DIFFRACTIVE PROCESSES AND QCD

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Diffractive processes are considered from both s–channel and t–channel points of view. Soft and hard diffractive processes are discussed. An importance of unitarity (shadowing) effects for a relation between hard diffraction at HERA and Tevatron is emphasized. A model, which takes into account unitarity effects is developed for to predict interaction of high–energy virtual photons with nucleons. It is shown that this model gives a good description of HERA data on both total $\gamma^* p$ total cross section and diffractive dissociation of virtual photons in a broad region of Q^2. Diffractive production of jets in hadronic interactions is investigated and an important role of unitarization effects is emphasized. It is shown how to describe the CDF data on diffractive jet production at Tevatron using an information on distribution of partons in the pomeron from HERA experiments. Double–pomeron production of jets at Tevatron shows a clear violation of factorization and is in a good agreement with our prediction.

Contents

1 Introduction

This review paper on diffractive processes is devoted to a memory of Misha Marinov. I has met Misha 40 years ago when I started to visit ITEP as a student of I.Ya.Pomeranchuk. I was always impressed by his very serious and deep style of thinking. He tried to understand all subtleties in the field of high–energy physics. His deep knowledge of many complicated problems in our field was well known and highly appreciated by physicists of ITEP. Misha had a great pedagogical talent which unfortunately was not used to full strength. Only once we asked him to give lectures on path integral methods. These lectures became a basis for his famous review on this subject. In 60-ies many people in ITEP including Misha were interested in high–energy hadronic interactions. The theoretical approach based on Regge theory was very popular at that time. Misha Marinov has investigated scattering of particles with spins at high energies. I learned a lot on this subject from his papers and discussions with him. At present there is a revival of interest to Regge approach mostly due to successful applications of this method to deep–inelastic scattering in the small–x region, studied experimentally at HERA and theoretical investigation of the role of Regge singularities in QCD.

Investigation of diffractive scattering of hadrons gives an important information on structure of hadrons and mechanisms of their interactions. There are two complementary points of view to these processes.

The s–channel view of diffraction

It is well known that an absorption of an initial wave due to many inelastic channels leads by unitarity to diffractive elastic scattering. At high energies lifetimes of hadronic fluctuations are large $\tau \sim E/m^2$ and hadronic constituents can undergo elastic scattering . This process leads to diffraction dissociation of hadrons. An elegant interpretation of diffraction dissociation in terms of eigenstates of diffractive part of S-matrix has been given by Good and Walker.[1] Consider the states ϕ_k which diagonalize the diffractive part of the T matrix. Such eigenstates of diffraction only undergo elastic scattering. Let us denote the orthogonal matrix which diagonalizes Im T by C, so that

$$\mathrm{Im}\, T \;=\; CFC^T \qquad \text{with} \qquad \langle\phi_k|F|\phi_j\rangle \;=\; F_j\,\delta_{jk}. \tag{1}$$

Now consider the diffractive dissociation of an arbitrary incoming state

$$|i\rangle \;=\; \sum_k C_{ik}\,|\phi_k\rangle. \tag{2}$$

The elastic scattering amplitude for this state satisfies

$$\langle i|\mathrm{Im}\ T|i\rangle \;=\; \sum_k |C_{ik}|^2\, F_k \;=\; \langle F\rangle, \tag{3}$$

where $F_k \equiv \langle\phi_k|F|\phi_k\rangle$ and where the brackets of $\langle F\rangle$ mean the average of F over the initial probability distribution of diffractive eigenstates. After the diffractive scattering described by T_{fi}, the final state $|f\rangle$ will, in general, be a different superposition of eigenstates than those of $|i\rangle$ shown in (2). At high energies one can neglect by the real parts of the diffractive amplitudes, then for cross sections at a given impact parameter b we have

$$\begin{aligned}
\frac{d\sigma_{\rm tot}}{d^2b} &= 2\,\mathrm{Im}\langle i|T|i\rangle \;=\; 2\sum_k |C_{ik}|^2\, F_k \;=\; 2\langle F\rangle \\
\frac{d\sigma_{\rm el}}{d^2b} &= |\langle i|T|i\rangle|^2 \;=\; \left(\sum_k |C_{ik}|^2\, F_k\right)^2 \;=\; \langle F\rangle^2 \\
\frac{d\sigma_{\rm el\,+\,SD}}{d^2b} &= \sum_k |\langle\phi_k|T|i\rangle|^2 \;=\; \sum_k |C_{ik}|^2\, F_k^2 \;=\; \langle F^2\rangle.
\end{aligned} \tag{4}$$

It follows that the cross section for the single diffractive dissociation of a proton,

$$\frac{d\sigma_{\rm SD}}{d^2b} \;=\; \langle F^2\rangle \;-\; \langle F\rangle^2, \tag{5}$$

is given by the statistical dispersion in the absorption probabilities of the diffractive eigenstates.

Note that if all the components ϕ_k of the incoming diffractive state $|i\rangle$ were absorbed equally then the diffracted superposition would be proportional to the incident one and the inelastic diffraction would be zero. Thus if, at very high energies, the amplitudes F_k at small impact parameters are equal to the black disk limit, $F_k = 1$, then diffractive production will be equal to zero in this impact parameter domain and so will only occur in the peripheral b region. This behaviour takes place in pp (and $p\bar{p}$) interactions at Tevatron energies.

