THEORY OF DIFFRACTIVE MESON PRODUCTION

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In this talk some recent developments in the theory of diffractive meson production are reviewed: (1) the issue of end point divergencies and diffractive factorisation; (2) elastic vector meson production in the context of the gluon saturation in the impact parameter plane; (3) diffractive vector meson photoproduction at large momentum transfer; (4) C-even meson photoproduction and the Odderon.

1. Introduction

Diffractive meson production offers unique opportunities to probe the colour singlet exchange in high energy scattering. In HERA experiments one measures properties of the scattering amplitude in a wide range of energy W, momentum transfer, t, and the incoming photon virtuality Q^2, for various mesons.[1] The multidimensional set of data is, with no doubts, a challenge for the theory. Standard methods of QCD, like the fixed order perturbative calculus and the hard factorisation, are not always sufficient to get a correct picture, and the use of resummation schemes, like the BFKL formalism[2] or phenomenological models as the saturation model,[3] is necessary. Therefore a lot of effort is directed towards improving the phenomenological models and rooting them in the fundamental theory. The main goal is to obtain a robust global description of the diffractive meson production data, in a simple conceptual framework grounded in QCD.

On the other hand, new measurements may be performed that probe some interesting mechanisms (e.g. the Odderon exchange) which are not fully understood yet. Furthermore, some quantities which are not calculable from first principles may be determined or constrained. The hadron structure is a good example here.

In this talk I am presenting short notes on four selected topics in which substantially new results have recently appeared and further progress is expected in the close future.

2. Diffractive factorisation

The theory of exclusive meson electroproduction is based on the diffractive factorisation theorem.[4] The proof holds for a longitudinally polarized incoming photon that carries a large virtuality. In this case the scattering amplitude is related to parton distributions in the target. To be specific, for the diffraction at high energies the dominant contribution comes from the gluon evaluated at the scale of the order of the photon virtuality or of the meson mass. It should be added, however, that the photon – meson transition requires a non-zero energy-momentum transfer from the proton to the photon. Hence, the operators defining the gluon distribution in the proton are evaluated in a non-forward kinematics. and the conventional gluon distribution has to be replaced by its off-diagonal counterpart.[5] The evolution of the off-diagonal partons is known in QCD,[5] and it was applied to construct models of these quantities.

In the case of transversely polarized photons, though, one finds some troublesome pieces in the amplitude which seem to spoil the diffractive factorisation. The hard part of the amplitude at the leading twist gives $1/[z(1-z)]^2$ dependence on the longitudinal momentum fraction z of the quark. The singular behaviour of this part at the end-points, $z \to 0$ and $z \to 1$, is not compensated by the asymptotic form of the meson distribution amplitude $\phi(z) = 6z(1-z)$, and infrared logarithmic divergencies appear in the integrals over z, breaking down the hard factorisation. Speaking in terms of the colour dipole model the photon wave function does not suppress large colour dipoles with an asymmetric sharing of the longitudinal momentum and those dipoles may interact with the target exchanging soft gluons.

An interesting concept was put forward[6] to use the singular behaviour of the hard amplitude to explain some peculiarities of the high-t diffractive meson photoproduction. Namely, the hard physics dictates a $1/|t|^4$ dependence of the amplitude without a helicity flip between the photon and the meson and a $1/|t|^3$ dependence of the helicity flip amplitude. Indeed, the data for ρ and ϕ exhibit approximately $1/|t|^3$ behaviour. Surprisingly, the analysis of angular distribution of the vector meson decay products, points to a strong dominance of the no-flip amplitude for all the range of momentum transfer.[7]

