Evidence for a low-lying unrenormalized vacuum trajectory from NN scattering

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The apparently anomalous energy dependences of the pp and pn elastic polarizations at low energies are used to motivate the existence of a low-lying vacuum trajectory σ with intercept $\alpha_\sigma < 0$ coupling more strongly to baryons than to mesons. The σ is identified with the medium-range attractive NN force in one-boson-exchange models at lower energies. We conjecture that the σ trajectory undergoes some nondiffractive renormalization at $p_{\text{lab}} \gtrsim 30$ GeV/c in analogy with recent models of the Pomeron. The possible presence of other low-lying trajectories is briefly discussed.

I. POLARIZATIONS

Recent measurements[1] of the pp and pn elastic polarizations between 2 and 6 GeV/c have indicated an unexpectedly rapid decrease with energy of the $n=1$, $I=0$ NN amplitude (we denote net s-channel helicity flip by n). We propose that this effect is associated with a low-lying vacuum trajectory σ. We shall relate the σ to a Reggeized continuation of the 0^+ exchange needed in models of low-energy NN scattering.[2] There, the σ provides a crucial medium-range attractive force with a large coupling. In quark language, the σ can be considered as a leading $2q2\bar{q}$ exotic state, coupling primarily to baryons.[3] The σ coupling to mesons is suppressed in this view since it cannot be planar and thus violates the Okubo-Zweig-Iizuka rule.[4] All this is consistent with meson-nucleon scattering, which exhibits a more "normal" behavior than does NN scattering.[5] While previous phenomenological fits[6,7] have invoked a low-lying vacuum trajectory which could be interpreted as the σ coupling to $\pi\pi$ (as distinct from the f' coupling mainly to $K\bar{K}$), the resulting couplings are smaller than those we shall invoke for NN scattering.

More recent work involving low-energy πN polarizations indicates that the σ as parametrized here does in fact contribute to πN scattering consistent with a small $\sigma\pi\pi$ to nonflip $\sigma N\bar{N}$ coupling ratio.[8]

At this point a theoretical remark is in order. As in the case of the leading vacuum trajectory, renormalization effects could be expected to exist for the σ, transforming it into some renormalized trajectory σ_R. In the schemes of Refs. 7 and 9 it is necessary to imagine, e.g., $K\bar{K}$ and $B\bar{B}$ nondiffractive inelastic thresholds renormalizing the leading vacuum trajectory with an intercept below 1 (the f generated by cylinder corrections to the planar bootstrap in Ref. 9 or the bare Pomeron $\hat{P}$ in Ref. 7) into the bare Pomeron of the Gribov calculus with an intercept above 1.[10] Such renormalization is expected above 30 GeV/c on the basis of inelastic $K\bar{K}$,$B\bar{B}$ production data;[11] the same conclusion is reached via detailed two-body phenomenology within this context.[7] Some related high-energy phenomenology for Pomeron renormalization is currently being pursued.[12] The physics behind trajectory renormalization by nondiffractive and diffractive thresholds via unitarity is described in great detail in Ref. 10, to which we refer the interested reader.

Although we would expect the σ to undergo renormalization via this same mechanism, the phenomenological details are undoubtedly complex and will not be broached here. That these effects are non-negligible is indirectly implied by pp polarization data at Serpukhov,[13] which implies a change in energy behavior. This transition cannot be fixed from pp polarization data since none exists between 17.5 and 45 GeV/c; however, we will fit data to 17.5 GeV/c.

In Sec. III we present a canonical renormalization model[7,10] applied to the σ trajectory. Some detailed analysis there leads to the following important point. That is, the association of the mass m_σ phenomenologically determined in low-energy NN scattering with the σ trajectory depends on whether or not one believes that the σ determines a pole in the S matrix at $t=m_\sigma^2$. If it does, m_σ^2 is to be determined by the renormalized trajectory α_σ^R. If it is only regarded as corresponding to a term in an effective Lagrangian, it is to be determined by the unrenormalized trajectory α_σ. The argument generalizes one originally made by Chew[14] in another context.

The unrenormalized trajectory $\alpha_\sigma(t)=-0.4+t$ that we will be using corresponds to a mass parameter $m_\sigma=600$ MeV, sensibly close to values of m_σ derived in NN potentials. A renormalized trajectory $\alpha_\sigma^R(t)=-0.2+t$ would yield a mass parameter $m_\sigma=450$ MeV, corresponding to the true mass of the σ meson if it exists.

