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Charles B. Chiu and Jerome Finkelstein

January 10, 1967

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ABSTRACT

With some simple assumptions on the behavior of the discontinuity of the cut in the angular momentum plane (associated with the simultaneous exchange of the ρ and the P trajectories), we show that the high energy πN charge-exchange polarization is expected to have the following features: It has a zero near the forward direction [say with $|t| < 1\ (\mathrm{GeV}/c)^2$], whose position depends on the energy. In the t region between $t = 0$ and the position of this zero, the polarization decreases with energy and beyond this region the polarization increases with energy. Consistent with all the other available πN data, we show that the present charge-exchange polarization and differential cross section data can be fitted by a small ρ-P cut contribution in addition to the ρ contribution, with the ρ trajectory choosing either sense or nonsense states at $\alpha_\rho = 0$. If the ρ trajectory chooses the nonsense state, the crossover effect observed in the high energy $\pi^\pm p$ differential cross section can be explained naturally without introducing a zero in the helicity nonflip amplitude of the ρ, which is needed in the present pure Regge pole model.

It is perhaps well-known by now that the high energy $\pi^- p$ charge-exchange differential cross section near the forward direction can be explained successfully in terms of the t-channel $(\pi\pi \rightarrow N\bar{N})$ ρ trajectory exchange.[1] If the $\pi^- p$ charge-exchange amplitude near the forward direction is indeed dominated by the ρ amplitude, then the phase of the helicity-flip and nonflip amplitudes should be about the same, which implies that the charge-exchange polarization should be small. The latest available data on charge-exchange polarization by Bonamy et al. at CERN[2] is shown in Fig. 1. The average polarization for all points with $0.24 \geqslant |t| \geqslant 0.04$ $(\mathrm{GeV/c})^2$ is $16 \pm 3.5\%$ at 5.9 GeV/c and is $14 \pm 4.5\%$ at 11.2 GeV/c. The polarization data are certainly not consistent with zero. To explain the observed magnitude, one has to explicitly include, in addition to the ρ amplitudes, some background contributions. Various possibilities within the Regge model have been proposed ever since the report of the preliminary results on nonzero charge-exchange polarization at 5.9 GeV/c became available.[3] Among these proposals, there are, for example, (a) interference between the ρ amplitude and the tail of direct channel resonances,[4] (b) interference between the ρ amplitude and the contribution of another trajectory,[5] (c) the possibility that the ρ trajectory is complex for negative t,[6] and (d) interference between the ρ and the contribution of the cut associated with ρ and Pomeranchon exchange.[7,8]

Let us examine the energy dependence of the charge-exchange polarization that would be expected from these various possibilities. Define the t-channel helicity nonflip and the flip amplitude A' and B such that at high energy one has[9] for the differential cross section:

$$\frac{d\sigma}{dt} = \frac{1}{16\pi E^2} \left(k_a |A'|^2 - t\, k_b |BE|^2\right) , \tag{1}$$

and for the polarization

$$P = \frac{-2(k_a k_b |t|)^{\frac{1}{2}} \quad \mathrm{Im}(BE)^* A'}{\left[k_a |A'|^2 - t\, k_b |BE|^2\right]} ; \tag{2}$$

where E is the total lab energy of incident pions. Factors k_a and k_b are given in Ref. 9. For $E \gg M$ and $|t| \ll 4M^2$, $k_a \approx 1$ and $k_b \approx 1/4M^2$. If a single trajectory α is exchanged, $A' \propto E^{\alpha}$ and $B \propto E^{\alpha-1}$.

(a) If we assume that the tails of direct channel resonances behave at large E as x/E, then the polarization due to interference between these tails and the ρ falls at high energy as $E^{-(\alpha_\rho+1)}$. If the elasticities x are not constant but are themselves decreasing functions of E , then the polarization disappears even faster.

(b) Let us now consider the polarization produced by the exchange of two t-channel trajectories of $\alpha_1(t)$ and $\alpha_2(t)$, where we take $\alpha_1 > \alpha_2$. If the contribution to the amplitude of α_1 is larger than that of α_2, the decrease in the polarization is given by $P \propto E^{\alpha_2-\alpha_1}$. On the other hand, if there is an energy region in which α_2 dominates, then in that region the polarization will increase: $P \propto E^{\alpha_1-\alpha_2}$. Note also that Bose statistics require any trajectory which can contribute to $\pi^- p$ charge-exchange scattering to have odd signature, which implies for small $|t|$, $P \propto t^{\frac{1}{2}} \sin \frac{\pi}{2} (\alpha_2 - \alpha_1)$.

(c) The proposal that the ρ trajectory may be complex for certain negative values of t may be thought of as an extension of the idea of having two trajectories, which would require $\alpha_1(t)$ and $\alpha_2(t)$ instead of being real, to be complex conjugates of each other for negative t. It is possible to get the polarization to be independent of energy with this proposal.

