Azimuthal Dependence of Forward-Jet Production in DIS in the High-Energy Limit

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Abstract:

As a signal for the BFKL Pomeron in small-x deep inelastic $ep$ scattering, we calculate the azimuthal dependence of the inclusive cross section of forward jets relative to the outgoing electron. For not very large differences in rapidity between the current jet and the forward jet the cross section peaks at $\pi/2$. For increasing rapidity BFKL dynamics predicts a decorrelation in the azimuthal dependence between the electron and the forward jet.

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1 Introduction

Forward jet production in Deep Inelastic Scattering is of great value for testing QCD at large energies where the probability of having multi-gluon emission is strongly increased and the need arises of resumming diagrams to all orders in $\alpha_s$. The resummation can only be performed to leading log accuracy and is technically done by solving the BFKL-equation [1]. The virtuality $Q^2$ and the jet transverse momentum $k_t^2$, both of them being larger than $5 GeV^2$, provide enough hardness in the process to justify the use of perturbative QCD excluding soft contributions. At fixed $Q^2$ the maximum parton subenergy $\hat{s}$ is achieved by taking the lowest possible $x_{bj}$-values and the largest parton or jet energy $xp$ with $p$ being the proton momentum.

In hadron-hadron collisions, kinematic configurations of this kind may be realized by selecting two-jet events at large rapidity intervals $\eta \simeq \ln(\hat{s}/Q^2)$, the so-called Mueller-Navelet jets [2]. The main contribution at the parton level comes from gluon exchange in the crossed channel. Then the BFKL theory dresses the exchanged gluon with the multiple emission of gluons, that uniformly fill the rapidity interval between the two extreme jets, and resums the leading powers in $\eta$, including both real and virtual corrections. The signature of the BFKL dynamics in this context is an exponential rise in $\eta$ of the inclusive two-jet cross section. However, since $\hat{s} = x_1x_2s$, with $x_1$ and $x_2$ the momentum fractions of the incoming partons, the rapidity interval may be increased by keeping $x_1$ and $x_2$ fixed and by raising $s$, which may be fulfilled only at a variable-energy collider, or by increasing $x_1$ and $x_2$ at a fixed-energy collider. The second option, though, introduces a damping in the cross section, due to the falling parton luminosity as $x \to 1$, which
conceals the dynamic rise of the partonic cross section, induced by the BFKL ladder [3].

In the case of ep colliders [4], [5], [6] the evolution parameter $\eta$ of the BFKL ladder is $\eta = \ln(\hat{s}/Q^2) = \ln(x/x_{bj})$, with $x$ the momentum fraction of the parton initiating the hard scattering. Since a fixed-energy ep collider is a variable-energy collider in the photon-proton frame, it is possible to increase $\eta$ by decreasing $x_{bj}$ while keeping $x$ fixed, thus avoiding the kinematical limitations noticed in the hadron colliders. The lowest-order process featuring gluon exchange in the crossed channel is three-parton production at $O(\alpha^2 \alpha_s^2)$. The exchanged gluon is then dressed with a BFKL ladder (Fig.1). In sect. 2 we reproduce the cross section for inclusive forward-jet production. As detailed in the next paragraph, we work out the cross section in the lab frame, however our result coincides with the one in the photon-proton frame [5]† since the forward-jet production rate is invariant under boosts between the two frames in the high-energy limit.

Another signature of the BFKL dynamics in Mueller-Navelet jet production in hadron-hadron collisions is the correlation in the azimuthal angle $\phi$ between the two tagging jets. The correlation has a maximum at $\phi = \pi$ and is expected to decrease as $\eta$ grows, because of the multiple gluon emission between the tagging jets induced by the BFKL ladder [3, 8]. The analogous process in DIS is the correlation in $\phi$ between the electron and the forward jet. Transverse-momentum and azimuthal-angle distributions in DIS have been considered previously in the parton model [3], [10] and in perturbative QCD at $O(\alpha_s)$ [11], [12]. In the parton model process $e + q \to e + q$, i.e. at $x = x_{bj}$, the simple two-body kinematics constrain the jet and the electron to be produced back-to-back,

†The relevant theoretical formulae may also be obtained by taking the massless limit of heavy-flavor production in DIS [7].
thus at the parton level it is $\phi = \pi$, with a smearing induced by the hadronic corrections \cite{8}, \cite{10}. For the $O(\alpha_s)$ corrections to the parton-model jet production, $e + g \rightarrow e + q\bar{q}$ or $e + g \rightarrow e + q + g$, the distribution in $\Phi$ in the photon-proton frame has been considered, with $\Phi$ the azimuthal angle between the jet and the incoming lepton. It has the functional form,

$$
\frac{d\sigma}{d\Phi} \sim A + B \cos \Phi + C \cos(2\Phi).
$$

(1)

