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**CONSTRAINTS ON POMERON COUPLINGS
FROM MULTIPLE DIFFRACTIVE DISSOCIATION**

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Since the calculation of Kajantie and Finkelstein ⁽¹⁾, which showed that multiple exchange of a factorizable Pomeron trajectory with unit intercept is inconsistent with single Pomeron exchange for the total cross section, and indeed violates the Froissart bound, much interest has centered on ways of circumventing this result. In order to get a violation of the Froissart bound, or any inconsistency at all, from the Kajantie-Finkelstein calculation one must include increasingly many Pomeron exchanges as the energy increases. A modern and elegant way of getting an inconsistency, but not a violation of the Froissart bound, was obtained from single Pomeron exchange by Abarbanel, Chew, Goldberger, and Saunders ⁽²⁾ in terms of inclusive cross sections. These authors found that an integration of single particle inclusive cross sections over the "triple-Pomeron region" gives rise to an energy dependence

$$\sigma(s) \propto \frac{g_p(0)}{2\alpha'} \ln \ln s$$

at large energy, where g_p is the triple-Pomeron coupling and $\alpha_p(t) = 1 + \alpha' t + \dots$. Consistency requires that $g_p(0) = 0$.

An extension of these results, obtained by Jones, Low, Tye, Veneziano and Young ⁽³⁾ shows that

$$\sigma(s) \propto \frac{g_{PRR}(t) V_{PR}^2(t)}{2\alpha'} \ln \ln s .$$

Here $g_{PRR}(t)$ is the coupling of a Pomeron at zero mass to two Regge trajectories of mass t , and $V_{PR}(t)$ is the coupling of a Pomeron of zero mass, a Regge trajectory of mass t , and a physical particle. Consistency now requires that either $g_{PRR}(t)$ or $V_{PR}(t)$ vanish. Brower and Weis ⁽⁴⁾ have given strong arguments showing that the

vanishing of $g_{PRR}(t)$ or $V_{PR}(t)$ requires, by continuation in t , that the Pomeron decouple from physical particles. No really satisfactory way out of this difficulty, which keeps a constant total cross section, has yet been achieved.

In this note we shall give a modest strengthening of the above results. We shall show that a single Pomeron exchange gives rise to a term like

$$\sigma(s) \propto \frac{[g_P(0)]^2}{2\alpha'} \ln s$$

while N Pomeron exchange gives rise to a term like

$$\sigma(s) \propto \left[\frac{g_P^2(0) \ln s}{2\alpha' N} \right]^N.$$

These results may be of interest for several reasons. (1) $\ln \ln s$ is never large for all practical purposes, but $\ln s$ can be large at accelerator energies. (2) A violation of the Froissart bound is obtained by considering only three Pomeron exchanges. (3) Attempts to overcome the decoupling problem by setting $\alpha' = 0$ do not seem to avoid our strengthened result.

We begin by considering those events, contributing to the total cross section, σ , for which a gap size greater than Δ overlaps the center of the rapidity plot as shown in Fig. 1. We have labelled the incoming momenta, p_A and p_B , by

$$p_A = (M_A, 0, 0, 0)$$

$$p_B = M_B(\text{ch } Y, 0, 0, \text{sh } Y),$$

Each event which contributes to σ has at most one gap greater than Δ which overlaps $Y/2$. Since we are neglecting a subset of the events contributing to the total cross section we have the rigorous inequality

$$\sigma \geq \int_{\Delta}^{Y/2} d\mu_1 \int_{\Delta}^{\min\{Y/2, Y - \mu_1 + \Delta\}} d\mu_2 \int_{-\infty}^0 dt \frac{d\sigma}{dt d\mu_1 d\mu_2} \quad (1)$$

The quantity $\frac{d\sigma}{dt d\mu_1 d\mu_2}$ is the double diffractive dissociation cross section shown in Fig. 2. The clusters have masses M_1 and M_2 , and the momentum transfer between particle A and one of the clusters is t . Here, also,

$$\begin{aligned} \mu_1 &= \ln M_1^2 \\ \mu_2 &= \ln M_2^2 \end{aligned}$$

The sum over the left and right hand clusters is entirely unrestricted. We include all states with $n = 2, 3, 4, \dots$ particles and any distribution of gaps within the clusters.

