Pomeranchuk-Exchange Contribution to Forward Photon Processes*

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Assuming pure Regge-pole behavior and the crossing relations for helicity amplitudes, we show that (1) the Pomeranchuk trajectory does not contribute to Compton scattering from spin-0 or spin-$\frac{1}{2}$ particles in the forward direction; (2) Pomeranchuk exchange does contribute to forward vector-meson photoproduction, and at high energies the production of vector mesons with helicity ±1 dominates over the production of those with zero helicity. Some interesting properties of the crossing relations are derived and used to obtain these results. We briefly discuss the difficulties associated with conclusion (1).

I. INTRODUCTION

WHEN the Pomeranchuk trajectory can be exchanged, it provides the dominant Regge-pole contribution to a scattering process at high energy. If it contributes to the forward helicity nonflip amplitudes for an elastic-scattering process, the total cross section for those incident particles approaches a constant at high energies; indeed, historically this was the reason for introducing the Pomeranchuk trajectory. When Pomeranchuk exchange does not contribute to the forward nonflip elastic amplitudes of some process, the Regge-pole model predicts that the corresponding total cross section approaches zero asymptotically as $s\rightarrow\infty$.

In the present paper we shall investigate the contribution of the Pomeranchuk trajectory to forward vector-meson photoproduction and Compton scattering. We shall assume pure Regge-pole behavior, i.e., partial-wave amplitudes with only moving poles in the complex angular-momentum plane, although in our final conclusions we shall be forced to question this assumption.

Our results follow directly from: (a) the crossing relations between helicity amplitudes,[1] (b) the Reggeization of kinematical-singularity-free helicity amplitudes (including the case of unequal masses) according to the usual prescriptions,[2–5] and (c) the assumption of a Pomeranchuk trajectory with positive signature and $\alpha_P(t=0)=1$.[6] The principal results are: (1) The Pomeranchuk exchange contribution to Compton scattering from spin-0 or -$\frac{1}{2}$ particles vanishes in the forward direction[7]; (2) Pomeranchuk exchange does contribute to forward vector-meson photoproduction and at high energies the amplitude for producing vector mesons with helicity ±1 (equal to the incident photon helicity) is a factor (s/s_0) larger than the amplitude for producing helicity 0 vector mesons.

Combining (1) and (2) we can show that for sufficiently high energies $\sigma^{\text{inel}}(\gamma N)>\sigma^{\text{tot}}(\gamma N)$; indeed our analysis of Compton scattering alone directly implies (for large s) $\sigma^{\text{el}}(\gamma N)>\sigma^{\text{tot}}(\gamma N)$. Clearly this conclusion is unacceptable and tells us that one must modify the assumption of "pure" Regge behavior. We shall not attempt to resolve this question in the present paper but only suggest some possibilities to be explored.

Section II contains a review of the relevant kinematics and notes some interesting consequences of the helicity crossing relations for massless particles or special scattering angles and high energies. In Sec. III we briefly summarize the pertinent Reggeization procedure and then in Sec. IV we use the results given in the previous sections to analyze the special cases of forward Compton scattering and vector-meson photoproduction. We discuss the interpretation of our results in Sec. V.

II. KINEMATICS AND THE CROSSING RELATIONS

We define the direct (or s) channel to be the reaction

$$a+b\rightarrow c+d;$$

the crossed (t) channel is then

$$D'+b'\rightarrow c'+A'.$$

The helicity amplitudes describing these reactions and the crossing relations between them are defined exactly as in Ref. 3. In all of our applications we choose b to be the incident photon and d to be the final photon or vector meson. The helicity crossing relations can be written[3]

$$f^s{}_{cd,ab}(s,t)=\sum_{c'A'D'b'} d_{A'a}{}^{J_a}(X_a)d_{C'c}{}^{J_c}(X_c)d_{b'b}{}^{J_d}(X_b)\times d_{D'd}{}^{J_d}(X_d)f^t{}_{c'A',D'b'}(s,t). \quad (1)$$

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[3] Ling-Lie Wang, Phys. Rev. **142**, 1187 (1966).
[4] Ling-Lie Wang, Phys. Rev. Letters **16**, 756 (1966).
[5] Ling-Lie Wang, Phys. Rev. **153**, 1664 (1967).
[6] There have been some recent suggestions that $\alpha_P(0)<1$, e.g., N. Cabibbo, J. Kokedee, L. Hurwitz, and Y. Ne'eman, Nuovo Cimento **45**, 275 (1966); but in the present work we shall always assume $\alpha_P(0)=1$.
[7] This conclusion has also been reached independently by V. D. Murr, Zh. Eksperim. i Teor. Fiz. **44**, 2173 (1963) [English transl.: Soviet Phys.—JETP **17**, 1458 (1963)], and by H. D. I. Abarbanel and S. Nussinov, Phys. Rev. **158**, 1462 (1967). The latter authors, whose work is very similar in spirit and methods to our own, have also considered the contribution of the Pomeranchuk to vector-meson photoproduction.

