Flavoring of Pomeron and Diffractive Production at Tevatron Energies

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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 \( \text{Im} \ T_{a,b}(s,t) \approx \beta_a(t)s^{\alpha_P(t)}\beta_b(t) \). The Pomeron trajectory has two important features. First, its zero-energy intercept is greater than one, \( \alpha_P(0) \equiv 1 + \epsilon, \epsilon \approx 0.08 \sim 0.12 \), leading to rising \( \sigma_{\text{tot}}(s) \). Second, its Regge slope is approximately \( \alpha'_P \approx 0.25 \sim 0.3 \ \text{GeV}^{-2} \), 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 leads to a “classical triple-Pomeron” formula, \( \frac{d\sigma}{dtd\xi} \to \frac{d\sigma_{\text{classical}}}{dtd\xi} \equiv F_{P/P}^{\text{cl}}(\xi, t)\sigma_{P/P}^{\text{cl}}(M^2, t), \) where \( M^2 \) is the missing mass variable and \( \xi \equiv M^2/s \). The first term, \( F_{P/P}^{\text{cl}}(\xi, t) \), is the so-called “Pomeron flux”, and the second term is the “Pomeron-particle” total cross section. With \( \epsilon \sim 0.1 \), it has been observed that the extrapolated \( p\bar{p} \) single diffraction dissociation cross section, \( \sigma_{\text{sd}} \), based on the standard triple-Pomeron formula is too large at Tevatron energies by as much as a factor of \( 5 \sim 10 \) and it could become larger than the total cross section.

Let us denote the singly diffractive cross section as a product of a “renormalization” factor and the classical formula,

\[
\frac{d\sigma}{dtd\xi} = Z(\xi, t; s)\frac{d\sigma_{\text{classical}}}{dtd\xi}.
\]  

It was argued by K. Goulianos in Ref. 3 that agreement with data could be achieved by having an energy-dependent suppression factor, \( Z(\xi, t; s) \to Z_G(s) \equiv N(s)^{-1} \leq 1 \), so that the new “Pomeron flux”, \( F_N(s, \xi, t) \equiv N(s)^{-1}F_{P/P}^{\text{cl}}(\xi, t) \), is normalized to unity for \( s \geq \bar{s}, \sqrt{\bar{s}} \approx 22 \ \text{GeV} \). An alternative suggestion has been made recently by P. Schlein, where \( Z \to Z_S(\xi) \).

In view of the factorization property for total and elastic cross sections, the “flux renormalization” procedure appears paradoxical. We shall refer to this as “Dino’s paradox”. Finding a resolution that is consistent with Pomeron pole dominance for elastic and total cross sections at Tevatron energies will be the main focus of this study. In particular, we want to maintain the following factorization property, \( \frac{d\sigma}{dtd\xi} \to \sum_k D_k(\xi, t)\Sigma_k(M^2) \), when \( \xi^{-1} \) and \( M^2 \) become large.

It is generally expected that the resolution to the paradox should lie in a proper implementation of screening corrections to the classical triple-Pomeron formula. Our treatment lies in a proper implementation of final-state screening correction, (or final-state unitarization), with “flavoring” for Pomeron as the primary
dynamical mechanism for setting the relevant energy scale. In our treatment, 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 only on the flavoring whereas the treatment of final-state screening can be found in the Ref. 1 and the effect turns out to be small.

2 Dynamics for Soft Pomeron and Flavoring

Although we have shown in Ref. 1 that final-state screening would automatically avoid unitarity violation, the primary source of high energy suppression actually comes from a proper treatment of scale-dependence for Pomeron couplings. Consider for the moment the following scenario where one has two different fits to hadronic total cross sections:

• (a) “High energy fit”: \( \sigma_{ab}(Y) \simeq \beta_a \beta_b e^{\epsilon Y} \) for \( Y >> y_f \),

• (b) “Low energy fit”: \( \sigma_{ab}(Y) \simeq \beta_{a}^{\text{low}} \beta_{b}^{\text{low}} \) for \( 0 < Y << y_f \).

