Spin Dependence of Deep Inelastic Diffraction: an Outline

J. Bartels[a], T. Gehrmann[b], M.G. Ryskin[c]

[a] *II. Institut für Theoretische Physik, Universität Hamburg, D-22761 Hamburg, Germany*

[b] *DESY Theory Group, DESY, D-22603 Hamburg, Germany*

[c] *Petersburg Nuclear Physics Institute, 188350, Gatchina, St.Petersburg, Russia*

Abstract

We argue that the relative magnitude of perturbative to non-perturbative contributions to the polarized deep inelastic diffractive cross section is more favourable than in the unpolarized case, since a large part of the polarized cross section could be calculated in perturbative QCD. We briefly outline this argument and describe which calculations need to be done for a determination of the perturbative contribution.

1 Introduction

The occurrence of large rapidity gaps between current jet direction and proton remnant direction in electron-proton collisions at HERA [1] represents one of the most puzzling phenomena in the physics of deep inelastic scattering. The first observation of these diffractive deep inelastic scattering events four years ago [1] has triggered much theoretical effort towards an understanding of this phenomenon. Although much progress both in the theoretical description and in the experimental study [2] of diffraction in deep inelastic scattering has been made, it is still fair to say that this phenomenon is not unambiguously understood at present, since it contains both perturbative and nonperturbative components.

To summarize our present understanding of (unpolarized) diffraction in deep inelastic scattering, the diffractive final states can be attributed either to the nonperturbative soft Pomeron or to the 'hard Pomeron', represented by the (slightly nonforward) gluon structure function. Examples of the former class are the aligned jet configuration of diffractively produced quark-antiquark pairs or the analogous quark-antiquark gluon configuration where the gluon is rather soft (in particular, has a small transverse momentum). For the latter class of final states, we mention the diffractive production of longitudinal vector particles or quark-antiquark jets with large transverse momenta. One of the most prominent signatures of these distinct classes is the energy dependence: for the nonperturbative part one expects to see the energy dependence of the soft Pomeron, whereas the perturbative part should be characterized by the stronger increase with energy as observed in the gluon structure function at small x. The diffractive contribution to deep inelastic structure functions can not be easily attributed to one of both classes. Experimental data on it lie indeed between the two extremes, thus indicating the existence of sizable contributions from both the hard and the soft Pomeron.

Attempts to find a theoretical description of this complicated situation can naturally be grouped into several classes, each emphasizing either the nonperturbative or the hard part of the cross section. The soft pomeron approach [3] interprets this process as being due to the emission of a soft pomeron off the proton. The incoming electron then scatters deep inelastically off the pomeron, thus probing its partonic content. The diffractive structure function in this approach is expected to factorize into a pomeron emission factor, which can be predicted from Regge theory and the structure function of the pomeron, which has to be determined from experiment. Correspondingly, the value of the intercept $\alpha(0)$ which describes the energy dependence of the diffractive cross section at low values of x_P is expected to be close to 1.08, as determined from soft hadron-hadron interaction phenomena. The semi-classical approach to diffraction [4] assumes that the underlying dynamics are governed by the scattering of partons in the incoming photon off the (classical) colour field of the proton. A diffractive final state is obtained if the partons leave the proton field in a color singlet configuration. This approach leads to diffractive distributions [5] which can be expressed as averages over particular configurations of the soft colour field inside the proton, which are completely non-perturbative. The energy dependence corresponds to a non-perturbative Pomeron intercept of $\alpha_P(0) = 1$. Finally, the perturbative approach [6] describes diffraction as the exchange of a system of two perturbative gluons (gluon structure function), leading to a diffractive cross section which rises like $[x_P g(x_P, \bar{Q}^2)]^2$. This approach will be outlined in more detail below, since it provides the basis of our proposal.

The fact that none of these approaches covers the full set of final states, i.e. both nonperturbative and perturbative dynamics, makes it impossible to decide how much of the cross section belongs to either side. In particular, so far there is no reliable theoretical estimate of how much of the diffractive cross section is calculable perturbatively. Qualitatively, the connection between the perturbative and the nonperturbative part can be understood most easily by considering the perturbative part of the cross section (in diffractive $q\bar{q}$ production: final states with large transverse momenta) and then trying to integrate over the full k_t region: the integral over the transverse momentum has the form $\int \frac{dk_t^2}{k_t^4}[x_P g(x_P, \bar{Q}^2)]^2$, and thus the main contribution comes from the region of small k_t where perturbation theory is inapplicable. We will return to this point further below.

