Diffraction Association and Flavoring of Pomeron

Chung-I Tan

(1)Department of Physics, Brown University, Providence, RI 02912, USA

Abstract

The most important consequence of Pomeron being a pole is the factorization property. However, due to Pomeron intercept being greater than 1, the extrapolated single diffraction dissociation cross section based on a classical triple-Pomeron formula is too large leading to a potential unitarity violation at Tevatron energies. It is our desire here to point out that the “flavoring” of Pomeron plays the dominant role in resolving this apparent “paradox”.

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1 Introduction

One of the more interesting developments from recent collider experiments is the finding that hadronic total cross sections as well as elastic cross sections in the near-forward limit can be described by the exchange of a “soft Pomeron” pole. i.e., the absorptive part of the elastic amplitudes can be approximated by \( Im T_{a,b}(s, t) \simeq \beta_a(t)s^{\alpha_P(t)}\beta_b(t) \). The Pomeron trajectory has two important features. First, its zero-energy intercept is greater than one, \( \alpha_P(0) \equiv 1 + \epsilon, \epsilon \simeq 0.08 \sim 0.12 \), leading to rising \( \sigma^{tot}(s) \). Second, its Regge slope is approximately \( \alpha'_P \simeq 0.25 \sim 0.3 GeV^{-2} \), leading to the observed shrinkage effect for elastic peaks. The most important consequence of Pomeron being a pole is factorization. For a singly diffractive dissociation process, factorization

\[
\frac{d\sigma}{d\xi_t} \simeq (s/\sqrt{s})^{-\epsilon} \lesssim 1 \quad \text{for} \quad 1 \lesssim s/\sqrt{s} \lesssim 25
\]

Several phenomenological suggestions for this modification factor have been made:

- It was argued by K. Goulianos [3] that agreement with data could be achieved by having an energy-dependent suppression factor,

\[
Z(\xi, t; s) \rightarrow N(s)^{-1} \leq 1,
\]

so that a new “Pomeron flux”, \( F_N(s, \xi, t) \equiv N(s)^{-1}F_{P/a}^{cl}(\xi, t) \), is always normalized to be less than unity: \( N(s) \simeq (s/\bar{s})^{2\epsilon} > 1 \) for \( s \geq \bar{s} \), and \( N(s) = 1 \) for \( s < \bar{s} \), where \( \sqrt{s} \simeq 22 GeV \).

- An alternative suggestion has been made recently by P. Schlein, [4] by introducing a “flux damping” factor,

\[
Z(\xi, t; s) \rightarrow D(\xi).
\]

This factor has the following features: (i) \( D(\xi) \sim 0(1) \), for \( 1 > \xi > \xi_1 \), (ii) \( D(\xi) \) drops by a factor of 2 as \( \xi \) decreases from \( \xi_1 \) to \( \xi_2 \), and (iii) \( D(\xi) \rightarrow 0 \) rapidly for \( 0 \leq \xi < \xi_2 \), with \( \xi_1 \sim 0.015 \) and \( \xi_2 \sim 10^{-4} \).

In view of the factorization property for total and elastic cross sections, the “flux renormalization” procedure, which breaks factorization, appears paradoxical. On the other hand, the occurrence of unusual scales for the Schlein damping factor, e.g., \( \xi_2 \sim 10^{-4} \), appears equally mysterious. We shall refer to this as “Dino’s paradox”. Finding a resolution consistent with Pomeron pole dominance for elastic and total cross sections at Tevatron energies will be the main focus of this study. [4]

A natural expectation for the resolution to this paradox lies in implementing a large screening correction to the classical triple-Pomeron formula. However, this appears too simplistic. In the absence of a new energy scale, a screening factor of the order \( 5 \sim 10 \), if obtained, would apply both at Tevatron energies and at ISR energies. This indeed is the case for the eikonalization analysis by Gotsman, Levin, and Maor, [5] as pointed out by Goulianos. Since a successful triple-Pomeron phenomenology exists up to ISR energies, a subtler explanation is required. We shall assume that
any screening effect can supply at the most a 10 \sim 20\% suppression and it cannot serve as the primary mechanism for explaining the paradox.