A. Elastic Scattering: $m_a=m_c$, $m_b=m_d=0$

For forward scattering, $\theta_s=0$ $(z_s=\cos\theta_s=1)$ and $t=0$. We have for z_t, the cosine of the t-channel scattering angle, as $s\to\infty$,

$$z_t\to\frac{s}{2p_{D'b'}p_{c'A'}},\tag{2}$$

where $p_{D'b'}$ and $p_{c'A'}$ are the momenta of the initial and final particles in the t-channel cm system. Evaluating the angles[8] in the crossing relations referring to the initial and final photons, we find $\cos X_b=+1$, $\cos X_d=-1$; hence, $X_b=0$, $X_d=\pi$ and these values hold for all s and t. Since

$$d^J{}_{\lambda\lambda'}(\theta=0)=\delta_{\lambda\lambda'}$$

and

$$d^J{}_{\lambda\lambda'}(\pi)=(-1)^{J-\lambda'}\delta_{\lambda,-\lambda'},$$

we seen that we must have

$$b'=b,\quad D'=-d,$$

where the letters now refer to the helicities of the particles. With a little more calculation we can derive the general result: Under crossing, the helicity of a photon flips if the line gets crossed [i.e., a particle (antiparticle) in s becomes an antiparticle (particle) in t], and the photon helicity remains the same if the line is uncrossed. As noted above, this result is true for all s and t and, in fact, is also independent of the values of the other masses or any relations between them.[9]

B. Inelastic Scattering: $m_a=m_c\neq0$, $m_b\neq m_d$

For unequal masses, the kinematics becomes considerably more complicated. In the forward direction, $z_s=1$, and expanding for large s with $m_a=m_c$, $m_b\neq m_d$,[10]

$$-t\xrightarrow[z_s=1,s\to\infty]{}\frac{m_a^2(m_d^2-m_b^2)^2}{s^2}+\frac{(m_a^2+m_b^2+m_c^2+m_d^2)m_a^2(m_d^2-m_b^2)^2}{s^3}+\cdots.\tag{3}$$

Because of the unequal masses, we find[11] for $z_s=1$

$$z_t=-1.\tag{4}$$

The crossing angles behave as follows[12] $(m_b=0)$:

$$\cos X_b=+1\quad\text{for all } s,\,t,\tag{5a}$$

$$\cos X_d\xrightarrow[z_s=1]{}1,\tag{5b}$$

$$\cos X_a\xrightarrow[z_s=1]{}-1,\tag{5c}$$

$$\cos X_c\xrightarrow[z_s=1]{}-1.\tag{5d}$$

These results have several interesting consequences:

(1) $X_b=0$ and hence $b'=b$ as in case A.

(2) As $s\to\infty$, $X_d\to0$ and thus $D'=d$, the vector-meson helicity does not flip. Recall that if d were a photon we would have $X_d=\pi$ and $D'=-d$. Thus the vector meson "d" does *not* behave (under crossing) like a photon as $s\to\infty$ and the crossing property for $m_d=0$ is crucially different from that when $m_d>0$ and $s\to\infty$. In the final Sec. V we shall make a more general statement on how the helicities behave under crossing for inelastic scattering.

(3) As $s\to\infty$, the dominant terms in the crossing relation have $A'=-a$, $c'=-c$,[13] i.e., the helicities of particles a and c both flip.[14] (We would find $A'=+a$, and $c'=+c$ if $m_b>m_d$ instead of $m_b=0<m_d$.)