Let us consider as an example just two diffractive channels[2,3,4] (say, p, N^*), and assume, for simplicity, that the elastic scattering amplitudes for these two channels are equal. Then the T matrix has the form

$$\mathrm{Im}\ T \;=\; 1 \;-\; e^{-\Omega/2}, \tag{6}$$

where the eikonal matrix Ω has elements

$$\Omega^{fi}_{f'i'} \;=\; \Omega_0\, \omega^{fi}\, \omega_{f'i'}. \tag{7}$$

The individual ω matrices, which correspond to transitions from the two incoming hadrons, each have the form

$$\omega = \begin{pmatrix} 1 & \gamma \\ \gamma & 1 \end{pmatrix}. \tag{8}$$

The parameter $\gamma(s, b)$ determines the ratio of the inelastic to elastic transitions. The overall coupling Ω_0 is also a function of the energy $\sqrt{s}$ and the impact parameter b.

With the above form of ω, the diffractive eigenstates are

$$|\phi_1\rangle = \frac{1}{\sqrt{2}} \left(|p\rangle + |N^*\rangle\right), \qquad |\phi_2\rangle = \frac{1}{\sqrt{2}} \left(|p\rangle - |N^*\rangle\right). \tag{9}$$

In this basis, the eikonal has the diagonal form

$$\Omega^{mn}_{m'n'} = \Omega_0 \, r^{mn} \, r_{m'n'}, \tag{10}$$

where $m, n = \phi_1, \phi_2$ and

$$r = \begin{pmatrix} 1+\gamma & 0 \\ 0 & 1-\gamma \end{pmatrix}. \tag{11}$$

In the case where γ is close to unity, $\gamma = 1 - \varepsilon$, one of the eigenvalues is small.

The t–channel point of view on diffraction

The t–channel approach is based on the Regge model for diffractive processes. Regge poles have been introduced in particle physics in the beginning of 60-ies[5,6] and up to present time are widely used for description of high-energy interactions of hadrons and nuclei. Regge approach establishes an important connection between high energy scattering and spectrum of particles and resonances. In this approach diffractive processes are mediated by an exchange by the pomeron (P) – the leading Regge pole with vacuum quantum numbers. The pomeron plays a role of an exchanged "particle" and gives factorizable contributions to scattering amplitudes. In impact parameter space Regge amplitude has asymptotically gaussian form. This contradicts to a peripheral form for inelastic diffraction, which follows from unitarity in the s-channel picture. However in the Regge theory it is necessary to take into account not only Regge poles but also Regge cuts, which corresspond to exchange by several Regge poles in the t-channel. These contributions restore unitarity of the theory and for inelastic diffraction they lead to a peripheral form of impact

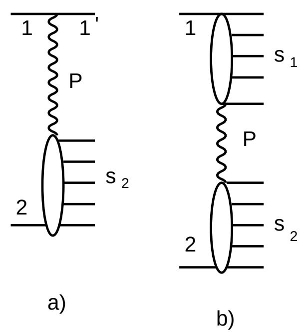

Figure 1: Diagrams for diffractive production of hadrons in the Regge pole model.

parameter distributions and allow to understand an energy dependence of the corresponding cross sections.

For all diffractive processes (Fig. 1) there is a large rapidity gap between groups of produced particles. For example for single diffraction dissociation there is a gap between the particle $1'$ and the rest system of hadrons. This rapidity gap $\Delta y \approx ln(1/1-x)$, where x the x_F for hadron $1'$. A mass of diffractively excited state at large s can be large. The only condition for diffraction dissociation is $s_i \ll s$.

The cross section for inclusive single diffraction dissociation in the Regge pole model can be written in the following form

$$\frac{d^2\sigma}{d\xi_2 dt} = \frac{(g_{11}(t))^2}{16\pi}|G_p(\xi',t)|^2\sigma_{P2}^{(tot)}(\xi_2,t) \tag{12}$$

where $\xi_2 \equiv ln(s_2/s_0), \xi' = ln(s/s_2)$ and $G_p(\xi',t) = \eta(\alpha_p(t))exp[(\alpha_p(t)-1)\xi']$ is the pomeron Green function. The quantity $\sigma_{P2}^{(tot)}(\xi_2,t)$ can be considered as the pomeron–particle total interaction cross section.[7] Note that this quantity is not directly observable one and it is defined by its relation to the diffraction production cross section, Eq. (12). This definition is useful however because at large s_2 this cross section has the same Regge behavior as usual cross sections

$$\sigma_{P2}^{(tot)}(s_2,t) = \sum_k g_{22}^k(0) r_{PP}^{\alpha_k}(t) \left(\frac{s_2}{s_0}\right)^{\alpha_k(0)-1} \tag{13}$$

where the $r_{PP}^{\alpha_k}(t)$ is the triple–reggeon vertex , which describes coupling of two pomerons to reggeon α_k.

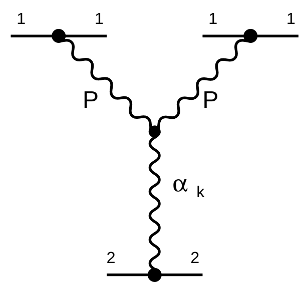

Figure 2: Triple–Regge diagram.

