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Parity Rule and Helicity Conservation in ρ Photoproduction

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The parity rule for natural- and unnatural-parity exchange is proved on the basis of the L-S scheme. The parity rule is shown to be a general feature of t-channel exchange processes. Experimental consequences of the parity rule are discussed. The s-channel helicity conservation in ρ^0 photoproduction is considered on the basis of the Regge-pole model with L-S coupling and general arguments from the complex-angular-momentum theory. It is found that the s-channel helicity conservation in ρ^0 photoproduction can be described by the coupling $L=\alpha-2$, $S=2$ with a multiplicative fixed pole at the nonsense wrong-signature point $\alpha=1$.

I. INTRODUCTION

One of the difficulties encountered by the Regge-pole theory[1-4] is the description of the coupling between the exchanged Regge pole and the external particles with spin. Considerable progress has been made, among other things, by the introduction of the concept of parity-conserving amplitude[5] and the investigation of its kinematical singularities.[6]

From a phenomenological analysis of resonance-production data,[7] the author has earlier proposed a model[8] to describe the Regge-pole couplings, which is based on the L-S scheme. This model has been applied to simple processes.[9] One cannot hope to be able to explain detailed aspects, such as polarization, from such a simple model. For the explanation of polarization, additional correction terms to the Regge-pole contribution will be required.[10] As a first approximation, nevertheless, this model is still of interest, since the residues are characterized by the total spins of the external particles as a manifestation of the residue dependence on the external spins and quantum numbers. This might be useful, since one of the sources of ambiguities in Regge-pole models is the parametrization of the residues.

The L-S scheme itself has been shown to be useful in the study of the kinematical factors and the threshold condition by Jackson and Hite.[11]

Recently, more data suitable for the study of these couplings became available. These data are the data summarized in the Morrison empirical rule[12] and the data on the spin-density matrices for ρ^0 photoproduction,[13] which indicate s-channel helicity conservation.[14,15]

The Morrison rule[12] gives the condition for the appearance of the diffraction scattering stated in terms of spins and parities of the external particles involved in the vertices of the corresponding t-channel exchange process. The rule is stated in Sec. III. This rule is an attempt to classify the dynamics of the diffraction phenomena. The rule as it is stated suggests its intimate relationship with the nature of the coupling at the t-channel vertices. Leader[12,16] has shown that the rule for a spin-0-induced reaction is implied by the kinematics of the reaction. Considering that this rule has so far not been established for diffraction processes induced by particles with nonzero spin, it is possible that this rule has its origin in the kinematics alone. On the other hand, the s-channel helicity conservation might be related to the dynamics of the diffraction phenomenon, e.g., due to the nature of the above-mentioned coupling. In this case the Morrison rule would be valid only for a spin-0-induced reaction and the approximate s-channel helicity conservation is a characteristic of the diffraction phenomenon. These considerations are the motivation for the present investigation.[17,18]

Note that the evidence for s-channel helicity conservation in ρ photoproduction comes from the measurement of the spin-density matrix of the produced ρ^0.

The density matrices in general, roughly speaking, depend on ratios of bilinear combinations of the helicity amplitudes. Unlike the polarization, the density matrices are less sensitive to small phase differences between the components of the helicity amplitude.[9] If only one particular exchange is dominant, the main features of the density matrix are determined by the *ratios* between the components of this dominant exchanged amplitude. In ordinary Regge-pole exchange, these amplitude ratios reduce to residue ratios, which are characteristic of the coupling, and the trajectory dependence cancels out. Consequently, the spin-density matrix becomes less sensitive to the details of the trajectory at least in a small-momentum-transfer t region, like the diffraction region[13] ($|t|<0.4$ GeV^2).

Experiments show that natural-parity exchange is dominant in ρ^0 photoproduction.[13,19] This dominant exchange is the Pomeranchon exchange. Presumably there is also a contribution from Regge cuts.[20] If these cuts are important, then the above arguments have to be modified because the amplitude ratio does not depend only on the residue of the pole but also on the relative magnitude of the pole and the strength of the cut. It will be assumed that the shielding cuts either are negligible or do not seriously affect the description of the density matrices in terms of the Regge pole only.

In connection with the s-channel helicity conservation, two problems have been raised. First, it is possible that s-channel helicity conservation alone is sufficient to exclude the t-channel exchange description. It is shown that a certain t-channel coupling picture has been ruled out by s-channel helicity conservation in ρ photoproduction.[21] Reference 14 has shown that s-channel helicity conservation is compatible with a t-channel exchange description with factorizable couplings in the processes $\pi\pi\to\pi\pi$, $\pi N\to\pi N$ and $NN\to NN$ and $\pi\pi\to\pi A_1$, $\pi\pi\to A_1A_1$. Moreover, it is conjectured[14] that it is possible to describe the diffraction phenomenon with a t-channel exchange process with factorizable couplings in a special way to result in s-channel helicity conservation. If one wishes to maintain the t-channel picture, the second question which remains is then the dynamical origin for the special coupling.

Preliminary results of the present investigation have been reported.[22] It is found that the Morrison rule for spin-0-induced reactions can be shown from the kinematics alone, on the basis of the L-S scheme. In the general case a modification of the rule, referred to as "the parity rule" for natural-parity exchange, based on the same scheme has been proposed. Concerning the s-channel helicity conservation in ρ photoproduction, it is found that the t-channel description for this process is compatible provided the coupling chooses a special form expressed in terms of the residue ratios.