Triple-Regge phenomenology has had a long history. It has enjoyed many successes since early seventies, and it should emerge as a feature of any realistic representation of non-perturbative QCD for high energy scattering. In particular, it should be recognized that, up to ISR energies, triple-Pomeron phenomenology has provided a successful description for the phenomenon of diffractive dissociation. A distinguishing feature of the successful low-energy triple-Pomeron analyses is the value of the Pomeron intercept. It has traditionally been taken to be near 1, which would lead to total cross sections having constant “asymptotic values”. In contrast, the current paradox centers around the Pomeron having an intercept greater than 1, e.g., $\epsilon \simeq 0.1$.

Instead of trying to ask “how can one obtain a large suppression factor at Tevatron energies”, an alternative approach can be adopted. We could first determine the “triple-Pomeron” coupling by matching the diffractive cross section at the highest Tevatron energy. A naive extrapolation to lower energies via a standard triple-Pomeron formula would of course lead to too small a cross section at ISR energies. We next ask the question:

- Are there physics which might have been overlooked by others in moving down in energies?
- In particular, how can a high energy fit be smoothly interpolated with the successful low energy triple-Pomeron analysis using a Pomeron with intercept at 1, i.e., $\epsilon \simeq 0$.

A key observation which will help in understanding our proposed resolution concerns the fact that, even at Tevatron energies, various “subenergies”, e.g., the missing mass squared, $M^2$, and the diffractive “gap”, $\xi^{-1}$, can remain relatively small, comparable to the situation of ISR energies for the total cross sections. Our analysis has identified the “flavoring” of Pomeron \cite{6,7} as the primary dynamical mechanism for resolving the paradox. A proper implementation of final-state screening correction, (or final-state unitarization), assures a unitarized “gap distribution”, with flavoring setting the relevant energy scale. We find that initial-state screening remains unimportant, consistent with the pole dominance picture for elastic and total cross section hypothesis at Tevatron energies. In fact, we shall concentrate in the present discussion mostly on the flavoring and will comment only briefly on screening. (A more complete treatment of final-state screening can be found in the Ref. 1 , and we find that the effect turns out to be small.)

2 Soft Pomeron at Low Energies

In order to be able to answer the questions we have posed, it is necessary to first provide a dynamical picture for a soft Pomeron and to briefly review the notion of “Harari-Freund” duality.

2.1 Harari-Freund Duality

Although Regge phenomenology pre-dated QCD, it is important to recognize that it can be understood as a phenomenological realization of non-perturbative QCD in a “topological expansion”, e.g., the large-$N_c$ expansion. In particular, an important feature of a large-$N_c$ expectation is emergence of the Harari-Freund two-component picture. \cite{8}

For $P_{lab} \leq 20$ GeV/c, it was recognized that the imaginary part of any hadron two-body amplitude can be expressed approximately as the sum of two terms:

$$Im A(s, t) = R(s, t) + P(s, t).$$
From the s-channel point of view, $R(s,t)$ represents the contribution of s-channel resonance while $P(s,t)$ represents the non-resonance background. From the t-channel point of view, $R(s,t)$ represents the contribution of “ordinary” t-channel Regge exchanges and $P(s,t)$ represents the diffractive part of the amplitude given by the Pomeron exchange. Three immediate consequences of this picture are:

- (a) Imaginary parts of amplitudes which show no resonances should be dominated by Pomeron exchange, ($R \simeq 0$, and $P \simeq \text{constant}$).
- (b) Imaginary parts of $A(s,t)$ which have no Pomeron term should be dominated by s-channel resonances,
- (c) Imaginary parts of amplitudes which do not allow Pomeron exchange and show no resonances should vanish,

Point (b) can best be illustrated by partial-wave projections of $\pi N \to \pi N$ scattering amplitudes from well-defined t-channel isospin exchanges. Point (c) is best illustrated by examining the $K^+ p \to K^0 n$, where, by optical theroyem, $\text{Im}A(K^+ p \to K^0 n) \propto \sigma_{tot}(K^+ p) - \sigma_{tot}(K^0 n)$. The near-equality of these two cross sections, from the t-channel exchange view point, reflects the interesting feature of exchange degeneracy for secondary Reggeons. Finally, let us come to the point (a). From the behavior of $\sigma_{\pi^\pm p}$, $\sigma_{K^\pm p}$, $\sigma_{pp}$ and $\sigma_{\bar{p}p}$, one finds that the near-constancy for the $P$-contribution corresponds to having an effective “low-energy” Pomeron intercept at 1, i.e.,

$$\alpha_P^{\text{low}}(0) \simeq 1.$$ 

### 2.2 Shadow Picture and Inelastic Production

A complementary treatment of Pomeron at low energies is through the analysis of inelastic production, which is responsible for the non-resonance background mentioned earlier. Diffraction scattering as the shadow of inelastic production has been a well established mechanism for the occurrence of a forward peak. Analyses of data up to ISR energies have revealed that the essential feature of nondiffractive particle production can be understood in terms of a multipertipheral cluster-production mechanism. In such a picture, the forward amplitude at high energies is predominantly absorptive and is dominated by the exchange of a “bare Pomeron”.