(d) The energy dependence of the polarization expected from interference between the ρ Regge pole and the ρ-P cut has been studied recently in two papers, by Delany and Gross[7] and by Muzinich and Teplitz.[8] In this note we wish to examine in further detail the dependence of this polarization upon both energy and momentum transfer. We will point out that the simplest assumptions about the discontinuity of the cut, such as were made in Ref. 7 and Ref. 8 imply that the polarization should decrease as the energy increases (contrary to the expectation of these authors), but not so fast as to be in conflict with the experimental data.[2]

We will first derive a general expression for the polarization due to the interference between the ρ pole and the ρ-P cut; we ignore the contributions of higher-order cuts as hopefully negligible and certainly complicated. Since in our fits the ρ-P cut contribution to the differential cross section is small, it is plausible that the contributions of higher order cuts are even smaller. We then discuss in some detail the expected high-energy behavior of this polarization. Finally, we present two sample fits to the existing charge-exchange polarization data consistent with all other relevant πN data.

Associated with the simultaneous exchange of the ρ and the Pomeranchuk Regge poles in the t-channel, there will be a cut in the angular momentum plane, whose rightmost extension α_c is given by[10]

$$\alpha_c(t) = \max_{\substack{\sqrt{|t'|} + \sqrt{|t''|} \geq \sqrt{|t|} \\ t', t'' \leq 0}} [\alpha_\rho(t') + \alpha_P(t'') - 1] \ . \qquad (3)$$

Assuming that the trajectories are linear (for small t) and that $\alpha_P(0) = 1$, this gives

$$\alpha_c(t) = \alpha_\rho(0) + \frac{\alpha_\rho' \, \alpha_P'}{\alpha_\rho' + \alpha_P'} \, t \ . \qquad (4)$$

The value of α_P' is at present not well determined, although earlier fits[11] to the πN data suggest that α_P' lies between 0 and $0.4 (\mathrm{GeV}/c)^{-2}$. The fits which will be discussed later have $\alpha_P' = 0.35$, which together with the value $\alpha_\rho' = 1.0$ gives $\alpha_c' = 0.26$. From Eq. (4) $\alpha_c' < \alpha_\rho'$, for any value of α_ρ'. The conclusion of this paper is essentially independent of the value of α_P' used. For instance α_P' could be zero, or it could be as large as α_ρ'.

Performing the Mandelstam-Watson-Sommerfeld transformation on the partial wave expansion for the helicity nonflip and flip amplitudes and keeping only the ρ pole term and the contribution of the cut, we get for large E ,

$$A'(E,t) = \beta^+(t)\left(\tan \frac{\pi\alpha_\rho}{2} + i\right) \left(\frac{E}{E_0}\right)^{\alpha_\rho(t)} + \int_{-L}^{\alpha_c} dJ \, D^+(J,t)\left(\tan \frac{\pi J}{2} + i\right)\left(\frac{E}{E_0}\right)^J , \qquad (5)$$

and a similar expression for $B(E,t)$ with β^+ and D^+ replaced by β^- and D^-. The symbols $\beta^\pm$ are the reduced residue functions for the pole terms and $D^\pm$ the reduced discontinuities across the cut in the J plane of the t-channel helicity amplitudes. We assume that the background integral can be pushed at least to $J = -\frac{1}{2}$, i.e $L \geqslant \frac{1}{2}$. For simplicity, we will ignore contributions to the integral in Eq. (5) from the lower limit, these contributions being of background size. This is equivalent to evaluating the integral as if $L = \infty$. From hermitian analyticity, $D(J,t)$ is real for $t < 4m_\pi^2$.

We assume that the discontinuites $D^\pm(J,t)$ can be written as

$$D^\pm(J,t)(\tan\frac{\pi J}{2} + i) = \frac{D^\pm(J,t)}{\cos\frac{\pi J}{2}} \exp\left[-\frac{i\pi}{2}(J-1)\right]$$

$$= \sum_{n=0}^{\infty} C_n^\pm(\alpha_c - J)^{\gamma+n} \exp\left[-\frac{i\pi}{2}(J-1)\right]. \tag{6}$$

The $C_n^\pm$'s are given by

$$C_n^\pm(t) = \frac{(-)^n}{n!} \frac{\partial^n}{\partial J^n} \left[\frac{D^\pm(J,t)}{\cos\frac{\pi J}{2}(\alpha_c - J)^\gamma}\right]_{J=\alpha_c}, \quad \text{and are real.}$$

If the cut is logarithmic as expected,[12] then $\gamma = 0$. From (5) and (6), the contribution A_c' of the cut to A' is

$$A_c'(E,t) = \left(\frac{E}{E_0}\right)^{\alpha_c}(\tan\frac{\pi\alpha_c}{2} + i) \sum_{n=0}^{\infty} \frac{\Gamma(n+\gamma+1)\,\overline{C}_n^{\,+}}{\left[\ln\left(\frac{E}{E_0}\right) - i\frac{\pi}{2}\right]^{\gamma+n+1}}, \tag{7}$$

and a similar expression for B_c with $\bar{C}_n^+$ replaced by $\bar{C}_n^-$, where $\bar{C}_n^{\pm} = C_n^{\pm} \cos \pi\alpha_c/2$. A formula similar to Eq. (7) has been derived in Ref. 7. To order $1/\ln E$ the phase of A_c' is given by the first (n = 0) term:

$$\left[\text{phase of } A_c'\right] = \frac{\pi}{2}(1 - \alpha_c) + (\gamma + 1)\tan^{-1}\frac{\pi}{2\ \ln\frac{E}{E_0}} + O\left(\frac{1}{\ln^2\frac{E}{E_0}}\right). \tag{8}$$

It is crucial for the discussion below that to order $1/\ln E$ the phase of A_c' does not depend on the details of the discontinuity of the cut (the set of C_n); thus to this order the phase of A_c' and B_c are the same. Since for small t, α_ρ and α_c are nearly equal, the term

$$(\gamma + 1)\tan^{-1}\frac{\pi}{2\ \ln\frac{E}{E_0}}$$

is really the leading term of the phase difference between the cut and the pole, and so must not be neglected; a similar point has been made in Ref. 7.

