Exchange-degeneracy-violating A_1-Z Regge amplitude and dibaryon resonances in NN system

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The measured large value of $\Delta\sigma_L$ for pp scattering must be interpreted as a violation of strong exchange degeneracy of A_1 and Z. A strong-exchange-degeneracy–violating A_1-Z contribution to $\mathrm{Im}U_0$ provides a fine representation of $\Delta\sigma_L$ for $p_L \geq 3$ GeV/c. The Legendre coefficients of this amplitude draw the loops on the Argand diagram. We discuss a similarity between the top positions of them and the possible dibaryon resonances, and a duality scheme due to low-lying trajectories.

The last two years have witnessed a considerable change in the philosophy underlying the discussion of the resonance problem. It was believed that every resonance is built up of qqq (baryons) or $q\bar{q}$ (mesons). Other combinations of quarks and antiquarks are called exotic states and their existence was scarcely believed. However, the dual string model,[1] the bag model,[2] and the naive quark model[3] approaches predict exotic states called baryoniums, and candidates for them have been reported experimentally.

Recently, de Boer *et al.*, Biegert *et al.*, and Auer *et al.* measured pp total cross sections in pure spin states with polarized beam and target at $p_L = 1$–6 GeV/c. They have observed an unexpected energy dependence in the total-cross-section differences, $\Delta\sigma_T = \sigma^{\mathrm{tot}}(\uparrow\downarrow) - \sigma^{\mathrm{tot}}(\uparrow\uparrow)$ (Ref. 4) and $\Delta\sigma_L = \sigma^{\mathrm{tot}}(\rightleftarrows) - \sigma^{\mathrm{tot}}(\rightrightarrows)$ (Ref. 5) for pp scatterings in initial transverse and longitudinal spin states, respectively. These theoretical and experimental results shed a new light on the subject of dibaryon resonances. Their data show two bumps at $p_L = 1.5$–2 GeV/c, and suggest that these bumps may be candidates for the dibaryon resonances. In order to study the spin states of these bumps, Grein and Kroll presented some combinations of σ^{tot}, $\Delta\sigma_T$, and $\Delta\sigma_L$ which project out certain groups of partial waves. They argued that a very pronounced bump in $\Delta\sigma_L$ at around $p_L = 1.5$ GeV/c is an uncoupled triplet state ($L = J$) and another bump in $\Delta\sigma_T$ at around 2 GeV/c is a spin singlet state.[6] It has been suggested that the former bump is in the $^3F_3$ wave on the bases of p-p phase-shift analysis[7] and the energy dependence of the coefficients of a Legendre-polynomial expansion of $Pd\sigma/d\Omega$ and the structure appearing in the data of $\Delta\sigma_L$.[8] Another candidate is $^1D_2(2.16)$ which is expected to be a resonance from phase-shift analysis.[9]

It has been generally thought that the pp scattering was an exotic channel and its amplitudes, except for the Pomeron contribution, were dominantly real at high energies. However, the imaginary parts proportional to $\Delta\sigma_T$ and $\Delta\sigma_L$ are not zero at high energies. Thus, it is true that the exchange degeneracy (EXD) is not complete. It is a very interesting temptation to speculate that these imaginary parts are dual to the pp resonances, i.e., the dibaryon resonances.

In this paper the imaginary parts are fitted by Regge poles at high energies,[10] and we project these Regge amplitudes with a signature factor into a Legendre-polynomial expansion. Then we can observe that the coefficients draw the loops[11] on the Argand diagrams which are closely connected to the low-energy region. We discuss a similarity between the top and bottom positions of the loops and the possible dibaryon resonances. However, we cannot prove that these positions correspond exactly to the poles of the S-matrix elements. We also give an application to some other channels and comment on a duality scheme due to low-lying Regge trajectories.

