†On leave of absence from the Deutches Elektronen Synchrotron.

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[7]By integrating the theoretical cross section over the $N^*$, $\rho^0$ mass region, partial waves other than the $p$ waves from $N^*$ and $\rho^0$, respectively, are contributing, too. In the case of the $N^*$ region their effect on the off-shell corrections is supposed to be negligible. In the case of the $\rho^0$ meson, the $T=0,2$ $S$-wave contributions to $\sigma_{\pi^+\pi^-}(m)$ have been estimated from a $\pi\pi$ phase-shift analysis (Ref. 5). Instead of (3), the following correction formula was used:

$$q_t\sigma_{\pi^+\pi^-}(m,t)=\left[c_0+c_1\left(\frac{q_t}{q}\right)^2\frac{1+q^2R_\rho^{\ 2}}{1+q_t^{\ 2}R_\rho^{\ 2}}\right]q\sigma_{\pi^+\pi^-}(m) \quad (3a)$$

with $c_1=0.8+\frac{1}{2}(m-0.6)$, $m$ in units of GeV, and $c_0=1-C_1$.

[8]N. Gelfand, Columbia University Report No. NEVIS-137, 1965 (unpublished) (2.35 GeV/$c$). The data at 3-4 and 6.95 GeV/$c$ have been taken from Ref. 1, and from P. Slatterly, H. Kraybill, B. Forman, and T. Ferbel, University of Rochester Report No. UR-875-153, 1966 (unpublished) (6.95 GeV/$c$); Aachen-Berlin-Birmingham-Bonn-Hamburg-London (I.C.)-München Collaboration, Nuovo Cimento 35, 659 (1965); Phys. Rev. 138, B897 (1965) (4.0 GeV/$c$).

[9]H. C. Dehne, E. Raubold, P. Soding, M. W. Teucher, G. Wolf, and E. Lohrmann, Phys. Letters 9, 185 (1964); H. C. Dehne, thesis, University of Hamburg, 1964 (unpublished) (3.6 GeV/$c$); V. Alles-Borelli, B. French, A. Frisk, and L. Michejda, Nuovo Cimento 48A, 360 (1967) (5.7 GeV/$c$).

# GHOST-ELIMINATING ZEROS IN THE REGGE-POLE MODEL

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The superconvergent sum rules for helicity-flip and -nonflip amplitudes controlled by only the $R$ trajectory are used to clarify the ghost-eliminating mechanisms.

The diffraction shrinkage at high energy for the reaction $\pi^-+p\rightarrow\pi^0+n$ has been successfully explained by the Regge-pole model based on a single $\rho$-meson exchange.[1] In addition, the dip phenomenon observed in the above reaction around $t\approx-0.6$ GeV$^2$ has also been clearly explained in the Regge-pole model with a vanishing helicity-flip amplitude at $\alpha=0$.[2,3] Shrinkage is also seen in another reaction $\pi^-+p\rightarrow\eta+n$ which is supposed to be controlled by the $R$ (or $A_2$) trajectory with even signature.[4] In this case, however, there remain the following unsettled questions[5]:

(i) Experimentally, it is not yet known whether the $R$ trajectory passes through spin zero.[4] If it does, we have a ghost problem. Then to eliminate this ghost[6] the residue function of the helicity-nonflip amplitude will have to vanish.

(ii) It is an open question whether the exchange of the $R$ trajectory can produce helicity flip at $\alpha=0$.[3] If the trajectory simply "chooses nonsense" as suggested by Gell-Mann,[6] the helicity-flip amplitude comes out finite but nonvanishing at $\alpha=0$. On the other hand, if the trajectory "chooses sense," an additional zero occurs at $\alpha=0$ in the residue of the helicity-flip amplitude as implied by the ghost-eliminating mechanism of Chew.[7] Then the helicity-flip term will actually vanish at $\alpha=0$. Therefore, it will be of great interest for the models of ghost elimination whether the helicity-flip amplitude due to the $R$ exchange indeed vanishes at $\alpha=0$. Phenomenologically, the discrimination between these models by Arbab, Bali, and Dash[8] has not yet been convincing since two trajectories contribute, and the theory involves a number of parameters.