An assumption[6] that the distribution amplitudes do not vanish at $z \to 0$ and $z \to 1$ leads to an appearance of a power divergence at the end-points that replaces the usual logarithmic divergence. The necessary physical infrared cut-off is naturally provided by some soft scale $\sim \Lambda_{QCD}$. Thus, the end-point contributions would have a different scaling with t than it follows from the perturbative analysis. This is how one could explain the observed t-dependence $\sim 1/|t|^3$ of the dominant no-flip amplitude. Another consequence of this scenario is that the helicity need not be conserved for the quark–anti-quark configurations which contain a very slow parton and which are sensitive to non-perturbative effects, e.g. the instantons. Thus, this model could explain the qualitative features of the data for high-t meson production but no quantitative predictions have been provided so far. Although the idea is interesting, the assumptions behind it are controversial. For instance, the invariant mass $M_{q\bar{q}}$ of the end-point configurations at the quark level, given by $M^2_{q\bar{q}} \simeq (m_q^2 + k_T^2)/[z(1-z)]$, is much larger than the meson mass M_V (with m_q and k_T denoting quark mass and transverse momentum respectively). Thus, such the configurations should be strongly suppressed. The reason is that in an exclusive process, the invariant mass at the parton level should be close to the meson mass if the local parton-hadron duality is to hold.

Recently, a thorough analysis of the factorisation at the end-points beyond the leading twist contribution has been carried out.[8] The momentum dependent meson wave function is modelled in the following way. The starting point is the perturbative photon wave function, next its point-like component is removed using the Borel transform. The obtained meson wave function $\Psi_V(z, k_T) \sim \exp\left(-\frac{k_T^2 + m_q^2}{M_V^2\, z(1-z)}\right)$ contains contributions from all twists. The resulting helicity structure is the same as the helicity structure of the photon. Note that this wave function gives the asymptotic form of the distribution amplitudes after the integration over the transverse momentum, and that it naturally suppresses the configurations with large invariant mass at the parton level.

This model permits a detailed analysis of the end-point contributions at all twists. An apparent singular behaviour of the integrals at the end-points is found to be only an effect of inappropriate extrapolations at the leading twist. In the complete model of the wave function the end-point divergences do not appear. Instead, at large photon virtuality Q^2 one finds logarithms of Q^2 which can be fully absorbed into the generalized evolution of t-channel $gg \to q\bar{q}$ exchange. Thus the end point contributions are separated into two factorised terms, containing the gluon distribution and the an effective $q\bar{q}$ distribution. The latter term may be neglected at small x and the gluon dominates.[8]

3. Elastic production of vector mesons

It was shown long time ago[9,10] that elastic production of vector mesons may be used as a probe of the collinear parton densities in the proton, giving a direct access to the off-diagonal gluon. In fact, a detailed measurement of the forward peak provides even more information. The t-dependence of the scattering is related to the profile of the scattering amplitude in the impact parameter plane by the Fourier transform. Furthermore, in collisions at high energies the transverse position of the dipole is conserved in the scattering, thus, in the b-representation the S-matrix is diagonal and much easier to determine. There is one difficulty left, however. Experiments measure cross sections, and the information about relative complex phases of the scattering matrix elements is partly lost. A procedure that was recently proposed for S-matrix determination[11] is based on a simplifying assumption that this matrix is real. In this scheme a closed formula was derived for the S-matrix elements in a dipole scattering off the proton, at a given impact parameter. The necessary experimental input was the differential cross section $d\sigma/dt$ for the elastic electroproduction of longitudinally polarised ρ meson.[11] The differential cross section of this process $d\sigma/dt$ is known for a few values of Q^2, in each case covering about a decade in x and the range of $-t < 0.6$ GeV^2. The results obtained were parameterised in a form following from the hypothesis of b-dependent saturation, and the saturation scale was determined as a function of the impact parameter.

In the colour dipole representation the dependence of dipole scattering amplitude on the impact parameter is influenced by the matter (mostly gluons) distribution in the proton. Thus, in principle, the cross section for vector meson production could be used to provide some clues on the three dimensional structure of the proton.

