

Balitskiĭ-Fadin-Kuraev-Lipatov Pomeron manifestation in inclusive single jet production in high-energy hadron collisions

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A discrepancy between new data on inclusive single jet production at the Fermilab Tevatron and perturbative QCD is discussed. It is shown that the discrepancy may be accounted for by the BFKL Pomeron. [S0556-2821(98)50403-1]

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The description of the inclusive production of hadron jets is one of the successes of perturbative QCD (PQCD). Quantitative agreement between data and theory has been achieved for jets produced over a wide kinematical range. In particular, data from the Collider Detector at Fermilab (CDF) and *D0* Collaborations at the Fermilab Tevatron on inclusive single jet production for $\sqrt{s}=1800$ GeV [1,2] is in agreement with PQCD for jet transverse energy ranging approximately from 15 to 400 GeV. Over this transverse energy range the cross section falls by seven orders of magnitude.

With this good quantitative description of the dependence on the parameters of the produced jets at fixed total energy of the collision, the natural next task is to examine the dependence of the production cross section on the total energy. Dimensional analysis and scaling hypothesis predetermines this energy dependence. PQCD (see, e.g., review [3]) dictates a particular mechanism of scaling violation involving a hadronic scale, Λ_{QCD} , which yields a specific non-trivial energy dependence. Any deviation from this prediction would manifest an inadequacy of the PQCD framework for managing nonperturbative physics (soft hadronic radiation).

We will show that existing data already contain evidence for additional non-PQCD effects which are consistent with the Balitskiĭ-Fadin-Kuraev-Lipatov (BFKL) Pomeron [4] framework. We can begin with a more general observation that a potential mechanism for such non-PQCD scaling violation is implied by resummation of the leading energy logarithms of QCD (for a recent review, see Ref. [5]). In this case, the dependence of the cross section on the QCD running coupling constant α_s differs from the simple power dependence of PQCD. This change away from a power dependence yields an altered energy dependence of the cross section.

There are preliminary data from CDF [6] on the cross section for inclusive single jet production at $\sqrt{s}=1800$ GeV

and $\sqrt{s}=630$ GeV. In particular, we examine the ratio of cross sections scaled by the jet transverse energy E_\perp , taken at the same values of $x_\perp = 2E_\perp / \sqrt{s}$ and the rapidity of the jet η , but at different total energies [7,8]:

$$R(x_\perp, \eta) = \frac{(E_\perp^4 Ed^3\sigma/d^3k)|_{\sqrt{s_1}, x_\perp, \eta}}{(E_\perp^4 Ed^3\sigma/d^3k)|_{\sqrt{s_2}, x_\perp, \eta}}. \quad (1)$$

Without scaling violation the scaled ratio is unity, regardless of the dynamics. On the PQCD leading order predictions for the ratio see [9].

Comparison of the next-to-leading order (NLO) PQCD prediction [10] with the data [6] shows a noticeable discrepancy at small x_\perp (Fig. 1). This problem has been already seen in previous CDF data [8]. It is unlikely that complications connected with jet algorithms and various uncertainties of PQCD approach [11] may account for the discrepancy.

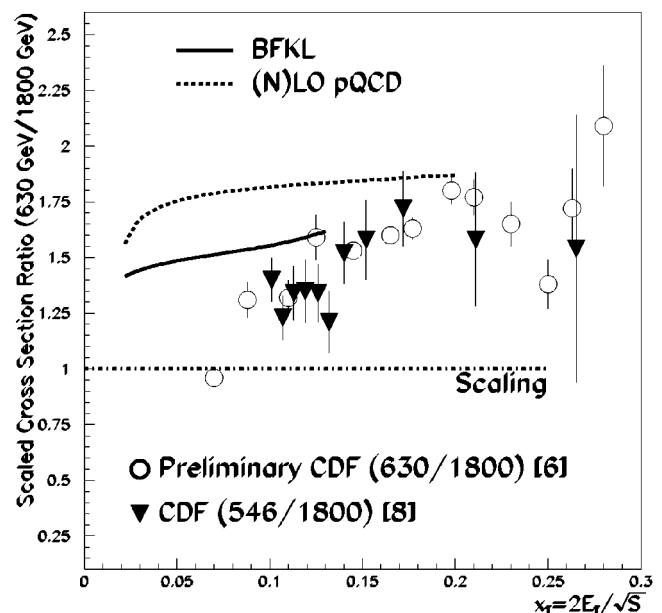


FIG. 1. The x_\perp -dependence of the scaled cross-section ratio (630 GeV)/(1800 GeV). Only statistical errors are shown.