The cross sections are expressed in terms of the conventional s-channel helicity amplitudes[10] as

$$\sigma^{\mathrm{tot}} = \frac{1}{2pW}\mathrm{Im}\tfrac{1}{2}(\phi_1+\phi_3)\Big|_{t=0} = \frac{1}{2pW}\mathrm{Im}N_0\Big|_{t=0}, \tag{1}$$

$$\Delta\sigma_T = -\frac{1}{2pW}\mathrm{Im}\phi_2\Big|_{t=0} = \frac{1}{2pW}\mathrm{Im}(N_2 - U_2)\Big|_{t=0} \tag{2}$$

$$\Delta\sigma_L = \frac{1}{2pW}\mathrm{Im}(\phi_1 - \phi_3)\Big|_{t=0} = \frac{1}{2pW}\mathrm{Im}2U_0\Big|_{t=0}, \tag{3}$$

where p and W are the c.m. momentum and energy, respectively. N and U correspond to t-channel natural- and unnatural-parity exchange, respectively. Their subscripts 0 and 2 refer to the total s-channel helicity flip.

At first we will analyze U_0 which contains the contribution of trajectories associated with the

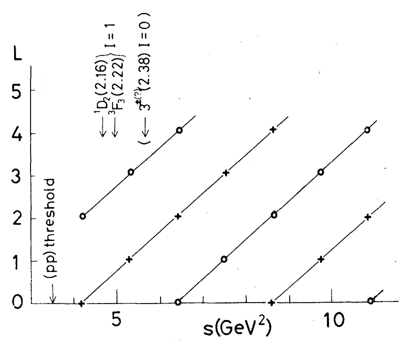

FIG. 1. Top (open circle) and bottom (cross) positions of $2U_0^{A_1\text{-}Z} = \phi_1^{A_1\text{-}Z} - \phi_3^{A_1\text{-}Z}$ for NN scattering vs s.

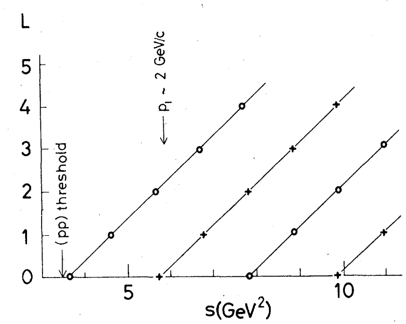

FIG. 2. Top and bottom positions for $U_2^{\rho\text{-}A_2} - N_2^{\rho\text{-}A_2} = \phi_2^{\rho\text{-}A_2}$ for NN scattering vs s.

$A_1(J^{PC}=1^{++})$ and conjectured $Z(2^{--})$ nonets. Since pp is an exotic channel, the contribution of EXD $(A_1 - Z)$ must produce a real amplitude. However, the measured value of $\Delta\sigma_L$ is nearly equal to -1 mb at $p_L = 6$ GeV/c. This must be interpreted as a violation of EXD. Here we assume the weak EXD and parametrize the s-channel helicity amplitudes as

$$\phi_1^{A_1} = -\phi_3^{A_1} = \tfrac{1}{2}[\gamma_{A_1 p\bar{p}}{}^2(t)](1 - e^{-i\pi\alpha(t)})(\alpha' s)^{\alpha(t)}, \tag{4}$$

$$\phi_1^{Z} = -\phi_3^{Z} = \tfrac{1}{2}[\gamma_{Z p\bar{p}}{}^2(t)](1 + e^{-i\pi\alpha(t)})(\alpha' s)^{\alpha(t)}, \tag{5}$$

where

$$\alpha_{A_1}(t) = \alpha_Z(t) = \alpha(t) = \alpha_0 + \alpha' t.$$

Then the U_0 amplitude leads to

$$\begin{aligned} 2U_0^{A_1\text{-}Z} &= \phi_1^{A_1\text{-}Z} - \phi_3^{A_1\text{-}Z} \\ &= [\gamma^2(t) - \Delta(t) e^{-i\pi\alpha(t)}](\alpha' s)^{\alpha(t)}, \end{aligned} \tag{6}$$

where $\Delta(t)$ shows the magnitude of strong-exchange-degeneracy (SEXD) violation and we will parametrize

$$\Delta(t) = [\gamma_{A_1 p\bar{p}}{}^2(t) - \gamma_{Z p\bar{p}}{}^2(t)] = \Delta e^{bt}, \tag{7}$$

$$\gamma^2(t) = [\gamma_{A_1 p\bar{p}}{}^2(t) + \gamma_{Z p\bar{p}}{}^2(t)] = \gamma^2 e^{at}. \tag{8}$$