The b-dependent saturation of the elastic colour dipole – proton cross section is a phenomenon which follows directly from the unitarity constraints. A vector meson production process may probe dipole scattering in the kinematical domain where rescattering effects should be important. It opens a possibility to analyse in detail how the amplitude saturates in the impact parameter plane and for the increasing collision energy. In the classical saturation model proposed by Golec-Biernat and Wüsthoff (GBW)[3] it is assumed that the nucleon-dipole interaction is confined to a disc in the transverse plane. In a more realistic picture the scattering matrix elements should be characterised by a smooth distribution $S(b)$ in the impact parameter. The large-b behaviour is constrained by the fact that QCD is a theory with a mass gap, thus $S(b)$ should decrease faster than $\exp(-mb)$ for $b \to \infty$. It may be argued that the Gaussian tail $S(b) \sim \exp(-Ab^2)$ accounts better for the confining properties of QCD vacuum.[12]

Two recent extensions of the GBW model[13,14] use the Gaussian profile $S(b)$ of a single dipole scattering. The unitarity constraints are imposed by the Glauber-Muller prescription at each impact parameter[13] or by inserting[14] the b-dependence into solutions of Balitsky-Kovchegov equation.[15] Both the models were confronted with the data for elastic J/ψ photo- and electroproduction, giving a good fit. For the J/ψ though, the typical dipole size is related to the inverse c-quark mass, and the unitarity corrections are not large.[13] An important goal to achieve in the close future is to provide a global fit of all the data for exclusive vector meson production in the saturation model with the impact parameter dependence. Such an analysis, based on the present and future HERA data, would provide stringent experimental constraints of the dynamics of b-dependent saturation.

4. Diffractive production at large t

Recently, ZEUS[7] and H1[16] measured the differential cross section $d\sigma/dt$ for the ρ^0, ϕ and J/ψ diffractive production over a wide range of t, and also the corresponding spin density matrix elements. From those data estimates of the helicity amplitudes may be obtained. The amplitude without helicity flip dominates for all mesons and for the accessible t range. The double flip and single flip amplitudes have the magnitudes of one fifth and one tenth of the leading amplitude respectively.

High energy diffractive meson production in non-forward direction tests colour singlet exchange at the large momentum transfer t. Elastic scattering of a dipole with size r is strongly suppressed for $r^2|t| \gg 1$ because the dipole tends to break up. Also a gluon radiation that spoils the rapidity gap is important in this regime. Therefore the momentum transfer provides the hard scale for the problem and the perturbative approach is justified for the photon-meson transition, even for the light mesons and the photoproduction limit.

In the colour dipole representation at high energies, the dipole scattering amplitude is diagonal in the quark helicity basis. In the leading twist approximation the quark momenta are collinear and the vector meson helicity is a simple sum of quark helicities. Beyond this approximation, the angular momentum originating from a motion of quarks inside the meson and a possible appearance of additional partons in higher Fock states, has to be taken into account and the sum of quark helicities need not match the meson helicity. This explains how the double helicity flip occurs in the high-t photoproduction of ρ^0 and ϕ.[17,18,19]

The lowest order approximation to the C-even colour singlet in QCD is given by two gluons combined into a colour neutral state. This simple picture is not sufficient, though, when the rapidity gap $\hat{y}$ between the meson and the proton remnant is large. Perturbative corrections to the two-gluon exchange contain terms that behave like $(\hat{y}\alpha_s)^n$ at the n-th order of expansion. Since $\hat{y}\alpha_s \sim 1$, the fixed order perturbative expansion gives inaccurate results and a resummation of leading logarithms $(\hat{y}\alpha_s)^n$ (with $\hat{y} \sim \log s$) is necessary. Those contributions are diagrammatically represented as gluon ladders. The resummation is performed by the BFKL equation[2] providing a consistent framework to analyse the non-forward diffractive amplitudes.