(4) Putting all of these results together we see that for $m_a=m_c\neq0$, $m_b=0$, $m_d>0$, the crossing relation reduces to[14]

$$f^s{}_{cd,ab}\xrightarrow[z_s=1,s\to\infty]{}f^t{}_{-c\,-a,db},\tag{6}$$

where only the leading term has been retained.[12]

III. REGGEIZATION

Here we briefly summarize the results we shall need. For more details see Refs. 2–5. Define

$$\bar f^t{}_{c'A',D'b'}\equiv(1-z_t)^{-|\lambda_t-\mu_t|/2}\times(1+z_t)^{-|\lambda_t+\mu_t|/2}f^t{}_{c'A',D'b'},\tag{7}$$

where $\lambda_t=D'-b'$, $\mu_t=c'-A'$. The $\bar f^t$ are assumed to have no kinematical singularities or zeros in the s plane.[3] Other kinematical singularities in t can be analyzed according to the prescriptions of Ref. 3. It is easiest to Reggeize the so called parity-conserving

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[8] These are defined in Ref. 3.

[9] With $m_b=m_d$ and $m_a=m_c\neq0$, for forward scattering, one finds $X_a=X_c=\pi/2$ for all s, but this result will not be needed in what follows.

[10] If $m_a\neq m_c$ and $m_b\neq m_d$, one finds

$$-t\xrightarrow[z_s=1,s\to\infty]{}[(m_b^2-m_d^2)(m_a^2-m_c^2)s^{-1}]+[(m_a^2m_b^2-m_c^2m_d^2)\times(m_a^2+m_b^2-m_c^2-m_d^2)s^{-2}]+\cdots.$$

These expansions are easily derived using Appendix A of Ref. 3.

[11] This is for the case $m_a=m_c\neq0$, $m_b\neq m_d$. If both $m_a\neq m_c$ and $m_b\neq m_d$, we have $z_t=\mp1$ if $(m_d^2-m_b^2)(m_c^2-m_a^2)>0$ or <0. These results follow from the equation

$$(1-z_t^2)(T_{ac}T_{bd})^2/t=(1-z_s^2)(S_{ab}S_{cd})^2/s,$$

where

$$S_{ab}^2\equiv[s-(m_a+m_b)^2][s-(m_a-m_b)^2],\text{ etc.}$$

[12] In a conversation with the author, Dr. Ling-Lie Wang has correctly observed that for exactly forward (inelastic) scattering $|\cos X_i|=1$ for *all* s. Hence when the scattering is precisely in the forward direction, the simple helicity crossing properties we derive are true for all s and not just in the limit $s\to\infty$. However, we are especially interested in the large-s region, since it is there that we expect a Regge-pole expansion to be useful (and Pomeranchuk exchange to dominate). Experimentally, it is also more realistic to consider the limit of nearly forward scattering, rather than the single point $\theta=0$. I am grateful to Dr. Wang for her comments.

[13] This is only for the unequal-mass case, $m_b\neq m_d$. Recall for case A above $X_a=X_c=\pi/2$.

[14] For the moment we are neglecting any differences in the s dependence of the t-channel helicity amplitudes. See Sec. III.

helicity amplitudes[2]

$$\bar{f}^t{}_{c'A',D'b'} \pm \bar{f}^t{}_{-c'-A',D'b'},$$

and we shall assume that it is correct to write the Regge asymptotic behavior for such parity-conserving helicity amplitudes even in the case of unequal masses.[15] With this assumption, and employing the Reggeization procedure of Ref. 2, we have for the leading s dependence (ignoring other factors)

$$\bar{f}^t{}_{c'A',D'b'} \underset{s\to\infty}{\longrightarrow} (s/s_0)^{\alpha-\lambda_m}, \tag{8}$$

where

$$\lambda_m = \max(|\lambda_t|,|\mu_t|).$$

Inverting Eq. (7) and using Eqs. (2) and (4), this implies (in the forward direction)

$$f^t \underset{s\to\infty}{\longrightarrow} (s/s_0)^{\alpha}, \quad (m_a=m_c,\ m_b=m_d), \tag{9a}$$

$$f^t{}_{c'A',D'b'} \underset{s\to\infty}{\longrightarrow} (s/s_0)^{\alpha-\lambda_m}\ \lambda_{t'-\mu_t} \qquad (m_a=m_c,\ m_b\neq m_d), \tag{9b}$$

i.e., in the unequal mass case, $f^t \sim s^{\alpha-\lambda_m}$ for $\lambda_t=-\mu_t$ and $f^t=0$ otherwise.[16] Besides its asymptotic s dependence, we shall need the behavior of a Reggeized amplitude near a sense-nonsense value of the angular momentum.[17] The result can be stated as follows: $\bar{f}_{\mu,\lambda}$ vanishes with a factor $(\alpha-J_{sn})$ at those values of J_{sn} such that $|\mu| \le J_{sn} < |\lambda|$ or $|\lambda| \le J_{sn} < |\mu|$ and $(-1)^{J_{sn}}=-\tau$ where τ is the signature of the Regge trajectory being exchanged.[18]

