

# Proton dissociation into three jets

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## Abstract

We explore the possibility to observe hard exclusive three-jet production in early LHC runs, corresponding to diffractive dissociation of the incident proton into three jets with large but compensating transverse momenta. This process is sensitive to the proton unintegrated gluon distribution at small  $x$  and to the distribution of the three valence quarks in the proton at small transverse distances. The corresponding cross section is calculated using an approach based on  $k_t$  factorization. According to our estimates, observation of hard diffractive three-jet production at LHC is feasible for jet transverse momenta  $q_\perp \sim 5$  GeV.

**1.** The physics potential of forward detectors at LHC, within and beyond the standard model, is attracting a lot of attention, cf. [1, 2, 3, 4, 5]. In this Letter we explore the possibility to observe hard exclusive diffractive dissociation of a proton into three hard jets in proton-proton collisions

$$p(p_1) + p(p_2) \rightarrow jet(q_1) + jet(q_2) + jet(q_3) + p(p'_2), \quad (1)$$

cf. [6] In this process one proton stays intact and the other one dissociates into a system of three hard jets separated by a large rapidity gap from the recoil proton, see Fig. 1. The main aim of our study is to estimate the cross section of this reaction and the corresponding event rates at LHC and Tevatron.

Note that we are interested in *exclusive* three-jet production which constitutes a small fraction of the inclusive single diffraction cross section. The exclusive and inclusive mechanisms have different final state topologies and can be distinguished experimentally. A characteristic quantity is e.g. the ratio  $R_{jets}$  of the three-jet mass to the total invariant mass of the system produced in the diffractive interaction. Exclusive production corresponds to the region where  $R_{jets}$  is close to unity. This strategy was used recently at the Tevatron [7] where central exclusive dijet production,  $p\bar{p} \rightarrow p + jet + jet + \bar{p}$ , in double-Pomeron collisions was measured for the first time.

Exclusive dijet production in the central region has much in common with the exclusive Higgs boson production process,  $p\bar{p} \rightarrow p + H + \bar{p}$ . In [8] it was argued that studies of exclusive dijet production and other diffractive processes at the early data runs of the LHC can provide

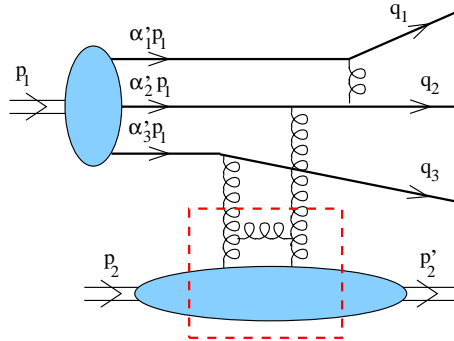


Figure 1: Proton dissociation into three jets. The unintegrated gluon distribution includes the hard gluon exchange as indicated by the dashed square.

valuable checks of the different components of the formalism. Indeed, this was the main motivation for Tevatron experiment. The exclusive 3-jets production in single diffraction (1) offers another interesting example since factorization of hard and soft interactions in this case is less complicated. In particular, the fluctuation of a proton projectile into a state with small transverse size, which is the underlining mechanism for (1), suppresses secondary soft interactions that may fill the rapidity gap. Thus one can get an access to the gluon distribution at small  $x$  in a cleaner environment, having no problems with gap survival probability and factorization breaking that introduce major conceptual theoretical uncertainties in the calculations of diffractive Higgs production.

Our approach to exclusive three-jet production derives from experience with coherent pion diffraction dissociation into a pair of jets with large transverse momenta which was measured by the E791 collaboration [9, 10]. The qualitative features of the E791 data have confirmed some earlier theoretical predictions [11, 12, 13]: a strong  $A$ -dependence which is a signature for color transparency, and a  $\sim 1/q_{\perp}^8$  dependence on the jet transverse momentum. These features suggest that the relevant transverse size of the pion  $r_{\perp}$  remains small, of the order of the inverse transverse momenta of the jets  $r_{\perp} \sim 1/q_{\perp}$ .