We now turn to our analysis of the data. We denote by P the diffractive component of Refs. 7 and 9, or alternatively the more usual sum of a "naive" Pomeron pole P_N with intercept at 1, along with an exchange-degenerate f_{exd}. We denote the $I=1$ Reggeons A_2-ρ by R. Neglecting $n=2$ terms (a point to which we shall return), the polarizations $P(pp), P(pn)$ for elastic pp and pn scattering are then

$$\begin{aligned} |P_{n=0}|^2 P(pp) &\cong -2\,\mathrm{Im}[(P+\sigma-\omega+R)_{n=0} \\ &\qquad \times (P+\sigma-\omega+R)^*_{n=1}], \\ |P_{n=0}|^2 P(pn) &\cong -2\,\mathrm{Im}[(P+\sigma-\omega-R)_{n=0} \\ &\qquad \times (P+\sigma-\omega-R)^*_{n=1}]. \end{aligned} \tag{1.1}$$

It is experimentally observed[1] that for $2 < P_{1\text{ab}} < 6$ GeV/c

$$S \equiv P(pp) + P(pn) \approx O(s^{\alpha_{\text{eff}}-1}), \tag{1.2}$$

where

$$\alpha_{\text{eff}}(t) \approx -0.3 + t, \tag{1.3}$$

with a magnitude of about 0.3 at $t=-0.3$, $P_{1\text{ab}}=4$ GeV/c.

We can easily obtain qualitative conclusions from Eq. (1.1). $\mathrm{Im}(P_{n=0}P^*_{n=1})$ vanishes if P is assigned to be a single power $-[\exp(-i\pi/2)s]^{\alpha_P}$; it is small if small cut corrections are added, but these would yield terms of $O(s^{\alpha_P-1})$ in S which disagrees with the experimental energy behavior. The more conventional decomposition $P=P_N+f_{\text{exd}}$ yields $(P_N)(f_{\text{exd}})^*$ cross terms, but these would be wrong since they are of $O(s^{-1/2})$. The same is true of P-ω cross terms. The $\omega_{n=0}\omega^*_{n=1}$ and $R_{n=0}R^*_{n=1}$ terms are of $O(s^{-1})$, which is the correct behavior, but both terms are much too small; the factors are mainly in phase and in any case are much smaller than $|P_{n=0}|^2$. The only term that can plausibly be assumed to have both the correct energy dependence and magnitude is the $P_{n=0}\sigma^*_{n=1}$ term. This is the only term we shall keep in S.

At this point it is worthwhile mentioning that a scheme involving the $I=0$, $C=-1$ Freund-Nambu O trajectory at $\alpha_0=\alpha' t$, arising in connection with a Pomeron singularity P_N at $\alpha=1+\alpha' t$ inserted by hand into the cylinder coupling,[15] fails immediately in describing the polarization since $\mathrm{Im}(P_N 0^*)=0$. An alternate scheme with a heavy ω' trajectory having $\alpha_{\omega'}(0)\sim -\frac{1}{2}$ in order to reproduce the energy dependence leads to the wrong zero structure in S.

We shall now describe a qualitative fit to the sum S of the pp and pn polarizations. Using the above arguments to discard all terms but $P_{n=0}$ and $\sigma_{n=1}$ we obtain

$$S \approx A(-t)^{1/2} \sin[\tfrac{1}{2}\pi(\alpha_P-\alpha_\sigma)](\nu/\nu_0)^{\alpha_\sigma-\alpha_P}, \tag{1.4}$$

where we have taken

$$\nu = \tfrac{1}{2}(s-u), \tag{1.5}$$

as an appropriate variable for continuing to small s. We shall ignore the t-dependence of ν so that $\nu = 2mE_{1\text{ab}}$. Choosing the unrenormalized Pomeron trajectory of Ref. 7