The purpose of this Letter is an attempt to answer the above questions (i) and (ii) in connection with the ghost-eliminating mechanisms, applying the previous techniques[9,10] of superconvergence sum rules to the helicity-nonflip and -flip amplitudes controlled by only the $R$

trajectory at high energies. This method allows us to deduce a clearer conclusion on the zeros of the $R$-pole residues. The essential point of this is to connect the $R$-pole parameters (as a function of $t$) with the low-energy integrals using the above technique, which enables us to investigate the zeros that are characteristic in the Regge formula.[11]

We shall start from question (i): whether the $R$ trajectory passes through spin zero or not. In order to answer this question, we first derive a superconvergent sum rule for the helicity-nonflip amplitude

$$A^{(+)}(\nu,t) \equiv \tfrac{1}{4}[A_{K^-p}(\nu,t)+A_{K^+p}(\nu,t) - A_{K^-n}(\nu,t)-A_{K^+n}(\nu,t)]$$

[which Singh[12] calls $A^{(+)\prime}(\nu,t)$] which will be controlled by only the $R$ pole at high energies. Here $\nu=\nu_L+t/4m_N$, with $\nu_L$ the incident $K$-meson energy in the lab system. Let us assume that there are no other singularities except for the $R$ pole in the complex $J$ plane for $\alpha \geqslant -2$ at $t=0$.[13] Then $A^{(+)}(\nu,t)$ can be separated as

$$A^{(+)}(\nu,t)=A_K^{R}(\nu,t)+A^{(+)\prime}(\nu,t) \tag{1}$$

with

$$\mathrm{Im}A_K^{R}(\nu,t)=\alpha_R(t)[\alpha_R(t)+1]\frac{C_R(t)}{m_K}\left(\frac{\nu}{m_K}\right)^{\alpha_R(t)} \tag{2}$$

in an asymptotic form at high energies.[14] Following the same procedure as in the previous papers,[9,10,15] we obtain the superconvergent sum rule

$$\int_0^\infty d\nu\,\nu\,\mathrm{Im}A^{(+)\prime}(\nu,t)=0. \tag{3}$$

If the Regge asymptotic behavior is assumed to be already established at high energies $\nu_L > \nu_A$ (we take[16] $\nu_A=2$ GeV for convenience), we can express the term involving the $R$-pole parameters in terms of the low-energy integrals as follows:

$$\frac{1}{4}\frac{g_\Lambda^2}{4\pi}\frac{1}{m_N^2}\nu_\Lambda(t)X_-(\Lambda,t)-\frac{1}{4}\frac{g_\Sigma^2}{4\pi}\frac{1}{m_N^2}\nu_\Sigma(t)X_-(\Sigma,t)+\frac{1}{4}\frac{g_{Y_0^*}^2}{4\pi}\frac{1}{m_N^2}\nu_{Y_0^*}(t)X_+(Y_0^*,t)$$
$$-\frac{1}{4}\frac{g_{Y_1^*}^2}{4\pi}\frac{1}{m_N^2}\nu_{Y_1^*}(t)Z(Y_1^*,t)+\frac{1}{2\pi^2m_N}\int_{m_K}^{\nu_A}d\nu_L\left(\nu_L+\frac{t}{4m_N}\right)\mathrm{Im}A^{(+)}\left(\nu_L+\frac{t}{4\pi_N},t\right)$$
$$=\frac{m_K}{2\pi^2m_N}\frac{\alpha_R(t)[\alpha_R(t)+1]}{\alpha_R(t)+2}C_R(t)\left(\frac{\nu_A+t/4m_N}{m_K}\right)^{\alpha_R(t)+2}, \tag{4}$$

where[17]