In this paper we extend the parity rule to unnatural-parity exchange. From the derivation of the parity rule it is suggestive that the parity rule is a *general feature of the t-channel exchange processes*. As to the s-channel helicity conservation, it is shown that exactly the above-mentioned constraint on the residue ratios can be naturally satisfied within the model with the L-S scheme and the general framework of complex angular momentum.[2-4] It is found that the data require a multiplicative fixed pole[23-25] at the wrong-signature point $\alpha=1$. The particular coupling in our model, which can accommodate this multiplicative fixed pole, has the residue ratios required for the s-channel helicity conservation. This result is compatible with the assumption stated above concerning the effect of cuts on the density matrices.[26]

In Sec. II we define the notation and review the essentials of the L-S scheme. In Sec. III the parity rule for the natural and unnatural parity is derived. Some illustrative examples are discussed. In Sec. IV the structure of t-channel helicity amplitudes within the L-S scheme is presented. In Sec. V we derive the relation between the ρ^0 spin-density matrices and the t-channel amplitudes in ρ^0 photoproduction by linearly polarized γ's. The expression for the spin-density matrices due to Regge exchange with a general coupling is also given in this section. In Sec. VI the spin-density matrices,

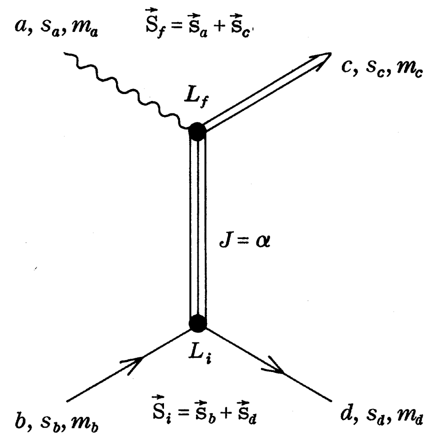

FIG. 1. The s-channel reaction is $a+b\to c+d$; the t-channel reaction is $\bar{d}+b\to c+\bar{a}$; m_x = the mass of x; s_x = the spin of x; ξ_x = the intrinsic parity of x. $L_{i(f)}$ = the orbital angular momentum in the initial (final) state t-channel reaction. $S_{i(f)}$ = the total spin of the external particles at initial (final) state t-channel vertices. In photoproduction the convention used is such that the t-channel process is $\bar{N}+N\to\rho^0+\bar{\gamma}$; μ = total helicity of the $\bar{N}N$ system and λ = total helicity of the $\rho^0,\bar{\gamma}$ system.

according to the s-channel helicity conservation, are expressed in the Gottfried-Jackson frame and the necessary coupling to produce s-channel helicity conservation is identified. The conclusions are summarized in Sec. VIII.

II. THE L-S COUPLING SCHEME IN THE REGGE-POLE MODEL

We are concerned with the reaction

$$a+b\rightarrow c+d. \tag{1}$$

The physical process (1) occurs in the s channel and the t channel will be defined by the process

$$\overline{d}+b\rightarrow c+\overline{a}. \tag{2}$$

The helicity of the particle x will be denoted by x, its mass by m_x, and its spin by s_x, where $x=a,\ b,\ c,\ d$. The helicity amplitude in the t channel will be denoted by $f_{c\overline{a},\overline{d}b}(s,t)$ or simply $f_{\lambda\mu}(s,t)$, where $\lambda=c-\overline{a}$ and $\mu=\overline{d}-b$; the s-channel amplitude will be denoted by $f^s_{cd,ab}$. For convenience, these notations are shown in Fig. 1. For the application to ρ^0 photoproduction it will be helpful to remember that λ is the helicity of the boson system $(\rho,\overline{\gamma})$ and μ is that of the $\overline{N}N$ system. The kinematical variables s and t are defined as usual: $s=(p_a+p_b)^2$ and $t=(p_a-p_c)^2$. The partial-wave expansion for $f_{\lambda\mu}$ is given by

$$f_{\lambda\mu}(s,t)=\sum(J+\tfrac{1}{2})f^J_{\lambda\mu}d^J_{\mu\lambda}(z_t), \tag{3}$$

where $z_t=\cos\theta_t$ and $\theta_t\equiv$production angle in the t channel. The general expression for z_t is

$$z_t=\frac{1}{T_{\overline{a}c}T_{b\overline{d}}}[2st+t^2-t\sum m^2+(m_b{}^2-m_d{}^2)(m_c{}^2-m_a{}^2)], \tag{4}$$

where

$$T_{ij}{}^2=[t-(m_i+m_j)^2][t-(m_i-m_j)^2],\quad i,j=\overline{a}c,\ b\overline{d}.$$

The total spins in the t-channel will be denoted by $\vec{S}_f=\vec{s}_{\overline{a}}+\vec{s}_c$ and $\vec{S}_i=\vec{s}_{\overline{d}}+\vec{s}_b$ and the addition of the spins follows the usual rule of addition of angular momenta (see Fig. 1). In the t channel, the Regge pole will be coupled to the particles in the initial (final) state with the orbital angular momenta $L_{i(f)}$ such that $\vec{J}=\vec{L}_{i(f)}+\vec{S}_{i(f)}$. Note that the mismatch between the total angular momenta J and the orbital angular momenta L, viz., $J-L$, is always an integer. The L-S scheme can be introduced in the partial-wave expansion (3); this leads to the relation

$$f^J_{c\overline{a},\overline{d}b}=\sum_{L,S}\langle J\lambda|L_fS_f\overline{a}c\rangle\langle L_fS_f|f^S|L_iS_i\rangle\langle L_iS_ib\overline{d}|J\mu\rangle, \tag{5}$$

where the coupling coefficient is given by

$$\langle J\lambda|L_fS_f\overline{a}c\rangle=\left(\frac{2L_f+1}{2J+1}\right)^{1/2}\langle L_fS_f0\lambda|J\lambda\rangle\langle s_{\overline{a}}s_c;\overline{a}c|S_f\lambda\rangle. \tag{6}$$

The general expression for the Regge-pole contribution is given by[8,11]

$$f_{\lambda\mu}=S_\pm(\alpha)K^\pm_{\lambda\mu}(t)R^\alpha_{\lambda\mu}(t)d^\alpha_{\lambda\mu}(z)\left(\frac{T_{\overline{a}c}T_{b\overline{d}}}{2s_0}\right)^{\alpha-M}, \tag{7}$$

where $M=\max(|\lambda|,|\mu|)$ and $S_\pm(\alpha)=[1\pm e^{-i\pi\alpha(t)}]/[2\sin\pi\alpha(t)]$ is the signature factor; $K^\pm_{\lambda\mu}(t)$ is the kinematic factor as derived from the crossing properties and $R^\alpha_{\lambda\mu}(t)$ is the Regge residue of the amplitude $f_{\lambda\mu}$. There are certain ambiguities in choosing the kinematic factors. We shall always choose the evasive solution,[16] which is also compatible with the L-S scheme.[11]

If the Sommerfeld-Watson transform is applied to the partial-wave expansion, with the L-S scheme, then we obtain the residue parametrization according to L_i, L_f, S_i, S_f. The relation between the residue $R_{\lambda\mu}$ and the residues in the L-S scheme is given by

$$R^\alpha_{\lambda\mu}(t)=\sum_{L,S}\langle\alpha\lambda|L_fS_f\overline{a}c\rangle\langle L_fS_f|\beta^\alpha|L_iS_i\rangle\langle L_iS_i,\ b\overline{d}|\alpha\mu\rangle. \tag{8}$$

For brevity we shall denote $\langle L_fS_f|\beta^\alpha|L_iS_i\rangle$ by $\beta^{L_fS_f,L_iS_i}(\alpha)$. We refer to Ref. 8 for the phenomenological arguments for introducing the L-S scheme.