The $\Phi$-dependence arises from interference of the photon helicity states $0, +, -$ (longitudinal and transverse polarization). The term $B \cos \Phi$ results from the mixing of longitudinal and transverse polarization and $C \cos(2\Phi)$ occurs when the helicities $\pm$ interfere \cite{11}. In the simple photon-gluon diagram $C$ turns out to be positive \cite{11}, \cite{12}.

In this paper, we consider the $O(\alpha_s^2)$ corrections to forward-jet production, $e + g \rightarrow e + q\bar{q} + g$ (Fig.1) or $e + q \rightarrow q\bar{q} + q$ as a function of the evolution parameter, $\eta = \ln(x/x_{bj})$, 

Figure 1: Forward-jet production in DIS.
of the BFKL ladder, and compute the distribution in $\phi$ in the lab frame, with $\phi$ the azimuthal angle between the electron and the forward jet. We find a correlation in $\phi$ of the form

$$\frac{d\sigma}{d\phi} \sim A' + C' \cos(2\phi),$$

with $C'$ being negative, now. Compared to the parton-model analysis of DIS with $O(\alpha_s)$ corrections one finds a maximum at $\phi = \pi/2$ rather than a minimum. The effect of the BFKL ladder is then as usual to flatten the $\phi$ distribution. The $\cos(\phi)$-term completely drops out due to antisymmetry in the polar angle distribution ($\theta \rightarrow \pi - \theta$) of the quark-antiquark pair at the top of the diagram. At $O(\alpha_s^2)$ these quarks are predominantly produced backwards, i.e they are not tagged as forward jet. After integration over their total phase space the antisymmetric contributions cancel out.

2 Forward-jet production

We work in the electron-proton lab frame and consider the lowest-order process featuring gluon exchange in the crossed channel. In order to achieve that, we must couple the virtual photon to the off-shell gluon via a quark box. Since we want to examine jet production near the proton fragmentation region, we couple then the off-shell gluon to a parton coming from the proton. The lowest-order diagrams with these features are the three-parton production one at $O(\alpha^2 \alpha_s^2)$, pictured in fig.11, and the one obtained from this by crossing the quark legs. All the other diagrams at the same order in $\alpha_s$ are subleading in the high-energy limit.

We label $p_e$ and $p_{e'}$ the momenta of the incoming and outgoing electrons, $p$ and $k$
the momenta of the incoming parton and the parton in the forward direction, \( p_q \) and \( p_{\bar{q}} \) the momenta of the quarks produced in the photon-gluon fusion, respectively. The jet-production rate may be then written in the high-energy limit as

\[
\frac{d\sigma}{d\hat{\sigma}} = \int dx f_{\text{eff}}(x, \mu) d\hat{\sigma},
\]

with the effective parton density, \( f_{\text{eff}}(x, \mu) \), given by \(^{13}\)

\[
f_{\text{eff}}(x, \mu) = G(x, \mu) + \frac{4}{9} \sum_f \left[ Q_f(x, \mu) + \bar{Q}_f(x, \mu) \right],
\]

with the sum over the quark flavors of the incoming parton. \( \mu \) denotes the factorization scale. The partonic cross section, \( d\hat{\sigma} \), is

\[
d\hat{\sigma} = \frac{(2\pi)^4 \delta^4(p_{e'} + p_q + p_{\bar{q}} + k - p_e - p)}{2xs} d\Pi_k d\Pi_{p_{e'}} d\Pi_{p_q} d\Pi_{p_{\bar{q}}} |\mathcal{M}|^2,
\]

with \( s \) the squared electron-proton center-of-mass energy, and with the phase space given in terms of rapidity \( \eta \) and transverse momentum \( k_\perp \) by,

\[
d\Pi = \frac{d\eta d^2k_\perp}{16\pi^3}.
\]