At energies so large that $\bar{C}_n/\bar{C}_0 \ll (\ln E)^n$, A_c' will be well approximated by the first term of Eq. (7). The energy dependence of the polarization will then be given by

$$P \sim \frac{\left(\frac{E}{E_0}\right)^{\alpha_c-\alpha_\rho} \sin\frac{\pi}{2}\left[\alpha_\rho - \alpha_c + \frac{2}{\pi}(\gamma + 1)\tan^{-1}\frac{\pi}{2\ \ln(E/E_0)}\right]}{\left[\ln^2\left(\frac{E}{E_0}\right) + \left(\frac{\pi}{2}\right)^2\right]^{(\gamma+1)/2}}. \tag{9}$$

In the small $|t|$ region of interest, for $\gamma = 0$, and a very wide range of E, the energy variation of the polarization is controlled primarily by the sine factor; see Fig. 2 for a graphical illustration of this. We should point out that a simplified version of Eq. (9) has been given by Delany and Gross in Ref. 7. However, when they discussed the energy dependence of the polarization in this small $|t|$ region of interest, they only considered the $E^{\alpha_c - \alpha_\rho}$ factor, neglecting the logarithmic denominator and the controlling sine factor, and so were led to conclude that the polarization should increase with energy.

We can define the effective position of the cut, α_{eff}, so that

$$\text{phase of } A_c' = \text{phase of } \left(\tan \frac{\pi\alpha_{eff}}{2} + i\right) . \tag{10}$$

From either Eq. (8) or Eq. (9), we can see that at large energy,

$$\alpha_{eff}(E,t) = \alpha_c(t) - \frac{2(\gamma + 1)}{\pi} \tan^{-1} \frac{\pi}{2\, \ln(E/E_0)} \sim \alpha_c(t) - \frac{\gamma + 1}{\ln(E/E_0)} . \tag{11}$$

We can now see what the energy dependence for the polarization will be, for E large enough so that Eq. (11) is valid. For given E and small enough $|t|$, α_{eff} will be less than α_ρ; as the energy increases α_{eff} will rise, the phase of the cut and the pole will get closer together, and the polarization will fall. At larger $|t|$, where $\alpha_{eff} > \alpha_\rho$, the polarization will increase. These remarks are illustrated in Fig. 3.

Between the region in t in which the polarization will rise, and the region in which it will fall, is the point $\alpha_{eff} = \alpha_\rho$; at this point the polarization is zero. That is, there is a moving zero in the polarization at the value of t satisfying

$$\alpha_c(t) - \alpha_\rho(t) = \frac{2(\gamma + 1)}{\pi} \tan^{-1} \frac{\pi}{2\,\ln(E/E_0)} \sim \frac{\gamma + 1}{\ln(E/E_0)} . \tag{12}$$

We turn now to the question of whether energies of 5.9 and 11.2 GeV are sufficiently large to warrant the approximation of keeping only the first term in Eq. (7). The answer to this question depends on the derivatives of the discontinuities (the C_n's), which we do not know. From Eq. (5), keeping the first term in Eq. (7) is equivalent to assuming that

$$R^{\pm}(J,t) = \frac{D^{\pm}(J,t)}{\cos\frac{\pi J}{2}\,(\alpha_c - J)^{\gamma}}$$

is independent of J, for those values of J for which $(E/E_0)^{\alpha_c - J}$ is significant. For the energies in question this is perhaps a full unit of angular momentum. We would expect that the contribution to the helicity-flip amplitude B_c should change sign at $J = 0$; however, since B_ρ is large in comparison with A_ρ', we assume that the polarization depends primarily on A_c',[13] and we know of no reason why $R^+(J,t)$ should be zero at $J = 0$. In our ignorance of any detailed explanation of the cut residues, we will make the simplest possible assumption, namely that $R^+(J,t)$ can be approximated by a constant (or at most an exponential) for $-\frac{1}{2} \leqslant J \leqslant \alpha_c$. The exponential can be included by adjusting E_0; for some choice of E_0, $C_n^+ \approx 0$ for $n > 0$, and $A_c{}'$ is given by the first term of Eq. (7). In fact in the fits we will take $E_0 = 1$ GeV and $\gamma = 0$, although these fits to the presently existing charge-exchange polarization data are insensitive to this particular choice. More explicitly the cut contribution is taken to be

$$A_c{}'(E,t) = \frac{\overline{C}_0{}^+(t)\left(\tan\frac{\pi\alpha_c}{2} + i\right)E^{\alpha_c}}{\ln E - i\frac{\pi}{2}} \qquad (13)$$

with

$$\overline{C}_0{}^+ = (\alpha_c + 1)a_0 \exp\left(a_1 t\right),$$

where a_0 and a_1 are the parameters. The factor $(\alpha_c + 1)$ is included here to be uniform with the parameterization for the ρ-amplitude of Case 1 below.