In a “shadow” scattering picture, the “minimum biased” events are predominantly “short-range ordered” in rapidity and the production amplitudes can be described by a multiperipheral cluster model. Under a such an approximation to production amplitudes for the right-hand side of an elastic unitary equation, $\text{Im}T(s,0) = \sum_{n} |T_{2,n}|^2$, one finds that the resulting elastic amplitude is dominated by the exchange of a Regge pole, which we shall provisionally refer to as the “bare Pomeron”. Next consider singly diffractive events. We assume that the “missing mass” component corresponds to no gap events, thus the distribution is again represented by a “bare Pomeron”. However, for the gap distribution, one would insert the “bare Pomeron” just generated into a production amplitude, thus leading to the classical triple-Pomeron formula.

Extension of this procedure leads to a “perturbative” expansion for the total cross section in the number of bare Pomeron exchanges along a multiperipheral chain. Such a framework was proposed long time ago, with the understanding that the picture can make sense at moderate energies, provided that the the intercept of the Pomeron is near one, $\alpha(0) \simeq 1$, or less.

However, with the acceptance of a Pomeron having an intercept greater than unity, this expansion must be embellished or modified. It is quite likely that the resolution for Dino’s paradox lies in understanding how such an effect can be accomodated within this framework, consistent with the Pomeron pole dominance hypothesis.
2.3 Bare Pomeron in Non-Perturbative QCD

In a non-perturbative QCD setting, the Pomeron intercept is directly related to the strength of the short-range order component of inelastic production and this can best be understood in a large-$N$ expansion. In such a scheme, particle production mostly involves emitting “low-mass pions”, and the basic energy scale of interactions is that of ordinary vector mesons, of the order of 1 GeV. In a one-dimensional multiperipheral realization for the “planar component” of the large-$N$ QCD expansion, the high energy behavior of a $n$-particle total cross section is primarily controlled by its longitudinal phase space, $\sigma_n \simeq (g^4 N^2/(n-2)!)(g^2 N \log s)^{n-2} s^{1/y}$, and, upon summing over $n$, one arrives at Regge behavior for the planar component of $\sigma^{\text{tot}}$ where

$$\alpha_R = (2\alpha_R - 1) + g^2 N. \tag{2}$$

At next level of cylinder topology, the contribution to partial cross section increases due to its topological twists, $\sigma_n \simeq (g^4/(n-2)!)(g^2 N \log s)^{n-2} s^{1/y}$, and, upon summing over $n$, one arrives at a total cross section governed by a Pomeron exchange, $\sigma^{\text{tot}}_0(Y) = g^4 e^{\alpha_P Y}$, where the Pomeron intercept is

$$\alpha_P = (2\alpha_R - 1) + 2g^2 N. \tag{3}$$

Combining Eq. (2) and Eq. (3), we arrive at an amazing “bootstrap” result, $\alpha_P \simeq 1$.

In a non-perturbative QCD setting, having a Pomeron intercept near 1 therefore depends crucially on the topological structure of large-$N$ non-Abelian gauge theories. In this picture, one has $\alpha_R \simeq .5 \sim .7$ and $g^2 N \simeq .3 \sim .5$. With $\alpha' \simeq 1 \text{ GeV}^{-2}$, one can also directly relate $\alpha_R$ to the average mass of typical vector mesons. Since vector meson masses are controlled by constituent mass for light quarks, and since constituent quark mass is a consequence of chiral symmetry breaking, the Pomeron and the Reggeon intercepts are directly related to fundamental issues in non-perturbative QCD. This picture is in accord with the Harari-Freund picture for low-energy Regge phenomenology.