In this kinematical region $s \gg s_2 \gg m^2$ the inclusive diffractive cross section is described by the triple–Regge diagrams (Fig. 2)and has the form

$$f^1 = \sum_k G_k(t)(1-x)^{\alpha_k(0)-2\alpha_P(t)}(\frac{s}{s_0})^{\alpha_k(0)-1} \tag{14}$$

The pomeron–proton total cross section and triple–Regge vertices r^P_{PP}, r^f_{PP} have been determined from analysis of experimental data on diffractive production of particles in hadronic collisions (see review[8]).

2 Small-x physics

An interesting class of diffractive processes is studied in deep inelastic scattering (DIS) at small values of the Bjorken variable x . They became especially actual due to recent experimental investigation of the small-x region at HERA accelerator.

Experiments at HERA have found two extremely important properties of small-x physics : a fast increase of parton densities as x decreases[9,10] and the existence of substantial diffractive production in deep inelastic scattering (DIS).[11,12]

In the paper[13] it was suggested that the increase of the effective intercept of the pomeron, $\alpha_{eff} = 1 + \Delta_{eff}$, as Q^2 increases from zero to several GeV^2 is mostly due to a decrease of shadowing effects with increasing Q^2. A parametrization of the Q^2 dependence of Δ_{eff} such that $\Delta_{eff} \approx 0.1$ for $Q^2 \approx 0$ (as in soft hadronic interactions) and $\Delta_{eff} \approx 0.2$ (bare pomeron intercept) for Q^2 of the order of a few GeV^2, gives a good description[13,14] of all existing data on $\gamma^* p$ total cross-sections in the region of $Q^2 \leq 5 \div 10$ GeV^2. At larger Q^2 effects due to QCD evolution become important. Using the above parametrization as the initial condition in the QCD evolution equation, it is possible to describe the data in the whole region of Q^2 studied at HERA.[13,15]

In the reggeon approach there are good reasons to believe that the fast increase of the $\sigma^{(tot)}_{\gamma^* p}$ with energy in the HERA energy range will change to a softer increase at much higher energies. This is due to multi–pomeron effects, which are related to shadowing in highly dense systems of partons - with eventual "saturation" of densities. This problem was investigated recently in our paper,[18] where reggeon approach was applied to the processes of diffractive $\gamma^* p$ interaction.

In the reggeon calculus[16] the amount of rescatterings is closely related to diffractive production. AGK-cutting rules[17] allow to calculate the cross-section of inelastic diffraction if contributions of multi–pomeron exchanges to the elastic scattering amplitude are known. Thus, it is very important for self-consistency of theoretical models to describe not only total cross sections, but, simultaneously, inelastic diffraction. In particular in the reggeon calculus the variation of Δ_{eff} with Q^2 is related to the corresponding variation of the ratio of diffractive to total cross sections.

In the paper[18] an explicit model for the contribution of rescatterings was constructed which leads to the pattern of energy behavior of $\sigma^{(tot)}_{\gamma^* p}(W^2, Q^2)$ for different Q^2 described above. Moreover, it allows to describe simultaneously diffraction production by real and virtual photons. In this model it is possible to study quantitatively a regime of "saturation" of parton densities.

Let us discuss briefly the qualitative picture of diffractive dissociation of a highly virtual photon at high energies. It is convenient to discuss this process in the lab. frame, where the quark–gluon fluctuations of a photon live a long time $\sim 1/x$ (Ioffe time[19]). A virtual photon fluctuates first to $q\bar{q}$ pair. There are two different types of configurations of such pair, depending on transverse distance between quarks (or $k_\perp$).

a) Small size configurations with $k_\perp^2 \sim Q^2$. These small dipoles ($r \sim 1/k_\perp \sim 1/Q$) have a small ($\sim r^2$) total interaction cross section with the proton.

b) Large size configurations with $r \sim 1/\Lambda_{QCD}$ and $k_\perp \sim \Lambda_{QCD} \ll Q$. They have a large total interaction cross section, but contribute with a small phase space at large Q^2, because these configurations are kinematically possible only if the fraction of longitudinal momentum carried by one of the quarks is very small $x_1 \sim k_\perp^2/Q^2 \ll 1$. This configuration corresponds to the "aligned jet", introduced by Bjorken and Kogut.[20]

Both configurations lead to the same behaviour of $\sigma^{(tot)}_{\gamma^* p} \sim 1/Q^2$, but they behave differently in the process of the diffraction dissociation of a virtual photon.[21,22] The cross section of such a process is proportional to a square of modulus of the corresponding diffractive amplitude and for a small size configuration it is small ($\sim 1/Q^4$). For large size configurations a smallness is only due to the phase space and the inclusive cross section for diffractive dissociation of a virtual photon decreases as $1/Q^2$, i.e. in the same way as the total cross section. This is true only for the total inclusive diffractive cross section, where characteristic masses of produced states are $M^2 \sim Q^2$. For exclusive channels with fixed mass (for example production of vector mesons) situation is different and these cross sections decrease faster than $1/Q^2$ at large Q^2.