In our numerical analysis, together with the charge-exchange polarization data given in Ref. 2, we include a sample of high energy πN data, such as the differential cross sections and polarizations of $\pi^{\pm}p$ elastic scattering, the $\pi^{\pm}p$ total cross sections and charge-exchange differential cross sections. The exact data included are enumerated in detail in Ref. 11 and 14. We add the ρ-P cut contribution of Eq. (13) to the Regge amplitudes of P, P' and ρ . The P and P' amplitudes are parameterized in the same manner as in Ref. 14. Two different parameterizations[14] for the ρ amplitudes, depending whether the ρ trajectory chooses sense or nonsense at $\alpha_\rho = 0$, are considered. We will proceed to discuss our fits for these two cases separately.

Case 1: The ρ trajectory chooses sense at $\alpha_\rho = 0$. We write

$$A_\rho' = (\alpha_\rho + 1)a_0 \left[(1 + G)\exp(a_1 t) - G\right](\tan\frac{\pi\alpha_\rho}{2} + i)E^{\alpha_\rho}$$

$$B_\rho = \alpha_\rho(\alpha_\rho + 1)b_0 \exp(b_1 t)(\tan\frac{\pi\alpha_\rho}{2} + i)E^{\alpha_\rho - 1} \tag{14}$$

where the term G is introduced here in order to explain the observed sign change in the difference

$$\frac{d\sigma}{dt}\Big|^{\pi^+ p} - \frac{d\sigma}{dt}\Big|^{\pi^- p}$$

in the small $|t|$ region (so-called the crossover effect).[15] With the inclusion of P, P' and ρ plus the cut contribution, we found that, having comparable fits to all the other data as in Ref. 11 and 14, the charge-exchange polarization data can be adequately fitted. The fit for the charge-exchange polarization is illustrated in Fig. 1. The parameters for the cut and the ρ contributions are tabulated in Table I. For comparison, we include in Table I also the corresponding parameters for the ρ amplitudes given in Ref. 11. Figure 4a shows the prediction of charge-exchange polarizations at 5.9, 11.2, 22, 50 and 100 GeV/c over a large t interval given by this solution. We illustrate in Fig. 5a our fit to the charge-exchange differential cross-section data[16] at 13.3 GeV/c and the contributions of the individual terms. Notice the dominance of the ρ contribution, except near the dip where the contribution due to the interference between the ρ and the cut is significant.

Case 2: The ρ trajectory chooses nonsense at $\alpha_\rho = 0$. We parameterize here

$$A_\rho' = \alpha_\rho(\alpha_\rho + 1)a_0 \exp(a_1 t)\left(\tan\frac{\pi\alpha_\rho}{2} + i\right)E^{\alpha_\rho}$$
$$B_\rho = \alpha_\rho(\alpha_\rho + 1)b_0 \exp(b_1 t)\left(\tan\frac{\pi\alpha_\rho}{2} + i\right)E^{\alpha_\rho - 1} . \tag{15}$$

In this case, the crossover effect described above between $\pi^\pm p$ differential cross sections is naturally explained by the change in sign in the sum of Im A_ρ' and Im A_c' . At $\alpha_\rho = 0$, $A_\rho' = B_\rho = 0$, the sizable charge-exchange differential cross sections observed near the dip is filled by the cut contribution. The cut and the ρ parameters for this case are also tabulated in Table I. To give a qualitative feeling how the crossover effect and the bottom of the dip in charge-exchange differential cross sections are fitted, we present in Table II a comparison in χ^2 associated with the high energy $\pi^\pm p$ and charge-exchange differential cross section data for Case 1 and 2 and the corresponding values in Ref. 11. The fit to the charge-exchange polarization data for this case is also illustrated in Fig. 1. Its prediction on charge-exchange polarization for a larger t interval at 5.9, 11.2 and higher energies is illustrated in Fig. 4b. Our fit to the charge-exchange data at 13.3 GeV/c and the contributions of the individual terms for this case are illustrated in Fig. 5b.

We summarize that, following the approach of Ref. 7, we have presented a general expression for the energy dependence of charge-exchange polarization due to ρ and ρ-P cut contribution. With the assumption that the discontinuity is constant and the charge-exchange polarization is mainly due to interference between A_c' and B_ρ, the charge-exchange polarization should behave as follows. There is a zero whose location in momentum transfer moves with energy. In the region where $|t|$ is less than the position of this zero, the polarization decreases with increase of energy, and in the larger $|t|$ region the polarization increases with energy. The cut contribution needed to explain the existing charge-exchange polarization data in the small $|t|$ region is quite small as illustrated in Fig. 5 (also for example at 11.2 GeV/c, $t = 0$, $A_c' = 0.13A_\rho'$ for Case 1), although its contribution in the large $|t|$ region, in particular near the dip region could be substantial. With the introduction of this cut contribution, the data can be consistently fitted with the ρ trajectory choosing either sense or nonsense states at $\alpha_\rho = 0$. With the ρ choosing nonsense at $\alpha_\rho = 0$, the crossover effect in $\pi^\pm p$ differential cross section can be explained naturally without introducing a zero in A_ρ'. The assumptions involved in this paper seem to be adequate in explaining the existing πN data. However, their validity remains to be further tested experimentally.