Finally we note that, in a Regge expansion, the relative importance of secondary trajectories to the Pomeron is controlled by the ratio $e^{\alpha_R y}/e^{\alpha_P y} = e^{-(\alpha_P - \alpha_R) y}$. It follows that there exists a natural scale in rapidity, $y_r$, $(\alpha_P - \alpha_R)^{-1} < y_r \simeq 3 \sim 5$. The importance of this scale $y_r$ is of course well known: When using a Regge expansion for total and two-body cross sections, secondary trajectory contributions become important and must be included whenever rapidity separations are below 3 \sim 5 units. This scale of course is also important for the triple-Regge region: There are two relevant rapidity regions: one associated with the “rapidity gap”, $y \equiv \log \xi^{-1}$, and the other for the missing mass, $y_m \equiv \log M^2$.

2.4 Conflict with Donnachie-Lanshoff Picture

It has become increasingly popular to use the Donnachie-Landshoff picture where Pomeron intercept above one, i.e., $\epsilon \sim 0.1$. Indeed, it is impressive that various cross sections can be fitted via Pomeron pole contribution over the entire currently available energy range. However, it should be pointed out that Donnachie-Lanshoff picture is not consistent with the Harari-Freund picture at low energies.

It can be argued that the difference between these two approaches should not be important at high energies. This is certainly correct for total cross sections. However, we would like to stress that this is not the true for diffractive dissociation, even at Tevatron energies. This can best be understood in terms of rapidity variables, $y$ and $y_m$. Since $y + y_m \simeq Y$, $Y \equiv \log s$, it follows that,
even at Tevatron energies, the rapidity range for either $y$ or $y_m$ is more like that for a total cross section at or below the ISR energies. Therefore, details of diffractive dissociation cross section at Tevatron would depend on how a Pomeron is treated at relatively low subenergies.

3 Soft Pomeron and Flavoring

Consider for the moment the following scenario where one has two different fits to hadronic total cross sections:

- **(a) “High energy fit”**: $\sigma_{ab}(y) \sim \beta_a \beta_b e^{\epsilon y}$ for $y >> y_f$,

- **(b) “Low energy fit”**: $\sigma_{ab}(y) \sim \beta_{a,\text{low}} \beta_{b,\text{low}}$ for $y << y_f$.

That is, we envisage a situation where the “effective Pomeron intercept”, $\epsilon_{\text{eff}}$, increases from 0 to $\epsilon \sim 0.1$ as one moves up in energies. In order to have a smooth interpolation between these two fits, one can obtain the following order of magnitude estimate $\beta_p \simeq e^{-\frac{\epsilon y_f}{\sqrt{s}}} \beta_{p,\text{low}}^{\text{low}}$. Modern parametrization for Pomeron residues typically leads to values of the order $(\beta_p)^2 \simeq 14 \sim 17$ mb. However, before the advent of the notion of a Pomeron with an intercept greater than 1, a typical parametrization would have a value $(\beta_p^{\text{low}})^2 \simeq 35 \sim 40$ mb, accounting for a near constant Pomeron contribution at low energies. This leads to an estimate of $y_f \sim 8$, corresponding to $\sqrt{s} \sim 50$ GeV. This is precisely the energy scale where a rising total cross section first becomes noticeable.

The scenario just described has been referred to as “flavoring”, the notion that the underlying effective degrees of freedom for Pomeron will increase as one moves to higher energies, and it has provided a dynamical basis for understanding the value of Pomeron intercept in a non-perturbative QCD setting. In this scheme, in order to extend a Regge phenomenology to low energies, both the Pomeron intercept and the Pomeron residues are scale-dependent. We shall review this mechanism shortly. However, we shall first introduce a scale-dependent formalism where the entire flavoring effect can be absorbed into a flavoring factor, $R(y)$, associated with each Pomeron propagator.

3.1 Effective intercept and Scale-Dependent Treatment

In order to be able to extend a Pomeron representation below the rapidity scale $y \sim y_f$, we propose the following **scale-dependent** scheme where we introduce a flavoring factor for each Pomeron propagator. Since each Pomeron exchange is always associated with energy variable $s$, (therefore a rapidity variable $y \equiv \log s$), we shall parametrize the Pomeron trajectory function as

$$\alpha_{\text{eff}}(t; y) \simeq 1 + \epsilon_{\text{eff}}(y) + \alpha' t,$$

where $\epsilon_{\text{eff}}(y)$ has the properties

- $\epsilon_{\text{eff}} \simeq \epsilon \simeq 0.1$ for $y >> y_f$,

- $\epsilon_{\text{eff}} \simeq \epsilon_0 \equiv \alpha_{p,\text{low}} - 1 \simeq 0$ for $y << y_f$.