Inclusive diffractive production of very large masses ($M^2 \gg Q^2$) can be described in the first approximation by triple–Regge diagrams.[23] From the point of view of the quark-gluon fluctuation of the fast photon triple–pomeron contribution corresponds to diffractive scattering of very slow (presumably gluonic) parton, which has a small virtuality.

The model[18] uses the picture of diffraction dissociation of a virtual photon outlined above and is a natural generalization of models used for the description of high-energy hadronic interactions. The interaction of the small size component in the wave function of a virtual photon is calculated using QCD perturbation theory.

The main parameter of the model – intercept of the pomeron was fixed from a phenomenological study of hadronic interactions[24] ($\Delta_P = 0.2$) and was found to give a good description of $\gamma^* p$-interactions in a broad range of Q^2 ($0 \leq Q^2 < 10\ GeV^2$). Another important parameter of the theory, the triple–pomeron vertex, obtained from a fit to the data ($r^{(0)}_{PPP}/g^P_{pp}(0) \approx 0.1$) is also in a reasonable agreement with the analysis of soft hadronic interactions.[24,23] The model describes experimental data on the structure function F_2 and diffractive structure function $F^{(3)}_{2D}$ quite well (Figs. 3, 4). It can be used to predict structure functions and partonic distributions at higher energies or smaller x.

Investigation of the Q^2-dependence of diffraction dissociation of a highly virtual photon gives a possibility to determine a distribution of gluons in the pomeron.[12,23,25] It turns out that contrary to the proton case both distributions of quarks and gluons inside the pomeron are rather hard. This indicates to an important contrubution of "valence" gluons in the pomeron. It should be noted however that the form of the distribution of gluons in the pomeron is not well known at present. This is partly due to systematic differences in Q^2 dependence of diffractive structure functions of H1 and ZEUS (for recent discussion of these points see[26]).

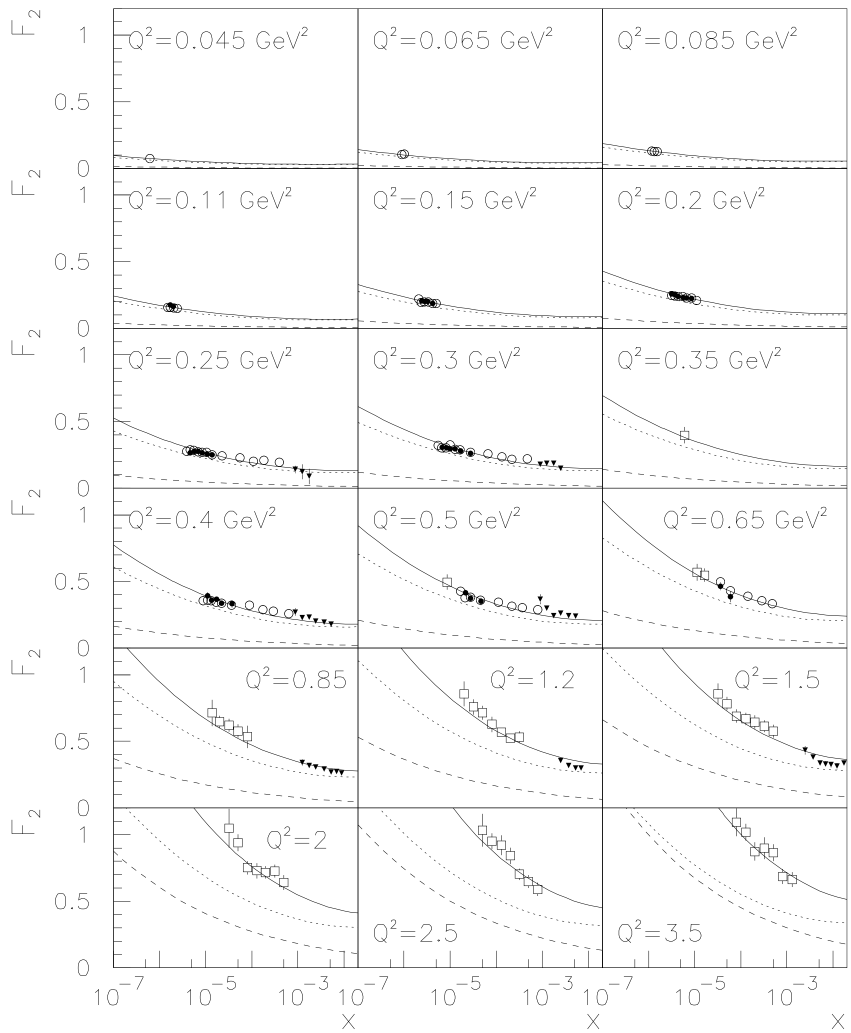

Figure 3: Structure function F_2 as a function of x for different values of Q^2 compared with experimental data. Dashed lines denote small distance contributions and dotted lines – large distance ones.