On a more quantitative level, we have shown [14, 15] that collinear factorization is violated in dijet production due to pinching of singularities between soft gluon (and quark) interactions in the initial and final state. However, the nonfactorizable contribution is suppressed compared to the leading contribution by a logarithm of energy so that in the double logarithmic approximation  $\ln q_{\perp}^2 \ln s/q_{\perp}^2$  collinear factorization is restored. Moreover, to this accuracy hard gluon exchange can be “hidden” in the unintegrated gluon distribution  $\mathcal{F}(x, q_{\perp})$ . Thus, in the true diffraction limit, for very large energies, hard exclusive dijet production can be considered as a probe of the hard component of the pomeron. The same interpretation was suggested earlier in [16] within the  $k_t$  factorization framework (see also [17]). The double logarithmic approximation turns out to be insufficient for the energy range of the E791 experiment, but might be adequate for the LHC. In this Letter we present an estimate for the cross section for the reaction (1) based on the generalization of these ideas.

**2.** At leading order the jets are formed by the three valence quarks of the proton, see Fig. 1. We require that all three jets have large transverse momenta which requires at least

two hard gluon exchanges. One of them can be effectively included in the high-momentum component of the unintegrated gluon density (the bottom blob) as indicated schematically by the dashed square, but the second one has to be added explicitly since the hard pomeron only couples to two of the three quarks of the proton.<sup>1</sup> The necessity for an additional hard gluon exchange makes calculation of proton diffraction dissociation more difficult as compared to the meson case.

Our notation for the momenta is explained in Fig. 1. We neglect power corrections in transverse momenta of the jets and also proton and jet masses so that  $p_1^2 = p_2^2 = p_3^2 = q_1^2 = q_2^2 = q_3^2 = 0$ . The jet momenta are decomposed in terms of momenta of the initial particles

$$q_k = \alpha_k p_1 + \beta_k p_2 + q_{k\perp}, \quad k = 1, 2, 3 \quad (2)$$

where

$$\vec{q}_{1\perp} + \vec{q}_{2\perp} + \vec{q}_{3\perp} = 0, \quad \alpha_1 + \alpha_2 + \alpha_3 = 1, \quad \beta_k = \vec{q}_{k\perp}^2 / (\alpha_k s), \quad (3)$$

The three-jet invariant mass is given by

$$M^2 = (q_1 + q_2 + q_3)^2 = \frac{\vec{q}_{1\perp}^2}{\alpha_1} + \frac{\vec{q}_{2\perp}^2}{\alpha_2} + \frac{\vec{q}_{3\perp}^2}{\alpha_3}, \quad \zeta = \frac{M^2}{s} = \beta_1 + \beta_2 + \beta_3. \quad (4)$$

where  $s = (p_1 + p_2)^2 = 2p_1 \cdot p_2$  is the invariant energy. Assuming that the relevant jet transverse momenta are of the order of 5 GeV, the typical values of the  $\zeta$  variable at LHC are in the range  $\zeta \sim 10^{-6} \div 10^{-5}$ .

At high energies, an amplitude is predominantly given by its discontinuity in the  $s$ -channel which usually implies that the amplitude is almost purely imaginary. In our case the situation is more complicated since in the physical region of (1) the amplitude develops a cut in the variable  $M^2$  as well, and it remains complex even after taking the  $s$ -channel discontinuity. Nevertheless, the  $s$ -channel discontinuity of the amplitude can be expressed, in our approximation, in terms of the unintegrated gluon distribution  $\mathcal{F}(x, k_\perp)$ .

