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THE DONNACHIE-LANDSHOFF POMERON VS. QCD†

by

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Abstract

The Donnachie-Landshoff model of diffractive scattering, which accounts for all confirmed diffractive data, is briefly reviewed. The picture emerging from this model is shown to directly contradict that coming from perturbative leading-log s calculations. The inclusion of non-perturbative effects brings the two pictures in closer agreement.

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The Donnachie-Landshoff model of diffractive scattering[1−5] has two basic features. Firstly, it is based on a simple physical picture, and thus contains only a few parameters (5 to be precise). Secondly, it works remarkably well, accounting for all diffractive data from 5 to 1800 GeV. In particular, it is the only model of diffractive scattering that *predicted* all of the following: the absence of a dip in $\bar{p}p$ elastic collisions at the ISR, the UA4(2) and E710 values of the ρ parameter, the total and the differential elastic cross sections as observed by E710, and most recently the total γ cross section observed at HERA. These features are very different from those of leading-log s QCD, and one of the challenges of QCD is to understand how such a simple structure might be the continuation of perturbative calculations to the diffractive region.

In the high-s limit, the simplest behaviour for the hadronic amplitude that one expects from Regge theory is a sum of simple poles, each pole corresponding to the exchange of the particles lying on a given Regge trajectory:

$$\mathcal{A}(s,t) = \sum_i \beta_i^2(t)\ s^{\alpha_i(t)} \xi(\alpha_i(t)) \tag{1}$$

with ξ the Regge signature factor. The first (degenerate) trajectories to be included in Eq. (1) are obviously those of the ρ and f mesons. They are the leading meson trajectories, clearly present when plotting the meson states in a J vs M^2 plane. When continued to negative values of M^2, these trajectories are responsible for the fall-off of the total cross section at small $\sqrt{s}$. At higher energies, the cross sections grow, and the most natural assumption is that another term in Eq. (1) is responsible for that rise although there is no observed trajectory to guide us. In fact, there is reasonable evidence that the exchange responsible for the rise with energy is close to a simple Regge pole, because of the factorizability of the pomeron vertex[6]: the ratio of the differential cross section for $pA \to XA$ to that for $pA \to pA$ does not depend on A. It is hard to imagine how a sum of poles (or a cut) with different couplings to p and A could satisfy this property. Furthermore, Eq. (1) leads, through the use of the optical theorem, to total cross sections that rise like a sum of two powers of s. The pp and $\bar{p}p$ total cross sections are indeed well fitted, from $\sqrt{s} = 5$ GeV to 1800 GeV, to the form[1]: $\sigma_{tot} = Xs^{0.0808} + Ys^{-0.4525}$, with X the same for pp and $\bar{p}p$, and Y changing because of the charge parity of the ρ trajectory. The two powers of s seem to be universal, and are again encountered in all the pX elastic/total cross section, for any target. Hence, it seems that the pomeron is a universal object, described by a simple Regge pole with intercept 1.08 and charge parity $C = +1$. It is this result that led to the predictions $\sigma_{tot}^{\bar{p}p}(\sqrt{s} = 1.8 \text{ TeV}) \approx 73$ mb, $\sigma_{tot}^{\gamma p}(\sqrt{s} = 200 \text{ GeV}) \approx 160\mu$b and $\rho(\sqrt{s} = 541 \text{ GeV}) \approx \rho(\sqrt{s} = 1.8 \text{ TeV}) \approx 0.13$. Note that the power of s describing the rise of the total cross section is insensitive to thresholds (charm, bottom, minijets...). This fact casts serious doubts on models assigning the rise of the total cross section to the opening of new inelastic channels.

If we now want to make a model for the function $\beta_P(t)$ of Eq. (1), we need to address the question of the pomeron coupling. Hadronic cross sections are observed to obey the quark counting rule, *i.e.* the cross sections are proportional to the number of valence quarks contained in the hadron. This is tested not only in πp vs pp cross sections, but also in cross sections involving strange quarks[1], *e.g.* $2\sigma(\Omega^- p) - \sigma(\Sigma^- p)$ can be predicted from $\sigma(pp)$, $\sigma(\pi p)$ and $\sigma(Kp)$. This property indicates that the pomeron is insensitive to the details of the quark wavefunction, as would be the case for a coupling to single quarks. The fact that the pomeron intercept is close to 1 then suggests that the pomeron-quark vertex has the structure γ_μ. This analogy with a photon coupling leads to the conclusion that $\beta_P(t)$ is proportional to the elastic Dirac form factor $F_1(t)$, which is measured in γp collisions: $\beta_P(t) = \beta_0 F_1(t)$. The fit to the total cross section fixes the value of β_0 to be $\beta_0 \approx (0.5 \text{ GeV})^{-1}$. Finally, the trajectory can be measured via low-t results for the differential elastic cross section. ISR results from experiment R211[7] clearly show that the trajectory is linear, $\alpha(t) = \alpha_0 + \alpha' t$, with $\alpha' =(2 \text{ GeV})^{-2}$. Note that the analogy pomeron-photon can be pushed further[2] in order to describe single-diffraction dissociation $pp \to pX$. The idea is that the unbroken proton couples to the pomeron as in elastic scattering, whereas the breaking of the other proton is described by the structure function observed in deep-inelastic scattering, and continued to the small-Q^2 region. This predicts that the differential cross section, $d\sigma/dtdM^2$, at small $x = -t/s$, will behave like $(s/M^2)^{2\alpha(t)-1}(1 - M^2/s)$, where the second factor comes from phase-space suppression. This prediction is realized, and, once refined[2] by the inclusion of the f trajectory, the model accounts very well for the data at all x.

