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A PERTURBATIVE DESCRIPTION OF THE POMERON SUGGESTED BY THE TWO-COMPONENT MODEL OF MULTIPARTICLE PRODUCTION*

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ABSTRACT

The two-component model of multiparticle production is interpreted as the zeroth order term in an expansion in powers of the triple pomeron coupling. This leads to an S-matrix perturbative description of the pomeron in terms of measurable quantities. The first order corrections are calculated and it is shown that the breaking of factorization can differ significantly from that expected in the two-component model. Rough numerical estimates are also given.

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The two-component model of multiparticle production suggested by Wilson [1] and developed by others [2], has had success in correlating the available NAL and ISR data, and in this letter we show how this model suggests a perturbative S-matrix approach to the pomeron. Each term in the expansion corresponds to a distinct, observable class of reactions; there are no renormalization ambiguities, and the ratio of successive terms is small.

The usual form of the two-component model describes the total cross section in terms of a "short-range correlation," or SRC, component (Fig. 1a) plus diffraction into low masses (Fig. 1b). The former may be defined [3] as a "fireball" with no large rapidity gap in the final state, and its contribution to the total cross section is assumed to be dominated at high energy by an isolated <u>factorizable</u> Regge pole, and is amenable to the Mueller analysis of inclusive spectra. Processes with a large gap in the final state, as in low-mass diffraction, are described by pomeron exchange and give cut contributions to the cross section. Phenomenologically [2], the SRC component represents ~ 80% of inelastic events at moderate energies. The contribution to the cross section of diffraction into high masses [4] (Fig. 2) or processes with more than one large gap (Fig. 3) are proportional to the triple pomeron coupling and are small, provided the latter is small. Processes involving several large gaps, as in Fig. 3b, will involve higher powers of the triple pomeron coupling or of the pomeron-pomeron-particle coupling. Furthermore each such process can be associated uniquely with a well-defined region of phase-space.

Since the two-component model includes only processes of <u>zeroth</u> order in the triple pomeron coupling, it is natural to investigate a systematic expansion in powers of this coupling [5]. This approach is closely related to Abarbanel's decomposition of multiparticle unitarity used in the study of the two-pomeron cut [6]. In the present paper we will analyze the first order processes of Figs. 2a and 2b and discuss how they "renormalize" the singularities of the two-component model, and the implications for factorization. At presently available energies this approximation should be excellent.

It is straightforward to extend our work to terms of any order, such as those associated with Figs. 3a and 3b, and we will give such an extension in a future publication [7]. We have a set of rules for the J-plane contribution of any diagram in terms of explicit Regge-poles and momentum transfer integrals. We will discuss in detail a solvable model in which pomeron-particle-pomeron vertices factorize in the pomeron masses.

We begin by writing the contribution of Fig. 1a, the SRC component, to Im A(s,0) [8] as

$$\frac{\beta_{\tilde{P}aa}(0)\,\beta_{\tilde{P}bb}(0)}{J - \alpha_{\tilde{P}}(0)} \tag{1}$$

where the singularity $\tilde{P}$ must differ from the true pomeron P, and satisfy $\alpha_{\tilde{P}}(0) < \alpha_P(0)$. This requirement follows because only a subset of the reactions which build P are allowed to contribute to $\tilde{P}$ [9].

Low-mass diffraction, Fig. 1b, leads to an AFS-like J-plane cut;

$$\frac{1}{16\pi s_0}\int dt\, \frac{\left(\sum_{a'} \beta^2_{Paa'}(t)\right)\left(\sum_{b'} \beta^2_{Pbb'}(t)\right)|\xi_P(t)|^2}{J - 2\alpha_P(t) + 1}\ . \tag{2}$$

We emphasize that we exchange the P trajectory, rather than $\tilde{P}$, since the complete elastic process appears in the sums; we assume the same singularity is responsible for all diffractive excitations. Also the cut contributions turn out to be small and not important for the subsequent discussion and therefore only the pole P is exchanged here.