$$\alpha_P(t) = 0.85 + \tfrac{1}{3}t, \tag{1.6}$$

the unrenormalized σ trajectory

$$\alpha_\sigma(t) = -0.4 + t, \tag{1.7}$$

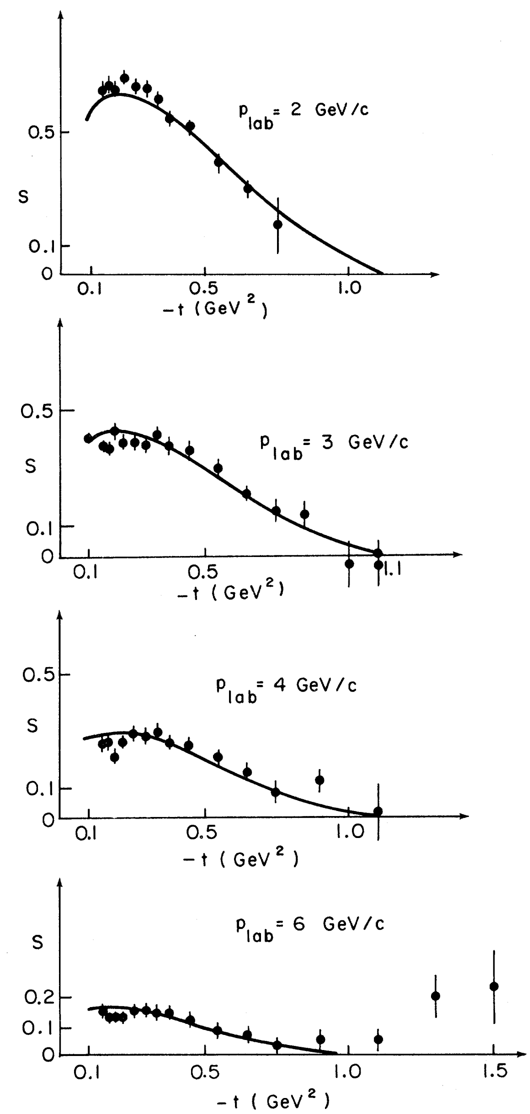

FIG. 1. $S=P(pp)+P(pn)$; $2 \le p_{\text{lab}} \le 6$ GeV/c. Data from Ref. 1. The continuous line corresponds to the parametrization of Eq. (1.4).

and

$$A(t)=A_0 e^{0.76t},$$
$$\nu_0=1\ \mathrm{GeV}^2, \qquad (1.8)$$
$$A_0=13.5\ \mathrm{GeV}^{-1},$$

we obtain the results shown in Fig. 1.

The agreement of the simple model with the data is surprisingly good. In particular, the zero at $t=-1.1$ predicted by the model at $\alpha_P-\alpha_\sigma=2$ seems to be borne out experimentally.

The sign of A_0 corresponds to taking

$$\frac{\sigma_{n=1}}{(-t)^{1/2}\sigma_{n=0}}<0, \qquad (1.9)$$

where

$$\sigma_{n=0}=-(e^{-i\pi/2}s)^{\alpha_\sigma}\beta_{\alpha_\sigma}/\sin(\tfrac{1}{2}\pi\alpha_\sigma). \qquad (1.10)$$

This form for the $n=0$ amplitude is consistent with the requirement that $\mathrm{Im}\,\sigma_{n=0}>0$, which holds for any vacuum Regge pole. The fact that the σ may undergo renormalization does not change this fact (cf. Sec. III).

The signs of the amplitudes $\sigma_{n=0}, \sigma_{n=1}$ are consistent with those of the attractive σ exchange parametrized in low-energy NN scattering.[2] (Note the sign in the elastic unitarity Eq. (3.33) of that reference.)

We may also fit the difference D of the pp and pn polarizations by making a simple assumption regarding the $I=1$ Reggeons. Taking $(A_2-\rho)_{n=1}=R$ as real, we obtain

$$D=B(-t)^{1/2}\sin\left(\frac{\pi\alpha_P}{2}\right)\left(\frac{\nu}{\nu_0}\right)^{\alpha_R-\alpha_P}. \qquad (1.11)$$

Choosing

$$B(t)=B_0 e^{0.076t},$$
$$B_0=0.619\ \mathrm{GeV}^{-1},$$
$$\alpha_R=0.56+0.85t$$