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 \( \beta_p \simeq e^{-\frac{\epsilon y_f}{2} \beta_{p}^{\text{low}}} \). Modern parametrization for Pomeron residues typically leads to values of the order \( (\beta_{p}^{\text{low}})^2 \simeq 14 \sim 17 \text{ mb} \). However, before the advent of the notion of a Pomeron with an intercept greater than 1, a typical parametrization would have a value \( (\beta_{p}^{\text{low}})^2 \simeq 35 \sim 40 \text{ mb} \), accounting for a near constant Pomeron contribution at low energies. This leads to an estimate of \( y_f \sim 8 \), corresponding to \( \sqrt{s} \sim 50 \text{ GeV} \). 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, \( y_f \) and it has provided a dynamical basis for understanding the value of Pomeron intercept in a non-perturbative QCD setting. In this scheme, both the Pomeron intercept and the Pomeron residues are scale-dependent. We shall briefly review this mechanism and introduce a scale-dependent formalism where the entire flavoring effect can be absorbed into a flavoring factor, \( R(y) \), associated with each Pomeron propagator.

2.1 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, \[ \sigma_n \simeq \frac{g^4 N^2}{(n-2)!}(g^2 N \log s)^{n-2}s^J_{eff}^{1/2}. \] Since there are only Reggeons at the planar diagram level, one has \( J_{eff} = 2\alpha_R - 1 \) and, after summing over $n$, one arrives at Regge behavior for the planar component of $\sigma_{tot}$ where

\[ \alpha_R = (2\alpha_R - 1) + g^2 N. \] (2)

At next level of cylinder topology, the contribution to partial cross section increases due to its topological twists, \( \sigma_n \simeq \frac{g^4}{(n-2)!}2^{n-2}(g^2 N \log s)^{n-2}s^J_{eff}^{1/2} \), and, upon summing over $n$, one arrives at a total cross section governed by a Pomeron exchange, \( \sigma_0^{tot}(Y) = g^4 e^{\alpha_P Y} \), where the Pomeron intercept is

\[ \alpha_P = (2\alpha_R - 1) + 2g^2 N. \] (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 \( \alpha' \simeq 1 \) $GeV^{-2}$, 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.

Finally we note that, in a Regge expansion, the relative importance of secondary trajectories to the Pomeron is controlled by the ratio

\[ e^{\alpha_R Y}/e^{\alpha P Y} = e^{-(\alpha_P - \alpha_R) Y}. \] (4)

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 “rapdity gap”, \( y \equiv \log \xi^{-1} \), and the other for the missing mass, \( y_m \equiv \log M^2 \).
2.2 Flavoring of Bare Pomeron

We have proposed sometime ago that “baryon pair” and “heavy flavor” production provides an additional energy scale, \( s_f = e^{\eta_f} \), 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. 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,

\[
\sigma_{p,q} \simeq (g^4/q!) 2^{p+q} (g^2 N \log s)^p (g_f^2 N \log s)^q s^{\eta_f - 1},
\]

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

\[
\alpha_P = \alpha_P^{old} + \epsilon,
\]

where \( \alpha_P^{old} \simeq 1 \) and \( \epsilon \simeq 2g_f^2 N \). 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.

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 \( N\bar{N} \) 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 \((\log(s/m_{eff}^2))^{p+1}\), where \( m_{eff} \) 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

\[
\sigma_{tot}^{q=1} \sim \sigma_0^{total}(Y - \delta)(2g_f^2 N) \log(Y - \delta) \theta(Y - \delta),
\]

where \( \alpha_P^{old} \simeq 1 \). 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)), \( \sigma^\text{total}_0 \). 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. (3).

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: \( y_f \geq 2\delta_0 + < q >_{\text{min}} \delta_f \).

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 \( y_f \simeq 8 \sim 10 \).