IV. APPLICATIONS: FORWARD PHOTON REACTIONS

Now we are ready to consider the Pomeranchuk contribution to forward Compton scattering and vector-meson photoproduction. In the forward direction, in order to conserve angular momentum we include only those s-channel helicity amplitudes $f^s{}_{cd,ab}$, which have

$$\lambda_s \equiv (a-b) = (c-d) \equiv \mu_s \quad \text{for} \quad z_s=1. \tag{10}$$

The total cross section for the scattering of particles 1 and 2 is given by

$$\sigma^{\text{tot}} = \left[\frac{-2\,\text{Im} f^s{}_{\lambda_1\lambda_2,\lambda_1\lambda_2}(\theta_s=0)}{p_{12}\sqrt{s}}\right], \tag{11}$$

where we would sum and average over the nonflip forward amplitudes to get the spin-averaged total cross section.

The differential cross section (for specified helicities) is

$$\left(\frac{d\sigma}{dt}\right)_{ab\to cd} \frac{|f^s{}_{cd,ab}|^2}{4\pi s p_{ab}{}^2}. \tag{12}$$

A. Compton Scattering

For the moment consider $m_a=m_c$, $J_a=J_c$ but let the spin J_a be general. The forward nonflip amplitudes are $f^s{}_{a1,a1}$ where a can be any of the allowed helicities corresponding to spin J_a. Equation (1) and the photon crossing property now give

$$f^s{}_{a1,a1} = \sum_{c'A'} d^{J_a}{}_{A'a} d^{J_a}{}_{c'a} f^t{}_{c'A',-11}. \tag{13}$$

Since $|\lambda_t|=2$ and if $|c'-A'|=|\mu_t|<2$, the value $J=1$ is a sense-nonsense point and hence Pomeranchuk exchange gives

$$f^t{}_{c'A',-11} \propto [\alpha_P(t)-1], \tag{14}$$

which vanishes at $t=0$, thus causing $f^s{}_{a1,a1}$ to vanish in the forward direction. If $J_a \ge 1$ we could have $|c'-A'| \ge 2$; for this case $f^s{}_{c'A',-11}$ is a nonsense-nonsense amplitude which vanishes when $\alpha_P=1$ if the Pomeranchuk is a "sense-choosing" trajectory.[2,18] Hence with this assumption (and rigorously for $J_a=J_c<1$) the contribution of the Pomeranchuk to *any* forward non-spin-flip elastic Compton amplitude vanishes, and thus by Eq. (11) the photon scattering total cross section vanishes as $s\to\infty$.[19] Next consider other possible nonzero forward amplitudes, i.e., amplitudes with $\lambda_s=\mu_s$ but now allowing helicity flip (these do not contribute to σ^{tot}). The only new possibility is an amplitude of the form $f^s{}_{c1,a-1}$ which after crossing has $|\lambda_t|=0$ and thus has no nonsense zero at $\alpha=1$ if $|c'-A'|<2$. Since $\lambda_s=\mu_s$ requires $|c-a|=2$ for these forward helicity flip amplitudes, we can rule out this

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[15] Since z_t does not become large as $s\to\infty$, when $z_s=1$ for unequal masses, it is very important whether we assume $\bar{f}^t$ or f^t has the usual (i.e., equal-mass) asymptotic behavior for its leading term as has been argued recently for the spinless case. We believe it is reasonable to assume that it is the kinematical-singularity-free amplitudes which Reggeize, and we also argue that it is correct to preserve the factors of $(1\pm z_t)$ which express some important physics, e.g., conservation of angular momentum. Wang, in Ref. 5, has made the same assumption. See also the comments in Refs. 16 and 21 below. For work on spinless unequal-mass scattering see D. Z. Freedman and J. M. Wang, Phys. Rev. Letters **17**, 569 (1966); Phys. Rev. **153**, 1596 (1967); D. Z. Freedman, C. E. Jones, and J. M. Wang, *ibid.* **155**, 1645 (1967); R. J. Oakes, Phys. Letters **24B**, 154 (1967); L. Durand, Phys. Rev. Letters **18**, 58 (1967); R. Omnes and E. Leader (to be published); G. Domokos (to be published).