For instance, exchanging such an effective Pomeron leads to a contribution to the elastic cross section $T_{ab}(s, t) \propto s^{\frac{1}{\epsilon_{\text{eff}}(y)} + \alpha' t}$. This representation can now be extended down to the region $y \sim y_r$. We shall adopt a particularly convenient parametrization for $\epsilon_{\text{eff}}(y)$ in the next Section when we discuss phenomenological concerns.

To complete the story, we need also to account for the scale dependence of Pomeron residues. What we need is an “interpolating” formula between the high energy and low energy sets. Once
a choice for $\epsilon_{eff}(y)$ has been made, it is easy to verify that a natural choice is simply $\beta^{eff}_a(y) = \beta_a e^{(1-\epsilon_{eff}(y))y}$. It follows that the total contribution from a “flavored” Pomeron corresponds to the following low-energy modification

$$T_{a,b}^{cl}(y, t) \rightarrow R(y) T_{a,b}^{cl}(y, t),$$

where $T_{a,b}^{cl}(y, t) \equiv \beta_a \beta_b e^{(1+\epsilon_{eff}(y))y}$ is the amplitude according to a “high energy” description with a fixed Pomeron intercept, and

$$R(y) \equiv e^{-\epsilon_{eff}(y)(y-y_f)},$$

is a “flavoring” factor. The effect of this modification can best be illustrated via Figure 1.

This flavoring factor should be consistently applied as part of each “Pomeron propagator”. With the normalization $R(\infty) = 1$, we can therefore leave the residues alone, once they have been determined by a “high energy” analysis. For our single-particle gap cross section, since there are three Pomeron propagators, the renormalization factor is given by the following product:

$$Z = R^2(y) R(y_m).$$

It is instructive to plot in Figure 2 this combination as a function of either $\xi$ or $M^2$ for various fixed values of total rapidity, $Y$.

### 3.2 Flavoring of Bare Pomeron

We have proposed sometime ago that “baryon pair”, together with other “heavy flavor” production, provides an additional energy scale, $s_f = e^y_f$, for soft Pomeron dynamics, and this effect can be responsible for the perturbative increase of the Pomeron intercept to be greater than unity, $\alpha_P(0) \sim 1 + \epsilon$, $\epsilon > 0$. One must bear this additional energy scale in mind in working with a soft Pomeron. [8] That is, to fully justify using a Pomeron with an intercept $\alpha_P(0) > 1$, one must restrict oneself to energies $s > s_f$ where heavy flavor production is no longer suppressed. Conversely, to extrapolate Pomeron exchange to low energies below $s_f$, a lowered “effective trajectory” must be used. This feature of course is unimportant for total and elastic cross sections at Tevatron energies. However, it is important for diffractive production since both $\xi^{-1}$ and $M^2$ will sweep right through this energy scale at Tevatron energies.

Flavoring becomes important whenever there is a further inclusion of effective degrees of freedom than that associated with light quarks. This can again be illustrated by a simple one-dimensional multiperipheral model. In addition to what is already contained in the Lee-Veneziano model, suppose that new particles can also be produced in a multiperipheral chain. Concentrating on the cylinder level, the partial cross sections will be labelled by two indices,

$$\sigma_{p,q} \simeq (g^4/p!)2^{p+q}(g^2 N \log s)^p(g_f^2 N \log s)^q s^{-1} \epsilon_{eff}^{-1},$$

where $q$ denotes the number of clusters of new particles produced. Upon summing over $p$ and $q$, we obtain a “renormalized” Pomeron trajectory

$$\alpha_p^{high} = \alpha_p^{low} + \epsilon,$$

where $\alpha_p^{low} \simeq 1$ and $\epsilon \simeq 2g_f^2N$. That is, in a non-perturbative QCD setting, the effective intercept of Pomeron is a dynamical quantity, reflecting the effective degrees of freedom involved in near-forward particle production.[9]