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FOOTNOTES AND REFERENCES

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13. Although the $\pi^{\pm}p$ elastic polarization in the same t and energy region is also observed to be ~15%, the situation there is much more complicated because P and P' trajectories can also be exchanged, and because the cut contribution to the polarization will presumably be coming from B_c rather than A_c'. It is thus conceivable that cuts would not contribute as significantly to elastic polarization as they might to charge-exchange polarization. It has been shown in Ref. 11 that the elastic polarization can be reasonably fitted with Regge pole amplitudes alone.

14. C. B. Chiu, S. Y. Chu, L. L. Wang, Regge Pole Model for the Secondary Maxima in πN and NN Scattering and the No-Compensation Mechanism, UCRL-17302 (in preparation).

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TABLI I: Parameters for the ρ and the cut

		Case 1	Case 2	Sol. a of Ref. 11
cut	a_0 mbGeV	- 0.63	- 0.71	--
	a_1 GeV^{-2}	1.27	0.04	--
ρ	a_0 mbGeV	1.69	$1.75\ \dot{\alpha}_\rho(0)^{-1}$	1.49
	a_1 GeV^{-2}	1.92	1.90	2.01
	G	1.21	--	1.79
	b_0 mb	28.9	28.5	29.2
	b_1 GeV^{-2}	0.07	0.07	0.12
α_i	$\alpha_\rho(t)$	0.57 + 1.01t	0.57 + 0.99t	0.576 + 1.02t
	$\alpha_P(t)$	1.0 + 0.35t	1.0 + 0.35t	1.0 + 0.34t
	$\alpha_c(t)$	0.57 + 0.26t	0.57 + 0.26t	--

TABLE II. A χ^2 comparison for the differential cross-section data

		Present work		Ref. 11	
Data	Total points	Case 1	Case 2	Sol. a	Sol. b
$\pi^\pm p$ dcs	141	158	164	133	161
$\pi^- p \to \pi^0 n$ dcs	56	86	98	87	87

FIGURE CAPTIONS

Fig. 1. $\pi^- p$ charge-exchange polarization at 5.9 and 11.2 GeV/c. Φ are from Ref. 2. The solid curves and the dashed curves are correspondingly the fits for Case 1 and Case 2.

Fig. 2. The energy dependence of factors in Eq. (9) for $\gamma = 0$, $E_0 = 1$ GeV at $t = -0.15$ $(\text{GeV/c})^2$.

Fig. 3. The functions $\alpha_{eff}(t)$ at 5.9 and 11.2 GeV/c, $\alpha_p(t)$, $\alpha_c(t)$ and $\alpha_\rho(t)$ for Case 1.

Fig. 4. Charge-exchange polarization at 5.9, 11.2, 22, 50 and 100 GeV/c as predicted by (a) Case 1 solution and (b) Case 2 solution.

Fig. 5. The fit to the charge-exchange differential cross section for (a) Case 1 and (b) Case 2 at 13.3 GeV/c. Data points Ξ are from Ref. 16a and Ξ from Ref. 16b. Notice the ρ trajectory contribution is dominating except near $t = -0.6$ $(\text{GeV/c})^2$ where B_ρ vanishes and the ρ-cut interference contribution is comparable to the ρ contribution.