The standard hypothesis of a factorized Pomeron gives

$$\frac{d\sigma}{dt d\mu_1 d\mu_2} = \frac{\beta_A(0)\beta_B(0)}{16\pi} [g_P(t)]^2 e^{2[\alpha_P(t) - 1](Y - \mu_1 - \mu_2) + (\alpha_P(0) - 1)t} \quad (2)$$

when μ_1, μ_2 , and $Y - \mu_1 - \mu_2$ are all larger than some fixed number Δ .

(Δ is chosen large enough to make secondary Regge exchanges negligible.) Toller angle dependence has been neglected since at $t = 0$ the Toller angle disappears.

(We shall be dealing with t values on the order of $1/\ln s$ so Toller angle dependence will not affect our "leading log" results.)

Using (2), with $\alpha(0) = 1$, (1) becomes

$$\sigma(Y) \geq \beta_A \beta_B \frac{[g_P(0)]^2}{32\pi\alpha'} \int_{\Delta}^{Y/2} d\mu_1 \int_{\Delta}^{\min\{Y/2, Y - \mu_1 + \Delta\}} d\mu_2 \frac{1}{Y - \mu_1 - \mu_2}$$

Call

$$\mu_1 = u_1 Y$$

$$\mu_2 = u_2 Y$$

then

$$\sigma(Y) \geq \beta_A \beta_B \eta_P Y \int_0^{\frac{1}{2}} du_1 \int_0^{\frac{1}{2}} du_2 \frac{1}{1 - u_1 - u_2}$$

where $\eta_P = \frac{[g_P(0)]^2}{32\pi\alpha'}$ is the triple Pomeron coupling. Thus, $\sigma(Y)$ grows at least as rapidly as $Y = \ln s$ unless $g_P(0) = 0$.

It should be emphasized that in Fig. 1 there may or may not be additional large gaps in rapidity in the two clusters. It is essential to include all events with a gap about $Y/2$ of a given width in order that the sum over states to the right and left build the full Pomerons in the M_1^2 and M_2^2 channels ^(a). We remark that if the u_1 and u_2 integrals are extended beyond $\frac{1}{2}$, the integral behaves as $\ln s \ln \ln s$. The $\ln s$ must now be associated with multiple counting and the $\ln \ln s$ with the non-vanishing triple-Pomeron coupling so the result is weakened.

It is a trivial matter to exhibit two additional particles, as illustrated in Fig. 3, with t_1 and t_2 representing secondary trajectories and t a Pomeron. Now, either $g_{PRR}(t)$ or $V_{PR}(t)$ must vanish or again $\sigma \propto Y$.

Now consider $N + 1$ cluster production as shown in Fig. 4a. Fig. 4b shows those events, contributing to σ , which have N gaps overlapping fixed positions $Y/2N, 3Y/2N, \dots, \frac{(2N-1)Y}{2N}$.

Then,

$$\frac{d\sigma}{dt_1 dt_2 \dots dt_N d\delta_1 d\delta_2 \dots d\delta_N d\mu_1 d\mu_2 \dots d\mu_N} =$$

$$= e^{\alpha(0)Y} \beta_A \beta_B \prod_{i=1}^N \frac{[g_p(t_i)]^2}{16\pi} e^{2[\alpha(t_i) - 1] \delta_i - \alpha(0)\delta_i} \quad (3)$$

Since each gap overlaps a fixed point, there is no multiple counting when one integrates over gap size. Thus,

$$\sigma \geq \int_{-\infty}^0 \prod_{i=1}^N dt_i \int_{\Delta}^{Y/2N - 2\Delta} \prod_{i=1}^N d\delta_i \prod_{i=1}^N \int_{\frac{2i-1}{2N} Y - \delta_i - \sum_{j=1}^{i-1} (\mu_j + \delta_j)}^{\frac{2i-1}{2N} Y - \sum_{j=1}^{i-1} (\mu_j + \delta_j)} d\mu_i \frac{d\sigma}{dt_1 \dots d\delta_1 \dots d\mu_1 \dots} \quad (4)$$

