembly of quasi-free partons. Only for small $|q^2|$ is this possible. To put it differently, the simple VMD picture we have been using is inadequate for the deep inelastic region.

APPENDIX

The t-channel helicity amplitudes $f_{Db}^{(t)}$ calculated from the M function $M_{\alpha\beta}$ in Eq. (2.11) are

$$\begin{aligned} 4m_d m_b f_{Db}^{(t)} = {} & 4m_d m_b \epsilon_D(p_D)\cdot\epsilon_b(p_b)A_1 - \delta_{0D}\delta_{0b}T_{db}{}^2 A_2 \\ & -\delta_{0b}\, T_{db} S_{cd} d^1_{0D}(\chi_d) A_3 \\ & +(-1)^b \delta_{0D} T_{db} S_{ab} d^1_{0b}(\chi_b) A_4 \\ & +(-1)^b S_{cd} S_{ab} d^1_{0D}(\chi_d) d^1_{0b}(\chi_b) A_5 . \end{aligned}$$

Here we have used the formulas

$$\begin{aligned} & 2m_d p_b \cdot \epsilon_D(p_D) = \delta_{0D} T_{db}, \\ & 2m_d p_c \cdot \epsilon_D(p_D) = S_{cd} d^1_{0D}(\chi_d), \\ & 2m_b p_D \cdot \epsilon_b(p_b) = \delta_{0b} T_{db}, \\ & 2m_b p_A \cdot \epsilon_b(p_b) = -(-1)^b S_{ab} d^1_{0b}(\chi_b). \end{aligned}$$

The double-flip amplitude is

$$f_{1,-1}^{(t)} = (\phi/2T_{db}{}^2)A_5$$

because

$$\epsilon_1(p_D)\cdot\epsilon_{-1}(p_b) = 0.$$

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