The relevant Feynman diagrams in leading order of perturbative QCD are shown in Fig. 2. They can be divided into three groups which differ by the attachments of the  $t$ -channel gluons to the quark lines (shown by crosses). In diagrams (f)-(j) the hard gluon exchange takes place between quark  $q_2$  and  $q_3$ . Hence the transverse momentum of one of the  $t$ -channel gluons coincides with that of the quark  $q_1$ . As a consequence, the contribution of this group of diagrams involves the unintegrated gluon distribution at the same scale,  $\mathcal{F}(\zeta, q_{1\perp})$ . In diagrams (k)-(o) the hard gluon exchange connects the quarks  $q_1$  and  $q_3$ , the transverse momentum of the  $t$ -channel gluon coincides with the momentum of the quark  $q_2$  and, therefore, the unintegrated gluon distribution enters at this scale,  $\mathcal{F}(\zeta, q_{2\perp})$ . Similarly, the contribution of diagrams (a)-(e) is proportional to  $\mathcal{F}(\zeta, q_{3\perp})$ .

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<sup>1</sup>Alternatively, one can consider double-pomeron exchange in the  $t$ -channel. This contribution is suppressed by a power of the jet transverse momentum so it is of higher twist.

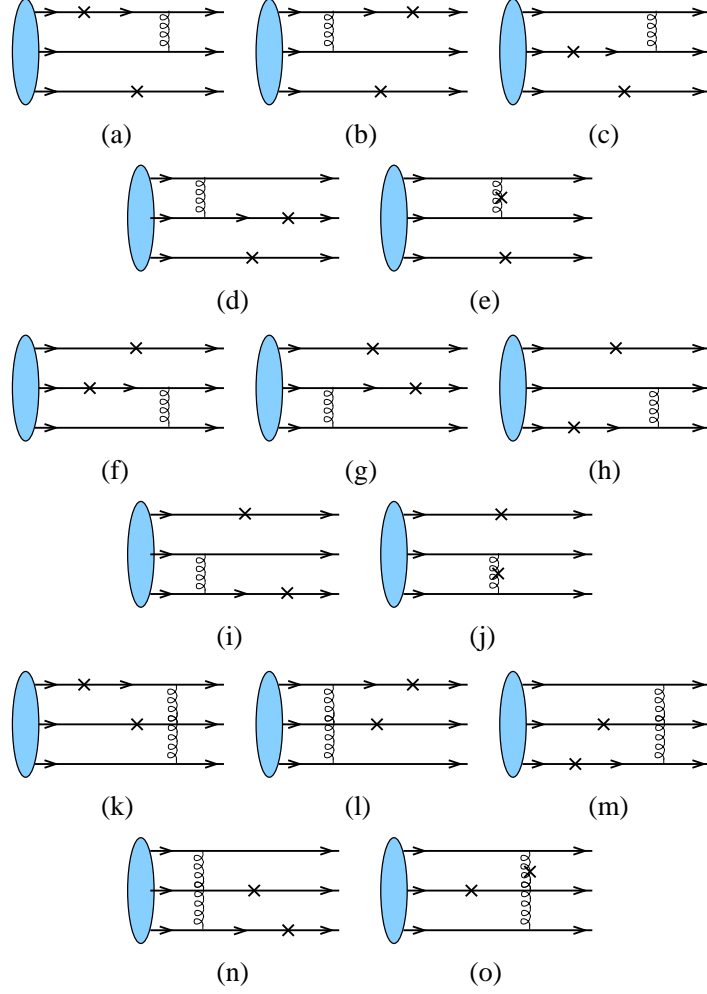


Figure 2: The leading-order contributions to proton disintegration into three jets within  $k_t$  factorization. The points where the t-channel gluons are attached to the quarks are shown by crosses.