2 The polarized diffractive cross section

Turning now to polarized deep inelastic scattering, the case of diffractive final states seems to be much better since the 'soft' (nonperturbative) part of the polarized cross section is expected to be rather small. To be definite, let us consider the longitudinal asymmetry A_{LL}. While the leading contribution to unpolarized soft diffraction arises simply from the soft pomeron pole exchange, in the polarized case one has to consider far more involved exchanges in order to obtain a non-vanishing contribution in soft diffraction, and they are expected to be small.

To see this in more detail, let us consider elastic hadron-hadron scattering $a + b \to a + b$, described by Reggeon (Pomeron) exchanges . In the scattering amplitude, the reggeon (Pomeron)-hadron (say, proton (V_{pRp}) or reggeon-quark (V_{qRq})) coupling has the general form [7]

$$V * (1 + \frac{a}{2m_N}\sqrt{k_t^2}(\vec{\sigma}\vec{n})) \tag{1}$$

where m_N is the nucleon mass (introduced in order to fix the normalization), k_t is the transverse momentum transfered through the reggeon, and the unit 3-vector $\vec{n}$ is perpendicular to the

scattering plane ($n_i \sim \epsilon_{ijl} k_{jt} q_l$ with $\vec{q}_l$ being the initial hadron momentum along the z-direction). The asymmetry we are looking for – the helicity-helicity asymmetry A_{LL} – must be proportional to $\sigma_z^a \cdot \sigma_z^b$ where the spin matrices σ^a and σ^b act on the hadron a and b respectively. As it is seen from eq.(1) there is no such term in a single reggeon (or Pomeron) exchange. The only possibility to obtain a nonvanishing contribution is to consider multi-reggeon cuts. Then the matrix σ_z^a may be generated as an antisymmetric product $[\sigma_x^a \cdot \sigma_y^a - \sigma_y^a \cdot \sigma_x^a]$. However, dealing with only one type of reggeons (say, multipomeron cuts) one never finds a σ_z this way, since the multireggeon vertex should be symmetric under permutations of the reggeons as a result of the Bose-Einstein statistics. Therefore, we have to include at least one secondary reggeon into our amplitude. Note that it is not enough to take the simplest reggeon-pomeron or other 2 reggeon cut at $t = 0$. In this case both spin-matrices σ will be directed along the same $\vec{n}$ axis and never give us σ_z. Thus one has to start from the 3-reggeon cut.

Let us first consider the RPP cut which is suppressed by the energy factor $1/\sqrt{s}$, compared to the Pomeron pole exchange in the unpolarized case. Because of the very small spin-flip part of the pomeron vertex this cut provides only a very small contribution. The spin-flip factor (a in eq.(1)) extracted from the old Regge phenomenology [7] is about $a_P = 0.2$ or less (consistent with zero). Thus the contribution of a RPP cut to the asymmetry A_{LL} is suppressed by at least the factor $(a_P \cdot a_R \frac{k_t^2}{4m_N^2})^2 \simeq (0.05 * 0.05)(a_P \cdot a_R)^2$ (here we have used $k_t^2 \simeq 0.2$ GeV^2) in comparison with the same RPP-cut contribution to the unpolarized cross section. Therefore it may be better to consider the RRP-cut: if $a_P = 0$, this will in fact be the only possibility. Here it might be possible to build the antisymmetric $\rho^+\rho^-$ configuration with a rather large $a_\rho \sim 6$, which almost compensates the suppression due to a small $k_t < 2m_N$. Now the weight factor (relative to the RRP-cut in the unpolarized case) is much larger: $(a_R \cdot a_R \frac{k_t^2}{4m_N^2})^2 \simeq (0.05 * 0.05)(a_R \cdot a_R)^2 \sim 3$. Nevertheless, the relative strength in comparison with the unpolarized simple Pomeron pole exchange remains small. There is the small numerical factor which, in general, suppresses the 3-reggeon cut contribution in comparison with the single Pomeron exchange [8]. In addition, there is a suppression in the spin-dependent part (of the RRP cut) due to the angular integration, and finally there may be cancelations between different configurations (e.g. between the RRP and the RPR contribution). As a result, one expects that the nonpertubative part of the asymmetry couples weaker than the Pomeron pole in the unpolarized cross section.