The explicit expression for $d^\alpha_{\lambda\mu}(z)$ for large z is given by[27]

$$d^\alpha_{\lambda\mu}(z)\sim e^{i\eta(\pi/2)|\lambda-\mu|}\phi_{\lambda\mu}(\alpha)[\tfrac{1}{2}(1+z)]^{|\lambda+\mu|/2}[\tfrac{1}{2}(1-z)]^{|\lambda-\mu|/2}(\tfrac{1}{2}z)^{\alpha-M}[1+O(1/z^2)], \tag{9}$$

where $\eta = \operatorname{sgn} \operatorname{Im} z$, and the normalization factor $\phi_{\lambda\mu}(\alpha)$ is given by

$$\phi_{\lambda\mu}(\alpha) = \frac{(2\alpha+1)!}{[(\alpha+M)!(\alpha-M)!(\alpha+N)!(\alpha-N)!]^{1/2}} , \tag{10}$$

where $M = \max(|\lambda|, |\mu|)$ and $N = \min(|\lambda|, |\mu|)$. For noninteger x, then, $x! = \Gamma(x+1)$.

III. THE PARITY RULES

In a recent paper[22] the parity rule for natural-parity exchange and its role in the diffraction scattering has been reported. In this section it is generalized to the case of unnatural-parity exchange. The parity rule can be stated as follows. In the reaction (1), the *natural-parity* exchange will be coupled to the t-channel helicity amplitudes $f_{0,\mu}(s, t)$ for *all* s and t, *only if* the parities $\xi_{\bar{a}}$ and ξ_c satisfy the relation

$$\xi_{\bar{a}}\xi_c = (-1)^{S_f}, \tag{11}$$

where $\vec{S}_f = \vec{s}_{\bar{a}} + \vec{s}_c$. The corresponding rule for the $b\bar{d}$ vertex can be obtained by replacing $\bar{a}$ by $\bar{d}$, c by b, and S_f by S_i. The reason for considering only the $f^t_{0\mu}$ amplitude has been discussed in Ref. 22.

In order to prove the parity rule for the natural-parity exchange, we need Eq. (8) and substitute $\lambda = 0$. Then

$$R^{\alpha}_{0\mu}(t) \propto \langle \alpha 0 | L_f S_f \bar{a} c \rangle \propto \langle L_f S_f; 00 | \alpha 0 \rangle .$$

Since $P = \xi_{\bar{a}}\xi_c(-1)^{-L}$, then it follows that $\eta = +1$ is equivalent to

$$\eta = \xi_{\bar{a}}\xi_c(-1)^{J-L} = +1. \tag{12}$$

Note that the angular-momentum mismatch $J - L$ is an integer. The Clebsch-Gordan coefficient

$$\langle L_f S_f; 00 | \alpha 0 \rangle \propto \begin{pmatrix} L_f & S_f & \alpha \\ 0 & 0 & 0 \end{pmatrix} \neq 0 \quad \text{if } (-1)^{J-L+S} = +1. \tag{13}$$

From Eqs. (12) and (13) one then obtains Eq. (11).

For the *unnatural parity,* condition (12) becomes

$$\eta = \xi_{\bar{a}}\xi_c(-1)^{J-L} = -1.$$

Together with (13), the condition for the coupling of unnatural parity to $f_{0\mu}$ becomes

$$\xi_{\bar{a}}\xi_c = -(-1)^{S_f}. \tag{14}$$

Let us consider some illustrative examples. The reaction $PN \to P'N$, where P, P' is a 0^- meson and N is the nucleon, satisfies (11) and does not satisfy (14). Therefore, only natural parity can be exchanged in this reaction. Similarly, according to these equations, for $PN \to SN$,[28] where S is a 0^+ meson, only unnatural parity can be exchanged in the t channel. These well-known results can also easily be obtained using the parity-conserving amplitude.[16]

Let us consider the application of the parity rules to the ρ^0 photoproduction $\gamma p \to \rho^0 p$. Possible total angular momenta in the boson vertex are $S_f = 0, 1, 2$. The total spin of the external particles in the coupling to the exchanged object depends on the C and P quantum numbers. Since for the neutral system $\bar{\gamma}\rho$, $C = (-1)^{L_f + S_f}$ and $P = \xi_\gamma \xi_\rho (-1)^{L_f}$, then $CP = (-1)^{S_f}\xi_\gamma\xi_\rho$. If the exchanged particle has $CP = +1$, then it is coupled to $\bar{\gamma}\rho^0$ with total spin $S_f = 0, 2$ and those with quantum number $CP = -1$ are coupled with $S_f = 1$. Then it follows from Eqs. (11) and (14) that the natural-parity exchange with $CP = +1$ will be coupled to $f_{0\mu}$ and the unnatural-parity exchange with $CP = +1$ (such as A_1) will be decoupled from the amplitude $f_{0\mu}$. On the other hand, unnatural-parity exchange with $CP = -1$ (e.g., π, η, B) according to Eq. (14) will be coupled to $f_{0\lambda}$. In Regge-pole exchange, this unnatural-parity contribution is suppressed by a factor $O(\sqrt{t})$.[11,29]

In the $\bar{N}N$ vertex, it can easily be shown that exchange of $CP = +1$ will be coupled to the $\bar{N}N$ system in the triplet state and exchange of $CP = -1$ will be coupled to $\bar{N}N$ in the singlet state.[8] According to (11), the natural parity will be coupled to $f_{\lambda 0}$, and according to (14), the unnatural parity with $CP = +1$ will be decoupled from $f_{\lambda 0}$. The exchange of unnatural parity with $CP = -1$ is only coupled to $f_{\lambda 0}$ due to $S_i = 0$. From these properties it follows that in any reaction $aN \to cN$ there can be no interference between unnatural-parity-exchange amplitudes with $CP = +1$ and $CP = -1$.