[16] If we had assumed that f^t rather than $\bar{f}^t$ has the usual asymptotic behavior for the unequal-mass case, we would have $f^t{}_{\mu_t,\lambda_t} \sim s^{\alpha}$ for all λ_t, μ_t.

[17] See Refs. 2, 4, and 5. Those $J<|\lambda_t|$ or $|\mu_t|$ are called "nonsense" values since the physical J of a state cannot be less than some component of the total spin of the state.

[18] If $(-1)^{J_{sn}}=+\tau$, the factor $(\alpha-J_{sn})$ merely removes the pole in $\bar{f}_{sn}$ leaving a (nonzero) constant for the amplitude at $\alpha=J_{sn}$. When both $J<|\lambda|$ and $J<|\mu|$, the amplitude vanishes or not depending on whether the trajectory "chooses sense" or "nonsense." See Ref. 2.

[19] For $s\to\infty$, Eq. (11) gives $\sigma^{\text{tot}} \sim (s/s_0)^{\alpha(\theta_s=0)-1}$, where the α is that of the leading contributing trajectory. Since we have shown the P does not contribute, $\alpha(0)<1$ and

$$\sigma^{\text{tot}} \underset{s\to\infty}{\longrightarrow} 0.$$

If the leading nonzero contribution comes from the P' with $\alpha(0) \cong 0.7$, then we have $\sigma^{\text{tot}} \sim (s/s_0)^{-0.3}$.

possibility for spin-0 or spin-$\frac{1}{2}$ particles. Thus $J_a=J_c=0$ or $\frac{1}{2}$ (e.g., pions or nucleons) the contribution of the Pomeranchuk trajectory to *any* s-channel helicity amplitude vanishes in the forward direction.

But for $J_a=J_c\geq 1$, the Pomeranchuk may contribute to some forward helicity flip amplitudes (when $|\lambda_c-\lambda_a|=2$). It is also possible to show that for the case $m_a\neq m_c$ (with b and d photons), the Pomeranchuk does not contribute in the forward direction no matter what spins we allow for particles a and c.

B. Photoproduction of Vector Mesons

We shall specifically consider the case of photoproduction from nucleons ($J_a=J_c=\frac{1}{2}$).

$$\gamma(b)+p(a)\rightarrow V^0(d)+p(c).$$

Of the 12 independent helicity amplitudes for this process, 3 are nonzero for $\theta_s=0$:

$$f^s{}_{\frac{1}{2}1,\frac{1}{2}1},\quad f^s{}_{-\frac{1}{2}1,-\frac{1}{2}1},\quad \text{and}\quad f^s{}_{-\frac{1}{2}0,\frac{1}{2}1}.$$

Keeping only the dominant term at high energies, Eq. (6) implies[12]

$$f^s{}_{\pm\frac{1}{2}1,\pm\frac{1}{2}1}\xrightarrow[z_s=1,s\to\infty]{} f^t{}_{\mp\frac{1}{2}\mp\frac{1}{2},11}, \tag{15a}$$

$$f^s{}_{-\frac{1}{2}0,\frac{1}{2}1}\xrightarrow[z_s=1,s\to\infty]{} f^t{}_{\frac{1}{2}-\frac{1}{2},01}. \tag{15b}$$

We first note that none of these amplitudes vanishes at $\alpha=1$ since we do not have $|\lambda_t|$ or $|\mu_t|>1$ and thus $\alpha=1$ is a sense-sense value for these amplitudes. [Also we do have $\lambda_t=-\mu_t$ as required by Eq. (9b).] Hence we conclude that Pomeranchuk exchange *does* contribute to forward vector-meson photoproduction.[20] Next we may compare the energy dependence of the amplitude for producing helicity-1 vector mesons ($|\lambda_V|=1$), Eq. (15a), with the amplitude for producing helicity-0 vector mesons ($|\lambda_V|=0$), Eq. (15b). Using Eq. (9b) we have (neglecting constants)

$$\frac{f^s{}_{|\lambda_V|=1}}{f^s{}_{|\lambda_V|=0}}\xrightarrow[z_s=1,s\to\infty]{}\frac{(s/s_0)^\alpha}{(s/s_0)^{\alpha-1}}=(s/s_0), \tag{16}$$

and thus

$$\left[\frac{d\sigma}{dt}(\theta_s=0)\right]_{|\lambda_V|=1}\Big/\left[\frac{d\sigma}{dt}(\theta_s=0)\right]_{|\lambda_V|=0}\sim\left(\frac{s}{s_0}\right)^2, \tag{17}$$

i.e., the production of helicity$=1$ vector mesons dominates as $s\rightarrow\infty$.[21] The energy dependence predicted by Eq. (16) or (17) can be tested experimentally, but present data are probably too low in energy to justify comparison with this asymptotic Regge behavior.[22]