Vector meson production data at $|t| > 1$ GeV^2 are dominated by events with the dissociated proton. Thus, in most theoretical analyses all the final states of the proton remnant are included, having a given transverse momentum the total colour charge equal to zero. Then, the hard colour singlet couples predominantly to individual partons in the nucleon[20] resolving them at the hard scale $\sim |t|$. This approximation improves with increasing $|t|$, and this allows to factorise the cross sections into parton densities in the proton and a partonic cross section for the elastic meson production.

The scattering amplitude A in the BFKL formalism is given by a convolution over the gluon transverse momenta, of the universal gluon $2 \to 2$ Green function G and of two specific impact factors Φ_i describing the coupling of two gluons to the systems at the two sides of the rapidity gap, $A \sim \Phi_1 \otimes G \otimes \Phi_2$. The analytic approach to the BFKL equation is based on its conformal invariance[21] which permits a decomposition of the Green function in the basis of its conformal Eigenfunctions, indexed by a discrete conformal spin n and a continuous variable, ν. This method was successfully applied to the J/ψ diffractive production. At first, only the leading conformal component with $n = 0$ [22,23] was taken into account, and then the calculation was extended to all the conformal spins.[24]

For the production of light vector mesons, the impact factor representing photon-meson transition depends substantially on the details of both the photon and the meson structure. For the photon, one may use the perturbative wave function with a constituent quark mass. An alternative approach is based on the Operator Product Expansion in which only the leading dependence on the dipole vector $\vec{r}$ is retained at $r \to 0$. The dependence on the longitudinal momentum fraction is parameterised by the distribution amplitudes.[25] The approximation may be systematically improved by including higher twist terms and the QCD evolution. The same method may be applied also for the vector mesons.[26]

Recently, a few attempts have been made,[17,23,24,18,19,27] to describe the high-t photoproduction data. In those papers important observations about the meson production dynamics are performed, but none of the approaches is completely successful. It was argued[17] that non-perturbative condensates in the QCD vacuum generate a sizeable chiral odd coupling of the photon to the quark anti-quark pair. This component significantly enhances the amplitude without helicity flip, and the discrepancy between the theory and the data decreases. Furthermore, an analysis was carried out[23,24,18,19] that went beyond the two-gluon exchange approximation[17] and the complete BFKL amplitudes for all helicities were determined. The differential cross sections for all the mesons were fitted very well[23,24,19] as well as the energy dependence and spin density matrix for J/ψ production. Although the inclusion of BFKL effects greatly improves the fit of the spin density matrix elements, this description is not satisfactory for light vector mesons. In particular, both the sign and the magnitude of the single helicity flip amplitude come out wrong when the most probable values of wave function parameters are used. The model is capable to provide a reasonable fit of the data with an unrealistically large value of the constituent quark mass in the photon wave function. This may indicate that the model lacks some physical mechanism that suppresses the scattering of large dipoles e.g. the Sudakov form-factor.[19]

Another important idea is to use the meson wave function beyond the fixed twist approximation.[27] In the model, the BFKL evolution is included in the doubly logarithmic approximation but the chiral odd components in the photon are not taken into account. As a consequence, the single helicity flip amplitude is found to dominate in a severe contradiction with the data.

The lesson learned from those three approaches may be summarised in the following. The chiral odd components of the photon play an important rôle in the vector meson photoproduction. Calculations based on the BFKL equation[23,24,18,19] give a better picture of all the observables than the two gluon approximation. An improvement is needed in the treatment of the large dipole scattering, perhaps an inclusion of the Sudakov form-factor. Another interesting task is to apply the BFKL formalism to compute the helicity amplitudes at a large photon virtuality and non-zero $|t|$.

5. The Odderon

The Odderon is the leading part of the colour singlet exchange at high energies with odd C-parity, a partner of the C-even Pomeron. The surprising feature of the Odderon is that only very little evidence of this exchange has been found so far.[28] HERA experiments may be able to shed some light on the puzzle.