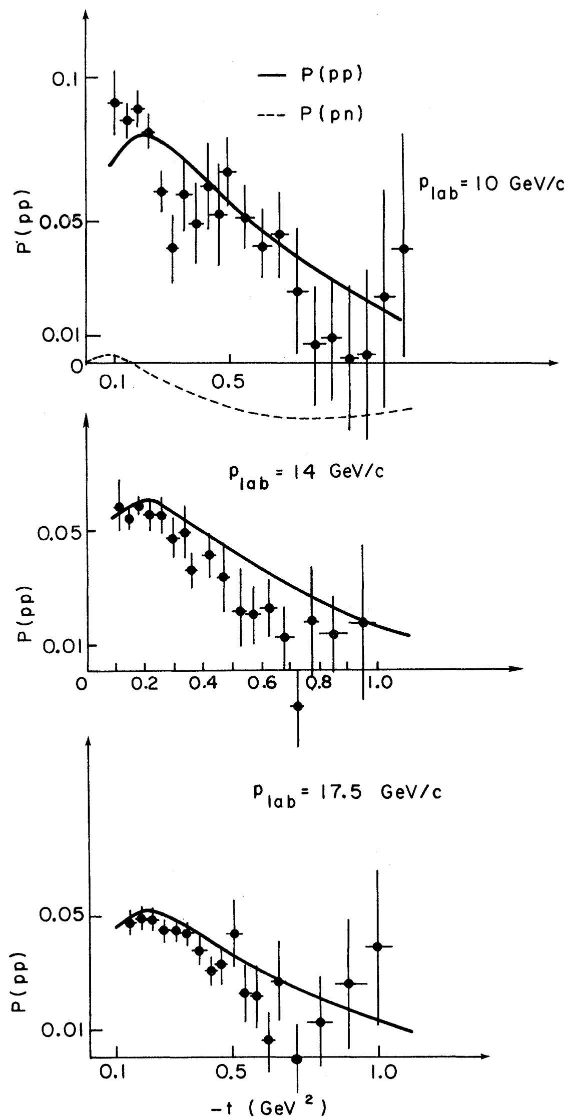

FIG. 3. pp polarization; $10\le p_{\mathrm{lab}}\le 17.5$ GeV/c. Data from Ref. 17. The continuous line corresponds to our prediction.

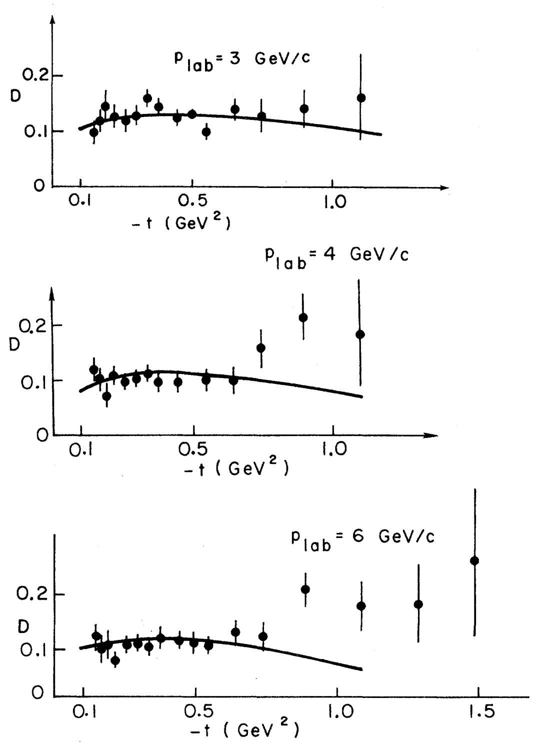

FIG. 2. $D=P(pp)-P(pn)$; $3\le p_{\mathrm{lab}}\le 6$ GeV/c. Data from Ref. 1; the curves are Eq. (1.11).

leads to the results shown in Fig. 2 for P_{lab} between 3–6. Again the agreement is quite reasonable.

We have not included the 2-GeV/c data. These data show a very anomalous behavior, being *smaller* than at higher energies. This may be related to the observation[16] that the difference of polarized-target–polarized-beam total pp cross sections $\Delta = -\sigma(\uparrow\uparrow) + \sigma(\uparrow\downarrow)$ is large at 2 GeV/c and decreases very rapidly, since Δ is determined by the $n=2$ amplitude which we have ignored up to now.