If the new degree of freedom involves particle production with high mass, the longitudinal phase space factor, instead of $(\log s)^q$, must be modified. Consider the situation of producing one $NN$ bound state together with pions, i.e., $p$ arbitrary and $q = 1$ in Eq. (5). Instead of $(\log s)^{p+1}$,
each factor should be replaced by \((\log(s/m_{\text{eff}}^2))^p+1\), where \(m_{\text{eff}}\) is an effective mass for the \(N\bar{N}\) cluster. In terms of rapidity, the longitudinal phase space factor becomes \((Y - \delta)^p+1\), where \(\delta\) can be thought of as a one-dimensional “excluded volume” effect. For heavy particle production, there will be an energy range over which keeping up to \(q = 1\) remains a valid approximation. Upon summing over \(p\), one finds that the additional contribution to the total cross section due to the production of one heavy-particle cluster is \(\sigma_q^\text{eff} = \sigma_0^\text{total}(Y - \delta)(2p_f^2 N)\log(Y - \delta)\theta(Y - \delta)\), where \(\alpha^\text{low}_P \simeq 1\). Note the effective longitudinal phase space “threshold factor”, \(\theta(Y - \delta)\), and, initially, this term represents a small perturbation to the total cross section obtained previously, (corresponding to \(q = 0\) in Eq. (3)), \(\sigma_0^\text{total}\). Over a rapidity range, \([\delta, \delta + \delta_f]\), where \(\delta_f\) is the average rapidity required for producing another heavy-mass cluster, this is the only term needed for incorporating this new degree of freedom. As one moves to higher energies, “longitudinal phase space suppression” becomes less important and more and more heavy particle clusters will be produced. Upon summing over \(q\), we would obtain a new total cross section, described by a renormalized Pomeron, with a new intercept given by Eq. (1).

We assume that, at Tevatron, the energy is high enough so that this kind of “threshold” effects is no longer important. How low an energy do we have to go before one encounter these effects? Let us try to answer this question by starting out from low energies. As we have stated earlier, for \(Y > 3 \sim 5\), secondary trajectories become unimportant and using a Pomeron with \(\alpha \simeq 1\) becomes a useful approximation. However, as new flavor production becomes effective, the Pomeron trajectory will have to be renormalized. We can estimate for the relevant rapidity range when this becomes important as follows: \(y_f > 2\delta_0 + < q >_{\text{min}} \delta_f\). The first factor \(\delta_0\) is associated with leading particle effect, i.e., for proton, this is primarily due to pion exchange. \(\delta_f\) is the minimum gap associated with one heavy-mass cluster production, e.g., nucleon-antinucleon pair production. We estimate \(\delta_0 \simeq 2\) and \(\delta_f \simeq 2 \sim 3\), so that, with \(< q >_{\text{min}} \simeq 2\), we expect the relevant flavoring rapidity scale to be \(y_f \simeq 8 \sim 10\).

4 A Caricature of High Energy Diffractive Dissociation

Both the screening function and the flavoring function depend on the effective Pomeron intercept, and we shall adopt the following simple parametrization. The transition from \(\alpha^\text{low}_P(0) = 1 + \epsilon_o\) to \(\alpha^\text{high}_P(0) = 1 + \epsilon\) will occur over a rapidity range, \((y_f^{(1)}, y_f^{(2)})\). Let \(y_f \equiv \frac{1}{2}(y_f^{(1)} + y_f^{(2)})\) and \(\lambda_f^{-1} \equiv \frac{1}{2}y_f^{(2)} - y_f^{(1)}\). Similarly, we also define \(\epsilon \equiv \frac{1}{2}(\epsilon + \epsilon_o)\) and \(\Delta \equiv \frac{1}{2}(\epsilon - \epsilon_o)\). A convenient parametrization for \(\epsilon_{\text{eff}}\) we shall adopt is \(\epsilon_{\text{eff}}(y) = \frac{\epsilon}{\Delta} \tan\lambda_f(y - \tilde{y}_f)\). The combination \([\epsilon - \epsilon_{\text{eff}}(y)]\) can be written as \((2\bar{\epsilon})(1 + (s/s_f)^{2\lambda_f})^{-1}\) where \(s_f = e^{\gamma_f}\). We arrive at a simple parametrization for our flavoring function

\[
R(s) \equiv (\frac{S_f}{s})^{(2\bar{\epsilon})(1 + (\frac{s}{s_f})^{2\lambda_f})^{-1}}.
\]

With \(\alpha^\text{low}_P \simeq 1\), we have \(\epsilon_o \simeq 0, \bar{\epsilon} \simeq \Delta \simeq \epsilon/2\), and we expect that \(\lambda_f \simeq 1\) and \(y_f \simeq 8 \sim 10\) are reasonable range for these parameters. 