$$\nu_Y(t)\equiv\frac{m_Y^2-m_N^2-m_K^2}{2m_N}+\frac{t}{4m_N}, \tag{5}$$

$$X_\pm(Y,t)\equiv\pm m_Y+m_N+\frac{\nu_Y(t)}{1-t/4m_N^2}, \tag{6}$$

$$Z(Y_1^*,t)\equiv-(m_N+m_{Y_1^*})\frac{t}{2}$$
$$+\frac{[(m_N+m_{Y_1^*})^2-m_K^2][(m_N+m_{Y_1^*})\{(m_N-m_{Y_1^*})^2-m_K^2\}-m_{Y_1^*}(m_N^2+m_K^2-m_{Y_1^*}^2)]}{6mY_1^{*2}}$$
$$+\frac{\nu_{Y_1^*}(t)}{1-t/4m_N^2}\left[\frac{t}{2}+\frac{\{(m_N+m_{Y_1^*})^2-m_K^2\}\{(m_N-m_{Y_1^*})^2-m_K^2-2m_Nm_{Y_1^*}\}}{6m_{Y_1^*}^2}\right], \tag{7}$$

and $m_A$ denotes the mass of a particle $A$. In the narrow-width approximation,

$$\frac{1}{2\pi^2 m_N}\int_{m_K}^{\nu_A} d\nu_L(\nu_L+t/4m)\,\mathrm{Im}A^{(+)}(\nu_L+t/4m_N,t)=\sum_{I,l_\pm}\pm C_I\frac{\Gamma_{l_\pm}^{\mathrm{el}}}{8m_N^2q_{l_\pm}^3}\left(\nu_{l_\pm}+\frac{t}{4m_N}\right)[P_{l\pm1}'(Z_{l_\pm})\{(M_{l_\pm}-m_N)^2$$
$$-m_K^2\}X_+(M_{l_\pm},t)-P_l'(Z_{l_\pm})\{(M_{l_\pm}+m_N)^2-m_K^2\}X_-(M_{l_\pm},t)], \qquad (8)$$

with the same notation as in II and III except that $\mu$ should be replaced by $m_K$ and $C_I=\frac{1}{2}$ $(-\frac{1}{2})$ for $I=0$ (1), respectively. Here $Z_{l_\pm}=1+t/2q_{l_\pm}^2$. In order to investigate possible zeros in the Regge amplitude we shall search for the zeros of the left-hand side of Eq. (4) with the use of Eq. (8). For evaluating the left-hand side as a function of $t$ the following data are used: For the $KN\Lambda$ and $KN\Sigma$ coupling constants, nothing definite is known. Therefore, we considered the following two cases: case (a), values obtained by Zovko,[18]

$$g_\Lambda^2/4\pi=6.8\pm2.9 \text{ and } g_\Sigma^2/4\pi=2.1\pm0.9;$$

case (b), values[19] obtained from SU(3),

$$\frac{g_\Lambda^2}{4\pi}=\frac{g_N^2}{4\pi}\frac{(1+2\alpha)^2}{3} \text{ and}$$

$$\frac{g_\Sigma^2}{4\pi}=\frac{g_N^2}{4\pi}(1-2\alpha)^2 \text{ with } \frac{g_N^2}{4\pi}=14.5.$$

The coupling constants $g_{Y_0^*}$ and $g_{Y_1^*}$, according to the estimates of Warnock and Frye,[20] are

$$g_{Y_0^*}^2/4\pi=0.32 \text{ and } g_{Y_1^*}^2/4\pi=1.9/m_N^2.$$

As direct-channel resonance parameters, those tabulated in the Rosenfeld table[21] have been used. In Fig. 1 we plot the left-hand side of Eq. (4) for various values of $t$. We find two zeros at $t\approx-0.11$ GeV$^2$ and $t\approx-0.39$ GeV$^2$ (see Fig. 1). Of course, there are some ambiguities arising from the above data. Notice, however, that these ambiguities are not large enough to change the above results qualitatively.[22]

Therefore, we can conclude as follows:

(1) Within the present approximation, the left-hand side of Eq. (4) gives us two zeros around $t=-0.11$ GeV$^2$ and $t=-0.39$ GeV$^2$, which suggests[23] that $\alpha_R(t)=0$ at $t\approx-0.39$ GeV$^2$, as well as that $C_R(t)$ also vanishes at $t\approx-0.11$ GeV$^2$ (see Fig. 1).