Berger *et al.*[10] succeeded in fitting the high-energy data by fixing the parameters $\alpha_0 = -0.19$, $\alpha' = 0.9$ (GeV/c)$^{-2}$, $b = 4$ (GeV/c)$^{-2}$, and $\Delta \sin\pi\alpha_0 = -52$. We will use these values in the following.

Now, after symmetrizing the amplitude (6) as

$$2U_0^{\pm}(z) = U_0(z) \pm U_0(-z),$$

we define the coefficients of a Legendre expansion of the amplitude as

$$\begin{aligned} A_L^{\pm} &= \int_{-1}^{1} 2U_0^{\pm}(z) P_L(z)\, dz \\ &= \tfrac{1}{2}[1 \pm (-1)^L] A_L, \\ A_L &= \int_{-1}^{1} 2U_0(z) P_L(z)\, dz. \end{aligned} \tag{9}$$

A_L is expressed in terms of many partial waves, e.g.,

$$A_2 = \tfrac{2}{3}\sqrt{6}\, {}^3(P,F)_2 + {}^1D_2 + \tfrac{1}{3}\, {}^3F_2 - \tfrac{7}{12}\, {}^3F_3 + \tfrac{1}{4}\, {}^3F_4 - \tfrac{1}{5}\sqrt{5}\, {}^3(F,H)_4 + \tfrac{1}{5}\, {}^3H_4 - \tfrac{11}{30}\, {}^3H_5 + \cdots, \tag{10}$$

where ${}^{2S+1}L_J$ represents a partial-wave amplitude. Here we concentrate on the weak EXD A_1-Z Regge-pole singularity and use the expression (6) for $U_0(Z)$. In doing so, we ignore possible EXD breaking contributions to, e.g., ϕ_4, which due to identical-particle symmetry could generate a sizable backward peak. Substituting the expressions (6)–(8) into Eq. (9), we can obtain

$$\begin{aligned} A_L^{A_1\text{-}Z} = {} & \gamma^2(\alpha' s)^{\alpha_0} \exp\{[a + \alpha' \ln(\alpha' s)](t_{\max} - 2p_i p_f)\}\{\pi/2[a + \alpha' \ln(\alpha' s)]p_i p_f\}^{1/2} I_{L+1/2}(2[a + \alpha' \ln(\alpha' s)]p_i p_f) \\ & + \Delta(\alpha' s)^{\alpha_0} \exp\{[b + \alpha' \ln(\alpha' s)](t_{\max} - 2p_i p_f)\} \left| j_L(2\pi\alpha' p_i p_f + i2[b + \alpha' \ln(\alpha' s)]p_i p_f) \right| \exp[i\Phi(s)], \end{aligned} \tag{11}$$

where p_i and p_f are the initial and final state c.m. momenta, respectively, and $t_{\max}$ is the maximum value of momentum transfer squared. The phase factor Φ is defined by

$$\Phi(s) = \pi[\alpha'(2p_i p_f - t_{\max}) - \alpha_0 - \tfrac{1}{2}L + 1] + \arg j_L + 2m\pi, \tag{12}$$

where m is an arbitrary integer. It can be shown that the contribution of $\arg j_L$ to Φ is not so large at small s that we can safely discard this term. Note that the factor $2p_i p_f$ is roughly equal to $s/2$. Then it is clear that the second factor of A_L generates counterclockwise circular motion (Schmid circle)[11] on the

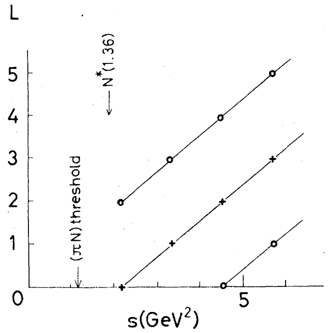

FIG. 3. Top and bottom positions for $P+N\to\pi+N$ scattering amplitude vs s.