Armed with the above low-energy results, we may now predict the pp and pn polarizations at all energies up to 30 GeV/c. Past those energies at least the Pomeron we have used is renormalized;[10] qualitative estimates based on Serpukhov data indicate that the σ intercept is renormalized upward by $\alpha_\sigma^R(0) - \alpha_\sigma(0) \gtrsim 0.2$. If it were not renormalized, the model $P(pp)$ would be too small at high energies.

The results for $P(pp)$ from 10 to 17.5 GeV/c are shown in Fig. 3. The agreement with the data[17] is quite good. In Fig. 4 we show predictions for $P(pn)$. Notice that mirror symmetry is approached as the energy is increased, as expected in Regge models with only $I=1$ flip amplitudes.

Next, again invoking exchange degeneracy to write $(A_2+\rho)_{n=1} = Re^{-i\pi\alpha}$, we easily find the polarization for $p\bar{p}$ scattering as

$$P(p\bar{p}) = \tfrac{1}{2}[P(pp)+P(pn)] + \tfrac{1}{2}[P(pp)-P(pn)]\frac{\sin[\frac{1}{2}\pi(\alpha_P - 2\alpha)]}{\sin(\frac{1}{2}\pi\alpha_P)}. \tag{1.13}$$

The results at 6 GeV/c are shown in Fig. 5.

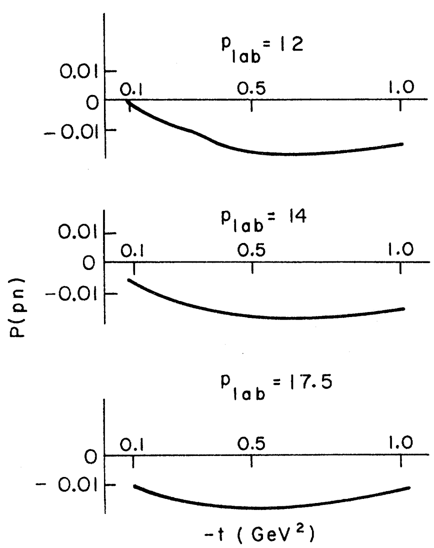

FIG. 4. Prediction for pn polarization at $12 \leq p_{\text{lab}} \leq 17.5$ GeV/c.

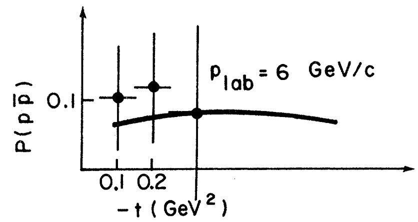

FIG. 5. Prediction for $p\bar{p}$ polarization at 6 GeV/c. Data from Ref. 17.

Again sensible agreement with the data[17] is obtained.

We may go even further and notice that since the polarization $P(K^+p)$ in K^+p elastic scattering[17] is given by an interference between the $0^-\frac{1}{2}^+$ $P_{n=0}$ and $R_{n=1}$ amplitudes (the σ is excluded by virtue of its supposedly small cylindrical $K\bar{K}$ coupling), there could *a priori* be some similarity between $P(K^+p)$ and $P(pp) - P(pn)$. That this is indeed the case is shown in Fig. 6. The prediction for the similar behavior of $P(K^-p)$ and $[P(p\bar{p}) - P(p\bar{n})]$ is also given.

II. THE n=0,2 AMPLITUDES AND OTHER LOW-LYING EXCHANGES

In this section we shall comment on other low-energy anomalies and possible low-lying exchanges