Accordingly, we have three different contributions to the amplitude:

$$\begin{aligned}
\mathcal{M} = & -i 2^7 \pi^5 s \alpha_s^2 \left[ \frac{e^{ijk} \left(\frac{1+N}{N}\right)^2}{4N!(N^2-1)} \right] \int D\alpha' \\
& \times \left( \mathcal{L}_{f \div j} \frac{\delta(\alpha_1 - \alpha'_1)}{q_{1\perp}^4} \mathcal{F}(\zeta, q_{1\perp}) + \mathcal{L}_{k \div o} \frac{\delta(\alpha_2 - \alpha'_2)}{q_{2\perp}^4} \mathcal{F}(\zeta, q_{2\perp}) + \mathcal{L}_{a \div e} \frac{\delta(\alpha_3 - \alpha'_3)}{q_{3\perp}^4} \mathcal{F}(\zeta, q_{3\perp}) \right), \tag{5}
\end{aligned}$$

where  $\int D\alpha' = \int_0^1 d\alpha'_1 d\alpha'_2 d\alpha'_3 \delta(1 - \sum \alpha'_i)$  corresponds to the integration over the quark momentum fractions in the incident proton,  $e^{ijk}$  describes the color state of the final quarks,  $N = 3$  is the number of colors. The dimensionless quantities  $\mathcal{L}_i$  are expressed in terms of different Dirac structures where for convenience we introduce a "positron-like" Dirac spinor  $v$ ,

$$(\bar{u}(q_2))^T = C v(q_2), \quad (\bar{u}(q_2) \gamma_\mu)^T = -C \gamma_\mu v(q_2). \tag{6}$$

Here  $C$  is the charge-conjugation matrix. The diagrams with the 3-gluon vertex do not contribute due to vanishing color factors. Therefore we need to calculate in total 12 nontrivial diagrams (4 for each group).

The calculation is straightforward, though rather tedious. Here we present the final results only:

$$\begin{aligned}
\mathcal{L}_{f \div j} = & [\mathcal{V} \bar{u}(q_1) \not{p}_1 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_1 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1)] \\
& \times \left( \frac{\alpha_1(\alpha_3 + \alpha'_3)}{[-\alpha'_2 \alpha_3 \beta_2 s^3][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{-\alpha_1(\alpha_3 + \alpha'_3)}{[-\alpha'_3 \alpha_3 \beta_3 s^3][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right) \\
& + [\mathcal{V} \bar{u}(q_1) \not{p}_2 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_2 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1)] \\
& \times \left( \frac{\alpha'_2 \beta_2 - \alpha_1 \beta_1}{[-\alpha'_3 \alpha_3 \beta_3 s^3][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{\alpha_1 \beta_1 + \alpha_3 \beta_3 + (\alpha_3 - \alpha'_2) \beta_2}{[-\alpha'_2 \alpha_3 \beta_2 s^3][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right. \\
& \quad \left. + \frac{\alpha'_2 \beta_2 + \alpha'_3 \beta_3}{\alpha'_2 \alpha'_3 \beta_2 \beta_3 (\alpha_2 + \alpha_3) s^3} \right) \\
& + [\mathcal{V} \bar{u}(q_1) \not{p}_1 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_1 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1)] \\
& \times \left( \frac{-\alpha_1(\alpha_2 + \alpha'_2)}{[-\alpha'_2 \alpha_3 \beta_2 s^3][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{\alpha_1(\alpha_2 + \alpha'_2)}{[-\alpha'_3 \alpha_3 \beta_3 s^3][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right) \\
& + [\mathcal{V} \bar{u}(q_1) \not{p}_2 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_2 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1)] \\
& \times \left( \frac{\alpha'_2 \beta_2 + \alpha_1 \beta_1}{[-\alpha'_3 \alpha_3 \beta_3 s^3][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{-\alpha_1 \beta_1 + \alpha_3 \beta_3 + (\alpha_3 - \alpha'_2) \beta_2}{[-\alpha'_2 \alpha_3 \beta_2 s^3][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right. \\
& \quad \left. + \frac{\alpha'_2 \beta_2 + \alpha'_3 \beta_3}{\alpha'_2 \alpha'_3 \beta_2 \beta_3 (\alpha_2 + \alpha_3) s^3} \right) \\
& + \mathcal{T} [\bar{u}(q_1) v(q_2) \bar{u}(q_3) \gamma_5 N(p_1) - \bar{u}(q_1) \gamma_5 v(q_2) \bar{u}(q_3) N(p_1)] \\
& \times \left( \frac{\alpha'_3 - \alpha_1}{[-\alpha'_2 \beta_2 s^2][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{(\alpha_2 - \alpha_1) \beta_2 + (\alpha_2 - \alpha'_3) \beta_3}{[-\alpha'_3 \beta_3^2 s^2][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right. \\
& \quad \left. + \frac{(\alpha_2 - \alpha_1)(\alpha'_2 \beta_2 + \alpha'_3 \beta_3)}{\alpha'_2 \alpha'_3 \beta_2 \beta_3^2 (\alpha_2 + \alpha_3) s^2} \right) \\
& + \mathcal{T} [\bar{u}(q_1) (ip_{1\mu} p_{2\nu} \sigma_{\mu\nu}) v(q_2) \bar{u}(q_3) \gamma_5 N(p_1) - \bar{u}(q_1) (ip_{1\mu} p_{2\nu} \sigma_{\mu\nu}) \gamma_5 v(q_2) \bar{u}(q_3) N(p_1)] \\
& \times \left( \frac{2(1 - \alpha'_2)}{[-\alpha'_2 \beta_2 s^3][\alpha'_3(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} + \frac{2((\alpha_1 + \alpha_2) \beta_2 + (\alpha_2 - \alpha'_3) \beta_3)}{[-\alpha'_3 \beta_3^2 s^3][\alpha'_2(\beta_2 + \beta_3) - \alpha_1 \beta_1 + i\epsilon]} \right. \\
& \quad \left. + \frac{2(\alpha_1 + \alpha_2)(\alpha'_2 \beta_2 + \alpha'_3 \beta_3)}{\alpha'_2 \alpha'_3 \beta_2 \beta_3^2 (\alpha_2 + \alpha_3) s^3} \right), \tag{7}
\end{aligned}$$