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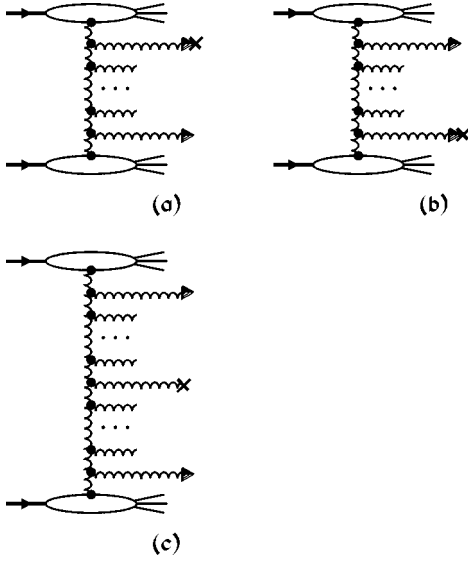


FIG. 2. Diagrams for inclusive single jet production. \times denotes tagged jet, \triangleright marks most forward (backward) jet.

In this paper we show that the discrepancy in the scaled ratio at small x_{\perp} is accounted for by the BFKL Pomeron [4].

In our approach, the dependence of the inclusive jet cross section on α_S may be expressed as a multiple integration and summation over conformal dimensions and conformal spins of the BFKL Pomerons [12]. Each term of this ‘‘sum’’ depends on α_S as $x_{\perp}^{-\alpha_S \beta(\{\nu, n\})}$, where $\{\nu, n\}$ is a set of conformal dimensions, ν_i , and conformal spins, n_i . To get the weights with which these contributions enter the ‘‘sum’’ as well as the $\beta(\{\nu, n\})$, we use the effective Feynman-like rules defined in Ref. [12]. Substituting the running α_S into the BFKL formulas for the inclusive cross sections, one obtains a new x_{\perp} -dependence for the scaled ratio R . The calculation of the scaled ratio along this line gives the result presented in Fig. 1.

As was pointed out in Ref. [13], it is important to keep track of the most forward and backward jets of the events for the BFKL kinematics. In particular [12], the inclusive single jet cross section is a sum of three terms. The first (second) term comes from the processes with untagged most forward (most backward) jet and corresponds to the diagram of Fig. 2(a) [Fig. 2(b)]. The third term [Fig. 2(c)] corresponds to the processes with both most forward and most backward jets untagged. The analytic expressions corresponding to the diagrams of Fig. 2 read as follows:

$$\frac{\alpha_S N_c}{2\pi^2} \vartheta(x^+ - x^-) F_B(e^{-(x^- - x_B^-)}, \mu_H) \times \int_{-\infty}^{\infty} d\nu \tilde{W}_A(x_A^+ - x^+, x_A^+ - x^-, \nu, \mu_H), \quad (2)$$

$$\frac{\alpha_S N_c}{2\pi^2} \vartheta(x^+ - x^-) F_A(e^{-(x_A^+ - x^+)}, \mu_H) \times \int_{-\infty}^{\infty} d\nu \tilde{W}_B(x^+ - x_B^-, x^+ - x_B^-, \nu, \mu_H), \quad (3)$$

$$\frac{\alpha_S N_c}{2\pi^2} \int \int_{-\infty}^{\infty} d\nu_1 d\nu_2 W_A(x_A^+ - x^+, x_A^+ - x^-, \nu_1, \mu_H) \times R_{\varphi}(0, -\nu_1 - \nu_2) W_B(x^+ - x_B^-, x^+ - x_B^-, \nu_2, \mu_H), \quad (4)$$