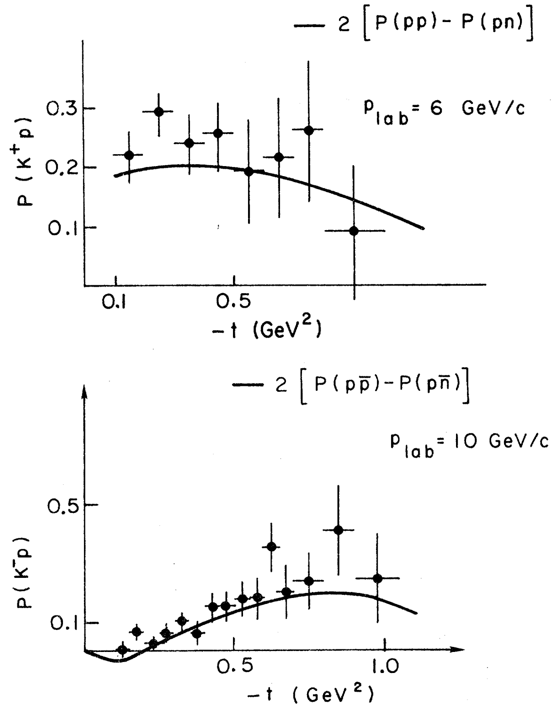

FIG. 6. Exhibits the similarity between $P(K^+p)$ and D (data from Ref. 17), and the similarity between $P(K^-p)$ and $P(p\bar{p}) - P(p\bar{n})$.

other than the σ, which we regard as established from Sec. I. Our conclusions will be much less definite.

We first consider NN and $N\bar{N}$ total cross sections. By taking suitable linear combinations of these to isolate definite quantum numbers, it is easily established that low-lying exchanges coupling strongly to NN must exist. For example, between 3 and 6 GeV/c the fact that[5] $\sigma_{pp}+\sigma_{p\bar{p}}\sim O(s^{-0.35})$ while $\sigma_{K^+p}+\sigma_{K^-p}\sim O(s^{-0.15})$ indicates the presence of $\sigma_{n=0}$ as confirmed by $\sigma(pd)+\sigma(\bar{p}d)$ (see Ref. 18) between 2 and 6 GeV/c. Similarly, $\sigma_{p\bar{p}}-\sigma_{pp}\sim O(s^{-1})$ implies the existence of some low-lying $C=-1$, $I=0$ contribution to the $n=0$ amplitude (though probably not in the $n=1$ amplitude as we have already observed in Sec. III). This is supported by the energy behavior of $\sigma(pd)-\sigma(\bar{p}d)$. Also $\sigma_{pp}-\sigma_{pn}\sim O(s^{-2})$ implies some $C=-1$, $I=1$ object. The energy dependence of the quantity $\tilde{\Delta}=(d\sigma/dt)(pp)-(d\sigma/dt)(p\bar{p})^5$ at low energies indicates the presence of low-lying contributions as well, at least at $t\approx 0$. (Past $-t\sim 0.4$ the energy dependence of $\tilde{\Delta}$ seems to become more "normal," indicating a strong t-damping of these contributions.) A third anomaly, already mentioned, is the difference[16] $\Delta=-\sigma_{pp}(\uparrow\uparrow)+\sigma_{pp}(\uparrow\downarrow)\sim O(s^{-3})$ which strongly indicates the presence of a low-lying $n=2$ cut, a conspiracy of some low-lying trajectories, or both. The type-2 conspiracy of Ref. 19 leads to an energy behavior

$$\Delta\sim\frac{1}{q(s)^{1/2}}(s-2m_N{}^2)^{\alpha^d(0)}$$

where $\alpha^d(0)$ is the intercept of the first daughter of the unnatural parity trajectory: $\alpha^d(0)=\alpha_u(0)-1$. The rapid falloff of Δ might be consistent with $\alpha_u(0)=-\frac{1}{2}$ as shown in Fig. 7.

It is also worth mentioning the presence of anomalies in meson-nucleon scattering at low energies. A variety of effects, such as exchange-degeneracy breaking in KN and $K\Delta$ charge-exchange reactions,[20] the different energy dependences of $\pi^-p\to\eta n(\eta\Delta)$,[21] and the need of a low-lying singularity for the description of the charge-exchange reaction $\pi^-p\to\pi^0 n$,[22] all point to the influence of such effects. Further study of these points is presently under investigation.

To summarize, the study of the energy dependences of σ_{tot} and other data at low energies (2 GeV/$c\lesssim P_{lab}\lesssim 6$ GeV/c) seems to imply that a trajectory one unit below the common Regge trajectory of intercept $\frac{1}{2}$ might be needed to reproduce the experimental data.

Figure 8 indicates a possible pattern for these low-lying singularities in a first approximation.