We introduce light-cone, or Sudakov, variables, with light-cone directions taken to be the ones of the proton and the electron, which we call + and − respectively. In the high-energy limit, for which the momenta \( p_{e'}^+, p_q^+ \) and \( p_{\bar{q}}^+ \) are negligible with respect to the one of the forward jet, \( k^+ \), the momentum conservation, \( \delta(\sum p^+) \), fixes \( x \) in eq.(3),

\[
x = \frac{k_\perp e^{\eta_k}}{2E_P},
\]

with \( E_P \) the proton energy and \( \eta_k \) the jet rapidity. With obvious modifications of the DIS standard conventions we may write the square of the invariant amplitude, integrated
over the phase space of the outgoing quarks produced in the photon-gluon fusion, as

$$\int d\Pi p d\Pi p |\mathcal{M}|^2 \delta(\sum p^-) \delta^2(\sum p_\perp) = \frac{e^4}{Q^4} L_{\mu\nu} W_{\mu\nu} \sum_\mathcal{q} e_q^2,$$

with $L_{\mu\nu}$ and $W_{\mu\nu}$ the tensors which describe the leptonic and hadronic structure respectively, and include the average (sum) over the initial-state (final-state) spin and color degrees of freedom. On the right-hand side we have singled the sum over the flavors of the final-state quarks out of $W_{\mu\nu}$. The leptonic tensor has the form,

$$L_{\mu\nu} = 2(p^\mu p'^\nu + p'^\mu p^\nu - g_{\mu\nu} p_e \cdot p_{e'}),$$

where the contribution of the Z-boson exchange has been neglected. The hadronic tensor $W_{\mu\nu}$ depends on the momenta of the proton $P$, the virtual-photon $q$ and the jet $k$. Thus for an unpolarized cross section it may be expressed in terms of four structure functions [10]. We have found convenient, though, just to determine the contractions $W_{\mu\nu} g_{\mu\nu}$ and $W_{\mu\nu} p^\mu p'^\nu$, with $W_{\mu\nu} p^\mu p'^\nu = W_{\mu\nu} p^\mu p'^\nu$ because of gauge invariance. For the gluon coupling to the gluon exchanged in the crossed channel (fig.1), we use the high-energy limit, i.e. we retain only the helicity-conserving term [1]. This entails that the component $v^-$ of the gluon exchanged in the crossed channel is neglected with respect to $p_q^-$ and $p_{\bar{q}}^-$. The calculation is long and tedious, and follows the lines of the ones performed in ref. [5], [6] and [7]. We obtain the forward-jet production cross section,

$$\frac{d\sigma}{d\eta_e dk^2_\perp d\eta_e dq^2_\perp d\phi} = x f_{eff}(x, \mu) \frac{d\hat{\sigma}}{d\eta_e dk^2_\perp d\eta_e dq^2_\perp d\phi},$$

with $\eta_e$ the electron rapidity, $q_\perp = -p_{e\perp}$, $\phi$ the azimuthal angle between the jet and the electron, $x$ fixed by eq.(7), and $f_{eff}$ taken from eq.(4). Introducing the electron energy loss $y$, satisfying the relation $q^2_\perp = (1-y)Q^2$, eq.(10) reads:

$$\frac{d\sigma}{dx dk^2_\perp dy dQ^2 d\phi} = f_{eff}(x, \mu) \frac{d\hat{\sigma}}{d\eta_e dk^2_\perp d\eta_e dq^2_\perp d\phi}. $$

(11)
The partonic cross section in eq. (11) is
\[
\frac{d\hat{\sigma}}{dk_2^\perp dy dQ^2 d\phi} = \frac{N_c \alpha_s^2 \alpha_s^2}{2\pi^2} \sum_q e_q^2 \frac{1}{y(Q^2 k_2^\perp)^2} F(k_2^\perp, Q^2, \phi, y), \tag{12}
\]
with \( N_c = 3 \) the number of colors. The impact factor for the final-state quarks, \( F(k_2^\perp, Q^2, \phi, y) \), is
\[
F(k_2^\perp, Q^2, \phi, y) = k_2^\perp Q^2 \int_0^1 d\alpha \int_0^1 dz \left[ \frac{1}{2} - \alpha(1-\alpha) - z(1-z) + 2\alpha(1-\alpha)z(1-z) \right] y^2 \\
+ \left[ 1 - 2z(1-z) - 2\alpha(1-\alpha) + 12\alpha(1-\alpha)z(1-z) \right] (1-y) \\
- 4\alpha(1-\alpha)z(1-z)(1-y) \cos(2\phi) \right). \tag{13}
\]