A simple and robust signal of the Odderon could be provided by a diffractive photoproduction of a C-even meson. In the perturbative QCD, three t-channel gluons is the simplest system which can take the quantum numbers of the Odderon and that contributes at asymptotic energies. The necessary inclusion of the leading logarithmic corrections leads to the BKP evolution equation[29] for the Odderon. The equation yields amplitudes with the intercept equal to one,[30] that is the diffractive cross sections constant with the energy.

Diffractive η_c photoproduction was studied within the perturbative framework. The charmed quark mass sets the scale for the problem, and the $\gamma \to \eta_c$ impact factor may be calculated in perturbative QCD. For the coupling of the Odderon to the proton, one is forced to use a model, and thus, to introduce some uncertainties. The first estimates of the cross section based on the three gluon approximation gave $\sigma(\gamma p \to \eta_c p) \simeq 10$ pb,[31] a signal probably too weak to detect given the low acceptance for η_c. The QCD evolution effects increased the prediction significantly[32] to yield $\sigma(\gamma p \to \eta_c p \simeq 50)$ pb. In addition, the cross section for a similar diffractive process with the proton dissociation was found to be even larger,[33] $\sigma(\gamma p \to \eta_c X \simeq 60)$ pb. Thus, the combined cross section for the diffractive η_c production exceeds earlier estimates by one order of magnitude and may turn out to be large enough for HERA experiments to discover the perturbative Odderon exchange.

In search of the Odderon, other promising observables have been also proposed. Estimates of the diffractive photoproduction cross-section of π^0, f_2, a_2^0 were obtained in the Stochastic Vacuum Model.[34] H1 collaboration set the upper limits for the corresponding cross sections[35] and, unfortunately, the results contradict the prediction for π^0 production. Besides that, an important idea was developed to use the Pomeron-Odderon interference effects in diffractive processes in the Odderon searches. The charge and spin asymmetries in $\pi^+\pi^-$ diffractive electroproduction were shown to be sensitive to the Odderon amplitude.[36]

6. Conclusions

Theory of diffractive meson production has been developing rapidly since about ten years. It touches some central problems of QCD, transition between the soft and the hard regime of the scattering, the resummation schemes at large energies and the saturation phenomenon. In addition, the analysis of exclusive processes showed importance of the non-forward parton distributions and stimulated studies of their properties and their QCD evolution. Analysis of the high-t photoproduction of vector mesons clearly indicates that the photon and light vector mesons contain sizeable chiral odd components. Diffractive production of C-even mesons may shed some light on the Odderon puzzle. All of this shows how much inspiration originates from this field.

Recent advances of the theory narrow the gap between the first principles of QCD and some successful phenomenological models. The mechanism of diffractive factorisation which is the basis for perturbative approach is getting better understood. The Balitsky-Kovchegov equation provides a QCD description of the perturbative parton saturation phenomenon within a well defined approximation scheme. Exchange amplitudes of the perturbative Pomeron and Odderon are computed for increasing number of processes at the leading logarithmic accuracy, and estimates of non-leading corrections are also known in some cases. The good news is, that together with the improving theoretical foundations of the phenomenological analysis, the data become better described. This may indicate that the crucial mechanisms that govern the diffractive meson production are correctly identified. Still, the quantitative description of the data is rather crude.

Wealth of HERA data, to be supplemented by the new data from HERA2 will certainly keep stimulating efforts to obtain a more complete QCD description of the diffractive meson production. With an improved statistics, the constraints on the theoretical estimates will become more stringent. We will learn more on some important phenomena, like the unitarisation of QCD amplitudes in the impact parameter plane. Finally, there is a lot of room left for measurements of new observables, for surprising results and for new theoretical ideas.