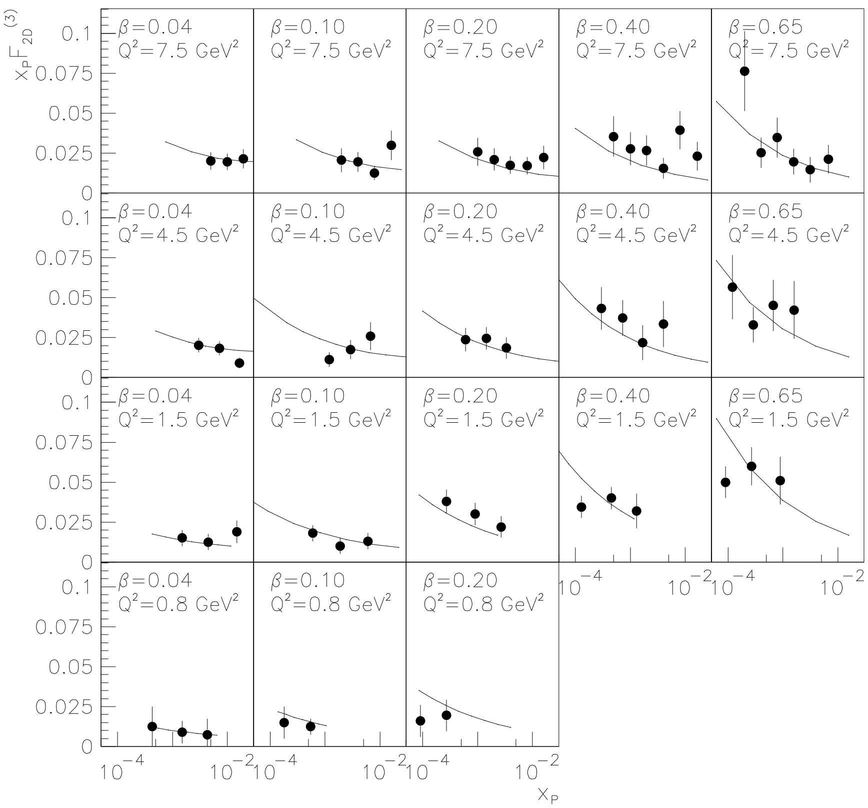

Figure 4: The diffractive structure function $x_P F_2^{D(3)}$ as a function of $x_P = x/\beta$ for fixed values of Q^2 and $\beta = Q^2/(Q^2 + M^2)$.

Shadowing effects for nuclear structure functions

A study of the shadowing effects for structure functions of nuclei in the small x region provides a stringent test of the reggeon approach to small–x problem. The shadowing effects are enhanced for nuclei ($\sim A^{1/3}$) and lead to deviations from A^1 behaviour for structure functions of nuclei. Glauber–Gribov[27,28] approach to interactions of particles with nuclei gives a possibility to calculate rescattering corrections for interaction of a virtual photon with a nucleus in terms of diffractive interaction of a photon with a nucleon, which was discussed above.

A contribution of a double rescattering term to the $\sigma_{\gamma^* A}$ is directly expressed in terms of the differential cross section for a diffraction dissociation of a virtual photon in $\gamma^* N$-interactions

$$\sigma^{(2)} = -4\pi \int d^2b \, T_A^2(b) \int dM^2 \, \frac{d\sigma^{DD}_{\gamma^* N}(t=0)}{dM^2 dt} \, F_A(t_{min}) \tag{15}$$

where $F_A(t_{min}) = \exp(R_A^2 t_{min}/3)$, $t_{min} \approx -m_N^2 x_P^2$, and $T_A(b)$ is the nuclear profile function ($\int d^2b T_A(b) = A$).

Higher order rescatterings are model dependent and in the generalized Schwimmer model[29] we obtain in the region of small x

$$F_{2A}/F_{2N} = \int d^2b \, \frac{T_A(b)}{1 + F(x,Q^2) T_A(b)} \tag{16}$$

with

$$F(x,Q^2) = 4\pi \int dM^2 \, \frac{d\sigma^{DD}_{\gamma^* N}(t=0)}{dM^2 dt} \, \frac{F_A(t_{min})}{\sigma_{\gamma^* N}(x,Q^2)} \, .$$

Theoretical predictions,[30] based on Eq. (16) and the model for diffraction dissociation of Ref. 23 are in a very good agreement with NMC-data on nuclear structure functions at very small x.[31] I believe that this appraoch gives reliable predictions for nuclear shadowing effects in the region of smaller x not yet studied experimentally. This region will play an important role in dynamics of heavy ions collisions at super–high (LHC) energies.

3 Hard diffraction in hadronic interactions

Diffractive production of large mass states, which include hard subprocesses, has been studied recently in high–energy hadronic interactions. Most of the results have been obtained at Tevatron where diffractive production of jets, W-bosons, heavy quarks and heavy quarkonia was observed.[32,33,34]

The most interesting results were obtained for diffractive dijet production. These observations, coupled with the diffractive measurements by H1[12] and ZEUS[11] at HERA, offer the opportunity to explore the relation between hard diffraction at HERA and Tevatron. The processes are shown schematically in Fig. 5, in the absence of rescattering corrections. Without rescattering

