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I. POLARIZATIONS

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Recent measurements 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 schannel 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" C σ can be considered as a leading 2q2q 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 Iizuka-Okubo-Zweig rule.⁴ All this is consistent with meson-nucleon scattering, 🥧 which exhibits a more "normal" behavior than does NN scattering.² 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_{e 10} distinct from the f' coupling mainly to $K\overline{K}$), the resulting couplings are smaller than those we shall invoke for NN scattering.

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 $\sigma_{\rm R}$. In the schemes of Ref. 7 and 8 it is necessary to imagine, e.g., $K\bar{K}$ and $E\bar{E}$ nondiffractive inelastic thresholds renormalizing the leading vacuum trajectory with intercept below one (the f generated by cylinder corrections to the planar bootstrap in Ref. 8 or the bare pomeron \hat{P} in Ref. 7) into the bare pomeron of the Gribov calculus with intercept above one.⁹ Such renormalization is

EVIDENCE FOR A LOW-LYING UNRENORMALIZED VACUUM

TRAJECTORY FROM NN SCATTERING

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ABSTRACT

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 $p_{lab} \gtrsim 30$ GeV/c in analogy with recent models of the pomeron. The possible presence of other low-lying trajectories is briefly discussed.

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expected above 30 GeV/c on the basis of inelastic $K\overline{K}$, $B\overline{B}$ production data;¹⁰ the same conclusion is reached via detailed two-body phenomenology within this context.⁷ It is also indirectly implied by pp polarization data at Serpukhov¹¹ 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.

It is important to note that the association of the σ mass m_{σ} phenomenologically determined in low-energy NN scattering with the σ trajectory depends on whether or not one belives 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} . We shall illustrate the point with a mathematical example in Section III.

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 $\alpha_{\sigma}(t) = -0.2 + t$ would yield a mass parameter $m_{\sigma} = 450$ MeV, corresponding to the true mass of the σ meson if it exists.

We turn now to our analysis of the data. We denote by P the diffractive component of Refs. 7 and 8, 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 $A_2 - \rho$ 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

$$|\mathbf{P}_{\mathbf{n}=\mathbf{0}}|^2 \mathbf{P} \begin{pmatrix} \mathbf{p}\mathbf{p} \\ \mathbf{p}\mathbf{n} \end{pmatrix} \cong -2 \operatorname{Im} \left[(\mathbf{P} + \boldsymbol{\sigma} - \boldsymbol{\omega} \pm \mathbf{R})_{\mathbf{n}=\mathbf{0}} \left(\mathbf{P} + \boldsymbol{\sigma} - \boldsymbol{\omega} \pm \mathbf{R} \right)_{\mathbf{n}=\mathbf{1}}^* \right].$$

$$(\mathbf{1},\mathbf{1})$$

It is experimentally observed¹ that for $2 < P_{lab} < 6 \text{ GeV/c}$

 $S \equiv P(pp) + P(pn) \approx O(s^{\alpha} eff^{-1})$ (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, $P_{lab} = 4 \text{ GeV/c}$. We can easily obtain qualitative conclusions from Eq. (1.1) $Im(P_{n=0} P_{n=1}^{*})$ vanishes if P is assigned to be a single power $-[exp(-i\pi/2)s]^{*}$; it is small if small cut corrections are added, but these would yield terms of $O(s^{P})$ in S which disagrees with the experimental energy behavior. The more conventional decomposition $P = P_N + f_{EXD}$ yields $(P_N) (f_{EXD})^{*}$ cross terms, but these would be wrong since they are of $O(s^{-1/2})$. The same is true of $P - \omega$ 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 being mainly in phase and in any case 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 $P_{_{_{I\!I}}}$ at $\alpha = 1 + \alpha't$ inserted by hand into the cylinder coupling¹² fails immediately in describing the polarization since $\operatorname{Im}(P_N^{\sigma}) = 0$. An alternate scheme with a heavy ω ' trajectory having $\alpha_{\omega}, (0) \sim -1/2$ 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

$$5 \approx A (-t)^{\frac{1}{2}} \sin \left[\frac{\pi}{2} (\alpha_{\rm P} - \alpha_{\sigma}) \right] (\nu/\nu_{\rm O})^{\alpha_{\sigma} - \alpha_{\rm P}}$$
 (1.4)

where we have taken

$$= \frac{1}{2}(s - u)$$
 (1.5)

as an appropriate variable for continuing to small s. We shall ignore the t-dependence of ν so that $\nu = 2m E_{lab}$. Choosing the unrenormalized pomeron trajectory of Ref. 7

$$_{\rm P}(t) = 0.85 + t/3$$
 (1.6)

the unrenormalized σ trajectory

 α

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

and

$$A(t) = A_0 e^{0.76t}$$

$$v_0 = 1 \text{ GeV}^2$$

$$A_0 = 13.5 \text{ GeV}^{-1}$$
(1.8)