$$\begin{aligned}
\mathcal{L}_{a\dot{+}e} = & [\mathcal{V} \bar{u}(q_1) \not{p}_1 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_1 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1)] \\
& \times \left( \frac{2(\alpha'_2 - \alpha_1)}{[-\alpha'_2 \beta_2 s^3][\alpha'_1(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} + \frac{2(\alpha'_1 - \alpha_2)}{[-\alpha'_1 \beta_1 s^3][\alpha'_2(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} \right) \\
& + [\mathcal{V} \bar{u}(q_1) \not{p}_2 v(q_2) \bar{u}(q_3) \not{p}_2 \gamma_5 N(p_1) - \mathcal{A} \bar{u}(q_1) \not{p}_2 \gamma_5 v(q_2) \bar{u}(q_3) \not{p}_2 N(p_1)] \\
& \times \left( \frac{2\beta_1}{[-\alpha'_2 \beta_2 s^3][\alpha'_1(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} + \frac{2\beta_2}{[-\alpha'_1 \beta_1 s^3][\alpha'_2(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} \right. \\
& \quad \left. + \frac{2(\alpha'_1 \beta_1 + \alpha'_2 \beta_2)}{\alpha'_1 \alpha'_2 \beta_1 \beta_2 (\alpha_1 + \alpha_2) s^3} \right) \\
& + \mathcal{T} [\bar{u}(q_1) v(q_2) \bar{u}(q_3) \gamma_5 N(p_1) - \bar{u}(q_1) \gamma_5 v(q_2) \bar{u}(q_3) N(p_1)] \\
& \times \left( \frac{-2\alpha_3}{[-\alpha'_2 \beta_2 s^2][\alpha'_1(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} + \frac{2\alpha_3}{[-\alpha'_1 \beta_1 s^2][\alpha'_2(\beta_1 + \beta_2) - \alpha_3 \beta_3 + i\epsilon]} \right) \quad (8)
\end{aligned}$$

