

The effect of NLO conformal spins in azimuthal angle decorrelation of jet pairs

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Abstract

Azimuthal angle decorrelation in inclusive dijet cross sections is studied analytically to take into account the next-to-leading corrections to the BFKL kernel while keeping the jet vertices at leading order. The spectral representation on the basis of leading order eigenfunctions is generalized to include the dependence on conformal spins. With this procedure running coupling effects and angular dependences are both included. It is shown how the angular decorrelation for jets with a wide relative separation in rapidity largely decreases at this higher order in the resummation.

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1. Introduction

Among the many relevant questions still open in quantum chromodynamics a very interesting one is how to describe scattering amplitudes in the so-called Regge limit. In Regge asymptotic the center-of-mass energy, s , is much larger than all other Mandelstam invariants and mass scales present in the process under investigation. It is possible to perturbatively keep track of the different contributions to the amplitude if some hard scale is present so that the strong coupling remains small. It is then needed to resum logarithmically enhanced contributions of the form $(\alpha_s \ln s)^n$ to all orders. This is achieved using the leading-logarithmic (LL) Balitsky–Fadin–Kuraev–Lipatov (BFKL) evolution equation [1].

The BFKL approach predicts a power-like rise in s of total cross sections. A lot of attention has been given to the search for BFKL effects in deep inelastic scattering (DIS) due to the rapid growth of structure functions at small values of Bjorken x . However, the golden process where

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the resummation of $\ln s$ is most important is the total cross section of two photons with large and similar virtualities. In this configuration the small x resummation, which includes ordering in rapidity and not in transverse scales, should be the dominant contribution to the scattering. This is not necessarily the case in DIS where ordering in k_t is important given that parton evolution takes place between an object with large transverse size, the proton, and a small highly virtual photon.

Observables where BFKL effects should prevail then require of enough energy to build up the parton evolution, and the presence of two large and similar transverse scales. In this work an example of this kind is investigated in detail by analytic means: the inclusive hadroproduction of two jets with large and similar transverse momenta and a big relative separation in rapidity, the so-called Mueller–Navelet jets. When Y , the distance in rapidity between the most forward and backward jets, is not large a fixed order perturbative analysis should be enough to describe the cross section but when it increases a BFKL resummation of $(\alpha_s Y)^n$ terms is needed.

Mueller–Navelet jets were first proposed in Ref. [2] as a clean configuration to look for BFKL effects at hadron colliders. A typical power-like rise for the partonic cross section was predicted in agreement with the value of the asymptotic LL hard pomeron intercept. However, at hadronic level, forward and backward jets are produced in a region of fast falling of the parton distributions, reducing the rise of the cross section. A way to make small x resummation effects more explicit is to look into the azimuthal angle decorrelation of the pair of jets. The relevant subprocess is $\text{parton} + \text{parton} \rightarrow \text{jet} + \text{jet} + \text{any number of soft emissions inside the rapidity interval separating the two jets}$. BFKL enhances soft real emission as Y increases reducing in this way the amount of angular correlation originally present in the back-to-back in transverse plane Born configuration. The LL prediction for this azimuthal dependence was first investigated in Ref. [3].

The results at LL are known to overestimate the rate of decorrelation and to lie quite far from the experimental data [4] as obtained from the Tevatron and subleading higher order effects have been called for an explanation of this discrepancy. Running coupling effects and kinematic constraints have been considered in Ref. [5]. In the present work the main target is to analytically understand how to include the $\alpha_s(\alpha_s Y)^n$ next-to-leading logarithmic (NLL) corrections to the BFKL kernel [6]. The effects of this kernel were numerically investigated using an implementation [7] of the NLL iterative solution proposed in Ref. [8] (different reviews can be found in [9]). It is left for future analysis the inclusion of the next-to-leading order (NLO) jet vertex [10], the investigation of more convergent versions of the kernel [11] and a study of parton distributions effects in order to make reliable phenomenological predictions at a hadron collider. Mueller–Navelet jets should be an important test of our understanding of small x resummation to be performed at the large hadron collider at CERN.