(2) Thus, the $R$ trajectory passes through spin zero and the ghost is shown to be eliminated by a zero.

Secondly we come to question (ii): choice of "sense" or "nonsense". In order to answer this question, we derive a similar sum rule for

$$B^{(+)}(\nu,t)\equiv\tfrac{1}{4}[B_{K^-p}(\nu,t)$$
$$+B_{K^+p}(\nu,t)-B_{K^-n}(\nu,t)-B_{K^+n}(\nu,t)].$$

Under the weaker assumption that there are no other singularities except for the $R$ pole in the $J$ plane for $\alpha\geqslant0$ at $t=0$, separating

$$B^{(+)}(\nu,t)=B_K^R(\nu,t)+B^{(+)\prime}(\nu,t) \qquad (11)$$

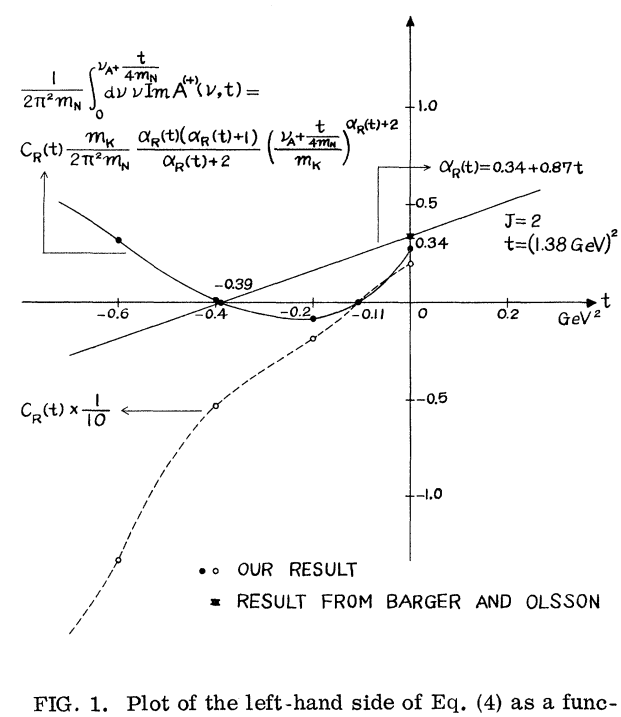


FIG. 1. Plot of the left-hand side of Eq. (4) as a function of $t$ in the narrow-width approximation. The helicity-nonflip residue function evaluated from Eq. (4) for $\alpha_R(t)=0.34+0.87t$ is also shown.

with

$$\mathrm{Im}B_K^{\;R}(\nu,t)=\alpha_R(t)[\alpha_R(t)+1]\,(\tilde{C}_R(t)/m_K^{\;2})(\nu/m_K)^{\alpha_R(t)-1}, \tag{12}$$

we obtain

$$\int_0^\infty d\nu\,\mathrm{Im}B^{(+)\prime}(\nu,t)=0. \tag{13}$$

Similarly as before, Eq. (13) reduces to

$$\frac{1}{4}\frac{g_\Lambda^{\;2}}{4\pi}-\frac{1}{4}\frac{g_\Sigma^{\;2}}{4\pi}+\frac{1}{4}\frac{g_{Y_0^*}^{\;2}}{4\pi}-\frac{1}{4}\frac{g_{Y_1^*}^{\;2}}{4\pi}\left[\frac{\{(m_N+m_{Y_1^*})^2-m_K^{\;2}\}\{(m_N-m_{Y_1^*})^2-m_K^{\;2}-2m_Nm_{Y_1^*}\}}{6m_{Y_1^*}^{\;2}}+\frac{t}{2}\right]$$