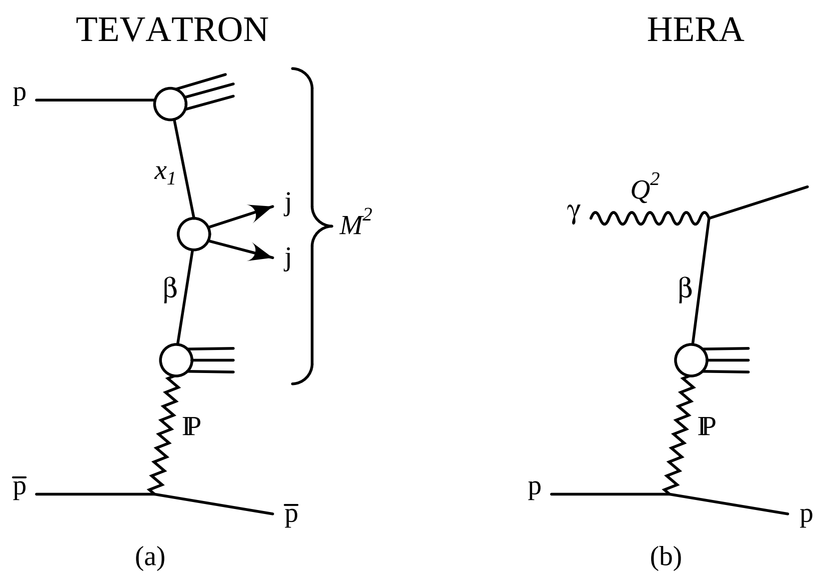

Figure 5: Schematic diagrams for diffractive dijet production at the Tevatron and for diffractive deep icelastic scattering at HERA.

corrections the cross section for diffractive dijet production , integrated over t, may be written as

$$\sigma = \sum_{i,j} \int F_P(\xi) f_i^P(\beta, E_T^2) f_j^p(x_1, E_T^2) \hat{\sigma} d\beta dx_1, \tag{17}$$

where $\hat{\sigma}$ is the cross section to produce dijets from partons carrying longitudinal momentum fractions x_1 and β of the proton and pomeron respectively. It was found[32,34] that calculation of the cross section, based on this factorized formula with diffractive structure functions obtained from HERA data, leads to a large discrepancy with the CDF measurements both in the normalization and in the shape of the observed distribution. Even taking into account uncertainties in gluonic distribution of the pomeron discussed above the calculation lies about a factor of 10 above the data. A fast increase of the cross section for $\beta \to 0$ observed at Tevatron differs strongly from predictions of models based on HERA data. The last result is difficult to understand in the models which take into account an extra suppression in hadronic interactions ("survival probability"[35]) due to rescatterings in the eikonal approximation.[36,37]

This problem was considered in Ref. 38. It was emphasized that a large cross section for single diffraction dissociation in hadronic collisions indicate to a large dispersion of eigen cross sections for diagonal states. It was assumed that the sea quarks and gluons mainly occur in large size configurations of the incident proton and thus have large cross sections, while the valence quarks occupy predominantly small–size configurations and are absorbed with substantially smaller cross sections.

In such a two-channel model the survival probability of the gaps can be written as follows

$$|S|^2 = \frac{\int d^2b \left(|M_v|^2 e^{-\Omega_v(s,b)} + |M_{sea}|^2 e^{-\Omega_{sea}(s,b)}\right)}{\int d^2b \left(|M_v|^2 + |M_{sea}|^2\right)}, \qquad (18)$$

where $M_{v,sea}$ are the probability amplitudes (in impact parameter space) of the hard diffractive process corresponding to the valence quark and to the sea quarks and gluons respectively. The functions Ω_i can be parametrized in the form

$$\Omega_i = K_i \frac{(g^P_{pp})^2 (s/s_0)^\Delta}{4\pi B} e^{-b^2/4B}, \qquad (19)$$

with $i = v, sea$, and where the slope of the pomeron amplitude is

$$B = \tfrac{1}{2} B_0 + \alpha' \ln(s/s_0), \qquad (20)$$

with $s_0 = 1$ GeV^2. We take $K_{\rm v} = 1 - \gamma$ and $K_{\rm sea} = 1 + \gamma$, consistent with the simple physical model introduced above. The value of the parameters g^P_{pp}, B_0, Δ and α' were determined in a two-channel global description of the total, elastic differential and soft diffraction cross sections,[4] in which the parameter γ was fixed to be 0.4.

In this model the soft rescattering effects ($\Omega_i \neq 0$) of the model based on (18) modify the β distribution of the dijet process in a characteristic way. First note that the CDF measurements cover a narrow ξ interval, $0.035 \le \xi \le 0.095$, and hence that the invariant mass squared of the diffractively produced state, $M^2 = \xi s$, remains close to the average value 2×10^5 GeV^2. Also the mass squared of the produced dijet system

$$M^2_{jj} \;=\; x_1 \beta M^2, \qquad (21)$$

does not change much compared to its average value of 1.2×10^3 GeV^2. Thus $x_1\beta \simeq 0.006$ and so for $\beta \geq 0.3$ we have $x_1 \leq 0.02$, whereas for $\beta \sim 0.03$ we have $x_1 \sim 0.2$. Therefore for large β (small x_1) sea quarks and gluons will give the dominant contribution, while for small β the valence quarks play an important role. Hence the survival probability should increase as x_1 increases and β decreases. Results of calculations based on this model (Fig. 6) are in a good agreement with CDF data both in magnitude and form of β distribution.[38] In particular an anomalously strong increase of this distribution as $\beta \to 0$ is reasonably described. This calculation of diffractive dijet production illus-