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References

1. see for instance P. Thompson, these Proceedings.
2. L. N. Lipatov, Sov. J. Nucl. Phys. **23** (1976) 338; E. A. Kuraev, L. N. Lipatov and V. S. Fadin, Sov. Phys. JETP **44**, 443 (1976); *ibid.* **45** (1977) 199; I. I. Balitsky and L. N. Lipatov, Sov. J. Nucl. Phys. **28** (1978) 822.
3. K. Golec-Biernat and M. Wüsthoff, Phys. Rev. D **59** (1999) 014017; Phys. Rev. D **60** (1999) 114023.
4. J. C. Collins, L. Frankfurt and M. Strikman, Phys. Rev. D **56** (1997) 2982.
5. A.V. Radyushkin, Phys. Lett. B **385** (1996) 333; Phys. Rev. D **56** (1997) 5524.
6. P. Hoyer, J. T. Lenaghan, K. Tuominen and C. Vogt, arXiv:hep-ph/0210124.
7. S. Chekanov *et al.* [ZEUS Collaboration], Eur. Phys. J. C **26** (2003) 389.
8. A. Ivanov and R. Kirschner, Eur. Phys. J. C **29** (2003) 353.
9. M. G. Ryskin, Z. Phys. C **57** (1993) 89.
10. S. J. Brodsky *et al.* Phys. Rev. D **50** (1994) 3134.
11. S. Munier, A. M. Staśto and A. H. Mueller, Nucl. Phys. B **603** (2001) 427.
12. L. Motyka, Acta Phys. Polon. B **34** (2003) 3069.
13. H. Kowalski and D. Teaney, arXiv:hep-ph/0304189.
14. E. Gotsman *et al.* Acta Phys. Polon. B **34** (2003) 3255.
15. I. Balitsky, Nucl. Phys. B **463** (1996) 99; Y. V. Kovchegov, Phys. Rev. D **60** (1999) 034008.
16. C. Adloff *et al.* [H1 Collaboration], Phys. Lett. B **539** (2002) 25; A. Aktas *et al.* [H1 Collaboration], Phys. Lett. B **568** (2003) 205.
17. D. Y. Ivanov *et al.* Phys. Lett. B **478** (2000) 101.
18. R. Enberg *et al.* JHEP **0309** (2003) 008;
19. G. G. Poludniowski *et al.* arXiv:hep-ph/0311017.
20. J. Bartels *et al.* Phys. Lett. B **348** (1995) 589.
21. L. N. Lipatov, Phys. Rept. **286** (1997) 131.
22. J. R. Forshaw and M. G. Ryskin, Z. Phys. C **68** (1995) 137.
23. J. R. Forshaw and G. Poludniowski, Eur. Phys. J. C **26** (2003) 411.
24. R. Enberg, L. Motyka and G. Poludniowski, Eur. Phys. J. C **26** (2002) 219.
25. P. Ball, V. M. Braun and N. Kivel, Nucl. Phys. B **649** (2003) 263.
26. P. Ball *et al.* Nucl. Phys. B **529** (1998) 323.
27. A. Ivanov and R. Kirschner, arXiv:hep-ph/0311077.
28. C. Ewerz, arXiv:hep-ph/0306137.
29. J. Bartels, Nucl. Phys. B **175** (1980) 365; J. Kwieciński and M. Praszałowicz, Phys. Lett. B **94** (1980) 413.
30. J. Bartels *et al.* Phys. Lett. B **477** (2000) 178.
31. J. Czyżewski *et al.* Phys. Lett. B **398** (1997) 400.
32. J. Bartels *et al.* Eur. Phys. J. C **20** (2001) 323.
33. J. Bartels, M. A. Braun and G. P. Vacca, arXiv:hep-ph/0304160.
34. E. R. Berger *et al.* Eur. Phys. J. C **9** (1999) 491; Eur. Phys. J. C **14** (2000) 673.
35. C. Adloff *et al.* [H1 Collaboration], Phys. Lett. B **544** (2002) 35.
36. P. Hagler *et al.* Phys. Lett. B **535** (2002) 117; Eur. Phys. J. C **26** (2002) 261.