$$+\frac{m_N}{2\pi^2}\int_{m_K}^{\nu_A}d\nu_L\,\mathrm{Im}B^{(+)}\left(\nu_L+\frac{t}{4m_N},t\right)=\frac{m_N}{2\pi^2m_K}[\alpha_R(t)+1]\tilde{C}_R(t)\left(\frac{\nu_A+t/4m_N}{m_K}\right)^{\alpha_R(t)}, \tag{14}$$

with

$$\frac{m_N}{2\pi^2}\int_{m_K}^{\nu_A}d\nu_L\,\mathrm{Im}B^{(+)}(\nu_L+t/4m_N,t)$$

$$=\sum_{I,l_\pm}\pm C_I\frac{\Gamma_{l_\pm}^{\;\mathrm{el}}}{8q_{l_\pm}^{\;3}}\left[\{(M_{l_\pm}-m_N)^2-m_K^{\;2}\}P_{l\pm1}{}'(Z_{l_\pm})-\{(M_{l_\pm}+m_N)^2-m_K^{\;2}\}P'(Z_{l_\pm})\right]. \tag{15}$$

In order to discriminate between "sense" or "nonsense," we shall investigate possible zeros of the left-hand side of Eq. (14). It should be noted, in this case, that the $\Lambda$ and $\Sigma$ make significant contributions to the left-hand side, but other ambiguities arising from higher resonances become smaller. We consider Case (a) and Case (b) separately for the $KN\Lambda$ and $KN\Sigma$ couplings. The left-hand side of Eq. (14) is plotted for Case (a) and for Case (b) in Fig. 2.

Thus, we would like to conclude as follows:

Case (a). If we use Zovko's values[18] for $g_\Lambda$ and $g_\Sigma$, the left-hand side of Eq. (14) becomes zero around $t\approx-0.48$ GeV$^2$ [24] (see Fig. 2). Therefore, there is a chance that the $R$ trajectory could choose "sense,"[7] within the errors of Zovko's values.[18]

Case (b). For values[19] obtained from SU(3), we plotted the left-hand side for three choices