the t integrals are easily done to give

$$\sigma \geq \int \prod_{i=1}^N \frac{d\delta_i}{\delta_i} \int \prod_{j=1}^N d\mu_j \beta_A \beta_B \eta_P^N,$$

which implies

$$\sigma \geq \beta_A \beta_B \eta_P^N \int_{\Delta}^{Y/2N - 2\Delta} \prod_{i=1}^N d\delta_i$$

or

$$\sigma \geq \beta_A \beta_B \left[\frac{\eta_P Y}{2N} \right]^N$$

We have chosen the gap positions and integrations in such a way as to obtain a simple result. The coefficient of $[\eta Y]^N$ is not the best possible ⁽⁵⁾, but no extra energy factors will emerge from a different choice of gaps.

It is easy to trace the Y^N behavior in (4). The multi-Regge phase space

contributes Y^N through the δ_i integrals, the integrals over the cluster masses μ_i give Y^{N+1} , the integrals over the t_i give Y^{-N} coming from the shrinkage of the diffraction peaks, and there is one factor Y^{-1} from the restriction $\sum_{i=1}^N \delta_i + \sum_{i=1}^{N+1} \mu_i = Y$. If the triple-Pomeron vertex has a linear zero in t , the t_i integrals give a factor Y^{-2N} so consistency is easily obtained ^(b). Another conceivable way to achieve consistency might be to increase the amount of shrinkage of the diffraction peaks ^{(6), (7)}. Roughly speaking, this can be done by taking trajectories of the form $\alpha_p(t) = 1 + \alpha(-t)^\gamma$ for which $\alpha'(0) = \infty$, and the above contributions vanish since they are proportional to $1/\alpha'$. Such increased shrinkage gives $Y^{-N/\gamma}$ for the t_i integrals. Thus consistency with constant cross sections would be obtained for $\gamma < \frac{1}{2}$ ^(c). However, at least for the elastic cross section shrinkage greater than Y^2 ($\gamma < \frac{1}{2}$) is forbidden by unitarity ⁽⁸⁾ so we believe that such schemes are probably unacceptable.

In conclusion, we emphasize that (4) provides a quite general constraint on multiple diffractive dissociation (i.e. approximately constant behavior for large δ_i and M_i^2) independent of the Pomeron fine structure (i.e. the number of powers of δ_i and μ_i by which the cross section differs from a constant). Thus, suppose (3) has the additional terms $(\delta_i)^{2a_i}$ and $(\mu_i)^{b_i}$ with diffraction peaks shrinking like $(\delta_i)^{-1/\gamma_i}$ and residues vanishing as $(t_i)^{2d_i}$. Then (4) becomes

$$\sigma_{\text{Tot}} \geq Y^{\sum_{i=1}^N (2a_i + 1) + \sum_{i=1}^{N+1} (b_i + 1) - \sum_{i=1}^N \frac{1}{\gamma_i} (2d_i + 1) - 1}.$$

By virtue of its generality the above constraint is rather untransparent. Nevertheless, it must be reckoned with in any model of diffractive production.

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FOOTNOTES

- a. Furthermore, the gap must have a fixed position in rapidity in order to avoid multiple counting.
- b. Ignoring shrinkage by replacing the Pomeron by a fixed pole gives a Y^{2N} behavior for the total cross section. The $\ln^2 s$ behavior of the two cluster decomposition may be of some phenomenological interest.
- c. The above leading log calculations seem to allow $\gamma = \frac{1}{2}$, but a more careful analysis shows that $\ln \ln s$ terms are present in this case.

FIGURE CAPTIONS

- Figure 1. A typical event with one gap overlapping $Y/2$.
- Figure 2. A double diffraction dissociation diagram.
- Figure 3. A double diffraction dissociation diagram with two particles singled out.
- Figure 4a. A multiple diffraction dissociation diagram.
- Figure 4b. A typical event with N gaps overlapping $Y/2N$,

$$\frac{3Y}{2N}, \dots, \frac{2N-1}{2N} Y.$$

FIGURE 1

FIGURE 2

FIGURE 3

FIGURE 4a

FIGURE 4b