V. DISCUSSION AND SUMMARY

In this section we wish to draw some consequences from our results and raise some questions for further study. In addition, we shall summarize the crossing behavior for particle helicities, since we find it quite interesting.

(1) For a massless particle it is easy to show from the crossing relations[1] that the helicity will flip or remain the same, depending on whether the particle line gets reversed or not. This statement is true for all s and t and whatever the values of the other masses.

For unequal-mass scattering (i.e., $a\neq c$ and/or $b\neq d$) one can derive the following general result, where in all cases it is understood we are referring to the behavior of the dominant term, in the limit[12] $\theta_s=0$, $s\rightarrow\infty$: (a) For the particles at a vertex (viz. a and c, or b and d) connecting *unequal* nonzero masses, both helicities flip (do not flip) if the mass of the particle whose line is reversed under crossing is less (greater) than the mass of the uncrossed particle. (b) At a vertex connecting *equal* nonzero masses, both helicities flip (do not flip) if (for the unequal mass pair at the other vertex) the mass of the particle which gets crossed is greater (less) than the mass of the uncrossed particle. In the present paper we shall not elaborate upon the implications of this general result for unequal mass scattering. It would also clearly be desirable to have a more transparent derivation of these simple crossing properties.

(2) By integrating the Pomeranchuk contribution to the elastic differential cross section for (nuclear or pion) Compton scattering we find $\sigma^{\text{elastic}}\sim 1/[\ln(s/s_0)]^2$. Since we found earlier $\sigma^{\text{tot}}\sim(s/s_0)^{-0.3}$, assuming the P' to give the leading nonzero contribution to σ^{tot},[19] we are led to conclude that for s sufficiently large, $\sigma^{\text{el}}>\sigma^{\text{tot}}$.[23] In order to avoid this unacceptable conclusion, it thus seems we are forced to alter our assumed "pure" Regge behavior in a manner which either allows the Pomeranchuk pole to contribute directly to forward Compton scattering or provides some other "large" contribution to $f_{\text{nonflip}}{}^{\text{elastic}}(\theta=0)$ and hence to σ^{tot}. We shall here not resolve this problem but only mention a few of the many possibilities which might be explored: (a) Perhaps the simplest alternative is to allow a fixed pole at $J=1$ in the Regge amplitude, as has been suggested recently for (isovector) photon-pion scattering.[24]

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[20] The difference between Eqs. (15a) and Eqs. (13) lies in the differing crossing properties for photons and vector mesons. See Eq. (5b) and comment (2) which follows it.

[21] If we had assumed that f^t and not $\bar{f}^t$ Reggeize for unequal masses we would find (see Ref. 16 above) both $f^s{}_{|d|=1}\sim(s/s_0)^\alpha$ and $f^s{}_{|d|=0}\sim(s/s_0)^\alpha$ and thus the ratio approaching a constant as $s\rightarrow\infty$.

[22] Present data (on $\gamma p\rightarrow\rho^0 p$) lie in the range $s\lesssim 12$ BeV^2. Experiments at Stanford Linear Accelerator Center may provide data up to $s\approx 40$ BeV^2 in a year or so.

[23] We also can show that the Pomeranchuk contribution to photoproduction gives $\sigma^{\text{inelastic}}\sim 1/\ln(s/s_0)$ and hence for $s\rightarrow\infty$ we have $\sigma^{\text{inelastic}}>\sigma^{\text{elastic}}>\sigma^{\text{total}}$.

[24] J. B. Bronzan, I. S. Gerstein, B. W. Lee, and F. E. Low, Phys. Rev. Letters **18**, 32 (1967); V. Singh, *ibid.* **18**, 36 (1967). Also see the comments of V. D. Murr, Zh. Eksperim. i Teor. Fiz. **44**, 2173 (1963) [English transl.: Soviet Physics—JETP **17**, 1458 (1963)].