In order to describe the data at large t, or at very large s, one must worry about n-pomeron exchanges. First of all, the fact that the data are well fitted to a single pomeron pole indicates that multiple exchanges are not going to play a major role in the total cross section, at least at present energies. This is fortunate, as no one knows how to evaluate precisely their contribution. Only their qualitative properties are known, *e.g.* the fact that they will eventually dampen the rise of the cross section and make it lower than the Froissart-Martin bound. Instead of using an ansatz for all n-pomeron exchanges, Donnachie and Landshoff have only considered two-pomeron exchange[3]. This contribution will turn out not to be too big, so that one may hope that further exchanges can be neglected. In the context of the Donnachie-Landshoff model, it is made of two terms: a planar one, for which the pomerons couple to the same quark, and a non-planar one for which they couple to two different quarks within the same hadron. The first process can be evaluated, for t/s sufficiently small, as two independent scatterings in impact-parameter space. The second one is assumed to be proportional to the first one, so that the sum is a constant λ times the second-order eikonal contribution. In order to determine λ, one demands that the imaginary parts of one- and two-pomeron exchanges cancel at the energy ($\sqrt{s} = 31$ GeV) and at the value of t where the dip is lowest in pp scattering. This leads to $\lambda = 0.43$, and to a contribution of about 10% to the total cross section at Sp$\bar{\text{p}}$S/Tevatron energies. Note that two-pomeron exchange has a slope $\alpha'/2$ and an intercept $2\alpha_0 - 1$. The simple power of s that comes from a fit to the total cross section must thus be considered as an *effective* power: the true one-pomeron intercept must be of the order of 0.085, and is lowered by the two-pomeron contribution, which accounts for about -7 mb at the Tevatron. Similarly, the meson pole will include contributions from meson-pomeron exchange[1].

Two-pomeron exchange is however not sufficient to reproduce the dip, as one still needs to cancel the real part of the amplitude. This is accomplished by making use of the Landshoff mechanism[4], which arises from 3-gluon exchange, each gluon being attached to a different quark. This process is dominant at high-t: when the three quarks are scattered by the same amount, there is no price to pay for their recombination into a proton, whereas *e.g.* one-pomeron exchange would be suppressed by $F_1(t)$. This $C = -1$, purely real contribution is calculated to be flat in s and to behave like t^{-8} for large $-t$. These features are confirmed by data from $\sqrt{s} = 27$ GeV to 53 GeV, and from $|t| \approx 3$ GeV^2 to ≈ 14 GeV^2. The magnitude of this term is fixed by demanding a cancellation with the left-over real part of one and two-pomeron exchange at the dip at $\sqrt{s} = 31$ GeV. Its contribution gets damped at small t by a regularization of the gluon propagator. As the cancellation in the dip region is largely due to the 3-gluon exchange term, and as that term is odd, the model predicted that there would not be a dip in $\bar{p}p$ scattering. That has now been verified at the ISR[8]. Note that this is the *only* evidence for an odderon contribution in hadronic scattering.

Finally, let us mention that this model also makes definite predictions concerning the photon and proton structure functions[5]. Indeed, in $\gamma^* p$ scattering, the same hadronic amplitude gets probed, the only change being that it gets attached to an upper photon line. The energy flowing through it is given by $s = -(k^2 + k_\perp^2)/x - k^2 + (1-x)m_p^2$, where k^2 is the off-shellness of the quark attached to the photon, $k_\perp$ its transverse momentum, and x the fraction of the proton momentum that it carries. Hence, at very small x, the effective s is very large, the process is dominated by pomeron exchange, and the structure functions are predicted to behave like $1/x^{1.08}$. One might point out that models that neglect the off-shellness of the quark and/or its transverse momentum, are bound to miss this behaviour. A recent parametrization of all the pre-HERA data consistent with the above ideas is now available[5].