The most important new parameter we have introduced for understanding high energy diffractive production is the flavoring scale, \(s_f = e^{\gamma_f}\). We have motivated by way of a simple model to show that a reasonable range for this scale is \(y_f \simeq 8 \sim 10\). Quite independent of our estimate, it is possible to treat our proposed resolution phenomenologically and determine this flavoring scale from experimental data.
It should be clear that one is not attempting to carry out a full-blown phenomenological analysis here. To do that, one must properly incorporate other triple-Regge contributions, e.g., the \(P\overline{P}PP\)-term for the low-\(y_m\) region, the \(\pi\pi\overline{P}\)-term and/or the \(R\overline{P}P\)-term for the low-\(y\) region, etc., particularly for \(\sqrt{s} \approx \sqrt{s_f} \sim 100\ \text{GeV}\). There are also “interference terms, e.g., \(R\overline{P}P\), to take into account. What we hope to achieve is to provide a “caricature” of the interesting physics involved in diffractive production at collider energies through our introduction of the flavoring factors.\(\uparrow\)

Let us begin by first examining what we should expect. Concentrate on the triple-Pomeron vertex \(g(0)\) measured at high energies. Let us for the moment assume that it has also been measured reliably at low energies, and let us denote it as \(g^{\text{low}}(0)\). Our flavoring analysis indicates that these two couplings are related by
\[ g^{\text{low}}(0) = \frac{1}{3\lambda y} g^{\text{PPP}\overline{P}}(0). \]
With \(\epsilon \approx 0.08 \sim 0.1\) and \(y_f \approx 8 \sim 10\), using the value \(g^{\text{low}}(0) = 0.364 \pm 0.025\ \text{mb}^2\), we expect a value of \(0.12 \sim 0.18\ \text{mb}^2\) for \(g^{\text{PPP}\overline{P}}(0)\). Denoting the overall multiplicative constant for our renormalized triple-Pomeron formula by \(K \equiv \beta_2^2(0) g^{\text{PPP}\overline{P}}(0) \beta_0(0)/16\pi\). With \(\beta_2^2 \approx 16\ \text{mb}\), we therefore expect \(K\) to lie between the range \(0.15 \sim 0.25\ \text{mb}^2\).

We begin testing our renormalized triple-Pomeron formula by first determining the overall multiplicative constant \(K\) by normalizing the integrated \(\sigma^{sd}\) to the measured CDF \(\sqrt{s} = 1800\ \text{GeV}\) value. With \(\epsilon = 0.1, \lambda_y = 1\), this is done for a series of values for \(y_f\) = 7, 8, 9, 10. We obtain respective values for \(K\) = 0.24, 0.21, 0.18, 0.15, consistent with our flavoring expectation. As a further check on the sensibility of these values for the flavoring scales, we find for the ratio \(\rho \equiv \sigma^{sd}(546)/\sigma^{sd}(1800)\) the values 0.63, 0.65, 0.68, 0.72 respectively. This should be compared with the CDF result of 0.834.

Having shown that our renormalized triple-Pomeron formula does lead to sensible predictions for \(\sigma^{sd}\) at Tevatron, we can improve the fit by enhancing the \(P\overline{P}P\)-term as well as \(R\overline{P}P\)-terms which can become important. Instead of introducing a more involved phenomenological analysis, we simulate the desired low energy effect by having \(\epsilon_0 \approx -0.06 \sim -0.08\). A remarkably good fit results with \(\epsilon = 0.08 \sim 0.09\) and \(y_f = 9\),\(\uparrow\) This is shown in Figure \(\uparrow\). The ratio \(\rho\) ranges from 0.78 \sim 0.90, which is quite reasonable. The prediction for \(\sigma^{sd}\) at LHC is \(12.6 \sim 14.8\ \text{mb}\).

Our fit leads to a triple-Pomeron coupling in the range of
\[ g^{\text{PPP}\overline{P}}(0) \approx 0.12 \sim 0.18\ \text{mb}^2, \] (8)

exactly as expected. Interestingly, the triple-Pomeron coupling quoted in Ref. 3 \((g(0) = 0.69\ \text{mb}^2)\) is actually a factor of 2 larger than the corresponding low energy value.\(\uparrow\) Note that this difference of a factor of 5 correlates almost precisely with the flux renormalization factor \(N(s) \approx 5\) at Tevatron energies.