The possible connection between the parity rule and the Morrison rule has been discussed in Ref. 22. The Morrison rule can be stated as follows: For the diffraction peak to be observed it is necessary that (1) Pomeranchon exchange be allowed, e.g., by G-parity conservation, (2) parities and spins of the particles involved in the reaction on each vertex satisfy the relation

$$\xi_a \xi_c = (-1)^{|S_a - S_c|} .$$

This rule is an attempt to classify the dynamics of the diffraction scattering.

For spin-0-induced reactions, this rule seems to be established. However, its validity for diffraction scattering induced by particles with nonzero spin still requires experimental support. For $S_a = 0$, the Morrison rule (2) is equivalent to the parity rule stated above. This observation raises the possibility that the Morrison rule might be equivalent to the parity rule, which arises from the kinematics alone. It is possible to distinguish these two alternatives by looking at the reactions, which the Morrison rule does not allow, but which is allowed by the parity rule. An example of this reaction is the production of $N^*(1535)$ with $J^P=\frac{1}{2}^-$. According to the Morrison rule this isobar cannot be produced diffractively, while according to the parity rule this production can be diffractive. The possibility of diffractive $N^*(1535)$ production is not yet excluded by recent experiments on $pp \to pN^*(1535)$.[30,31] Another example is the B^0 photoproduction. Assuming $J^P=1^+$ and $C_n=-1$ for B^0, then Pomeranchon exchange is allowed. According to the Morrison rule (2), B^0 photoproduction will not proceed diffractively, while by the parity rule, diffractive B^0 photoproduction is allowed. The parity rule requires that B^0 photoproduction proceeds with $S_f=1$.

The above-mentioned examples show that the parity rule is a general characteristic of t-channel exchange processes. The fact that diffraction scattering is compatible with the parity rule indicates that the diffraction scattering is also a t-channel exchange process.

IV. THE STRUCTURE OF THE HELICITY AMPLITUDE IN THE L-S SCHEME

For simplicity, in this section we shall restrict ourselves to the structure of the helicity amplitude in $\gamma p \to \rho^0 p$ for natural-parity exchange. The kinematical factors for this reaction are $K^+_{\lambda\mu}(t) \propto O(1)$ for $\mu=0$ and $K^+_{\lambda\mu}(t)=O(\sqrt{t})$ for $\mu\neq 0$,[29] and therefore the dominant amplitudes near the forward direction are $f_{\lambda 0}$ ($\lambda=0, 1, 2$). Using Eqs. (7) and (9), the Regge-pole contribution is given by

$$f_{\lambda 0}=S(\alpha)\beta_{\lambda 0}[\tfrac{1}{2}(1+z)]^{|\lambda+\mu|/2}[\tfrac{1}{2}(1-z)]^{|\lambda-\mu|/2}(s/s_0)^{\alpha(t)-M}, \tag{15}$$

where $\beta_{\lambda 0}=R_{\lambda 0}\phi_{\lambda 0}$ and $\phi_{\lambda 0}$ is given by Eq. (10). In the following the problem is to find the helicity dependence of the residue $\beta_{\lambda 0}$ for a particular coupling. We have shown in Sec. III that for natural parity the relevant couplings are $S=0$, $L=J$ and $S=2$, $L=J, J\pm 2$. When there is mixing between various couplings, we have the relation

$$\beta_{\lambda 0}=\sum_{L,S}\beta^{LS}_{\lambda 0}, \tag{16}$$

according to Eq. (8). Explicit calculation of the relevant Clebsch-Gordan coefficients[11] gives the following results. The orbital part of the coupling coefficient $\langle L_f S_f 0\lambda \mid \alpha\lambda\rangle$ in Eq. (6) is

$$\langle \alpha-2, 2; 0\lambda \mid \alpha\lambda\rangle=\left[\frac{(\alpha+\lambda)!(\alpha-\lambda)!}{(2+\lambda)!(2-\lambda)!}\right]^{1/2} a_-(\alpha), \tag{17}$$

$$\langle \alpha, 2; 0\lambda \mid \alpha\lambda\rangle=+[2(\alpha+1)-\lambda^2(2\alpha-1)]\left[\frac{(\alpha+\lambda)!(\alpha-\lambda)!}{(2+\lambda)!(2-\lambda)!}\right]^{1/2} c(\alpha), \tag{18}$$

$$\langle \alpha+2, 2; 0\lambda \mid \alpha\lambda\rangle=\frac{(-1)^\lambda a_+(\alpha)}{[(\alpha+\lambda)!(\alpha-\lambda)!(2+\lambda)!(2-\lambda)!]^{1/2}}, \tag{19}$$

and

$$\langle \alpha, 0; 0\lambda \mid \alpha\lambda\rangle=1, \tag{20}$$

where

$$a_-(\alpha)=\frac{1}{(\alpha-2)!}\left[\frac{(2\alpha-4)!4!}{(2\alpha)!}\right]^{1/2}, \tag{21}$$

$$c_-(\alpha)=-\frac{2}{(\alpha-1)!}\left[\frac{(2\alpha+1)(2\alpha-2)!}{(2\alpha+3)!}\right]^{1/2}, \tag{22}$$

$$a_+(\alpha)=(\alpha+2)!\left[\frac{(2\alpha+1)!4!}{(2\alpha+5)!}\right]^{1/2}. \tag{23}$$

The relevant spin part $\langle S_a S_c; \bar{a}c \mid S_f\lambda\rangle$ of the coupling coefficient in Eq. (6) is given by

$$\langle 1, 1; \bar{a}c \mid 2\lambda\rangle=\frac{1}{\sqrt{6}}\left[\frac{(2+\lambda)!(2-\lambda)!}{(1+\bar{a})!(1-\bar{a})!(1+c)!(1-c)!}\right]^{1/2} \tag{24}$$

and

$$\langle 1, 1; \bar{a}c \mid 0, 0\rangle=1/\sqrt{3}. \tag{25}$$

Using Eq. (8), we obtain the following results for the residues $\beta_{\lambda 0}^{L_f S_f}$:

$$\beta_{\lambda 0}^{\alpha-2,2}=\beta^{\alpha-2,2}\frac{A_-(\alpha)}{[(1+\bar{a})!(1-\bar{a})!(1+c)!(1-c)!]^{1/2}}, \tag{26}$$