III. THE σ MESON POTENTIAL AND THE σ TRAJECTORY

In this section we give a mathematical example of the nondiffractive renormalization of the σ trajectory α_σ into α_σ^R [cf. Eq. (1.7)]. In particular we shall show that it is consistent (1) to employ a one-boson-exchange (OBE) potential with a pole at $t=m_\sigma{}^2$ at low energies, (2) to not have the σ-meson pole exist in the S matrix, and (3) to determine the mass parameter $m_\sigma{}^2$ by using the value of t at which the *unrenormalized* trajectory $\alpha_\sigma(t)$ vanishes in the Reggeized σ amplitude used at higher (but not too high) energies. The argument generalizes that of Ref. 14. We write the $n=0$ σ-exchange

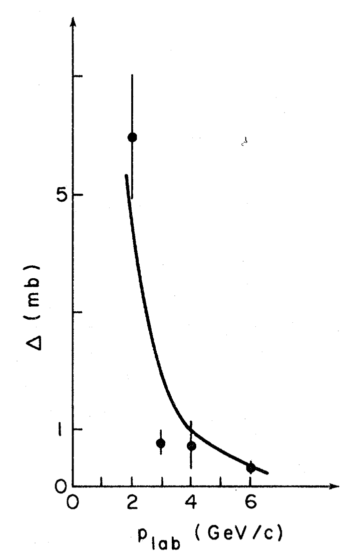

FIG. 7. $\Delta=-\sigma_{pp}(\uparrow\uparrow)+\sigma_{pp}(\uparrow\downarrow)$, $2\le p_{lab}\le 6$ GeV/c. Data from Ref. 16. The continuous line corresponds to an energy dependence

$$\Delta=A\frac{(s-2m_N{}^2)^{-1.5}}{2q(s)^{1/2}}.$$

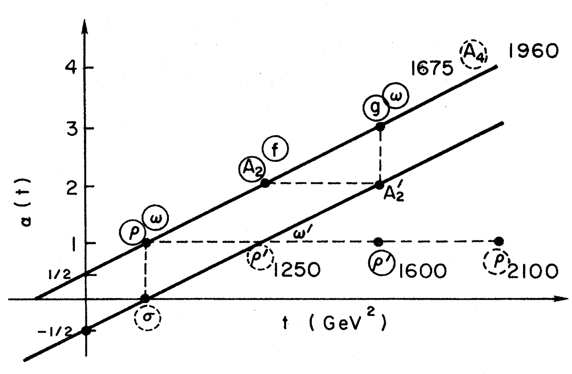

FIG. 8. Schematic pattern of the meson trajectory and its daughters. Full circles correspond to observed particles; dotted circles to questionable resonances.

amplitude in the form used previously to describe nondiffractive ("flavoring") renormalization of the bare Pomeron[7,9,10,12]

$$T_\sigma(s,t)=\int_{C-i\infty}^{C+i\infty}\frac{dj}{2\pi i}\left(\frac{s}{s_0}\right)^j\left(\frac{-e^{-i\pi j/2}}{\sin\pi j/2}\right)\frac{\tilde\beta_j e^{-bj}}{j-\alpha_\sigma-g^2e^{-bj}} \tag{3.1}$$

where $0>C>\alpha_\sigma^R,\ \alpha_\sigma$. Here g is the coupling which induces the renormalization and b is a parameter related to the threshold s_{th} of the renormalizing effect by $b\approx\frac{1}{2}\ln s_{th}$. If these effects are taken to be $K\bar K$ and $B\bar B$ inelastic production, $b\approx 2$.[10] As we shall see, b is also related to the lowest inelastic threshold $s_{th}^{(0)}$ by $b=\ln s_{th}^{(0)}$ in this example. (In principle we could have introduced another parameter b' for $\ln s_{th}^{(0)}$.) The renormalized trajectory $\alpha_\sigma^R(t)$ is determined by the leading zero of the denominator. The crucial point is contained in the possibility that $\tilde\beta_j$ vanishes at $j=\alpha_\sigma^R$. If it does *not* vanish, a pole in $T_\sigma(s,t)$ at $t=(m_\sigma^R)^2$ results corresponding to setting $\alpha_\sigma^R=0$. This is easily seen by moving the contour to the left past $j=\alpha_\sigma^R$. Now suppose that $\tilde\beta_j\sim(j-\alpha_\sigma^R)^2$. No pole in T_σ occurs at $t=(m_\sigma^R)^2$ and no simple result is found merely by moving the contour to the left. Instead we expand T_σ in a power series in g^2. For the $0(g^0)$ term, we move C to the left past $j=\alpha_\sigma$. We get the contribution [cf. Eq. (1.10)], assuming $\ln s>b$,

$$T_\sigma^{(0)}(s,t)=\frac{-(e^{-i\pi/2}s/s_0)^{\alpha_\sigma}}{\sin(\frac{1}{2}\pi\alpha_\sigma)}\tilde\beta_{\alpha_\sigma}e^{-b\alpha_\sigma}. \tag{3.2}$$