where x^{\pm} are connected with the light-cone components of the produced jet momentum: $x^{\pm} = \pm \log(k^{\pm}/\mu_R)$; $x_{A,B}^{\pm}$ are connected with the light-cone components of the momenta for the colliding hadrons A and B: $x_{A,B}^{\pm} = \pm \log(k_{A,B}^{\pm}/\mu_R)$; $F_{A,B}$ are the effective parton densities of the colliding hadrons A and B; μ_H is the normalization point for both the parton densities and the running coupling constant α_S (subscript H stands for ‘‘hard’’); μ_R is the normalization point for the energy logarithms which are resummed by the BFKL Pomeron (subscript R stands for ‘‘Regge’’). Integrations over conformal dimensions ν_i [4] are to account for different patterns of gluon radiation occupying the rapidity intervals spanned by the tagged jet and the most forward (backward) jet. Note that the radiation involves infinite number of radiated gluons. The following analytic expressions for $W_{A,B}$, $\tilde{W}_{A,B}$ may be obtained from the diagram technique of Ref. [12] by integration over parameters of the untagged most forward or (and) untagged most backward jet(s):

$$W_{A,B}(x_1, x_2, \nu, \mu_H) = \tilde{W}_{A,B}(x_1, x_2, \nu, \mu_H) \frac{\Gamma(\frac{1}{2} - i\nu)}{\Gamma(\frac{1}{2} + i\nu)}, \quad (5)$$

$$\begin{aligned} \tilde{W}_{A,B}(x_1, x_2, \nu, \mu_H) &= \frac{\alpha_S N_c}{2\pi^2} i \frac{e^{(-ix_1\{-\nu+i[1+\omega(0,\nu)]/2\})}}{\nu+i[1+\omega(0,\nu)]/2} \\ &\times \left[e^{(-ix_2\{-\nu+i[1+\omega(0,\nu)]/2\})} \right. \\ &\times M_{A,B}[1+\omega(0,\nu), \mu_H, m(x_1, x_2)] \\ &\left. - M_{A,B}\left(\frac{1+\omega(0,\nu)}{2} + i\nu, \mu_H, m(x_1, x_2)\right) \right], \quad (6) \end{aligned}$$

where $\omega(0, \nu)$ is the Lipatov eigenvalue $\omega(n, \nu) = (2\alpha_S N_c/\pi)\{\psi(1) - \text{Re} \psi[(|n|+1)/2 + i\nu]\}$ [4] taken at zero value of the conformal spin n (ψ here is the logarithmic derivative of the Euler Gamma-function);

$$M_{A,B}(\lambda, \mu_H, \gamma) = \int_{e^{-\gamma}}^1 dz z^{\lambda-1} F_{A,B}(z, \mu_H) \quad (7)$$

are incomplete moments of the parton densities; and $m(x_1, x_2) = \min\{x_1, x_2\}$.

One more object, $R_{\varphi}(0, \nu)$, entering Eq. (4) is an element of the diagram technique of [12],

$$R_{\varphi}(n, \nu) = \frac{i^{|n|} e^{i|n|\varphi}}{(|n|/2) - i(\nu + i\varepsilon)} \frac{\Gamma[(|n|/2) + 1 - i\nu]}{\Gamma[|n|/2 + 1 + i\nu]}, \quad (8)$$

taken at $n=0$.

The sum of Eqs. (2)–(4) gives the cross section $(s/\pi^4)d^2\sigma/dx^+dx^-$ which is easily connected to the invariant cross sections entering Eq. (1).

The first calculational task is to compute the incomplete momenta of the parton densities entering Eq. (6). For the effective parton densities, $F=g+4/9(q+\bar{q})$ [14], we use the following parametrization:

$$F(z, \mu_H) = a(\mu_H) z^{-b(\mu_H)} (1-z)^{c(\mu_H)}, \quad (9)$$

where we omitted the subscripts A and B , since, in our case, both colliding hadrons (p and \bar{p}) have the same effective parton density. With this parametrization the incomplete moments are analytically calculable as

$$M(\lambda, \mu_H, \gamma) = a(\mu_H) \left(B(\lambda - b(\mu_H), c(\mu_H) + 1) - \frac{e^{\gamma b(\mu_H) - \lambda}}{\lambda - b(\mu_H)} \Phi(\lambda - b(\mu_H), -c(\mu_H), \lambda - b(\mu_H) + 1, e^{-\gamma}) \right), \quad (10)$$

where B is the Euler Beta-function and Φ is the hypergeometric function. Parameters $a(\mu_H), b(\mu_H), c(\mu_H)$ were found by fitting the CTEQ4L [15] parton densities.