Finally, the contribution of the RRP-cut, compared to the Pomeron pole exchange in the unpolarized case, is suppressed by the power of energy ($1/s$). This is the same suppression factor x_P that one finds also in the pQCD Born approximation to the spin dependent amplitude in comparison with the Born approximation to the unpolarized amplitude, as we shall discuss below. In this approximation, therefore, the energy dependence of the perturbative and nonperturbative parts seem to be comparable. However, below we shall argue that the pQCD contribution is *enhanced by large logarithms*: in each order of perturbation theory one finds terms of the form $\alpha_s^n \ln^{2n}(1/x)$ in the polarized amplitude. It is the appearance of these logarithms which, together with the weak coupling of the 3-reggeon cut, gives rise to the hope that the perturbative part of the asymmetry will be considerably larger than the nonperturbative part.

In order to compare this nonperturbative part of the longitudinal asymmetry with the perturbative contribution in more detail, let us briefly go back to the unpolarized case [6]. For quark-antiquark production one has to start with the kinematic configuration where the use of perturbation theory can be justified, namely the region of large relative transverse quark momenta k_t. In the double-leading logarithmic approximation the amplitude is described by coupling the two gluons of the gluon structure function to the $q\bar{q}$ system. An intuitive description of inclusive and diffractive amplitudes at small x in this approach can be given in the wave

function formalism elaborated by Wüsthoff [9]. Viewed in the target rest frame, the scattering process can be described as follows. The incoming virtual photon (virtuality Q^2) splits into a partonic system of transverse size $1/Q$ well ahead of the target, the parton content of this system can be described by probabilistic wave functions. The simplest partonic configurations in deep inelastic diffraction are a $q\bar{q}$ dipole and a $q\bar{q}g$ dipole, where the latter only contributes at leading twist if quark and antiquark are close in impact parameter space, i.e. form a single source of color charge. Explicit wave functions for these dipoles, including already the dependence on the incoming photon helicity for the $q\bar{q}$ dipole, are given in [9]. The diffractive amplitude is modeled by the exchange of two gluons in a color singlet state. The coupling of the two gluons to the proton is described by the unintegrated gluon structure function $\mathcal{F}_g(x, Q^2, l_t^2)$, where l_t^2 is the transverse momentum circulating in the gluon loop, which is integrated over. In the double leading log limit [10], the diffractive cross section for $q\bar{q}$ production can be expressed through the gluon structure function, and it is proportional to $[\frac{1}{k_t^2} x_P g(x_P, k_t^2/(1-\beta))]^2$.

The competition of this perturbative final state with the nonperturbative region is easily seen (as it was mentioned already at the end of the introduction) when one integrates over k_t: because of the k_t^2-factor in the denominator, the integral receives its main contribution from the small-k_t region (more precisely, the aligned jet configuration) where the soft Pomeron comes into play, and the perturbative calculation becomes invalid. Therefore, the (integrated) diffractive $q\bar{q}$ cross section contains two competing pieces: the aligned-jet region with the soft Pomeron and the large-k_t region with the strongly rising gluon structure function. Based upon the experience with hadron scattering, one might have expected that the aligned jet configuration would dominate the diffractive cross section. However, because of the rise of the gluon structure function at small x_P, the dominant k_t-value is not so small, and the hard region gives a larger contribution. To put it in a more formal way, at small x_P the anomalous dimension γ ($g(x_P, k_t^2/(1-\beta)) \propto (k_t^2)^\gamma$) becomes large, and as soon as γ exceeds the value 1/2 the maximum in k_t starts to move towards larger values. For the HERA domain it lies around $k_t \sim k_0 \sim 1$ GeV. This value of k_0 increases with energy, but at present it is not yet large enough to allow the conclusion that, in the perturbative formulae, the main contribution comes completely from small distances. That is why one has to take into account the non-perturbative contribution as well.

Returning now to the polarized case of diffractive quark-antiquark production, the enhancement of the perturbative region relative to the nonperturbative region is expected to be even stronger. The perturbative calculation [11] of the polarized cross section again starts from the large k_t-region, and instead of the (unpolarized) gluon structure function $g(x_P)$ we now expect the polarized quark and gluon structure functions $\Delta q(x_P)$ and $\Delta g(x_P)$ to determine the cross section. The occurrence of $\Delta q(x_P)$ is due to the fact the polarized quark and gluon distributions still mix at small x, which is not the case for the unpolarized distributions.