$$\beta_{\lambda 0}^{\alpha 2}=\beta^{\alpha 2}\frac{[2(\alpha+1)-\lambda^2(2\alpha-1)]C(\alpha)}{[(1+\bar{a})!(1-\bar{a})!(1+c)!(1-c)!]^{1/2}}, \tag{27}$$

$$\beta_{\lambda 0}^{\alpha+2,2}=\beta^{\alpha+2,2}\frac{A_+(\alpha)}{[(1+\bar{a})!(1-\bar{a})!(1+c)!(1-c)!]^{1/2}(\alpha+\lambda)!(\alpha-\lambda)!}, \tag{28}$$

where

$$A_-(\alpha)=\frac{1}{\sqrt{6}}\left[\frac{(2\alpha-3)!4!}{(2\alpha+1)!}\right]^{1/2}\frac{(2\alpha+1)!}{\alpha!(\alpha-2)!}, \tag{29}$$

$$C(\alpha)=-\frac{2}{\sqrt{6}}\,\frac{(2\alpha+1)!}{\alpha!(\alpha-1)!}\left[\frac{(2\alpha+1)(2\alpha-2)!}{(2\alpha+3)!}\right]^{1/2}, \tag{30}$$

$$A_+(\alpha)=\frac{(\alpha+2)!(2\alpha+1)!}{\alpha!}\left[\frac{(2\alpha)!4!}{(2\alpha+4)!}\right]^{1/2}, \tag{31}$$

$$\beta_{\lambda 0}^{\alpha 0}=\beta^{\alpha 0}\delta_{0\lambda}. \tag{32}$$

As is well known,[23–25] the helicity amplitudes have fixed poles at nonsense wrong-signature points. These poles are compatible with the continued unitarity conditions[23,25] in the t channel, because of the presence of the shielding cuts. As far as the high-energy behavior in the s channel is concerned, the relevant question is whether a fixed pole and a Regge pole appear in additive or multiplicative form. In the first case,[25] the fixed pole does not influence the high-energy behavior, but in the second case,[4] it gives rise to a factor $[\alpha(t)-1]^{-1}$ in the Regge-pole residue and generally prevents the decoupling of the trajectory at the point $\alpha(t)=1$. The decision as to whether a fixed pole and a Regge pole appear in additive or multiplicative form can only come from the dynamics of the model.[4] Fixed poles at wrong-signature points, multiplicative or additive, require shielding cuts[4] which are moving branch points $\alpha_c(t)$ in the complex angular momentum, with $\alpha_c(0)=1$. For later discussion we shall assume that these shielding cuts are either negligible or do not seriously affect the description of the density matrices in terms of the poles only. This assumption will be justified when comparison with the data is made.

Consider now the special case of ρ^0 photoproduction. For the Pomeranchon exchange, the point $\alpha=1$ is a wrong-signature point. The amplitude f_{20} corresponds to a nonsense-sense transition (n, s) and the amplitudes f_{10} and f_{00} describe the sense-sense transition (s, s). Note that in the residue $\beta_{\lambda 0}$ of the amplitude we have included the coefficients $\phi_{\lambda 0}$ coming from the asymptotic form of the rotation functions. For brevity, in this section we use the index ns for $\lambda=2$ and $\mu=0$, and the subscript ss for $\lambda=0, 1$ and $\mu=0$, where λ is the helicity associated to the boson vertex and μ is the helicity associated with the $\bar{N}N$ vertex.

Let us consider the behavior of the residues for each coupling L-S. The residue $\beta_{\lambda 0}^{\alpha-2,2}$ according to Eq. (26) behaves like

$$\beta_{ns}^{\alpha-2,2}\propto A_-(\alpha),\quad \beta_{ss}^{\alpha-2,2}\propto A_-(\alpha). \tag{33}$$

It follows from Eq. (29) that the coefficient A_- near $\alpha=1$ behaves like $A_-(\alpha)\propto(\alpha-1)$, unless the residue has a pole at $\alpha=1$. Therefore, we have the following possibilities for the coupling with $L=\alpha-2$ and $S=2$:

(a) The residue is regular near $\alpha=1$; then

$$\beta_{ns}^{\alpha-2,2}\propto(\alpha-1),\quad \beta_{ss}^{\alpha-2,2}\propto(\alpha-1).$$

(b) The residue has a pole at $\alpha=1$; then

$$\beta_{ns}^{\alpha-2,2}=O(1),\quad \beta_{ss}^{\alpha-2,\alpha}=O(1).$$

In the usual terminology[32] case (a) corresponds to the Gell-Mann nonsense-choosing mechanism and case (b) corresponds to a multiplicative fixed pole at the wrong-signature point $\alpha=1$ mentioned above. These two alternatives can be distinguished by looking at the contribution of nonsense-sense transition amplitude f_{ns} (i.e., f_{20} in this case) to the observed spin-density matrix. It is helpful to keep these remarks in mind since they will be essential for our discussions in the following sections.

The other coupling does not enter into our discussions in the later sections. For completeness we shall indicate their behavior near the wrong-signature point $\alpha=1$. The coupling with $L=\alpha$, $\alpha+2$, and $S=2$, satisfy the following property:

$$\beta_{ns}^{LS}\propto(\alpha-1)\ \text{ and }\ \beta_{ss}^{LS}=O(1),$$

which correspond to the sense-choosing mechanism in the usual terminology. Finally, the coupling with $L=\alpha$, $S=0$ corresponds to the coupling in which t-channel helicity is conserved.[33] Note that this behavior changes as we go into wrong-signature points $\alpha=n$, where n is a negative integer. However, this is irrelevant for the present problem.

The ratios of the residues which will be needed later are as follows:

$$\frac{\beta_{20}^{\alpha-2,2}}{\beta_{00}^{\alpha-2,2}}=1,\quad \frac{\beta_{10}^{\alpha-2,2}}{\beta_{00}^{\alpha-2,2}}=\sqrt{2}\,, \tag{34}$$

$$\frac{\beta_{20}^{\alpha 2}}{\beta_{00}^{\alpha 2}}=0,\quad \frac{\beta_{10}^{\alpha 2}}{\beta_{00}^{\alpha 2}}=\frac{1}{\sqrt{2}\,(\alpha+1)}, \tag{35}$$

$$\frac{\beta_{20}^{\alpha+2,2}}{\beta_{20}^{\alpha+2,2}}=0,\quad \frac{\beta_{10}^{\alpha+2,2}}{\beta_{00}^{\alpha+2,2}}=\sqrt{2}\,\frac{\alpha}{\alpha+1}. \tag{36}$$

V. DECAY DENSITY MATRIX IN ρ PHOTOPRODUCTION

For the sake of completeness, the expression for the measured spin-density matrices, in terms of the s-channel helicity amplitudes,[34] will be reviewed briefly. Then, these quantities will be expressed in terms of the t-channel helicity amplitudes.