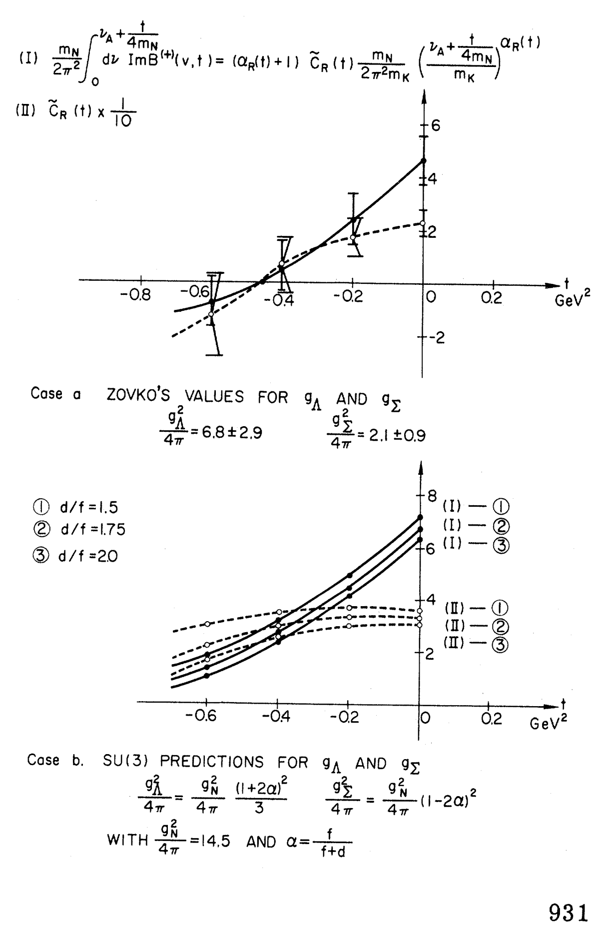


FIG. 2. For Case (a) and Case (b) the left-hand side of Eq. (14) is plotted as a function of $t$ in the narrow-width approximation. The helicity-flip residue functions evaluated from Eq. (14) for $\alpha_R(t)=0.34+0.87t$ are also shown. For Case (a) the bars show the variation within the errors of Zovko's values for $g_\Lambda$ and $g_\Sigma$. For Case (b) the results from three choices of the $d/f$ ratios are shown.

of $d/f$ ratios: $d/f = 1.5$, 1.75, and 2.0. We can find no zero in the interval $-0.8\ \mathrm{GeV}^2 \leq t \leq 0$ for any values between $2.0 \geq d/f \geq 1.5$ (see Fig. 2). Therefore, if Kim's analysis[19] that $g_\Lambda$ and $g_\Sigma$ are consistent with SU(3) predictions with $d/f \approx 1.5$ is correct, the $R$ trajectory would favor "nonsense."[6]

In conclusion, we would like to point out that the superconvergent sum rule for an amplitude controlled by only one Regge pole is particularly useful to pin down the zeros which are characteristic in the Regge formula and for tracing out the trajectory. In the present paper we have calculated the low-energy integrals of Eqs. (4) and (14) in the narrow-width approximation. If the phase shifts of $KN$ and $\bar{K}N$ scattering become available up to high energies (a few GeV) in the near future, then the superconvergent sum rules Eqs. (4) and (14) will allow us to argue more precisely on the $R$ trajectory and the ghost-eliminating zeros in the Regge amplitudes. We hope that careful investigation of the $\Lambda$ and $\Sigma$ couplings as well as detailed analysis of low-energy resonance parameters and phase shifts will be made as early as possible.

Detailed analysis including further applications will be published elsewhere.

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[8]F. Arbab, N. F. Bali, and J. Dash, University of California Lawrence Radiation Laboratory Report No. UCRL-17325, 1967 (unpublished). These authors attempted to discriminate between these ghost-eliminating mechanisms by fitting the charge-exchange reactions involving the $\rho$ and $R$ trajectories (like $\pi^- + p \to \pi^0 + n$, $\pi^- + p \to \eta + n$, and $K^- + p \to \bar{K}^0 + n$) in terms of the $\rho$ and $R$ Regge poles. However, it has not been fully successful because the theory involves a number of parameters.

[9]K. Igi and S. Matsuda, Phys. Rev. Letters 18, 625 (1967); "Existence of Other $\rho$-Regge Singularities in the Complex $J$-Plane" (to be published). Hereafter we refer to these papers as I and II.

[10]K. Igi and S. Matsuda, Phys. Rev. (to be published), referred to as III hereafter.

[11]R. Dolen, D. Horn, and C. Schmid, Phys. Rev. Letters 19, 402 (1967).

[12]V. Singh, Phys. Rev. 129, 1889 (1963). See also K. Igi, Phys. Rev. 130, 820 (1963).

[13]As was discussed in III, experiments are consistent with weak cut discontinuities even if they exist for the $S = 0$ and $I = 0$ quantum numbers. We will, therefore, assume $R$ dominance for the present case.

[14]The factor $\alpha_R(t)$ in Eq. (2) is introduced to eliminate the ghost and $(\alpha_R + 1)$ is included to emphasize the zero at $\alpha_R = -1$. A nonflip residue $C_R(t)$ is therefore defined to be dimensionless and is regular except for additional zeros at $\alpha_R = -2, -3, -4, \cdots$.

[15]A. A. Logunov, L. D. Soloviev, and A. N. Tavkhelidze, Phys. Letters 24B, 181 (1967).