5 Final Remarks:

In Ref. 1, a more elaborated treatment has been carried out where both the flavoring and the final-state screening effects were considered. We have shown, given Pomeron as a pole, the total Pomeron contribution to a singly diffractive dissociation cross section can in principle be expressed as
\[ \frac{d\sigma}{d\alpha (t)} = [S_i(s,t)][D_{\alpha\overline{P}}(\xi, t)][\sigma_{\overline{P}P}(M^2, t)], \]
and \(D_{\alpha\overline{P}}(\xi, t) = S_f(\xi, t)F_{\overline{P}/a}(\xi, t)\).

- The first term, \(S_i\), represents initial-state screening correction. We have demonstrated that, with a Pomeron intercept greater than unity and with a pole approximation for total and...
elastic cross sections remaining valid, initial-state absorption cannot be large. We therefore can justify setting $S_i \simeq 1$ at Tevatron energies.

- The first crucial step in our alternative resolution to the Dino's paradox lies in properly treating the final-state screening, $S_f(\xi, t)$. We have explained in an expanding disk framework that the final-state screening sets in at a rapidity scale determined by the flavoring scale, $y_f$, which correlates well with the mysterious scale, $\xi_2$, of Schlein.

- We have stressed that the dynamics of a soft Pomeron in a non-perturbative QCD scheme requires taking into account the effect of “flavoring”, the notion that the effective degrees of freedom for Pomeron is suppressed at low energies. As a consequence, we find that $F_{P/a}(\xi, t) = R^2(\xi^{-1})F_{P/a}^{cl}(\xi, t)$ and $\sigma_{Pb}(M^2, t) = R(M^2)\sigma_{Pb}^{cl}(M^2, t)$ where $R$ is the “flavoring” factor discussed in this paper.

It should be stressed that our discussion depends crucially on the notion of soft Pomeron being a factorizable Regge pole. This notion has always been controversial. Introduced more than thirty years ago, Pomeron was identified as the leading Regge trajectory with quantum numbers of the vacuum with $\alpha(0) \simeq 1$ in order to account for the near constancy of the low energy hadronic total cross sections. However, as a Regge trajectory, it was unlike others which can be identified by the particles they interolate. With the advent of QCD, the situation has improved, at least conceptually. Through large-$N_c$ analyses and through other non-perturbative studies, it is natural to expect Regge trajectories in QCD as manifestations of “string-like” excitations for bound states and resonances of quarks and gluons due to their long-range confining forces. Whereas ordinary meson trajectories can be thought of as “open strings” interpolating $qq$ bound states, Pomeron corresponds to a “closed string” configuration associated with glueballs. However, the difficulty of identification, presumably due to strong mixing with multi-quark states, has not helped the situation in practice. In a simplified one-dimensional multiperipheral realization of large-$N_c$ QCD, the non-Abelian gauge nature nevertheless managed to re-emerge through its topological structure. [10]

The observation of “pole dominance” at collider energies has hastened the need to examine more closely various assumptions made for Regge hypothesis from a more fundamental viewpoint. It is our hope that by examining Dino’s paradox carefully and by finding an alternative resolution to the problem without deviating drastically from accepted guiding principles for hadron dynamics, Pomeron can continued to be understood as a Regge pole in a non-perturbative QCD setting. The resolution for this paradox could therefore lead to a re-examination of other interesting questions from a firmer theoretical basis. For instance, to be able to relate quantities such as the Pomeron intercept to non-perturbative physics of color confinement represents a theoretical challenge of great importance.

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Figure 1: Effect of flavoring factor $R(s)$ when applied to a standard rising cross section: $\sigma^{cl} = \beta^2 s^\epsilon$, $\epsilon = 0.1$ and $\beta^2 = 16 \text{ mb}$, given by the solid curve. With $R(y)$ given by Eq. (7), the dashed-dotted curve has $\epsilon_o = 0$, $\lambda_f = 1$, and flavoring scale $y_f = 9$, and the dotted curve corresponds to $\epsilon_o = -0.04$. 
Figure 2: Renormalization factor due to flavoring alone, $Z_f(\xi; s) \equiv R^2(\xi^{-1})R(M^2)$, as a function of rapidity $y = \log \xi^{-1}$ for various fixed center of mass energies. These curves correspond to parameters used for the solid line in Figure 3.
Figure 3: Various fits to representative single diffraction dissociation cross sections extracted from Ref. 3 from ISR to Tevatron. The solid line and the dotted curve correspond to $\epsilon = 0.08$, $\epsilon_0 = -0.07$, $\lambda_f = 1$, $y_f = 9$, with small amount of final-state screening. The dashed-dotted curve corresponds to no screening.