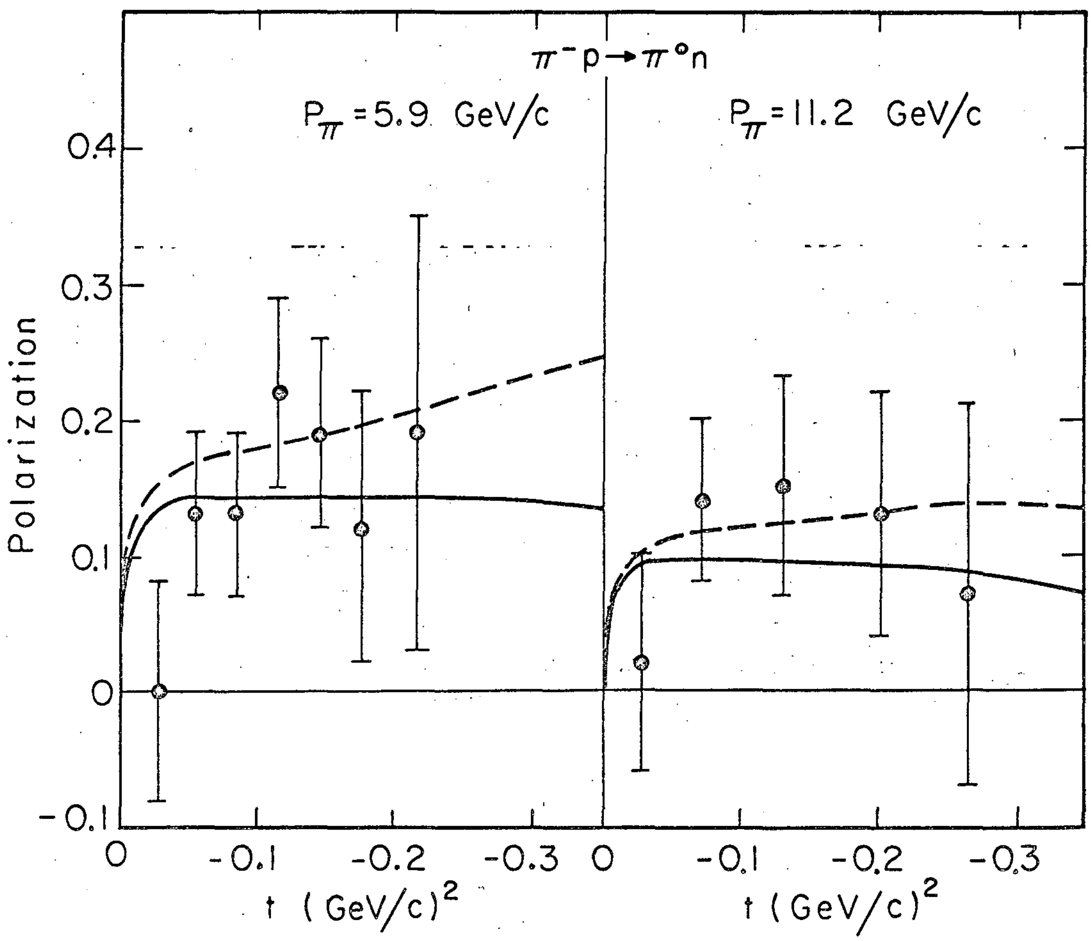

Fig. 1

XBL671-312

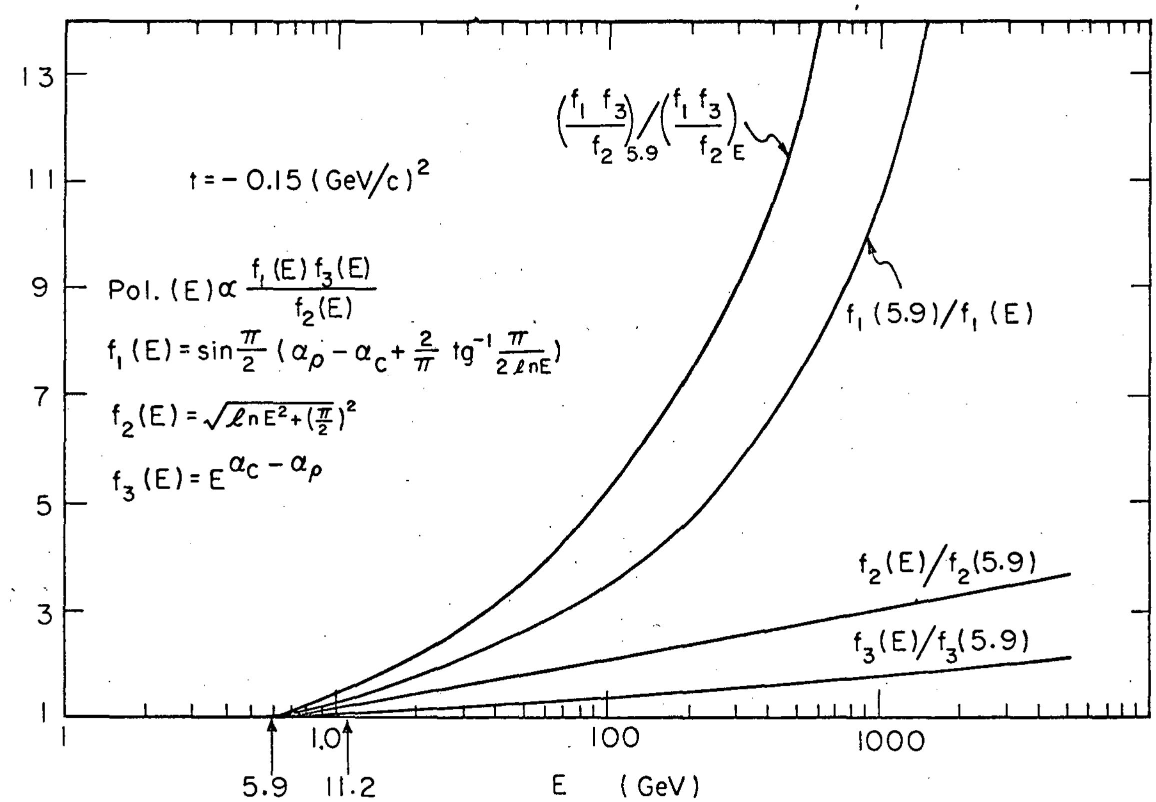

Fig. 2

XBL671-313

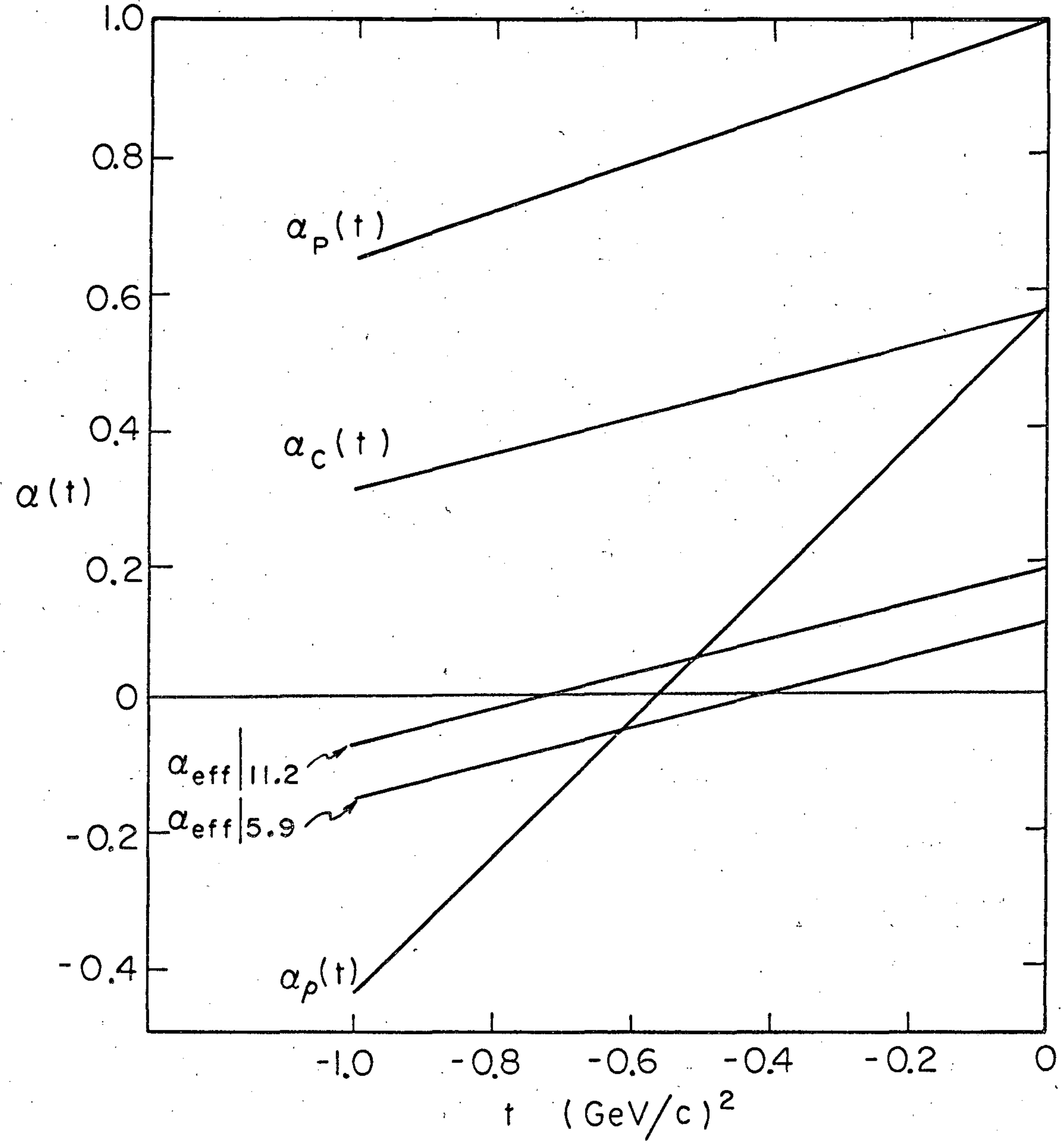

Fig. 3 XBL671-314