The small-x behaviour of $\Delta q(x_P)$ and $\Delta g(x_P)$ is closely related to the small-x behaviour of the inclusive structure function $g_1(x)$, which has been studied in [12]. A particular aspect of polarized quark and gluon distributions is the occurrence of terms proportional $\alpha_s^n \ln^{2n}(1/x)$, while the most singular terms in the unpolarized gluon distribution are $\alpha_s^n \ln^n(1/x)$. These leading logarithmic terms in the distribution functions can be resummed to all orders both in the unpolarized [13] as well as in the polarized [12] case. For not too small x, the results of [12] are in agreement with the small-x behaviour obtained from the polarized Altarelli-Parisi splitting functions, but at very small x a much stronger power-like rise is predicted. In particular, this rise is substantially stronger than that of the Regge-model which was discussed above. The difference between 'soft physics' and 'hard physics' seems to be larger than in the unpolarized case. *Therefore it is expected that the relative magnitude of perturbative to non-perturbative*

contributions is more favourable in the polarized diffractive cross section than in unpolarized diffraction.

3 Outlook

The calculation of the diffractive content of polarized structure functions at small x in the perturbative framework of a two gluon exchange model appears to be a feasible task. First steps in this direction have already been taken in [14], where the leading $\alpha_s^n \ln^n Q^2$ approximation to this process has been derived. A more elaborate study incorporating the resummation of double logarithmic terms $\alpha_s^n \ln^{2n}(1/x)$ is presently ongoing [11] and has briefly been motivated in this article. Since these double logarithmic terms are absent in the resummed unpolarized amplitudes, which contain only terms of order $\alpha_s^n \ln^n(1/x)$, one would expect that in polarized diffractive DIS at small x the perturbative part is even larger than in the unpolarized case.

The spin dependent diffractive cross section is therefore of particular interest, since one expects non-perturbative contributions to it to be comparably small. A measurement of this quantity could therefore provide an unambiguous test of the perturbative approach to diffractive deep inelastic scattering and of potential enhancement effects due to perturbative resummation. The size of these perturbative contributions will be derived in future work.

Acknowledgements

The work of MGR was supported by the INTAS grant 93-0079-EXT and by the Volkswagen-Stiftung.

References

[1] ZEUS Collaboration: M. Derrick et al., Phys. Lett. **B315** (1993) 481;
H1 Collaboration: T. Ahmed et al., Nucl. Phys. **B429** (1994) 477.

[2] H1 Collaboration: T. Ahmed et al., Phys. Lett. **B348** (1995) 681; C. Adloff et al., DESY preprint DESY-97-158 (hep-ex/9708016);
ZEUS Collaboration: M. Derrick et al., Z.Phys. **C68** (1995) 569; J. Breitweg et al., DESY preprint DESY-97-184 (hep-ex/9709021).

[3] G. Ingelman and P. Schlein, Phys. Lett. **B152** (1985) 256.

[4] W. Buchmüller, A. Hebecker and M.F. Mc Dermott, Nucl. Phys. **B476** (1996) 203; **B487** (1997) 283.

[5] A. Berera and D. Soper, Phys. Rev. **D53** (1996) 6162;
A. Hebecker, Nucl. Phys. **B505** (1997) 349.

[6] A.H. Mueller, Nucl. Phys. **B335** (1990) 115;
M.G. Ryskin, Sov. J. Nucl. Phys. **52** (1990) 529;
M.G. Ryskin and M. Besancon, Proceedings of the Workshop "Physics at HERA", eds. W. Buchmüller and G. Ingelman, DESY (Hamburg, 1991), p.215;
N.N.Nikolaev and B.G.Zakharov, Z. Phys. **C49** (1991) 607;
J. Bartels and M. Wüsthoff, Z. Phys. **C66** (1995) 157.

[7] K.G. Boresov, A.M. Lapidus, S.T. Sukhorukov and K.A. Ter-Martirosian, Sov. J. Nucl. Phys. **14** (1972) 457;
G. Kane and A. Seidl, Rev. Mod. Phys. **48** (1976) 305.

[8] A. Donnachie and P.V. Landshoff, Nucl. Phys. **B267** (1986) 690.

[9] M. Wüsthoff, Phys. Rev. **D56** (1997) 4311.

[10] J. Bartels, H. Lotter and M. Wüsthoff, Z. Phys. **C68** (1995) 121; Phys. Lett. **B379** (1996) 239.

[11] J. Bartels, T. Gehrmann and M.G. Ryskin, work in progress.

[12] J. Bartels, B. Ermolaev and M.G. Ryskin, Z. Phys. **C70** (1996) 273; **C72** (1996) 627.

[13] V.S. Fadin, E.A. Kuraev and L.N. Lipatov, Sov. Phys. JETP **44** (1976) 433;
Ya.Ya. Balitski and L.N. Lipatov, Sov. J. Nucl. Phys. **28** (1978) 822.

[14] M.G. Ryskin, hep-ph/9706505.