Note that we do not thereby violate unitarity, since it is a *fixed* pole only to lowest order in the photon-nucleon coupling and presumably becomes a moving pole (or essential singularity) in higher orders. (b) The existence of moving cuts in the angular-momentum plane (in this case generated by multiple Pomeranchuk exchange) could provide a contribution to $f_{ab,ab}(\theta=0)$, making σ^{tot} at least comparable to σ^{el} as $s\to\infty$. (c) Another possibility is some "conspiracy" of moving poles which are related at $t=0$ in such a way as to allow the Pomeranchuk to contribute.[25] (d) If $\alpha_P(0)<1$, we avoid the vanishing of the forward nonflip elastic amplitudes.

(3) Finally, it should be noted that the prediction given by Eq. (16) is essentially independent of the difficulty discussed above in (2) and will provide a further experimental test of the Regge-pole hypothesis.

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[25] Such a mechanism has been suggested in NN scattering. For the present case see V. D. Mur, Zh. Eksperim. i Teor. Fiz. **45**, 1051 (1963) [English transl.: Soviet Phys.—JETP **18**, 727 (1964)].

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Boson Masses. II

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The baryon-antibaryon model is employed to attempt a complete empirical assignment for the boson nonets with orbitals $^1S_0$, $^3S_1$, $^1P_1$, $^3P_0$, $^3P_1$, $^3P_2$. Only one state seems to be missing experimentally, and a search area is indicated by a simple degeneracy in the model. Some substantial level shifts are proposed to result from open channels for multiboson decay; with these can be associated the possibility of reduced ω-φ mixing. The octet SU_3 mass formula appears to be valid for only the $^1S_0$ nonet; a generalized formula reflects major R_7 and minor G_2 mixing with the basic SU_3. For K- and η-type mesons the triplet orbitals display strong tensor as well as spin-orbit splitting. As a consequence the $K^*(1400)$ should be a mixture of $2^+(^3P_2)$ and $1^-(^3D_1)$ resonances, with respective dominant modes $K\pi\pi$ and $K\pi$. Some remarks are added about D states and the validity of A parity.

I. INTRODUCTION AND SUMMARY

RECENTLY augmented data[1] allow us to extend and improve earlier considerations[2] on the baryon-antibaryon model for bosons. We first note that the model itself implies a likely degeneracy for bosons of a given ^3L_J nonet: between the charge singlet **1** and the $I=1$ member of the charge octet **8**. This is observed in the approximate degeneracies of ρ and ω, A_2 and f, and leads to the prediction of other resonances—in particular, a **1** state of $^3P_1$ at $\sim$1090 MeV to accompany the A_1.

Measured mass differences in these degenerate states can be semiquantitatively interpreted as level shifts due to open channels. This interpretation leads away from the idea that deviations from the SU_3 octet mass formula can be attributed to strong mixing of the ω-φ type. Accordingly, we attempt to analyze the four probably established nonets ^1S, ^3S, ^1P, ^3P in terms of clashing symmetries: $SU_3+gR_7+fG_2$. The empirical result is that f is very small, while variation in g is responsible for significant changes in pattern.

Comparison of the $^3P_{0,1,2}$ nonets suggest strong spin-orbit coupling throughout; in general there is also a tensor-type force, but this surprisingly vanishes just for the degenerate states **8** $I=1$ and **1**. Knowledge of this structure is sufficient to identify some D states among the fragmentary data at higher energies and to predict regions for other D states.

In conclusion are added a few remarks about A parity. Its validity is enhanced by our abandonment of ω-φ mixing. The chief experimental difficulty is then the comparable decay rates for $(K\pi)$ and $(K\pi\pi)$ modes of the $K^*(1400)$. We suggest that the reported state is in fact an accidental conjunction of $^3P_2$ and $^3D_1$ resonances induced by a strong tensor addition to spin-orbit splitting.

II. DEGENERACY ARGUMENT

The crux of the baryon-antibaryon model is avoidance of parastatistics in favor of ordinary Fermi sta-

[1] G. Goldhaber and R. H. Dalitz in, *Proceedings of the Thirteenth International Conference on High Energy Physics, Berkeley, 1966* (University of California Press, Berkeley, California, 1967).

[2] R. W. King and D. C. Peaslee, Phys. Rev. **143**, 1321 (1966). Equations (4) and (4′) of this reference are now seen to hold only for S orbitals. Exploration of D and F resonances will be needed to settle the question of even-odd alternation.