Diffraction into one high mass, as in Fig. 2a, gives [10]

$$\frac{d\sigma}{ds_b'\,dt} = \frac{1}{16\pi s^2}\left(\sum_{a'} \beta^2_{Paa'}(t)\right)|\xi_P(t)|^2$$

$$\times \left(\frac{s}{s_0}\right)^{2\alpha_P(t)}\left(\frac{s_b'}{s_0}\right)^{\alpha_{\tilde{P}}(0)-2\alpha_P(t)} g_{PP\tilde{P}}(t)\,\beta_{\tilde{P}bb}(0) \tag{3}$$

provided s_b'/s_0 and s/s_b' are large; we will require these ratios to be greater than e^{Δ} with $\Delta \sim 3$. We use the $\tilde{P}$ singularity in the fireball since there are no large gaps, and thus there is no difficulty with multiple counting [11]. The relation between g_{PPP} and $g_{PP\tilde{P}}$ will be discussed in Ref. 7; the latter is smaller for the same reason that $\alpha_{\tilde{P}}(0) < \alpha_P(0)$. To simplify the discussion we assume $\alpha_{\tilde{P}}(0) > 2\alpha_P(0) - 1$, which of course requires $\alpha_P(0) < 1$. The J-plane projection of Eq. (3) is then a pole

$$\frac{\beta^{(1)}_{\tilde{P}aa}(0)\;\beta_{\tilde{P}bb}(0)}{J-\alpha_{\tilde{P}}(0)}\;e^{2\Delta\left(\alpha_{\tilde{P}}(0)-J\right)} \tag{4}$$

plus a <u>cancelling</u> cut

$$-\frac{1}{16\pi s_0}\int dt\left(\sum_{a'}\beta^2_{Paa'}(t)\right)|\xi_P(t)|^2\;\frac{\gamma^{(1)}_{\tilde{P}bb}(t)}{J-2\alpha_P(t)+1}$$

$$\times\; e^{2\Delta\left(2\alpha_P(t)-1-J\right)} \tag{5}$$

where

$$\beta^{(1)}_{\tilde{P}aa}(0) = \frac{1}{16\pi}\int dt\left(\sum_{a'}\beta^2_{Paa'}(t)\right)|\xi_P(t)|^2$$

$$\times\;\frac{g_{PP\tilde{P}}(t)\;e^{\Delta\left(2\alpha_P(t)-1-\alpha_{\tilde{P}}(0)\right)}}{\alpha_{\tilde{P}}(0)-2\alpha_P(t)+1} \tag{6}$$

and

$$\gamma^{(1)}_{\tilde{P}bb} = s_0\,\frac{g_{PP\tilde{P}}(t)\;\beta_{\tilde{P}bb}(0)}{\alpha_{\tilde{P}}(0)-2\alpha_P(t)+1}\;e^{\Delta\left(\alpha_{\tilde{P}}(0)-2\alpha_P(t)+1\right)}\;. \tag{7}$$

The contribution of Fig. 2b may be obtained from these equations by interchanging a and b.

The correction $\beta^{(1)}_{\tilde{P}aa}$ is expected to be small compared to $\beta_{\tilde{P}aa}$ since it is proportional both to the small quantity $g_{PP\tilde{P}}$ and the moderately small quantity $\sum_{a'}\beta^2_{Paa'}$, which measures the amount of low mass diffraction. Combining (1), (4) and the analog of (4) with $a \longleftrightarrow b$, the $\tilde{P}$ contribution is $\left(\text{for } J \text{ near } \alpha_{\tilde{P}}(0)\right)$

$$\frac{\left(\beta_{\tilde{P}aa}(0)+\beta^{(1)}_{\tilde{P}aa}(0)\right)\left(\beta_{\tilde{P}bb}(0)+\beta^{(1)}_{\tilde{P}bb}(0)\right)}{J-\alpha_{\tilde{P}}(0)}\;. \tag{8}$$

We have taken the liberty of adding a very small higher order term so that (8) manifestly factorizes [12].