After this brief introduction, in Section 2 the normalization for the gluon Green's function is indicated together with the formulae for the partonic cross section. Then the operator formalism suggested by Ivanov and Papa in Ref. [12] is extended to introduce angular dependences. The form of the NLL kernel for all conformal spins calculated by Kotikov and Lipatov in Ref. [13] is also discussed. Towards the end of the section a compact expression for the angular differential cross section which includes the NLL contributions is derived. In Section 3 the numerical study of the previous formulæ is discussed in detail. Finally, several conclusions are drawn and different lines for future research highlighted.

2. Calculation of the dijet partonic cross section

As indicated in the Introduction, in this analysis the object of interest is the partonic cross section parton + parton \rightarrow jet + jet + soft emission, with the two jets having transverse momenta \vec{q}_1 and \vec{q}_2 and being produced at a large relative rapidity separation Y . This can be related to the external hadrons by its approximate relation to the longitudinal momentum fractions carried by the jets, i.e., $Y \sim \ln x_1 x_2 s / \sqrt{q_1^2 q_2^2}$, and a convolution with parton distribution functions whose analysis is left for a future work. In the present framework the resummed differential partonic cross section for the particular case of gluon–gluon scattering is

$$\frac{d\hat{\sigma}}{d^2\vec{q}_1 d^2\vec{q}_2} = \frac{\pi^2 \bar{\alpha}_s^2}{2} \frac{f(\vec{q}_1, \vec{q}_2, Y)}{q_1^2 q_2^2}, \quad (1)$$

with the usual notation $\bar{\alpha}_s = \alpha_s N_c / \pi$. One can now introduce the Mellin transform of the BFKL gluon Green's function in rapidity space:

$$f(\vec{q}_1, \vec{q}_2, Y) = \int \frac{d\omega}{2\pi i} e^{\omega Y} f_\omega(\vec{q}_1, \vec{q}_2). \quad (2)$$

The normalization for the BFKL integral equation, including for simplicity only the LL terms, then reads

$$\omega f_\omega(\vec{q}_1, \vec{q}_2) = \delta^{(2)}(\vec{q}_1 - \vec{q}_2) + \bar{\alpha}_s \int \frac{d^2\vec{k}}{\pi(\vec{q}_1 - \vec{k})^2} \left(f_\omega(\vec{k}, \vec{q}_2) - \frac{q_1^2 f_\omega(\vec{q}_1, \vec{q}_2)}{\vec{k}^2 + (\vec{q}_1 - \vec{k})^2} \right). \quad (3)$$

As it is well known, including the angular dependence on the transverse plane of \vec{q}_1 and \vec{q}_2 , the LL solution to Eq. (3) can be written as

$$f_\omega(\vec{q}_1, \vec{q}_2) = \frac{1}{2\pi^2} \sum_{n=-\infty}^{\infty} \int_{-\infty}^{\infty} d\nu (q_1^2)^{-i\nu-\frac{1}{2}} (q_2^2)^{i\nu-\frac{1}{2}} \frac{e^{in(\theta_1-\theta_2)}}{\omega - \bar{\alpha}_s \chi_0(|n|, \nu)}, \quad (4)$$

where the eigenvalue of the LL kernel,

$$\chi_0(n, \nu) = 2\psi(1) - \psi\left(\frac{1}{2} + i\nu + \frac{n}{2}\right) - \psi\left(\frac{1}{2} - i\nu + \frac{n}{2}\right), \quad (5)$$

is expressed in terms of the logarithmic derivative of the Euler gamma function. As it stands the n variable corresponds to a Fourier transform in the angular sector. A more sophisticated interpretation arises if Eq. (3) is considered for non-zero momentum transfer (see, e.g., [14]). In this case, if a representation in the complex plane for the transverse momenta of the form $\vec{q} = q_x + i q_y$ is introduced, it can be shown how the BFKL equation corresponds to a Schrödinger-like equation with a holomorphically separable Hamiltonian where $-iY$ is the time variable. Both the holomorphic and antiholomorphic sectors in the Hamiltonian are invariant under spin zero Möbius transformations with eigenfunctions carrying a conformal weight of the form $\gamma = \frac{1}{2} + i\nu + \frac{n}{2}$. In the principal series of the unitary representation ν is real and $|n|$ the integer conformal spin [15].