Other terms arising from the zeros of $\sin\pi j/2$ ($j<0$) vanish since $\beta_j=0$ there. The restriction $\ln s>b$ requires us to be above the lowest inelastic threshold. If $\ln s<b$, $T_\sigma^{(0)}$ is real and is approximately given by

$$T_\sigma^{(0)}(s,t)\cong-\frac{2}{\pi}\frac{\tilde\beta_0(t)}{\alpha_\sigma(t)}. \tag{3.3}$$

This can be seen by moving C to the right past $j=0$. Other terms exist in Eq. (3.3) but are small [cf. Eq. (3.5) with $k=0$, $N>0$]. Equation (3.3) is just the result obtained from Eq. (3.2) by setting $\alpha_\sigma\approx 0$.[14]

Let

$$T_\sigma(s,t)=T_\sigma^{(0)}(s,t)+\tilde T_\sigma(s,t). \tag{3.4}$$

Suppose for the moment that $\tilde T_\sigma$ is small for $t<0$. Then we may use $T_\sigma^{(0)}$ as a good approximation to the full amplitude T_σ. $T_\sigma^{(0)}$ has a pole at $\alpha_\sigma=0$ (recall we only assumed $\tilde\beta_j=0$ at $j=\alpha_\sigma^R$ so that $\tilde\beta_j$ will not in general vanish at $j=\alpha_\sigma$).

Under the conditions stated, it is clear that if the NN one-boson-exchange amplitude $T_\sigma^{\rm OBE}=-g_\sigma^2/(t-m_\sigma^2)$ is a good approximation to T_σ at low energies, we may identify it as the low-energy continuation of $T_\sigma^{(0)}$. (The nonflip amplitude in Ref. 2 is $-T_\sigma^{\rm OBE}$.) The pole at $t=m_\sigma^2$ in $T_\sigma^{\rm OBE}$ is then to be identified as the value at which $\alpha_\sigma(m_\sigma^2)=0$, and the coupling constant $g_\sigma^2=2\tilde\beta_0/\pi\alpha_\sigma'$. Since our $\tilde\beta_0$ depends exponentially on t instead of as an inverse power as in Ref. 2, we have not attempted a numerical comparison.

We now consider the term $\tilde T_\sigma(s,t)$ in Eq. (3.4). It arises from the zeros of $\sin\pi j/2$ for terms of order g^{2k}, $k\neq 0$. These terms are real and are easily shown to be small provided that we restrict our attention to "low" energies for renormalization; here this means $\ln s\ll 2b$. Using Cauchy's theorem in the right half j plane, we find, assuming $t\leq 0$,

$$\tilde T_\sigma(s,t)=\sum_{\substack{k>0\\ N\geq 0}}\frac{2g^{2k}\tilde\beta_{2N}}{\pi(2N-\alpha_\sigma)^{k+1}}e^{-2N[(k+1)b-\ln(s/s_0)]}, \tag{3.5}$$

where the Nth term results from the pole in $(\sin\pi j/2)^{-1}$ at $j=2N$. In particular the $N=0$ term is small if g^2 is small. This is the case if $\alpha_\sigma^R-\alpha_\sigma\approx 0.2$ ($g^2\approx 0.13$).

Although $\tilde T_\sigma(s,t)$ is not large at $t\leq 0$, $\ln s\ll 2b$, it is clear that it must have a pole at $t=m_\sigma^2$ since T_σ itself does not have it. Our association of $T_\sigma^{(0)}$ with the Reggeized $T_\sigma^{\rm OBE}$ at higher energies involves no contradiction since $T_\sigma^{\rm OBE}$ is of course never applied in analyzing data for $t>0$, and so need not be T_σ.

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