The BFKL ladder we wish to insert on the gluon exchanged in the crossed channel is [1],
\[
f(k_\perp, v_\perp, \phi, \bar{\phi}, \eta) = \frac{1}{(2\pi)^2} \frac{1}{(k_\perp^2 v_\perp^2)^{1/2}} \sum_{\nu=-\infty}^{\infty} e^{\text{in}(\tilde{\phi}-\phi)} \int_{-\infty}^{\infty} d\nu \ e^{\omega(\nu,n)\eta} e^{i\nu \ln(v_\perp^2/k_\perp^2)}, \tag{14}
\]
with \( v_\perp \) the transverse momentum of the off-shell gluon coupling to the quark box (Fig. 1), \( \tilde{\phi} \) the azimuthal angle between \( v_\perp \) and \( p_{e\perp} \), \( \eta = \ln(x/x_{bj}) \) and
\[
\omega(\nu, n) = -2\alpha_s N_c \frac{\alpha_s}{\pi} \text{Re} \left[ \psi \left( \left| n \right| + \frac{1}{2} + i\nu \right) - \psi(1) \right], \tag{15}
\]
with \( \psi \) the logarithmic derivative of the \( \Gamma \) function. The corrections of the BFKL ladder to the partonic cross section [12] are then given by the formula,
\[
\frac{d\hat{\sigma}}{dk_2^\perp dy dQ^2 d\phi} = \frac{N_c \alpha_s^2 \alpha_s^2}{\pi^2} \sum_q e_q^2 \frac{1}{y(Q^2 k_2^\perp)^2} \int \frac{d^2v_\perp}{v_\perp^2} f(v_\perp^2, k_\perp^2, \phi, \bar{\phi}, \eta) F(v_\perp^2, Q^2, \phi, y). \tag{16}
\]

In the \( \alpha_s \eta \to 0 \) limit, the BFKL ladder reduces to the gluon propagator exchanged in the crossed channel,
\[
\lim_{\alpha_s \eta \to 0} f(k_\perp, v_\perp, \bar{\phi}, \eta) = \frac{1}{2} \delta^2(k_\perp - v_\perp), \tag{17}
\]
and accordingly eq. (16) reduces to eq. (12). We term eq. (17) the Born approximation to the BFKL ladder. We substitute eq. (13) and (14) into eq. (16), make the measure \( d^2 v_\perp \) explicit as \( d^2 v_\perp = d\tilde{\phi} dv_\perp^2 / 2 \), and perform the integral over \( \tilde{\phi} \), which singles out the components \( n = 0, 2 \) of the eigenvalue (15) in the BFKL ladder (14). We can perform the integral over \( \alpha, z \) and \( v_\perp^2 \), by using the formula

\[
\frac{\int_0^1 d\alpha [\alpha (1-\alpha)]^{t_\alpha} \int_0^1 dz [z (1-z)]^{t_z} \int_0^\infty dv_\perp^2 \frac{(v_\perp^2)^{-1/2 + i\nu}}{\alpha (1-\alpha) Q^2 + z (1-z) v_\perp^2}}{(Q^2)^{-1/2 + i\nu} \frac{\pi}{\cosh(\pi \nu)} B \left( \frac{1}{2} + t_\alpha + i\nu, \frac{1}{2} + t_\alpha + i\nu \right) B \left( \frac{1}{2} + t_z - i\nu, \frac{1}{2} + t_z - i\nu \right)} =
\]

(18)

with \( t_\alpha = 0, 1 \) and \( t_z = 0, 1 \), and with \( B \) the Euler beta function. Performing then a bit of algebra, the partonic cross section eq. (16) becomes

\[
\frac{d\hat{\sigma}}{d^2 k_\perp^2 dy dQ^2 d\phi} =
\]

\[
N_c \alpha_s^2 \frac{1}{8\pi} \sum_q e_q^2 Q^2 (Q^2 k_\perp^2)^{3/2} y \int_0^\infty d\nu \cos \left( \nu \ln \frac{Q^2}{k_\perp^2} \right) \frac{\sinh(\pi \nu)}{\cosh^2(\pi \nu) \nu (1 + \nu^2)} \frac{1}{\nu (1 + \nu^2)} \left( e^{\omega(\nu,0)\eta} \left[ 3\nu^2 + \frac{11}{4} \right] (1-y) + \left( \nu^2 + \frac{9}{4} \right) \frac{y^2}{2} - e^{\omega(\nu,2)\eta} \cos(2\phi) \left( \nu^2 + \frac{1}{4} \right) (1-y) \right).
\]

(19)

Substituting eq. (13) into eq. (11), we obtain the forward-jet production rate, with the higher-order corrections of the BFKL ladder.