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 $\alpha_p - \alpha_q = 2$ seems to be borne out experimentally. The sign of A corresponds to taking

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

where

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

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

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

$$D = B \left(-t\right)^{\frac{1}{2}} \sin\left(\frac{\pi \alpha_{p}}{2}\right) \left(\nu/\nu_{0}\right)^{\alpha_{R}-\alpha_{p}}$$
(1.11)

Choosing

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

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 <u>smaller</u> than at higher energies. This may be related to the observation¹³ that the difference of polarized target-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. It is also worth noting that the 2 GeV/c polarizations were normalized with $d\sigma/dt$ from a different experiment.

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;⁹ qualitative estimates based on Serpukhov data indicate that the σ intercept is renormalized upward by $\alpha_{\sigma}(0) - \alpha_{\sigma}(0) \gtrsim 0.2$. If it were not renormalized the model P(pp) would be too small at high energies.

The results for P(pp) from 10-17.5 GeV/c are shown in Fig. 3. The agreement with the data¹⁴ 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 $(A_2 + \rho)_{n=1} = R e^{-i\pi\alpha}$ we easily find the polarization for $p\overline{p}$ scattering as

$$P(p\overline{p}) = \frac{1}{2} \left[P(pp) + P(pn) \right] + \frac{1}{2} \left[P(pp) - P(pn) \right] \frac{\sin \left[\frac{\pi}{2} (\alpha_p - 2\alpha) \right]}{\sin \left(\frac{\pi \alpha_p}{2} \right)}$$
(1.13)

The results at 6-14 GeV/c are shown in Fig. 5. Again sensible agreement with the data 14 is obtained.

We may go even further and notice that since the polarization $P(K^+p)$ in K^+p elastic scattering¹⁴ is given by an interference between the $0^{-\frac{1}{2}+} P_{n=0}$ and $R_{n=1}$ amplitudes (the σ being excluded by virtue of its supposedly small cylindrical $K\overline{K}$ 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 $\{P(p\overline{p}) - P(p\overline{n})\}$ 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 other than the σ , which we regard as established from Section I. Our conclusions will be much less definite.

We first consider NN and NN 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 3-6 GeV/c the fact that $\sigma_{pp} + \sigma_{p\bar{p}} \sim 0(s^{-0.35})$ while $\sigma_{p\bar{p}} + \sigma_{p} \sim 0(s^{-0.15})$ indicates the presence of $\sigma_{n=0}$ as confirmed by $\sigma(pd) + \sigma(\bar{p}d)^{-15}$ between 2 and 6 GeV/c. Similarly, $\sigma_{p\bar{p}} - \sigma_{pp} \sim 0(s^{-1})$ implies the existence of some low lying C = -1, I = 0 contribution to the n = 0amplitude (though probably not in the n = 1 amplitude as we have already observed in Section II). This is supported by the energy behavior of $\sigma(pd) - \sigma(\bar{p}d)$. Also $\sigma_{pp} - \sigma_{pn} \sim 0(s^{-2})$ implies some C = -1, I = 1 object. The energy dependence of the quantity $\widetilde{\Delta} = \frac{d\sigma}{dt} (pp) - \frac{d\sigma}{dt} (pp)^{5}$ at low energies indicates the presence of lowlying contributions as well, at least at $t \approx 0$. (Past $-t \sim 0.4$ the energy dependence of $\widetilde{\Delta}$ seems to become more "normal", indicating a strong t-damping of these contributions.) A third anomaly, already mentioned, is the difference $\Delta = \sigma_{pp}(\uparrow\uparrow) - \sigma_{pp}(\uparrow\downarrow) \sim 0(s^{-3})$ 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. 16 leads to an energy behavior

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

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 fall off of \triangle 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 mesonnucleon scattering at low energies. A variety of effects, such as severe exchange degeneracy breaking in KN and K \triangle charge exchange reactions, ¹⁷ the different energy dependences of π . $p \rightarrow \eta n(\eta \Delta)$, ¹⁹ and the need of a low lying singularity for the description of the charge exchange $\pi^- p \rightarrow \pi^0 n^{-18}$ 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 $\leq P_{lab} \leq 6$ GeV/c) deems to imply that a trajectory one unit below the common Regge trajectory of intercept 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} (c.f. 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) not have the σ meson pole exist in the S matrix, and (3) to determine the mass parameter m_{σ}^{2} by using the value of t at which the <u>unrenormalized</u> trajectory $\alpha_{\sigma}(t)$ vanishes in the reggeized σ amplitude used at higher (but not too high!) energies. The argument generalizes that of Ref. 20. Write the σ -exchange amplitude as (we take n = 0)

$$T_{\sigma}(s,t) = \int_{C-i\infty}^{C+i\infty} \frac{dj}{2\pi i} \left(\frac{s}{s_0}\right)^j \left(\frac{-e^{-i\pi j/2}}{\sin \pi j/2}\right) \frac{\tilde{\beta}_j e^{-bj}}{j - \alpha_{\sigma} - g^2 e^{-bj}} \quad (3j)$$

where $0 > C > \alpha_{\sigma}^{R}$, α_{σ} . Here g is the coupling inducing 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 thrshold $s_{th}^{(0)}$ 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 $\tilde{\beta}_{j} - (j - \alpha_{\sigma}^{R})^{2}$. 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}