The next step is to perform the integration of Eqs. (2)–(4) over conformal dimensions. This has been done numerically. The relative error of this numerical computation does not exceed 10% for the results presented, though, for most values of the parameters, it is less than few percents.

The only free parameters which we have is the normalization points μ_H and μ_R . They were taken to be proportional to the transverse energy: $\mu_{H,R} = \xi_{H,R} E_\perp$, as in PQCD [10]. Our main result is that with fixed values of $\xi_H=1$ and $\xi_R=0.5$, one has a good description of data [6] simultaneously for jet cross section at $\sqrt{s}=630$ and 1800 GeV (Fig. 3), and for the scaled ratio $(630 \text{ GeV})/(1800 \text{ GeV})$ (Fig. 1). As is seen from Fig. 1, PQCD fails to do this for the scaled ratio at low x_\perp . (Note, that CDF data [6] are still preliminary and require final analysis.)

Qualitatively, the above result is a consequence of the energy dependence of the gluon radiation (see Fig. 2) which is neglected in the finite order PQCD calculations: The radiation is more significant for higher energies where more ‘‘diffusion’’ of the transverse momenta is allowed [4,5].

Here we should comment on the applicability of the running coupling constant in conjunction with the BFKL Pomeron and on the normalization point dependence of the BFKL predictions.

There are no ultraviolet divergencies in our resummed leading energy logarithms and, hence, there is no running coupling constant and ultraviolet renormalization group invariance. Presumably, both should arise via resummation of subleading energy logarithms along with appropriate new factorization theorems. Connection between ‘‘hard’’ and ‘‘Regge’’ normalization points should also be established via analysis of subleading energy logarithms. Our work assumes the eventual success of such a program. For recent works on resummation of subleading energy logarithms see Ref. [16]

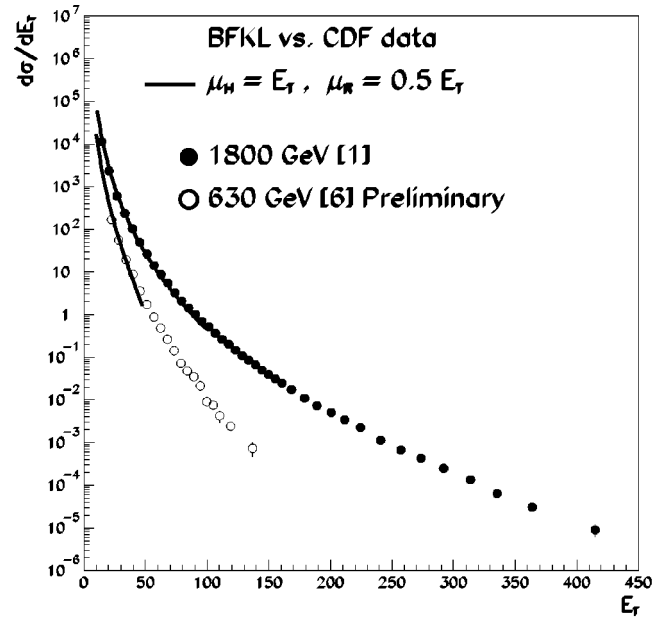


FIG. 3. Inclusive jet production cross sections for Tevatron energies.

and references therein. In addition, attempts to find the relevant versions of factorization theorems have appeared [17–19]. For more examples of the BFKL Pomeron in phenomenological applications encountering the problem of normalization point dependence see Refs. [13,19,20]. We note also that significant normalization point dependence and a deviation from the data [21] was observed in the calculation of forward jet production in NLO PQCD [22] for deep inelastic scattering at the HERA energies, while the BFKL calculation [23] agrees with the data [21].

Our conclusion is that the BFKL prediction has more capacity than the prediction of PQCD in fitting the data on inclusive single jet production for both $\sqrt{s}=630$ GeV and $\sqrt{s}=1800$ GeV. In both approaches, BFKL and PQCD, there is a fitting parameter μ_H ; BFKL approach has additional parameter, μ_R , which was also used to fit the data. In the case of PQCD the good fit for $\sqrt{s}=1800$ GeV is achieved at the expense of the fit for the lower energy at small x_\perp . However, the BFKL prediction is able to accommodate both energies.

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