The functions  $\mathcal{A}(\alpha'_1, \alpha'_2, \alpha'_3)$ ,  $\mathcal{V}(\alpha'_1, \alpha'_2, \alpha'_3)$  and  $\mathcal{T}(\alpha'_1, \alpha'_2, \alpha'_3)$  are the leading-twist light-cone nucleon distribution amplitudes defined as in [18]

$$\begin{aligned}
\langle 0 | \epsilon^{ijk} u_\alpha^i(a_1 z) u_\beta^j(a_2 z) d_\gamma^k(a_3 z) | N(p_1) \rangle = \\
= V(\not{p}_1 C)_{\alpha\beta} (\gamma_5 N(p_1))_\gamma + A(\not{p}_1 \gamma_5 C)_{\alpha\beta} (N(p_1))_\gamma + T(i\sigma_{\mu\nu} \not{p}_1 C)_{\alpha\beta} (\gamma^\mu \gamma_5 N(p_1))_\gamma, \quad (9)
\end{aligned}$$

where  $z^2 = 0$  and  $\sigma_{\mu\nu} = (i/2)[\gamma_\mu, \gamma_\nu]$ . Each invariant amplitude  $V, A, T$  depends on the scalar products  $a_i p_1 \cdot z$  and can be represented as the Fourier transform of the corresponding distribution amplitude, e.g.

$$V(a_i p_1 z) = \int D\alpha' e^{-ip_1 \cdot z \sum_i \alpha'_i a_i} \mathcal{V}(\alpha'_i). \quad (10)$$

Finally,  $\mathcal{L}_{k-o}$  is given by the expression similar to  $\mathcal{L}_{f-j}$  in Eq. (7) with the replacements  $\alpha_1 \leftrightarrow \alpha_2$ ,  $\alpha'_1 \leftrightarrow \alpha'_2$  and  $\beta_1 \leftrightarrow \beta_2$  everywhere except for the arguments of the distribution amplitudes.

One remark is in order. Another possible mechanism for the exclusive proton disintegration into three jets could be the exchange of three hard gluons in the  $t$ -channel. Such a contribution could involve two different color structures, proportional to  $f^{abc}$  and  $d^{abc}$ , which correspond to different  $C$ -parity in the  $t$ -channel. We found by explicit calculation that the color factor  $\sim f^{abc}$  corresponding to the  $C$ -parity-even exchange vanishes. The  $C$ -parity odd contribution  $\sim d^{abc}$  is related to odderon exchange and presumably small.

**3.** The differential cross section can be written as

$$d\sigma = \frac{|\mathcal{M}|^2}{2^5 (2\pi)^8 s^2} \frac{d\alpha_1 d\alpha_2 d\alpha_3 \delta(1 - \alpha_1 - \alpha_2 - \alpha_3)}{\alpha_1 \alpha_2 \alpha_3} d^2 \vec{q}_1 d^2 \vec{q}_2 dt d\phi_t \quad (11)$$

where  $t = (p_2 - p'_2)^2$  is the Mandelstam  $t$  variable of the  $pp$  scattering and  $\phi_t$  is the azimuthal angle of the final state proton. In our kinematics, for large transverse momenta of the jets and small  $t$ , one can neglect effects of azimuthal correlations between the jets and the final proton. Hence  $d\phi_t$  integration is trivial and gives a factor  $2\pi$ . For the  $t$  dependence we assume

a simple exponential form,  $d\sigma/dt \sim e^{bt}$ , and use  $b \sim 4 \div 5 \text{ GeV}^2$  for the slope parameter which is a typical value which describes HERA data for hard exclusive processes: DVCS and vector meson electroproduction at large  $Q^2$ . Thus, the integration over the proton recoil variables gives a factor  $\int dt d\phi_t \rightarrow \frac{2\pi}{b}$ .