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in a power series in g^2 . For the $O(g^0)$ term, we move C to the left past $j = \alpha_{\sigma}$. We get the contribution (c.f. Eq. (1.10)), assuming in s > b,

$$T_{\sigma}^{(0)}(s,t) = \frac{-\left(e^{-\frac{i\pi}{2}}\frac{s}{s_{0}}\right)^{\alpha}}{\sin\left(\frac{\pi}{2}\alpha_{\sigma}\right)} \tilde{\beta}_{\alpha_{\sigma}}e^{-b\alpha_{\sigma}}.$$
 (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 inelastic threshold. If ln s < b, $T_{\sigma}^{(0)}$ is real and is approximately given by

$$\mathbf{T}_{\sigma}^{(0)}(\mathbf{s},\mathbf{t}) \stackrel{\simeq}{=} -\frac{2}{\pi} \frac{\tilde{\beta}_{0}(\mathbf{t})}{\alpha_{-}(\mathbf{t})} . \qquad (3.3)$$

This can be seen by moving C to the right past j = 0. Other terms exist in Eq. (3.3) but are small (c.f. Eq. (3.14)). Equation (3.3) is just the result obtained from Eq. (3.2) by setting $\alpha_{\sigma} \approx 0^{(20)}$.

Let

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

Suppose for the moment that \tilde{T}_{σ} 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 $T_{\sigma}^{OBE} = -g_{\sigma}^2/(t - m_{\sigma}^2)$ 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 $-T_{\sigma}^{OBE}$.) The pole at $t = m_{\sigma}^2$ in T_{σ}^{OBE} is then to be identified as the value at which $\alpha_{\sigma}(m_{\sigma}^2) = 0$, and the coupling constant $g_{\sigma}^2 = 2\tilde{\beta}_0/\pi\alpha_{\sigma}^2$. Since our $\tilde{\beta}_0$ actually depends on t 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 $t \leq 0$,

$$\tilde{\mathbf{T}}_{\sigma}(\mathbf{s},\mathbf{t}) = \sum_{\substack{\mathbf{k}>0\\\mathbf{N}\geq 0}} \frac{2g^{2\mathbf{k}} \tilde{\boldsymbol{\beta}}_{2\mathbf{N}}}{\pi(2\mathbf{N}-\boldsymbol{\alpha}_{\sigma})^{\mathbf{k}+1}} e^{-2\mathbf{N}\left[(\mathbf{k}+1)\mathbf{b}-\mathbf{ln}(\mathbf{s}/\mathbf{s}_{0})\right]}, \quad (3.5)$$

where the <u>N</u>th 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 $T_{\sigma}(s,t)$ 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_{σ}^{OBE} at higher energies involves no contradiction since T_{σ}^{OBE} is of course never applied in analyzing data for t > 0, and so need not be T_{σ} .

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FIGURE CAPTIONS

Fig. 1. S = Pol(pp) + Pol(pn); 2 ≤ p_{lab} ≤ 6 GeV/c. Data from Ref. 1. The continuous line corresponds to the parametrization of Eq. (1.4)

Fig. 2.
$$D = Pol(pp) - Pol(pn);$$
 $3 \leq p_{lab} \leq 6 \text{ GeV/c}.$

Data from Ref. 1; the curves are Eq. (1.11).

Fig. 3. pp polarization; $10 \leq p_{lab} \leq 17.5 \text{ GeV/c}$.

Data from Ref. 14. The continuous line corresponds to our prediction.

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

Fig. 5. Prediction for pp polarization at 6 GeV/c. Data from Ref. 14.

- Fig. 6. Exhibits the similarity between $Pol(K^+p)$ and D. Data from Ref. 14 and the similarity between $P(K^-p)$ and $P(p\bar{p}) - P(p\bar{n})$.
- Fig. 7. $\triangle = -\sigma_{pp}(\uparrow \uparrow) + \sigma_{pp}(\uparrow \downarrow), 2 \le p_{lab} \le 6 \text{ GeV/c.}$ Data from Ref. 13. Continuous line corresponds to an energy dependence

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

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

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Fig. 3

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Fig. 5

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Fig. 6

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