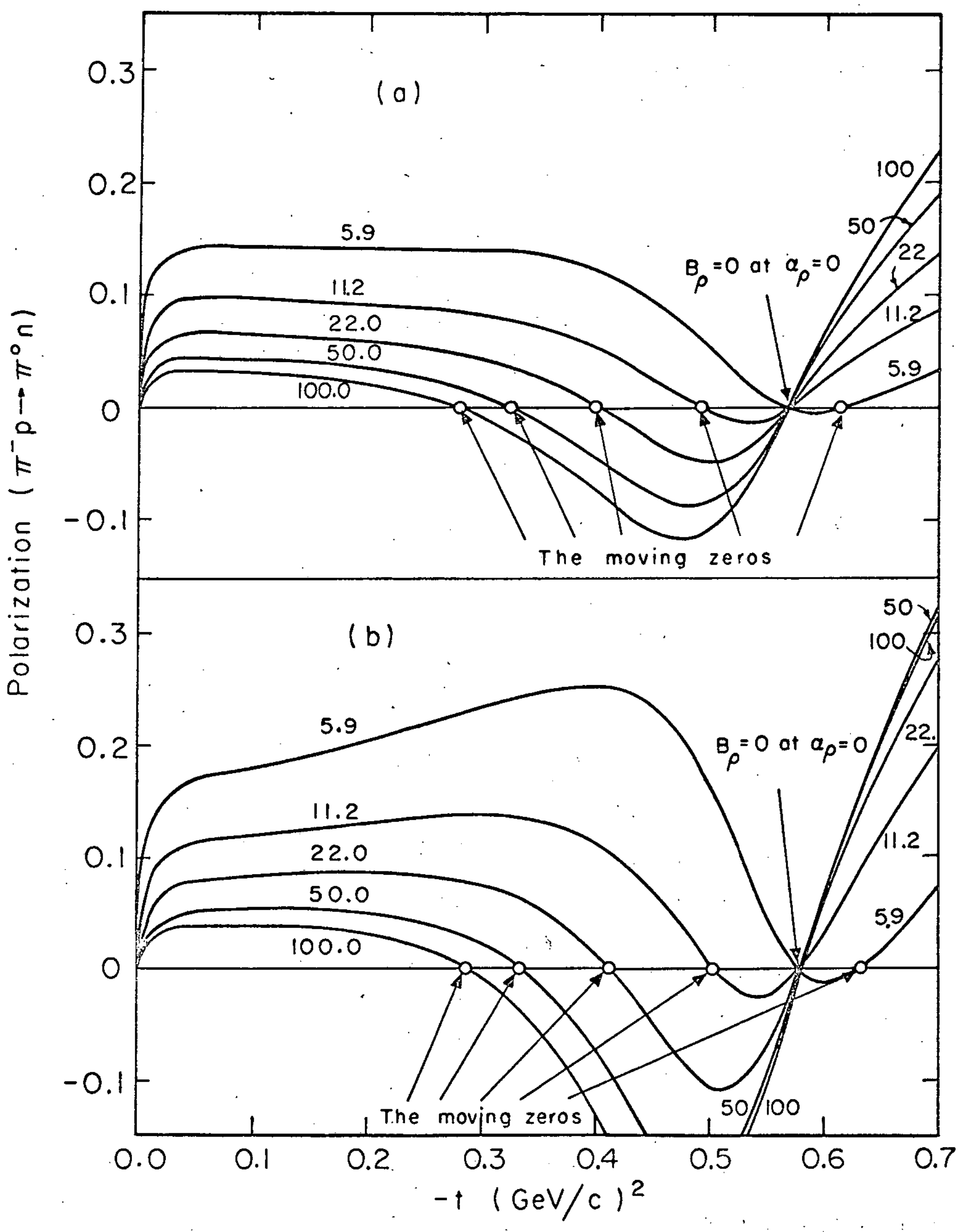

Fig. 4

XBL671-315

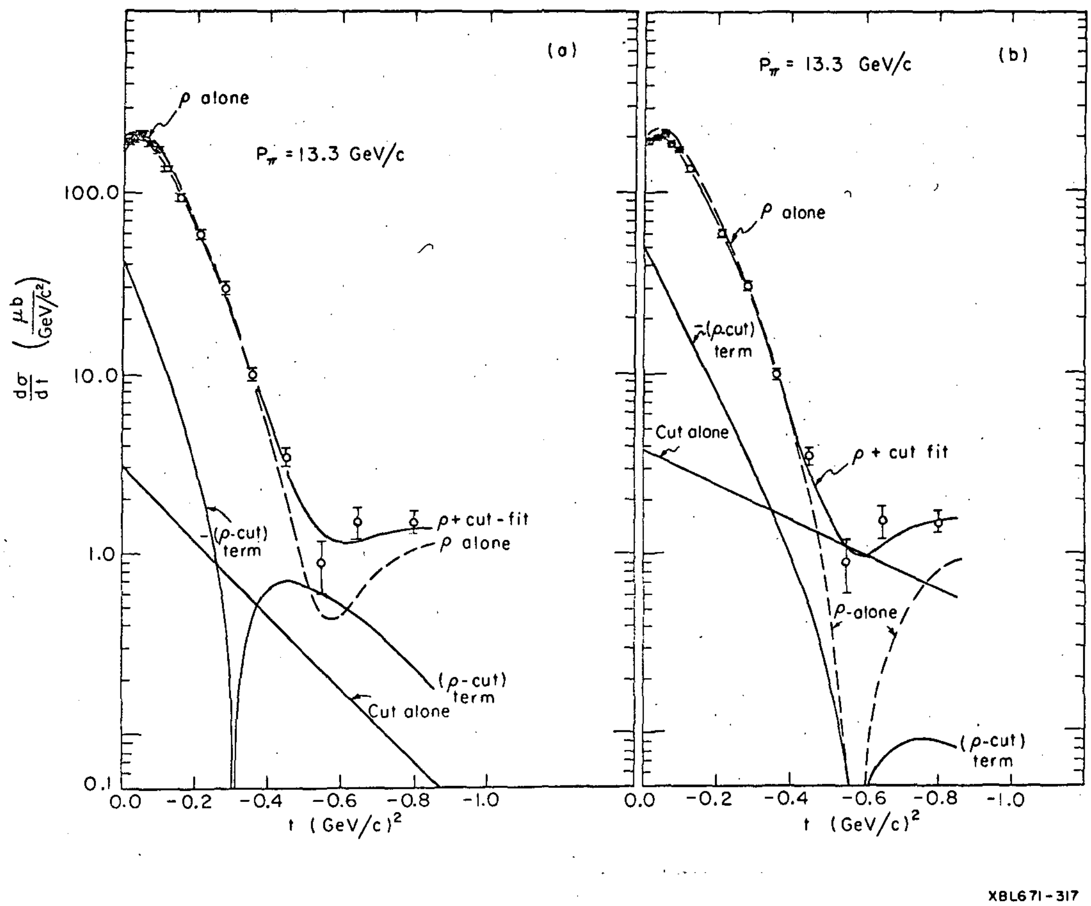
(a)
(b)
P_π = 13.3 GeV/c
ρ alone
ρ + cut - fit
ρ + cut fit
(ρ-cut) term
Cut alone
ρ-alone
$\frac{d\sigma}{dt}$ $\left(\frac{\mu b}{GeV/c^2}\right)$
t (GeV/c)^2
100.0
10.0
1.0
0.1
0.0
-0.2
-0.4
-0.6
-0.8
-1.0
XBL671-317

Fig. 5

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