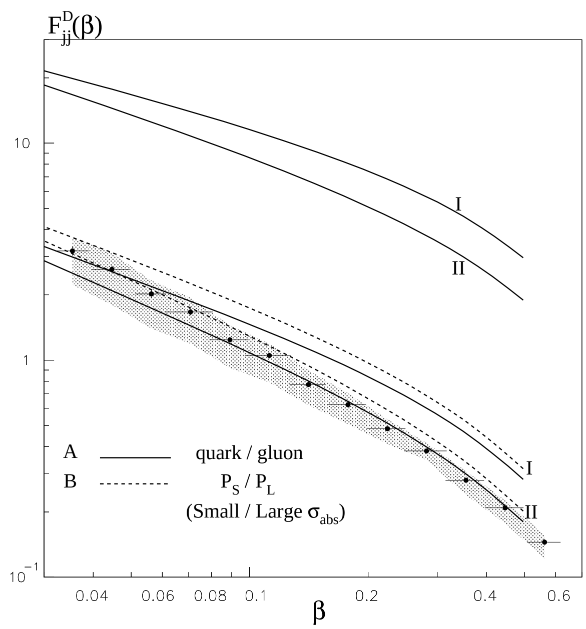

Figure 6: The predictions for diffractive dijet production at the Tevatron, obtained from two alternative sets of 'HERA'diffractive parton distributions , compared with the CDF data.[32] The upper two curves correspond to the neglect of rescattering corrections, whereas the lower four curves show the effect of including these corrections in two models A(continuous curves) and model B (dashed curves) for the diffractive eigenstates (see Ref. 38).

trates a crucial ingredient necessary in the description of rapidity gap processes. Namely that the survival probability of a gap can depend on x_1 of the partons of the proton . This leads to many experimental consequences for processes with rapidity gaps.[38]

Double–pomeron jet production

The process of particle production by double–pomeron exchange process (DPE) gives an important new information on mechanism of diffraction at high energies. Production of dijets with two large rapidity gaps gaps (x_p and $x_{\bar{p}}$ are close to 1) has been observed at Tevatron.[34] Comparison of cross section of this process (σ^{jj}_{DP}) with cross section of jet production in the process of single diffraction dissociation (σ^{jj}_{SD}) and in ordinary inelastic events (σ^{jj}_{in}) allows to study factorization properties of inelastic diffractive processes. The ratio $\sigma^{jj}_{SD}/\sigma^{jj}_{in}$ can be written as:

$$R_1 \equiv \frac{\sigma^{jj}_{SD}}{\sigma^{jj}_{in}} = \frac{F_P(\xi) f_g^P(\beta)}{f_g^p(x)} S_1^2, \tag{22}$$

where S_1^2 is the "survival probability" for single diffraction dissociation. For simplicity I consider only contribution of gluons. Account of quarks is straightforward.

The ratio of jet production in double–pomeron process to those for single one has the form

$$R_2 \equiv \frac{\sigma^{jj}_{DP}}{\sigma^{jj}_{SD}} = \frac{F_P(\xi_1) f_g^P(\beta_1)}{f_g^p(x_1)} \frac{S_2^2}{S_1^2}, \tag{23}$$

where S_2^2 is the "survival" probability for DPE.

So the ratio R_1/R_2 for the situation when $\xi = \xi_1$, $\beta = \beta_1$ ($x = x_1$) is equal to

$$R \equiv \frac{R_1}{R_2} = \frac{(S_1^2)^2}{S_2^2} \tag{24}$$

For a single Regge pole exchange ($S_i^2 = 1$) $R = 1$ and thus deviations of R from unity signals a breakdown of the factorization. With account of absorption $S_1^2 = 0.1$, $S_2^2 = 0.05$ at Tevatron energies.[4] So from Eq. (8) it follows that we can expect $R \approx 0.2$ at these energies. Recent results of CDF show that $R = 0.19 \pm 0.07$ in a good agreement with theoretical expectation.[39] This result shows that factorization is strongly violated in inelastic diffractive processes.

4 Conclusions

A study of diffractive processes provides an important information on interplay of soft and hard mechanisms in high energy interactions of hadrons. A new field of hard diffractive processes give possibility to study not only partonic content of the pomeron but also cross sections for absorption of different partonic configurations of the proton. The theory of supercritical pomeron with the intercept $\alpha_P(0) \approx 1.2$, which takes into account multi–pomeron exchanges gives a unified description of diffractive processes both for small–x DIS and hadronic interactions at high energies.