### 2.3 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

- (i) \( \epsilon_{\text{eff}} \simeq \epsilon \equiv \alpha_P^{\text{old}} - 1 \simeq 0 \) for \( y << y_f \).
- (ii) \( \epsilon_{\text{eff}} \simeq \epsilon \simeq 0.1 \) for \( y >> y_f \).

For instance, exchanging such an effective Pomeron leads to a contribution to the elastic cross section \( T_{ab}(s, t) \propto s^{1+\epsilon_{\text{eff}}(y)+\alpha't}. \) This representation can now be extended down to the region \( y \sim y_f \). 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” formul a between the high energy and low energy sets. Once a choice for $\epsilon_{\text{eff}}(y)$ has been made, it is easy to verify that a natural choice is simply $\beta_{\text{eff}}(y) = \beta_0 \epsilon^{(c-\epsilon_{\text{eff}}(y))y_f}$. It follows that the total contribution from a “flavored” Pomeron to a Pomeron amplitude is $T_{a,b}(y, t) = R(y) T_{a,b}^{\text{cl}}(y, t)$, where $T_{a,b}^{\text{cl}}(y, t) \equiv \beta_a \beta_b e^{(1+\epsilon+\alpha_{\text{eff}}^t) y}$ is the amplitude according to a “high energy” description with a fixed Pomeron intercept, and $R(y) \equiv e^{-[\epsilon-\epsilon_{\text{eff}}(y)](y-y_f)}$ is a “flavoring” factor. In terms of $s = e^y$, $R(s) \equiv (s_f/s)^{[c-\epsilon_{\text{eff}}(\log s)]}$.

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 instance, for the single-particle gap cross section, since there are three Pomeron propagators, one has for the renormalization factor: $Z = R^2(y)R(y_{\text{on}})$. It is instructive to plot in Figure 1 this combination as a function of either $\xi$ or $M^2$ for various fixed values of $Y$.

3 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 $\alpha^{\text{old}}(0) = 1+\epsilon_o$ to $\alpha^{\text{new}}(0) = 1+\epsilon$ 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 $\bar{\epsilon} \equiv \frac{1}{2}(\epsilon + \epsilon_o)$ and $\Delta \equiv \frac{1}{2}(\epsilon - \epsilon_o)$. A convenient parametrization for $\epsilon_{\text{eff}}$ we shall adopt is $\epsilon_{\text{eff}}(y) = [\bar{\epsilon} + \Delta \tanh \lambda_f(y - y_f)]$. The combination $[\epsilon - \epsilon_{\text{eff}}(y)]$ can be written as $(2\bar{\epsilon})[1 + (s/s_f)^{2\lambda_f}]^{-1}$ where $s_f = e^{y_f}$. We arrive at a simple parametrization for our flavoring function

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

With $\alpha^{\text{old}} \simeq 1$, we have $\epsilon_o \simeq 0$, $\bar{\epsilon} \simeq \Delta \simeq \epsilon/2$, and we expect that $\lambda_f \simeq 1 \sim 2$ 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, $s_f = e^{y_f}$. We have motivated by way of a simple model to show that a reasonable range for this scale is $y_f \simeq 8 \sim 10$. 

6
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 $\mathcal{P}\mathcal{P}\mathcal{R}$-term for the low-$y_m$ region, the $\pi\pi\mathcal{P}$-term and/or the $\mathcal{R}\mathcal{R}\mathcal{P}$-term for the low-$y$ region, etc., particularly for $\sqrt{s} \leq \sqrt{s_f} \sim 100 \text{ GeV}$. 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. [8]

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 $g_{\text{PPP}}(0) \simeq e^{-\frac{3}{2} \lambda_f y_f} g_{\text{PPP}}(0)$. With $\lambda_f = 0.1 \sim 0.1$ and $y_f = 8 \sim 10$, using the value $g_{\text{PPP}}(0) = 0.364 \pm 0.025$ $\text{mb}^\frac{1}{2}$, we expect a value of $0.12 \sim 0.18$ $\text{mb}^\frac{1}{2}$ for $g_{\text{PPP}}(0)$. Denoting the overall multiplicative constant for our renormalized triple-Pomeron formula by $K \equiv \beta_2^0(0) g_{\text{PPP}}(0) \beta_b(0) / 16\pi$. With $\beta_2^p \simeq 16$ $\text{mb}$, we therefore expect $K$ to lie between the range $.15 \sim .25$ $\text{mb}^2$.