Similarly from (2) and (5), the total cut contribution is

$$\frac{1}{16\pi s_0}\int dt\,\frac{|\xi_P(t)|^2}{J-2\alpha_P(t)+1}\left\{\left(\sum_{a'}\beta^2_{Paa'}(t)\right)\left(\sum_{b'}\beta^2_{Pbb'}(t)\right)\right.$$

$$\left.-\left(\sum_{a'}\beta^2_{Paa'}(t)\right)\gamma^{(1)}_{\tilde{P}bb}-\gamma^{(1)}_{\tilde{P}aa}\left(\sum_{b'}\beta^2_{Pbb'}(t)\right)\right\}\;. \tag{9}$$

If we follow the approximation made for the pole piece and add a higher order term [13] we get

$$\frac{1}{16\pi s_0}\int dt\;|\xi_P(t)|^2\,\frac{N_a(t)\,N_b(t)}{J-2\alpha_P(t)+1}\;, \tag{9'}$$

where

$$N_a(t) = \sum_{a'}\beta^2_{Paa'}(t)-\gamma^{(1)}_{\tilde{P}aa}(t)\;. \tag{10}$$

This is essentially the Gribov [14] formula for the cut, but with the positive sign expected from approach [6]. The N's are not quite the usual nonsense wrong signature fixed pole residues, since $\gamma^{(1)}$

depends in part on $\tilde{P}$, but the difference is a higher order correction. Going from (9) to (9') may not be as good an approximation, however, because in our numerical estimate $\gamma^{(1)}$ will turn out to be comparable to $\sum \beta^2$. This is important because (9), unlike (9'), is not positive definite (for $a = b$) and so there is a possibility of a negative cut in the energy range where only the first order terms in $g_{PP\tilde{P}}$ can contribute.

From Eqs. (6)-(10) we see that the effect of the first order corrections is to change the residues of the singularities of the two-component model [1],[2] by various averages of $g_{PP\tilde{P}}(t)$; the pole residue is increased and the strength of the cut decreased [3]. Thus the breaking of factorization in total cross sections should be less than the fraction of low-mass diffraction in the naive two-component model, and more in accord with experiment [14].

We mention at this point some of the results of the solvable factorizable model alluded to above [7]. The pomeron pole residue and two-P cut discontinuity are, to a good approximation, the results just derived in first order. The model will renormalize the pole at $\alpha_{\tilde{P}}(0)$ into a higher pole which should be identified as the real pomeron, the two being related by

$$\alpha_P(0) - \alpha_{\tilde{P}}(0) = e^{\Delta\left(\alpha_{\tilde{P}}(0)-\alpha_P(0)\right)} \int \frac{dt}{16\pi} \, g^2_{PP\tilde{P}}(t) \times \frac{e^{\Delta\left(2\alpha_P(t)-1-\alpha_P(0)\right)}}{\alpha_P(0) - 2\alpha_P(t) + 1} + O(g^4_{PP\tilde{P}}) \quad . \tag{11}$$

Finally, we wish to make a quantitative estimate. We must emphasize that some important input parameters, especially α'_P and $g_{PP\tilde{P}}(t)$ are not well known from experiment, and the numbers we obtain should be regarded with caution.