[16]We chose the value $\nu_A = 2$ GeV (which corresponds to the c.m. energy $W = 2.2$ GeV) for convenience, since the spin and parity are not known yet fo the $\Lambda(2340)$ and $\Sigma(2260)$. This value does not seem to be too low since, choosing $\nu_A = 2.44$ GeV, Yoshimura obtained reasonable values for $C_V(0)$ and $f/d$ ratio for the vector nonet [M. Yoshimura, "Vector Meson Regge Poles and $KN$ Superconvergence Sum Rules" (to be published)].

[17]In order to calculate the Born terms for the $Y_0{}^*$ and $Y_1{}^*$ we assumed the following effective Lagrangians for the $(Y_0{}^*KN)$ and $(Y_1{}^*KN)$ vertices:

$$L = g_{Y_0{}^*} \bar{Y}_0{}^* N \bar{K} + \mathrm{H.c.},$$

$$L = g_{Y_1{}^*} (\bar{Y}_1{}^*)_\mu N \partial_\mu \bar{K} + \mathrm{H.c.}$$

[18]N. Zovko, Phys. Letters 23, 143 (1966).

[19]J. K. Kim, Bull. Am. Phys. Soc. 12, 506 (1967). According to Kim's analysis, careful treatment of unphysical contributions leads to the result that $g_\Lambda$ and $g_\Sigma$ are consistent with SU(3) predictions with $d/f \approx 1.5$. The authors are thankful to Professor K. Nishijima for the informing us of Kim's work.

[20]R. L. Warnock and G. Frye, Phys. Rev. 138, B947 (1965).

[21]A. H. Rosenfeld et al., Rev. Mod. Phys. 39, 1 (1967).

[22]Fortunately, contributions from the $\Lambda$, $\Sigma$, $Y_0{}^*$, and $Y_1{}^*$ are small. The $\Lambda(1670)$ and $\Sigma(1660)$ whose elasticities are not known exactly at present also give a very small contribution to the integral due to a small kinematical factor for the $\Lambda(1670)$ and a preliminary small elasticity ($\sim 0.15$) for the $\Sigma(1660)$. Dominant contributions come from the $\Lambda(1520)$, $\Lambda(1820)$, $\Lambda(2100)$, $\Sigma(1770)$, and $\Sigma(2035)$.

[23]We can give plausible arguments that our choice of $\alpha_R(t = -0.39\ \mathrm{GeV}^2) = 0$ and $C_R(t = -0.11\ \mathrm{GeV}^2) = 0$ is preferable to other possible choices: (i) If we take a straight-line approximation $\alpha_R(t) = 0.34 + 0.87t$, deduced from $\alpha_R(t = -0.39\ \mathrm{GeV}^2) = 0$ and $\alpha_R(0) = 0.34 \pm 0.03$ [V. Barger and M. Olsson, Phys. Rev. Letters 18, 294 (1967)], it extrapolates to $\alpha = 2$ at $t = (1.38\ \mathrm{GeV})^2$ which is close to the $A_2$-meson mass, $m_{A_2}{}^2 = (1.30\ \mathrm{GeV})^2$.

(ii) If one is willing to take the "exchange degeneracy hypothesis" [R. C. Arnold, Phys. Rev. Letters 14, 657 (1965)], then it is reasonable to expect a zero of $C_R(t)$ around $t=-0.11$ GeV$^2$ just like the zero of $C_\rho(t)$, which is required to explain the crossover phenomenon of $\pi^{\pm}p$ differential cross sections [R. J. N. Phillips and W. Rarita, Phys. Rev. 139, B1336 (1965)]. (iii) Use together of an independent sum rule for $A^{(+)\prime}(\nu,t)/\nu$ [the type of Eq. (19) in III], with Eq. (4), will give us an independent determination of $\alpha_R(t)$ and $C_R(t)$. Discussions will be given elsewhere.