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 $\sqrt{s} = 1800$ GeV value. With $\epsilon = 0.1$, $\lambda_f = 1$, this is done for a series of values for $y_f = 7, 8, 9, 10$. We obtain respective values for $K = .24, .21, .18, .15$, 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 $\mathcal{P}\mathcal{P}\mathcal{R}$-term as well as $\mathcal{R}\mathcal{R}\mathcal{P}$-terms which can become important. Instead of introducing a more involved phenomenological analysis, we simulate the desired low energy effect by having $\epsilon_o \simeq -0.06 \sim -0.08$. A remarkably good fit results with $\epsilon = 0.08 \sim 0.09$ and $y_f = 9$. [8] This is shown in Figure 2. 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

$$g_{\text{PPP}}(0) \simeq .12 \sim .18 \text{ mb}^\frac{1}{2},$$

exactly as expected. Interestingly, the triple-Pomeron coupling quoted in Ref. 3 ($g(0) = 0.69 \text{mb}^\frac{1}{2}$) is actually a factor of 2 larger than the corresponding low energy
value. Note that this difference of a factor of 5 correlates almost precisely with the flux renormalization factor \( N(s) \approx 5 \) at Tevatron energies.

## 4 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

\[
\frac{d\sigma}{dt d\xi} = [S_i(s,t)][D_a,p(\xi,t)][\Sigma_pb(M^2)],
\]

(10)

\[
D_{a,p}(\xi,t) = S_f(\xi,t)F_{P/a}(\xi,t).
\]

(11)

- 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 \( S_i \approx 1 \) 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 setting why a final-state screening can set in relatively early when compared with that for elastic and total cross sections.

- 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 \( \Sigma_{Pb}(M^2) = R(M^2)\Sigma_{Pb}^{cl}(M^2) \) 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 \( \alpha(0) \approx 1 \) 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 interpolate. 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 $q\bar{q}$ 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 QCD, the non-Abelian gauge nature nevertheless managed to re-emerge through its topological structure. 

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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[1] For a more detailed discussion, see: Chung-I Tan, “Diffractive Production at Collider Energies: Soft Diffraction and Dino’s Paradox”, hep-ph/9706276.

[2] We use $s_0 = 1 \text{ GeV}^2$ as the basic energy scale through out this paper. These “classical” expressions are: $F_{P/a}^{cl}(\xi, t) = (1/16\pi)\beta_a(t)(\xi^{-1})^{2\alpha_P(t)-1}$, and $\sigma_{Pb}^{cl}(M^2, t) = g(t)(M^2)^{\alpha_P(0)-1}\beta_b(0)$. The triple-Pomeron coupling $g(t)$ can be in principle determined by data below $\sqrt{s} \leq 30 \text{ GeV}$ where cross sections are relatively insensitive to the choice of the Pomeron intercept value used.

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production as well as other effects. The effective degrees of freedom involved are “diquarks” and charm quarks respectively. For color counting, a baryon is considered as a bound state of a quark and diquark.

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[8] By choosing $\epsilon^{old} < 0$, it is possible to provide a global “average” description mimicking “secondary trajectory” contributions for various low energy regions. In Ref. 5, acceptable estimates are $\epsilon_0 \simeq -0.11 \sim -0.5$. 
Figure 1: Renormalization factor, $Z(\xi; s) \equiv R^2(\xi^{-1}) R(M^2)$, 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 2.
Figure 2: 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 $\epsilon = 0.08$, $\epsilon_o = -0.07$, $\lambda_f = 1$, $y_f = 9$, with small amount of final-state screening. The dashed-dotted curve corresponds to no screening.