We consider pp scattering, and take $\sigma_T \approx 38$ mb, $\sigma_{SRC} \approx 26$ mb, and $\sigma_D \approx 12$ mb [2]. From factorization, the single diffraction dissociation cross section is $\sigma_{s.d.d.} \approx 0.41\ \sigma_{e\ell}$. In making the estimates, the approximate relations $g_{PPP} \approx g_{PP\tilde{P}}$ and $\alpha_{\tilde{P}}(0) \approx \alpha_P(0) \approx 1$ are used. Since the correction terms depend only on integrals over $g_{PP\tilde{P}}(t)$, the results should not depend critically on the details of the parametrization. We take $\alpha'_P = 0.3\ \mathrm{GeV}^{-2}$ and $g_{PPP}(t) = -2ate^{1.5t}$ [16],[17], where $a = 1\ \mathrm{GeV}^{-4}$, and $\sum_{p'} \beta^2_{Ppp'}(t) \approx \beta^2_{Ppp}(0)(1 + \sigma_{s.d.d.}/\sigma_{e\ell})e^{bt}$ with $b \approx 4\ \mathrm{GeV}^{-2}$, assumed to be the same for all diffractive excitations of the proton. With $\Delta = 3.0$, we obtain [18] $\beta^{(1)}_{\tilde{P}pp}(0) \approx 1.25$ so $\beta_{\tilde{P}pp}(0) + \beta^{(1)}_{\tilde{P}pp}(0) \approx 9.31$. The factorizable piece of σ_T is then ≈ 35 mb, so the cut is ≈ 3 mb, --a considerable reduction. Using Eq. (11) for the shift in pole position gives $\alpha_P(0) - \alpha_{\tilde{P}}(0) \approx 0.0058$. Thus, since $\alpha_P(0)$ is very near one (as implied by the near constancy of the total cross section), so is $\alpha_{\tilde{P}}(0)$.

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FOOTNOTES AND REFERENCES

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3. G. F. Chew, An Estimate of the Magnitude of Triple Pomeron Coupling From The Observed Energy Dependence of Total and Elastic pp Cross Sections, Lawrence Berkeley Laboratory Report LBL-1556, Jan. 1973.
4. "Large mass" here means large enough for its contribution to the cross section to be given by the leading SRC Regge pole.
5. G. F. Chew, in Multiperipheral Dynamics, 1971 Coral Gables Conference Proceedings (Gordon and Breach, 1971).
6. H. D. I. Abarbanel, Phys. Rev. D6, 2788 (1972); G. F. Chew, Phys. Rev. D7, 934 (1973).
7. M. Bishari, G. F. Chew, and J. Koplik, in preparation.
8. Our normalization conventions are $\sigma_T = s^{-1}\,\mathrm{Im}\,A(s,0)$, where $A(s,t)$ is the elastic scattering amplitude. A reggeon R contributes

$$\frac{1}{s}\,\beta_{Raa}(0)\,\beta_{Rbb}(0)\,(s/s_0)^{\alpha_R(0)}$$

to σ_T and

$$\frac{1}{16\pi s^2}\,\beta^2_{Raa}(t)\,\beta^2_{Rbb}(t)\,|\xi_R(t)|^2(s/s_0)^{2\alpha_R(t)}$$

to $d\sigma/dt$, where $s_0 = 1\ \mathrm{GeV}^2$ is a scale energy and $\xi_R(t)$ is the signature factor.

9. For an explicit proof see J. Koplik, Phys. Rev. D7, 558 (1973).
10. The PPP' coupling can be shown to be unimportant in the present context, although it can be significant in diffractive inclusive spectra near the forward direction.
11. G. F. Chew and S. D. Ellis, Phys. Rev. D6, 3330 (1972).
12. In fact, Fig. 3a contributes such a term.
13. Figure 2c provides the extra term.
14. V. N. Gribov, Zh. Eksp. Teor. Fiz. 53, 654 (1967)[Sov. Phys. JETP 26, 414 (1968)].
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16. R. Rajaraman, Phys. Rev. Letters 27, 693 (1971).
17. A parametrization of $g_{PPP}(t)$ with a linear zero, based on the analysis of Ref. 14, has been used in connection with the P-P cut by I. J. Muzinich, F. E. Paige, T. L. Trueman, and L.-L. Wang, Phys. Rev. Letters 28, 850 (1972).
18. The incident energy is assumed sufficiently high that t_{min} effects can be neglected. We will discuss this point further in Ref. 7.