Argand diagram with angular velocity depending on α'. The momenta for the tops and bottoms of the loops are determined by the general expression

$$\alpha'(2p_ip_f - t_{\max}) - \alpha_0 - \tfrac{1}{2}L + 1 + 2m = \tfrac{1}{2} \text{ and } \tfrac{3}{2}, \tag{13}$$

respectively.

In the case of $NN\to NN$ (and $N\bar{N}\to N\bar{N}$) Eq. (13) is simplified to

$$s = 4m_N{}^2 + (1/\alpha')[L + 2\alpha_0 - 1 + \mathrm{mod}(4), \\ \text{and } L + 2\alpha_0 + 1 + \mathrm{mod}(4)]. \tag{14}$$

The top and bottom positions for $2U_0$ are illustrated in Fig. 1 by the open circles and crosses, respectively. We also show the positions of candidates for dibaryon resonances which have been reported experimentally.[7-9,12] If the $^1D_2$ (2.16) and $^3F_3$ (2.22) lie on the leading trajectory of dibaryon resonances, we can see from Fig. 1 that a spin-singlet and a spin-triplet state shift to the high- and the low-mass regions, respectively. If this hold in the case of $3^{\pm}$ $(2.38, I=0)$, this state is in $^1F_3$ or $^3G_3$ wave.

As for the amplitude ϕ_2 the situation is more complicated. This amplitude is dominated at small t by the $\pi+B$ cut contribution. At large t it has the important contribution from EXD $\rho+A_2$, which must produce a real amplitude. At intermediate value of t, however, it is possible for the imaginary part of ϕ_2 to remain an abnormally large value in contrast to the reduced real part due to the $\rho+A_2$ and $\pi+B$ cut cancellation.[10] Therefore we tentatively assume that $\phi_2^{\rho\text{-}A_2}$ is dominated by ρ contribution and we perform a procedure similar to $\phi_1^{A_1\text{-}Z}-\phi_3^{A_1\text{-}Z}$. Here we used ρ-trajectory parameters as $\alpha_0=0.57$ and $\alpha'=0.96$ $(\mathrm{GeV}/c)^{-2}$ in expression (12). The top and bottom positions of the loops are shown in Fig. 2 with the position of bump at near $p_L=2$ GeV/c. If this bump is a spin-singlet state,[6] we can guess that it is probably in $^1S_0$ or $^1D_2$ wave.

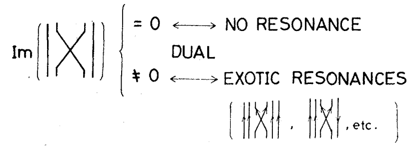

FIG. 4. Duality scheme of EXD-violating process.

Finally we will consider a diffraction-dissociation process such as $\pi+N\to\pi+(\pi N)_{\rm DD}$, i.e., a virtual process $P+N\to\pi+N$, where P is Pomeron. The leading exchange of this process is the s-channel helicity-nonflip and unnatural-parity exchange. If it is dominated by the SEXD-violating A_1+Z, we can also expect some resonances in this process. We show the top and bottom positions for this case in Fig. 3 together with an experimental candidate for an exotic baryon $(qqqq\bar{q})$ $N^*(1.36)$.[13-15] We expect that an experimental search for a new bump near the next top position in $\pi/N+N\to\pi/N+(\pi N)_{\rm DD}(1.83)$ will be performed.

In summary, the experiments with polarized beams and polarized targets have changed the situation and have revealed rather rare mechanisms such as those of (A_1+Z) exchange and $(\pi+B)$ cut-$(\rho+A_2)$ cancellation. We have shown that SEXD is possibly broken in these cases, where the exotic resonances are able to exist as is illustrated in the quark-line diagram (Fig. 4), and their existence is not inconsistent with the experimental data. From this point of view it is a very important problem to reform the duality scheme of high-energy physics.

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