Integrating over \( \phi \), the forward-jet production rate reduces to the one given in ref. [3]. In that context it was derived in the photon-proton frame, however in the high-energy limit the forward-jet production rate is invariant under boosts between the electron-proton and the photon-proton frames, i.e. in the notation of eq. (11) \( \Phi = \phi \).

Assuming \( k_\perp^2 \) to be of the order \( Q^2 \) and \( \eta \) to be very large the integration in eq. (19) can be performed approximately by means of the saddle point method. After expanding
\(\omega(\nu, 0)\) around \(\nu = 0\) one finds

\[
\frac{d\hat{\sigma}}{dk_\perp^2} \sim \frac{1}{(k_\perp^2)^{3/2}} \exp \left( 4 \ln 2 \frac{N_c \alpha_s}{\pi} \eta \right) \exp \left( -\frac{\ln^2(k_\perp^2/Q^2)}{4B\eta} \right) \quad \text{with} \quad B = 14\zeta(3) \frac{N_c \alpha_s}{\pi},
\]

where \(\zeta\) denotes the Riemann-\(\zeta\) function. Eq. (20) is regular as \(k_\perp \to 0\), however the approximation becomes invalid then. The limit \(k_\perp \to 0\) is of interest for the total, integrated cross section which is dominated by small \(k_\perp\). Assuming \(\log(Q^2/k_\perp^2) \to \infty\) the saddle point of eq. (19) is shifted towards the pole of \(\omega(\nu, 0)\) at \(\nu = -i/2\). This pole corresponds to the collinear emission of gluons, i.e. we pick up all collinear singularities of the gluon ladder. The gluon structure function according to its definition includes all collinear singularities, so that after integration over the jet momentum eq. (19) coincides at least in a formal sense with the inclusive small \(x_{bj}\) cross section.

Returning to the distribution in \(\phi\) of eq. (19) we recognize that it is periodic in \(\pi\) and has a maximum at \(\phi = \pi/2\). The electron-jet correlation in \(\phi\), though, quickly dies out as \(\eta = \ln(x/x_{bj})\) increases. Indeed, the main contribution to the integral over \(\nu\) in eq. (19) comes from the \(\nu \simeq 0\) region, and from eq. (13) we see that \(\omega(\nu = 0, n = 0) = 4 \ln 2 (\alpha_s N_c/\pi)\) while \(\omega(\nu = 0, n = 2) = 4(\ln 2 - 1)(\alpha_s N_c/\pi)\), thus the uncorrelated term in eq. (19) is enhanced while the correlated one is upset as \(\eta\) increases.

For a small \(\eta = \ln(x/x_{bj})\) the zeroth order parton configuration is relevant, and the usual correlation at \(\phi = \pi\), between the electron and the current-jet is obtained. As \(\eta\) grows the jet production is increasingly dominated by diagrams with two- and later on with three-final state partons and with gluon exchange in the crossed channel. The three parton final state has the functional form of eqs. (2) or (12). The higher-order corrections
to them induced by the BFKL ladder, eq. (19), dampen then the correlation in $\phi$. What said above does not suffice, though, to explain the correlation at $\phi = \pi/2$. In order to see that and the transition from the correlation at $\phi = \pi$ to the one at $\phi = \pi/2$, it is necessary to compare the asymptotic calculation of eq. (12) with the exact calculation at $O(\alpha_s^2)$ [14].

Finally, we note that because of the relation between the DIS cross section and the structure functions $F_1(2)$,

$$d\sigma \over dydQ^2 = {4\pi\alpha^2 \over yQ^4} \left[ (1 - y)F_2(x_{bj}, Q^2) + x_{bj}y^2F_1(x_{bj}, Q^2) \right], \tag{21}$$

with $F_1$ and $F_2$ related to the transverse and longitudinal polarizations of the virtual-photon total cross section by,

$$F_1(x_{bj}, Q^2) = {Q^2 \over 8\pi^2\alpha x_{bj}} \sigma_T(\gamma^* P), \tag{22}$$

$$F_2(x_{bj}, Q^2) = {Q^2 \over 4\pi^2\alpha} \left( \sigma_T(\gamma^* P) + \sigma_L(\gamma^* P) \right),$$

in the high-energy limit all the information about the correlation in $\phi$ is effectively contained in the expression for the longitudinal polarization of the virtual photon although it originates from the interference of the transverse helicities. Indeed, comparing the differential of eq. (21) in $\phi$ to eq. (11), with the partonic cross section given by eq. (12) or (19), we see that $dF_2/d\phi$ depends on $\phi$, while $dF_1/d\phi$ does not.

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