FIGURE CAPTIONS

Fig. 1. Processes in the two-component model.

Fig. 2. Diffraction into high mass.

Fig. 3. Processes with more than one large gap.

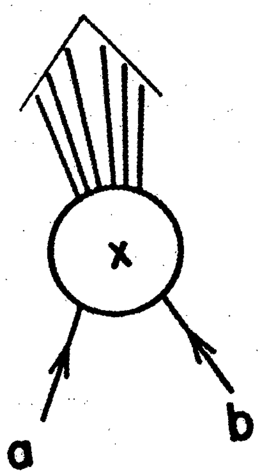

Fig. 1a

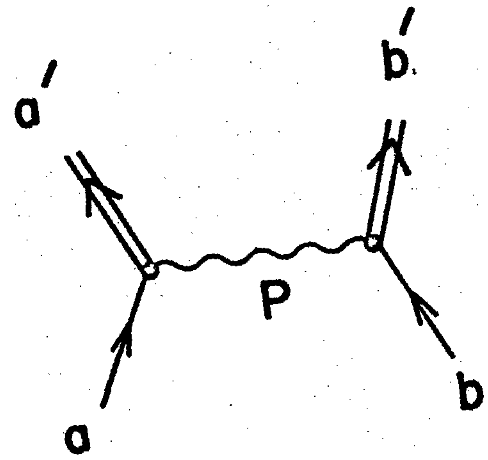

Fig. 1b

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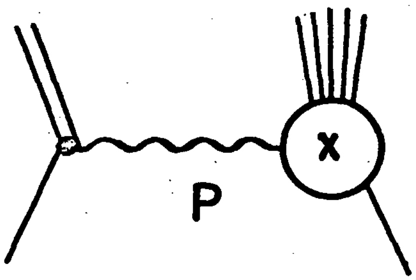

Fig. 2a

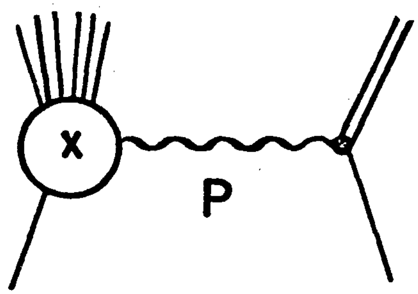

Fig. 2b

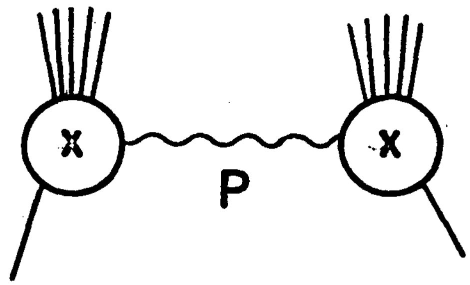

Fig. 2c

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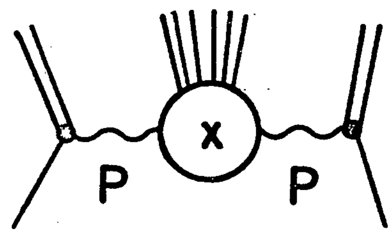

Fig. 3a

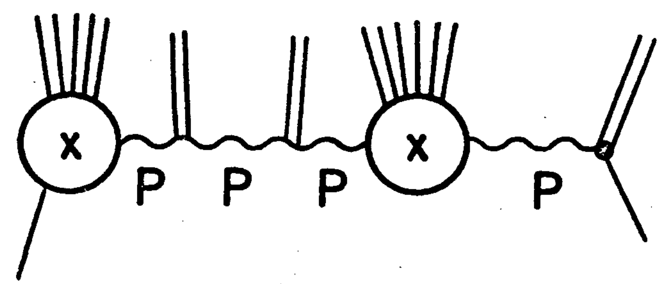

Fig. 3b

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