Let $\rho(\gamma)$ be the spin-density matrix of the polarized incident photon beam in the reaction $\gamma p \to \rho^0 p$. This matrix can be expanded in terms of the complete set of 2×2 matrix basis $I, \vec{\sigma}$, where $\vec{\sigma}$ are the Pauli spin matrices and I is the unit matrix. This expansion is given by

$$\rho(\gamma)=\tfrac{1}{2}I+\tfrac{1}{2}\vec{\sigma}\cdot\vec{\mathrm{P}}, \tag{37}$$

where $\vec{\mathrm{P}}$ is the degree of polarization. If the s-channel amplitude of the reaction $\gamma p \to \rho^0 p$ is f^s, then the spin-density matrix of the ρ^0 in the final state $\rho_{\lambda\mu}(V)$ will be given by

$$\rho_{\lambda\mu}(V)=[f^s\rho(\gamma)f^{s\dagger}]_{\lambda\mu}\,. \tag{38}$$

Since $\rho(\gamma)$ is linear in $\vec{\mathrm{P}}$, then $\rho(V)$ can be expressed as a linear function of $\vec{\mathrm{P}}$ as follows:

$$\rho_{\lambda\mu}(V)=\rho^0_{\lambda\mu}(V)+\sum_{\alpha=1}^{3}P^\alpha\rho^\alpha_{\lambda\mu}(V), \tag{39}$$

where

$$\rho^0_{\lambda\mu}(V)=\frac{1}{N}[f^sf^{s\dagger}]_{\lambda\mu}, \tag{40}$$

$$\rho^\alpha_{\lambda\mu}(V)=\frac{1}{N}[f^s\sigma^\alpha f^{s\dagger}]_{\lambda\mu}, \tag{41}$$

with $\alpha=1,2,3$. The normalization factor N is thus chosen such that $\mathrm{Tr}\rho_{\lambda\mu}(V)=1$, which gives

$$N=\sum_{\lambda\mu}|f^s_{\lambda\mu}|^2. \tag{42}$$

The spin-density matrices will then appear as parameters in the observed angular distributions of the ρ^0 decay. We refer to Ref. 34 for the details of these angular distributions. The information which can be extracted from the ρ^0-meson decay are as follows. The spin-density matrix $\rho^0_{\lambda\mu}$ can be obtained from the unpolarized-beam experiment. From the linearly polarized beam experiments one can extract the density matrices $\rho^1_{\lambda\mu}$ and $\rho^2_{\lambda\mu}$, and the circularly polarized experiments will give information on $\rho^3_{\lambda\mu}$. Since we are concerned with the experiments using linearly polarized γ, the discussion will be restricted to $\rho^0_{\lambda\mu}$, $\rho^1_{\lambda\mu}$, and $\rho^2_{\lambda\mu}$. The explicit expression of (39), (40), and (41) is

$$\rho^0_{\lambda\lambda'}=\frac{1}{N}\sum_{\gamma NN'}f^s_{\lambda N',\gamma N}f^{s*}_{\lambda' N',\gamma N}, \tag{43}$$

$$\rho^1_{\lambda\lambda'}=\frac{1}{N}\sum_{\gamma NN'}f^s_{\lambda N',-\gamma N}f^{s*}_{\lambda' N',\gamma N}, \tag{44}$$

$$\rho^2_{\lambda\lambda'}=\frac{1}{N}\sum_{\gamma NN'}\gamma f^s_{\lambda N',-\gamma N}f^{s*}_{\lambda' N',\gamma N}, \tag{45}$$

where γ, N, N' are the helicities of γ, N, N'. The expressions (43)–(45) give the density matrices, expressed in terms of the s-channel amplitudes, which are measured in the helicity frame. The helicity frame is the rest frame of ρ^0 with the z axis chosen opposite to the direction of the outgoing proton. For testing the t-exchange model it is more convenient to use the Gottfried-Jackson frame.[35] The Gottfried-Jackson (GJ) frame is defined to be the rest frame of the ρ^0 meson with the choice of the z axis along the direction of the incident photons.

It can easily be shown, using the orthogonality properties of the crossing matrices[35] of the nucleons N and N' and those of the photon crossing matrix, that in the GJ frame the density matrices are

$$\tilde{\rho}^0_{\lambda\lambda'}=\frac{1}{N}\sum_{\bar{\gamma}\bar{N}\bar{N}}f_{\lambda\bar{\gamma},\bar{N}N}f^*_{\lambda'\bar{\gamma},\bar{N}N}, \tag{46}$$

$$\tilde{\rho}^1_{\lambda\lambda'}=\frac{1}{N}\sum_{\bar{\gamma}\bar{N}N}f_{\lambda-\bar{\gamma},\bar{N}N}f^*_{\lambda'\bar{\gamma},\bar{N}N}, \tag{47}$$

$$\tilde{\rho}^2_{\lambda\lambda'}=\frac{1}{N}\sum_{\bar{\gamma}\bar{N}N}\bar{\gamma}f_{\lambda-\bar{\gamma},\bar{N}N}f^*_{\lambda'\bar{\gamma},\bar{N}N}\,. \tag{48}$$

The relevant property of the photon crossing matrix is that the crossing matrix is diagonal due to the zero mass of the photon.[36]