We have thus seen that the Donnachie-Landshoff model of diffractive scattering reproduces all data from $\sqrt{s} = 5$ GeV to present energies, for the scattering of protons onto $p, \bar{p}, \pi, K, \gamma, d$. It also provides a successful picture of diffractive dissociation, and, as explained elsewhere[9], of the pomeron structure function. Let us nevertheless mention that two recent measurements are in conflict with the model[10]: H1 observes a much steeper rise of the structure functions at small-x, and CDF measures a cross section which is larger than the E710 number predicted by the model. If confirmed, these data would indicate that something new is happening, maybe the long-awaited perturbative contribution. One would then have to decide how to subtract the non-perturbative contributions from the data, in order to make meaningful comparisons with theory. This question leads us naturally to the second part of this paper, which deals with the possibility of obtaining a description of non-perturbative scattering from QCD.

It is now well-known that the leading-log s resummation[11] cannot account for diffractive scattering. The first problem is that the amplitude is not factorizable: this is in fact a consequence of the infrared finiteness of the answer. Quark-quark scattering via gluon exchange diverges for massless gluons. Nevertheless, hadron-hadron scattering is infrared finite, as the colour of the hadron gets averaged for very long wavelength gluons. Hence there is a contribution that comes from the diagrams where gluons are exchanged between different quarks in the hadron. These diagrams feel the hadronic wavefunction, and hence their contribution depends on the target. A second problem is that the rise of the cross section predicted by perturbative QCD is entirely inappropriate to describe data, as its leading contribution to the hadronic amplitude goes like $s^{1+2.65\alpha_s}$. Finally, we have seen that hadronic data contain a scale of the order of $1/\sqrt{\alpha'} \approx 2$ GeV, which comes in the description of the t-dependence of the hadronic amplitude. No such scale is present in perturbative QCD, hence the differential elastic cross section has the wrong shape: its logarithmic slope at $t = 0$ is infinite and its curvature is too big.

One is thus led to the semantic distinction of a "soft pomeron", which describes the data at low momentum transfers, and a "hard pomeron", which is supposed to arise once α_s is small enough and s is big enough. No one knows how these objects would combine, and quite a few alternatives have been proposed in the literature[12]. We want here to address the question of the matching of a non-perturbative pomeron with a perturbative one at the level of quarks and gluons, *i.e.* we want to consider what ingredients would transform the perturbative calculation into a viable model of non-perturbative exchanges[13].

The first question concerns the factorizability of the amplitude. The n-gluon contribution to the hadronic amplitude $\mathcal{A}_h$ takes the form:

$$\mathcal{A}_h(s,t) = \sum_n \int \alpha_S^n \left[\prod_{i=1}^n dk_\perp^i\right] \mathcal{A}_{qq}(k_\perp^i) \ \left[\mathcal{E}_1^2(t) - \mathcal{E}_d(k_\perp^i)\right] \qquad (2)$$

where $\mathcal{A}_{qq}$ is the quark-quark scattering amplitude, $\mathcal{E}_1^2(t)$ is the form factor associated with the exchange when the gluons are attached to the same quark line $-$ one can show that $\mathcal{E}_1$ is equal to the Dirac form factor F_1 $-$, and $\mathcal{E}_d$ comes from the other diagrams, where the gluons are attached to different quarks. It is the latter that guarantees the infrared finiteness of the answer, and that produces violations of factorizability. Hence this term must be suppressed by non-perturbative effects. It contains explicitly the hadron wavefunction, and thus a scale R, the hadronic radius. Landshoff and Nachtmann[14] have observed that an infrared suppression of the gluon propagator would indeed generate a suppression. If the gluon propagator is smoother than a pole in the infrared region, then on dimensional grounds, it must contain a scale μ_0: $D(q^2) = (1/\mu_0^2)\ \mathcal{D}(q^2/\mu_0^2)$. If the scale μ_0 is big enough, then the propagator does not change much while the form factor $\mathcal{E}_d$ drops sharply. Hence, for $\mu_0 >> 1/R$, one gets a suppression proportional to $1/(\mu_0 R)^2$.

This idea that the gluon propagator should be smoother than the perturbative one has received some theoretical confirmation from lattice gauge theory[15], from studies of the Dyson-Schwinger equations[16] and from considerations related to the Gribov Horizon[17]. Changing the gluon propagator amounts to the inclusion of a subclass of diagrams (the gluon self-energy) which are resummed via non-perturbative methods. All these propagators give rise to a factorizable amplitude for μ_0 big enough.