Since our calculation is only done to double logarithmic accuracy, we use the simplest model for the unintegrated gluon distribution as given by the logarithmic derivative of the usual gluon parton distribution  $xg(x, Q^2)$

$$\mathcal{F}(x, q_\perp^2) = \frac{\partial}{\partial \ln q_\perp^2} x g(x, q_\perp^2). \quad (12)$$

The numerical estimates presented below are obtained using the CTEQ6L leading-order gluon distribution as provided by [19]. We also used the simplest, asymptotic model for the nucleon distribution amplitude

$$\mathcal{V}(\alpha'_i) = \mathcal{T}(\alpha'_i) = 120 f_N \alpha'_1 \alpha'_2 \alpha'_3, \quad \mathcal{A}(\alpha'_i) = 0. \quad (13)$$

The normalization parameter  $f_N$  is scale-dependent. To leading-logarithmic accuracy

$$f_N(\mu) = f_N(\mu_0) \left( \frac{\alpha_s(\mu)}{\alpha_s(\mu_0)} \right)^{\frac{2}{3\beta_0}}, \quad (14)$$

where  $\beta_0 = 11/3N - 2/3n_f$ . The existing QCD sum rule estimates

$$\begin{aligned} m_N f_N(\mu = 1 \text{ GeV}) &= (5.0 \pm 0.3) \times 10^{-3} \text{ GeV}^3 \quad [20], \\ m_N f_N(\mu = 1 \text{ GeV}) &= (5.1 \pm 0.4) \times 10^{-3} \text{ GeV}^3 \quad [21] \end{aligned} \quad (15)$$

are somewhat larger compared with a very recent  $n_f = 2$  (unquenched) lattice calculation

$$f_N(\mu = 2 \text{ GeV}) = (3.14 \pm 0.09) \times 10^{-3} \text{ GeV}^2 \quad [22]. \quad (16)$$

The given number corresponds to the lattice spacing  $a \simeq 0.067 \text{ fm}$ ; a continuum extrapolation was not attempted. For definiteness we use the value  $f_N = 5.0 \times 10^{-3} \text{ GeV}^2$  at 1 GeV as input, and evolve it to the relevant scale.

The integration over the phase space of the three jets was done numerically, restricting the longitudinal momentum fractions to the region<sup>2</sup>

$$0.1 \leq \alpha_1, \alpha_2, \alpha_3 \leq 0.8 \quad (17)$$

and requiring that the transverse momentum of *each* jet is larger than a given value  $q_0 = q_{\perp, \min}$ . For the value  $q_0 = 5 \text{ GeV}$  we obtain for the integrated three-jet cross section at the LHC energies

$$\sigma_{3-jets}^{\text{LHC}} = 4 \text{ pb} \cdot \left( \frac{f_N(q_0)}{4.7 \cdot 10^{-3} \text{ GeV}^2} \right)^2 \left( \frac{\alpha_s(q_0)}{0.21} \right)^4 \left( \frac{5 \text{ GeV}}{q_0} \right)^9. \quad (18)$$

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<sup>2</sup>We use this rather conservative cut condition to assure a clear three-jet event selection.

Assuming the integrated luminosity for the first LHC runs in the range  $100 \text{ pb}^{-1}$  to  $1 \text{ fb}^{-1}$  an observation of this process at LHC seems to be feasible. Note that the effective power  $\sigma \sim 1/q_0^9$  (fitted in the  $q_0 = 3 \div 8 \text{ GeV}$  range) is somewhat stronger than the naive power counting prediction  $\sigma \sim 1/q_0^8$ . This effect is due to the strong  $\zeta$  dependence of the unintegrated gluon distribution: larger values of  $q_0$  imply larger invariant masses  $M^2$  of the three-jet system (4) and consequently larger  $\zeta = M^2/s$ . The sizeable cross section for  $q_0 = 5 \text{ GeV}$  is in fact an implication of the expected rise of the LO gluon distribution more than two times as  $\zeta$  is decreasing by roughly a factor of 50 when going from Tevatron to LHC. The existing parameterizations of the LO gluon distribution at  $\zeta \sim 10^{-6}$  differ from each other by  $\sim 30\%$ . The unintegrated gluon distribution (12) enters as a square in the prediction for the cross section, therefore, the study of exclusive three-jet events at LHC may provide a valuable constraint for the gluon distribution at small momentum fractions.