Using the Regge-pole contribution given by Eq. (15), substituted into (46)–(48), we obtain

$$\tilde{\rho}^0_{00}\simeq\left(\frac{s_0}{2m_Nm_\rho}\right)\left|\frac{\beta_{10}}{\beta_{00}}\right|^2(-\tau), \tag{49}$$

$$\tilde{\rho}^0_{11}\simeq\frac{1}{2}\left(\frac{s_0}{2m_Nm_\rho}\right)\frac{\beta_{10}}{\beta_{00}}\left(\frac{-\tau}{2}\right)^{1/2}, \tag{50}$$

$$\tilde{\rho}^0_{1,-1}\simeq\left(\frac{s_0}{2m_Nm_\rho}\right)^2\frac{\beta_{20}}{\beta_{00}}(-\tau), \tag{51}$$

$$\tilde{\rho}^1_{00}\simeq\left(\frac{s_0}{2m_Nm_\rho}\right)^2\left|\frac{\beta_{10}}{\beta_{00}}\right|^2\tau, \tag{52}$$

$$\tilde{\rho}^1_{11}\simeq\left(\frac{s_0}{2m_Nm_\rho}\right)^2\frac{\beta_{20}}{\beta_{00}}(-\tau), \tag{53}$$

$$\tilde{\rho}^1_{10} \simeq -\tfrac{1}{2}\left(\frac{s_0}{2m_N m_\rho}\right)^2 \frac{\beta_{10}}{\beta_{00}}\sqrt{-\tau}\,, \tag{54}$$

$$\tilde{\rho}^1_{1,-1} \simeq \tfrac{1}{2} + O(\tau^2), \tag{55}$$

$$\mathrm{Im}\tilde{\rho}^2_{10} \simeq \left(\frac{s_0}{2m_N m_\rho}\right)\frac{\beta_{10}}{\beta_{00}}\sqrt{-\tau}, \tag{56}$$

$$\mathrm{Im}\tilde{\rho}^2_{1,-1} \simeq -\tfrac{1}{2} + O(\tau^2), \tag{57}$$

where $\tau = t/m_\rho^2$, taking into account only the dominant amplitude $f_{\lambda 0}$ ($\lambda = 0, 1, 2$). In obtaining these density matrices, the approximation $z \simeq s\sqrt{-t}/m_N m_\rho^2$ has been used. This is legitimate, because we are concerned only about the diffraction region in which $|t| < 0.4\ \mathrm{GeV}^2$.

VI. s-CHANNEL HELICITY CONSERVATION

In this section the prediction given by Eqs. (49)–(57) will be compared with the experimental data. In order to do this, the data will be expressed analytically near the forward direction, using the experimental fact that s-channel helicity is conserved. The helicity conservation can be expressed by demanding that the s-channel helicity amplitude be given by[21,34]

$$f^s_{\rho N',\gamma N} = g_{N\gamma}\delta_{N'N}\delta_{\rho\gamma}\,. \tag{58}$$

Substitution of this ansatz into Eqs. (43)–(45) leads to the following expressions for the density matrices:

$$\rho^0_{\lambda\lambda'} = \tfrac{1}{2}(\delta_{\lambda 1}\delta_{1\lambda'} + \delta_{\lambda,-1}\delta_{-1\lambda'}), \tag{59}$$

$$\rho^1_{\lambda\lambda'} = a(\delta_{\lambda,-1}\delta_{\lambda',+1} + \delta_{\lambda 1}\delta_{\lambda',-1}), \tag{60}$$

$$\mathrm{Im}\rho^2_{\lambda\lambda'} = a(\delta_{\lambda,-1}\delta_{\lambda' 1} - \delta_{\lambda 1}\delta_{\lambda,-1}). \tag{61}$$

In the helicity-conservation model[34] $a = \frac{1}{2}$ for $\eta = +1$ exchange. The density matrices in Eqs. (59), (60) are measured in the helicity frame. The density matrix in the GJ frame $\tilde{\rho}^i_{\lambda\lambda'}$ is related to the density matrix in the helicity frame $\rho^i_{\lambda\lambda'}$ by the equation[21,34]

$$\tilde{\rho}^i_{\lambda\lambda'} = d^1_{\lambda\mu}(\phi)\rho^i_{\mu\mu'}d^1_{\mu'\lambda'}(-\phi), \tag{62}$$

where

$$\cos\phi = \frac{1+\tau}{1-\tau}, \quad \sin\phi = \frac{2\sqrt{-\tau}}{1-\tau}, \tag{63}$$

$\tau = t/m_\rho^2$, and $d^1_{\lambda\mu}$ is the rotation function. Performing the transformation (62) up to the order τ, for the near-forward direction, the data can be represented well by

$$\tilde{\rho}^0_{00} \simeq -2\tau, \quad \tilde{\rho}^0_{10} \simeq \tfrac{1}{2}\sqrt{-2\tau}, \quad \tilde{\rho}^0_{1,-1} \simeq -\tau\,, \tag{64}$$

$$\begin{aligned}&\tilde{\rho}^1_{00} \simeq +2\tau, \quad \tilde{\rho}^1_{11} \simeq -\tau,\\ &\tilde{\rho}^1_{10} \simeq -\tfrac{1}{2}\sqrt{-2\tau}, \quad \tilde{\rho}^1_{1,-1} \simeq \tfrac{1}{2}(1+2\tau)\,,\end{aligned} \tag{65}$$

$$\mathrm{Im}\tilde{\rho}^2_{10} \simeq \tfrac{1}{2}\sqrt{-2\tau}, \quad \mathrm{Im}\tilde{\rho}^2_{1,-1} \simeq -\tfrac{1}{2}(1+2\tau). \tag{66}$$

Let us first of all consider the contributions of the sense-nonsense amplitude f_{20} into the density matrix in the first order. From Eqs. (51) and (53) it follows that the sense-nonsense amplitude f_{20} contributes to $\rho^0_{1,-1}$ and ρ^1_{11}, for the wrong-signature point $\alpha = 1$ of the Pomeranchon. The experimental points clearly exclude the possibility $\rho^0_{1,-1} = 0$ and $\rho^1_{11} = 0$ *in the diffraction region* $|t| < 0.4\ (\mathrm{GeV})^2$. Therefore, there are contributions from the sense-nonsense transition amplitude. In Sec. IV we have noted that there are two possibilities for the $L = \alpha - 2$, $S = 2$ coupling. First, the nonsense-chosing mechanism which *does not* contribute to the sense-nonsense transition amplitude and, secondly, the multiplicative fixed pole at the wrong-signature point, which *does* lead to a nonvanishing sense-nonsense transition amplitude. Since the data prefer a choice of $f_{ns} \neq 0$, it follows that the coupling $L = \alpha - 2$, $S = 2$ must be present and, from the general argument in ρ^0 photoproduction, this corresponds to the wrong-signature point multiplicative fixed pole. In addition to this, if one uses the ratio of the residue coupling given by (34) and takes $s_0 = 2m_N m_\rho$ in Eqs. (49)–(57), then one obtains Eqs. (64)–(66), with the exception of the equations for $\rho^1_{1,-1}$ and $\mathrm{Im}\rho^2_{1,-1}$ in the terms of order τ. This is due to an approximation where f_{01} is neglected with respect to f_{00}, since $f_{00} = O(1)$ and $f_{01} = O(\sqrt{t}\,)$ which comes from the angular factor $\frac{1}{2}(1 - z_t^2)^{1/2}$. While this gives a higher-order effect in τ, in other density matrices it will give a $O(\tau)$ correction in $\rho^1_{1,-1}$ and $\mathrm{Im}\rho^2_{1,-1}$.