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References

1. M.L. Good and W.D. Walker, *Phys. Rev.* **126**, 1857 (1960).
2. K.G. Boreskov, A.M. Lapidus, S.T. Sukhorukov and K.A. Ter-Martirosyan, *Nucl. Phys.* B **40**, 397 (1972).
3. E. Gotsman, E. Levin and U. Maor, *Phys. Lett.* B **452**, 387 (1999); *Phys. Rev.* D **60**, 094011 (1999).
4. V.A. Khoze, A.D. Martin and M.G. Ryskin, *Eur. Phys. J.* C **18**, 167 (2000).
5. V.N. Gribov, *ZhETF*, **41**, 667 (1961). .
6. G.F. Chew and S.C. Frautschi, *Phys. Rev. Lett.* **7**, 394 (1961).
   R. Blankenbecler and M. Goldberger, *Phys. Rev.* **126**, 766 (1962).
7. A.B. Kaidalov and K.A. Ter-Martirosyan, *Nucl. Phys.* B **75**, 471 (1974).
8. A.B. Kaidalov, *Phys. Rept.* **50**, 157 (1979).
9. T. Ahmed *et al.* (H1 Collaboration), *Phys. Lett.* B **299**, 374 (1992);
   C. Adloff *et al.* (H1 Collaboration), *Nucl. Phys.* B **497**, 3 (1997).
10. M. Derrick *et al.* (ZEUS Collaboration), *Phys. Lett.* B **293**, 465 (1992);
    J. Breitweg *et al.* (ZEUS Collaboration), *Phys. Lett.* B **407**, 432 (1997).
11. M. Derrick *et al.* (ZEUS Collaboration), *Zeit. Phys.* C **72**, 399 (1996).
12. C. Adloff *et al.* (H1 Collaboration), *Zeit, Phys.* C **76**, 613 (1997).
13. A. Capella, A. Kaidalov, C. Merino and J. Tran Thanh Van, *Phys. Lett.* B **337**, 358 (1994).
14. A. Kaidalov and C. Merino, *Eur. Phys. J.* C **10**, 153 (1999).
15. A. Kaidalov, C. Merino and D. Pertermann, *Eur. Phys. J.* C **20**, 301 (2001).
16. V.N. Gribov, *ZhETF*, **57**, 654 (1967).

17. V.A. Abramovsky, V.N. Gribov and O.V. Kancheli, *Sov. J. Nucl. Phys.* **18**, 308 (1974).
18. A. Capella, E. Ferreiro, A.B. Kaidalov and C.A. Salgado, *Nucl. Phys.* B **593**, 336 (2001); *Phys. Rev.* D **63**, 054010 (2001).
19. B.L. Ioffe, *Phys. Lett.* B **30**, 123 (1969).
20. J. D. Bjorken and J. B. Kogut, *Phys. Rev.* D **8**, 1341 (1973).
21. L.L. Frankfurt and M. Strikman, *Phys. Rept.* **160**, 235 (1988).
22. N.N. Nikolaev and B.G. Zakharov, *Zeit. Phys.* C **49**, 607 (1990).
23. A. Capella, A. Kaidalov, C. Merino and J. Tran Than Van, *Phys. Lett.* B **343**, 403 (1995).
    A. Capella, A. Kaidalov, C. Merino, D. Pertermann and J. Tran Than Van, *Phys. Rev.* D **53**, 2309 (1996).
24. A.B. Kaidalov, L.A. Ponomarev and K.A. Ter-Martirosyan, *Sov. J. Nucl. Phys.* **44**, 468 (1986).
25. K.J. Golec-Biernat and J. Kwiecinski, *Phys. Lett.0* B **353**, 329 (1995); T. Gehrmann and W.J. Stirling, *Zeit. Phys.* C **70**, 227 (1995).
26. C. Royon, L. Schoeffel, J. Bartels, H. Jung and R. Peschanski, *Phys. Rev.* D **63**, 074004 (2001).
27. R.J. Glauber,in*Lectures in Theoretical Physics*: Proceedings of Annual Summer Institute for Theoretical Physics, Boulder, Calorado, 16 Jun – 22 Aug 1958, eds. W.E. Brittin and L.G. Dunham (Interscience, 1959) Vol. 1, p. 315.
28. V.N. Gribov, *JETP* **56**, 892 (1969);
    V.N. Gribov, *JETP*, **57**, 1306 (1969).
29. A. Schwimmer, *Nucl. Phys.* B **94**, 445 (1975).
30. A. Capella, A. Kaidalov, C. Merino, D. Perterman and J. Tran Than Van, *Eur. Phys. J.* C **5**, 111 (1998).
31. M. Arneodo *et al.* (NMC Collaboration), *Nuovo Cim.* **107A**, 2141 (1994); P. Amaudruz *et al.* (NMC Collaboration), *Nucl. Phys.* B **441**, 3 (1995).
32. T. Affolder *et al.* (CDF Collaboration), *Phys. Rev. Lett.* **84**, 5043 (2000).
33. F. Abe *et al.* (CDF Collaboration), *Phys. Rev. Lett.* **79**, 2636 (1997).
34. K. Goulianos, *Nucl. Phys. Proc. Suppl.* **99**, 37 (2001).
35. J.D. Bjorken, *Phys. Rev.* D **47**, 101 (1993).
36. E.M. Levin, hep-ph/9912402 and references therein.
37. M.M. Block and F. Halzen, *Phys. Rev.* D **63**, 114004 (2001).
38. A.B. Kaidalov, V.A. Khoze, A.D. Martin and M.G. Ryskin, *Eur. Phys. J.* C **21**, 521 (2001).
39. A.B. Kaidalov, V.A. Khoze, A.D. Martin and M.G. Ryskin, (in preparation).