A comparison of the three-jet exclusive production at LHC and the Tevatron can be especially illuminating in this respect since other uncertainties do not have significant impact on the energy dependence. For Tevatron kinematics, assuming the value  $q_{\perp \text{min}} = 3 \text{ GeV}$ , our estimate for the cross section (fitted in the range  $q_0 = 2 \div 4.5 \text{ GeV}$ ) is

$$\sigma_{3\text{-jets}}^{\text{Tevatron}} = 50 \text{ pb} \cdot \left( \frac{f_N(q_0)}{4.7 \cdot 10^{-3} \text{ GeV}^2} \right)^2 \left( \frac{\alpha_s(q_0)}{0.255} \right)^4 \left( \frac{3 \text{ GeV}}{q_0} \right)^9. \quad (19)$$

Note that in this case  $M^2 \sim 100 \text{ GeV}^2$  and  $\zeta \sim 10^{-5} \div 10^{-4}$  where the gluon distribution is much better known: The typical difference between existing parameterizations is of order  $\sim 10\%$ .

In the most naive approximation, the longitudinal momentum fraction distribution of the jets is expected to follow that of the valence quarks in the proton: The momentum fraction distribution of a jet, arbitrary chosen in each event, is proportional to the proton distribution amplitude squared

$$\frac{d\sigma}{d\alpha} \sim \int D\alpha' \delta(\alpha - \alpha'_1) |\phi_N(\alpha')|^2, \quad (20)$$

where  $\phi_N = \mathcal{V} - \mathcal{A}$  [18]. In reality, a hard gluon exchange leads to a certain redistribution of the longitudinal momenta so that the resulting  $\alpha$ -dependence is more complex, see e.g. [14, 15]. To illustrate this effect, in Fig. 3 we compare our calculated normalized jet momentum fraction distribution (averaged over quark flavors) to the dependence in (20): The two distributions are similar, but the one resulting from the QCD calculation is shifted towards lower momentum fractions compared to the simple dependence in Eq. (20). Note that measurement of the valence quark momentum fraction distribution in a pion presented the main motivation for the E791 experiment [9, 10] which, in turn, triggered detailed studies of such reactions in pQCD.

**4.** To summarize, in this Letter we have studied the exclusive diffractive dissociation of a proton into three jets with large transverse momenta in the double-logarithmic approximation of perturbative QCD. This process is interesting in the broader context of diffractive processes at LHC, which will be studied using forward detectors, and in particular can be used to constrain the gluon distribution at very small values of Bjorken  $x$ . According to our estimates,

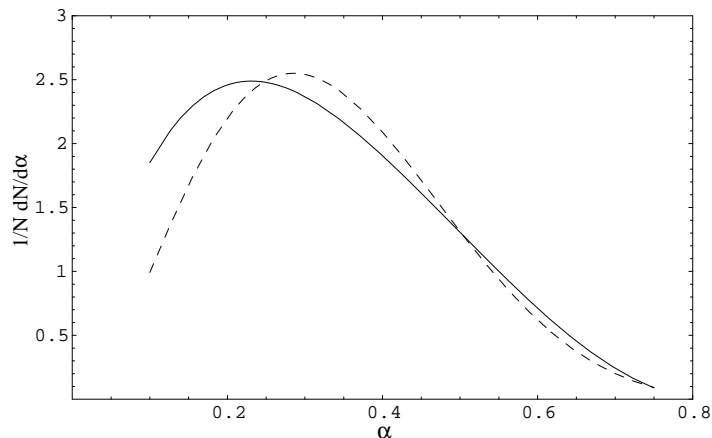


Figure 3: The normalized longitudinal momentum fraction distribution of the jet (solid curve). For comparison, the dependence corresponding to Eq. (20) is shown by dashes.

an observation of such processes in the early runs at LHC is feasible for jet transverse momenta of the order of 5 GeV.

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