The next question is whether the scale of the amplitude is correct. It is already known that purely perturbative 2-gluon exchange[18] gives reasonable numbers for the total cross section, but leads to the wrong shape for the differential elastic cross section. Although the various non-perturbative propagators are far from agreeing, for a scale μ_0 of the order of 0.5 to 1 GeV, each can give a good starting value for the total cross section, of the order of 20 mb. The logarithmic slope of the elastic cross section also gets cured by the introduction of μ_0, and numbers of the order of 10 GeV^{-2} can be obtained. At this order, as has now been observed by many authors[19], the inclusion of modified propagators in the calculation provides appreciable improvements and the improved order α_s^2 constitutes a good starting point for an expansion in $\log s$. Let us now see to which extent these improvements carry over to higher orders[13].

The complete order α_s^3 leads to a hadronic amplitude that has the following form:

$$\mathcal{A}(s,t) = \mathcal{A}_2 \left\{ i \left[1 + \log s \left(\epsilon_0 + \alpha' t + O(t^2) \right) \right] + f_{odd}(t) \right\} \tag{3}$$

It is of course tempting to see in this a first-order Taylor expansion in $\log s$ of a pomeron pole, plus a zeroth-order term from an odderon pole. As BFKL have shown[11], life is not so simple, and higher-order terms spoil the analogy. Hence, in the following, the terms "pomeron intercept" or "slope" must not be taken literally. At this order, the normalization of the cross section, $\mathcal{A}_2$, comes from two-gluon exchange; in the Feynman gauge, the coefficient of $\log s$, ϵ_0, is the ratio of two to three gluon exchange. The odd contribution $f_{odd}(t)$ comes entirely from 3-gluon exchange. Finally, the α' contribution comes from the diagrams involving the 3- and 4-gluon vertices and from the Taylor expansion of 3-gluon exchange. Note that one can show that, at least for 3-gluon exchange diagrams, the replacement of the perturbative gluon propagator by a non-perturbative counterpart does not violate gauge invariance and can be justified theoretically[13], provided that the propagators have a Hilbert transform.

Let us first consider the odd term f_{odd}. It is purely real, and proportional to s. As expected, it contains a contribution from the Landshoff mechanism. The structure of its form factor clearly shows that the Landshoff term is dominant at high-t whereas it contributes about 18% of the two-gluon exchange amplitude at $t = 0$. It is present only in baryon-baryon scattering, and not in meson-baryon scattering.

The two main problems of the perturbative expansion are already manifest at this order when one uses perturbative propagators: the pomeron intercept is $1 + 1.85\alpha_s$, and the slope turns out to be logarithmically divergent, although $\alpha' t \to 0$ when $t \to 0$. Non-perturbative gluon propagators bring the calculation in closer agreement with the data, but their effect is not sufficient: any substantial reduction of the order α_s^3 result also brings the normalization of the cross section down. For values of μ_0 favoured by our previous discussion on two-gluon exchange, we get values of the order of 2 for the intercept, and it seems impossible to get acceptable numbers both for two and three-gluon exchange. Similarly, the pomeron slope becomes finite once the infrared region is smoothed out. Values compatible with 0.25 GeV can again be achieved, but again for values of μ_0 or α_s that would suppress two-gluon exchange. Hence, one cannot accommodate a sizeable two-gluon exchange amplitude together with a slowly rising third-order amplitude by simply modifying the gluon propagator according to the results of Ref. 16.

Nevertheless, the inclusion of non-perturbative propagators inside the perturbative calculation has a rather large effect: qualitatively, it replaces the infinities of the perturbative answer ($B(0)$, α') by finite numbers, and gives rise to a factorizing amplitude. Quantitatively, it reduces ϵ_0 by a factor of the order of 2, and gives good starting values for the order α_S^2.

It is worth pointing out that the factorizability of the pomeron vertex, the smallness of the intercept and the pomeron slope point to rather large values of the scale μ_0. The problem would then be to find a mechanism to increase the two-gluon exchange term. It is conceivable that contributions from the color field would increase all orders by a common factor (as they would act at the form-factor level), and one would then need to use a larger value of μ_0, hence suppressing the higher orders more and more (as the number of propagators grows with the order).

It is also possible that the standard techniques perturbation theory (such as cutting rules) become modified in the non-perturbative regime. This could totally modify the estimate of higher orders. Indeed, one finds that the diagrams giving rise to ϵ_0 are made of two contributions: a cut through the quarks, and another cut through the quarks and one gluon. Were the latter to be diminished by a factor 2, ϵ_0 would become identically zero.

These problems are presently under further investigation.

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