The result of this section can be summarized as follows. The data prefer the contribution of sense-nonsense amplitude; hence there is the $L = \alpha - 2$, $S = 2$ coupling at the wrong-signature point $\alpha = 1$, with a multiplicative fixed pole. Subsequently, the choice of this coupling with multiplicative fixed pole gives the spin-density matrices which show the s-channel helicity conservation. Considering the fact that a shielding-cut contribution is present,[23,25] this result is consistent with the assumption that these shielding cuts are either negligible or do not seriously affect the description of the density matrices in terms of the multiplicative fixed pole only.

VII. CONCLUSION

The discussions and conclusions in the previous sections can be summarized as follows.

A. The Parity Rule

Consider the vertex $c\bar{a}$ in which the spins are s_c and $s_{\bar{a}}$. The total spin of the external particle is $\vec{S}_f = \vec{s}_c + \vec{s}_{\bar{a}}$. In the coupling to the Regge exchange, the system $c\bar{a}$ has orbital angular momentum L_f, such that the total angular momentum $\vec{J}=\vec{L}_f+\vec{S}_f$ (see Fig. 1). Denoting the Regge-exchange contribution to the t-channel amplitude by $f_{\lambda\mu}=f^{+}_{\lambda\mu}+f^{-}_{\lambda\mu}$, where $f^{\pm}_{\lambda\mu}$ is the contribution from natural- (unnatural-) parity exchange, the parity rule can be stated as follows: $f^{\pm}_{0\mu}\neq 0$ only if the condition $\xi_{\bar{a}}\xi_c=\pm(-1)^{S_f}$ is satisfied. If the above-stated condition is not satisfied, then $f^{\pm}_{0\mu}=0$ for all s and t. If the Regge exchange has a quantum number $C\times P$, evenness or oddness of S_f is determined by the relation

$$CP=\xi_{\bar{a}}\xi_c(-1)^{S_f},$$

if C conservation applies.

Remarks. In obtaining the parity rule, we have used only the L-S scheme and the well-known relations between the quantum numbers C and P of a system with spin and orbital angular momentum. The addition of spins at the vertices is specific to t-channel exchange processes. Therefore, this rule is a general feature of t-channel exchange. A consequence of this rule is that for Pomeranchon exchange, in order that the diffraction be observed, the intrinsic parities $\xi_{\bar{a}}$, ξ_c and the total spin must satisfy the relation $\xi_{\bar{a}}\xi_c=(-1)^{S_f}$. For a spin-0-induced reaction, this rule is equivalent to the Morrison rule. For a nonzero-spin-induced reaction, this rule leads to different consequences. For example, from the parity rule it is expected that the B^0 photoproduction and the N^*, with $J^P=\frac{1}{2}^-$, $\frac{3}{2}^+$ production, are expected to be produced diffractively. These possibilities are at present not yet excluded by the experiments. Since the parity rule seems to be satisfied by the diffraction processes it may suggest that the diffraction phenomenon is a t-channel exchange process.

B. Helicity Conservation in ρ^0 Photoproduction

The natural-parity exchange in ρ^0 photoproduction can couple to the boson vertex with $S_f=0,\ 2$. *A priori*, the coupling with $S_f=2$ and angular momenta $L=\alpha-2$ at the wrong-signature point $\alpha=1$ is consistent *either* with a nonsense mechanism *or* with a multiplicative fixed pole at the wrong-signature point. In the first case the amplitude does not contribute to the high-energy behavior, while in the second case, the amplitude does contribute to the high-energy behavior, in particular, the nonsense-sense amplitude gives a nonvanishing contribution. This is a well-known property in the Regge-pole theory. The other couplings with $S_f=2$, $L=\alpha$, $\alpha+2$, will lead to the nonsense-sense transition amplitude $f_{20}=0$ at $\alpha=1$. The experiment requires a nonvanishing amplitude f_{20} in the diffraction region, therefore it prefers the coupling $S_f=2$, $L=\alpha-2$. From the residue ratios it is found that the residue ratio of this coupling by itself gives the s-channel helicity-conserving density matrices in ρ^0 photoproduction by linearly polarized photons. Therefore we may conclude that the data indicate that the diffraction in ρ photoproduction occurs almost entirely through $l=\alpha-2$ wave with a multiplicative fixed pole. The implications of this result on the cut contribution in the framework of complex-angular-momentum theory[2-4] will be discussed below.

In the discussion of the s-channel helicity conservation we have used the classification of possible couplings according to the Regge-pole model with L-S coupling. In addition to these, we use general arguments from the theory of complex angular momentum[2-4]. According to the latter, we know that in general there are two types of fixed poles: i.e., the additive and multiplicative fixed pole at the wrong-signature point $\alpha=1$. The additive fixed pole does not contribute to the high-energy behavior. The multiplicative fixed pole at the wrong-signature point is a dynamical question.[4] Here the experimental evidence for this multiplicative fixed pole is indicated. However, within the framework of complex angular momentum, this fixed pole requires shielding cuts which are moving branch points $\alpha_c(t)$ in the complex angular-momentum plane, with $\alpha_c(0)=1$. The above result, that the residue ratios alone give the s-channel helicity conservation, is compatible with the assumption that these shielding cuts are either negligible or do not affect the description of the density matrices in terms of the multiplicative fixed pole.

For an ordinary Regge pole, the multiplicative fixed pole at wrong-signature points is not consistent with the exchange degeneracy.[37] But the Pomeranchon does not participate in exchange degeneracy, and therefore this fixed pole may be a specific feature of the Pomeranchon exchange in this process.

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