[24]Within the errors of Zovko's values (Ref. 18) we can find a zero in the interval $-0.36$ GeV$^2>t>-0.63$ GeV$^2$.

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# NEW TEST OF $T$ INVARIANCE IN $p$-$p$ SCATTERING*

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A test of $T$ invariance in $p$-$p$ triple scattering has been performed at 430 MeV. The difference in two polarizations which are identical if $T$ is valid was found to be 0.0006 ± 0.0028. The $T$-nonconserving amplitude was found to be less than $\frac{1}{2}$% of the $T$-conserving amplitude. The $T$-nonconserving phase was found to be less than 0.06 of its maximum value.

This Letter describes a test of time-reversal invariance in strong interactions using $p$-$p$ triple scattering at 430 MeV. The test is to compare the final polarization in Figs. 1(a) and 1(b) or, as seen in the laboratory, in Figs. 2(a) and 2(b). The two proton-spin and -momentum configurations shown in Fig. 1 in the $p$-$p$ center of mass are related by time reversal. The Lorentz transformation to the lab frame as shown in Fig. 2 leaves the incident spin directions unchanged, rotates the final momentum forward ($\bar{\theta}\to\theta$), and changes the angle between final spin and momentum [$\chi_f\to\chi_f+\theta$ for Fig. 1(a) and $\chi_i\to\chi_i-\theta$ for Fig. 1(b)].[1]

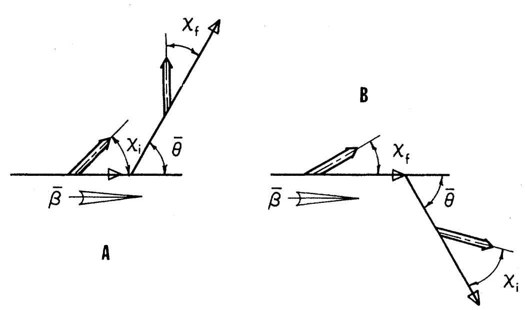


FIG. 1. Two configurations for initial- and final-state proton spin components (double arrows) and momenta for $p$-$p$ scattering through an angle $\theta$ in the center of mass. The angles $\chi_i$ and $\chi_f$ are arbitrary. The velocity of the center of mass is in the direction of $\bar{\beta}$. Reversing spins and momenta in $A$ gives $B$ after a space rotation.

Thus for arbitrary angles $\chi_i$ and $\chi_f$ the final-state polarization in Fig. 2(a) along the direction $\chi_f+\theta$ must be equal to the final-state polarization in Fig. 2(b) along the direction $\chi_i-\theta$ if time-reversal invariance is valid. The argument can be phrased in terms of the Wolf-

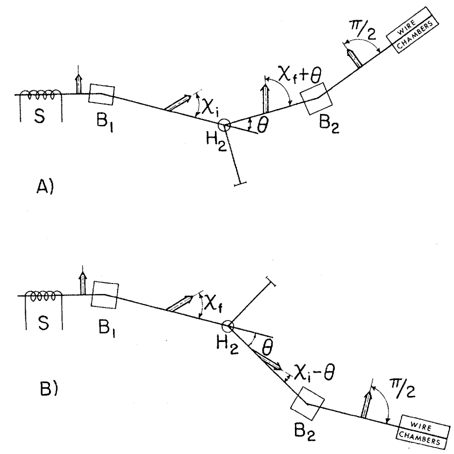


FIG. 2. Transforming the spins and momenta of Figs. 1(a) and 1(b) into the laboratory gives the spins and momenta shown in Figs. 2(a) and 2(b) respectively. The solenoid magnet $S$ and the first bending magnet $B_1$ prepared the initial spin direction $\chi_i=\chi_f=45°$. The scattering angle $\theta=30°$ was chosen. In geometry $A$ the magnet $B_2$ precessed the spin from 75° to 90° so that it could be analyzed by the wire chambers; in $B$ a precession from 15° to 90° was required. An "up-down" asymmetry resulted from this polarization component.