The partonic cross section is obtained by integration over the phase space of the two emitted gluons together with some general jet vertices, i.e.

$$\hat{\sigma}(\alpha_s, Y, p_{1,2}^2) = \int d^2\vec{q}_1 \int d^2\vec{q}_2 \Phi_{\text{jet}_1}(\vec{q}_1, p_1^2) \Phi_{\text{jet}_2}(\vec{q}_2, p_2^2) \frac{d\hat{\sigma}}{d^2\vec{q}_1 d^2\vec{q}_2}. \quad (6)$$

In the perturbative expansion of these jet vertices, $\Phi_{\text{jet}_i} = \Phi_{\text{jet}_i}^{(0)} + \bar{\alpha}_s \Phi_{\text{jet}_i}^{(1)} + \dots$, only leading-order terms are kept:

$$\Phi_{\text{jet}_i}^{(0)}(\vec{q}, p_i^2) = \theta(q^2 - p_i^2), \quad (7)$$

where p_i^2 corresponds to a resolution scale for the transverse momentum of the gluon jet. In this way a full NLO accuracy is not achieved but it is possible to pin down those effects stemming from the gluon Green's function. To extend this analysis it would be needed to calculate the Mellin transform of the NLO jet vertices in Ref. [10] where the definition of a jet is much more involved than here. Therefore one can proceed and write

$$\hat{\sigma}(\alpha_s, Y, p_{1,2}^2) = \frac{\pi^2 \bar{\alpha}_s^2}{2} \int d^2 \vec{q}_1 \int d^2 \vec{q}_2 \frac{\Phi_{\text{jet}_1}^{(0)}(\vec{q}_1, p_1^2)}{q_1^2} \frac{\Phi_{\text{jet}_2}^{(0)}(\vec{q}_2, p_2^2)}{q_2^2} f(\vec{q}_1, \vec{q}_2, Y). \quad (8)$$

At this stage it is very convenient to recall the work of Ref. [12] and to introduce the following transverse momenta operator representation:

$$\hat{q}|\vec{q}_i\rangle = \vec{q}_i|\vec{q}_i\rangle \quad (9)$$

with the normalization

$$\langle \vec{q}_1 | \hat{1} | \vec{q}_2 \rangle = \delta^{(2)}(\vec{q}_1 - \vec{q}_2). \quad (10)$$

In this notation the BFKL equation simply reads

$$(\omega - \hat{K})\hat{f}_\omega = \hat{1}, \quad (11)$$

where the kernel has the expansion

$$\hat{K} = \bar{\alpha}_s \hat{K}_0 + \bar{\alpha}_s^2 \hat{K}_1 + \dots \quad (12)$$

To NLO accuracy this implies that the solution can be written as

$$\hat{f}_\omega = (\omega - \bar{\alpha}_s \hat{K}_0)^{-1} + \bar{\alpha}_s^2 (\omega - \bar{\alpha}_s \hat{K}_0)^{-1} \hat{K}_1 (\omega - \bar{\alpha}_s \hat{K}_0)^{-1} + \mathcal{O}(\bar{\alpha}_s^3). \quad (13)$$

The next step is to define a basis on which to express the cross section. To generalize the study in Ref. [12] this basis should carry not only the dependence on the modulus of the transverse momenta but also the dependence on their angle on the transverse plane:

$$\langle \vec{q} | v, n \rangle = \frac{1}{\pi \sqrt{2}} (q^2)^{i\nu - \frac{1}{2}} e^{in\theta}. \quad (14)$$

The projection $\langle n, v | \vec{q} \rangle$ would be the complex conjugate of the previous expression. This basis has been chosen such that it is orthonormal:

$$\langle n', v' | v, n \rangle = \delta(v - v') \delta_{nn'}. \quad (15)$$

The action of the NLO kernel on this basis, which was calculated in Ref. [13], contains non-diagonal terms and can be written as

$$\hat{K} | v, n \rangle = \left\{ \bar{\alpha}_s \chi_0(|n|, v) + \bar{\alpha}_s^2 \chi_1(|n|, v) + \bar{\alpha}_s^2 \frac{\beta_0}{8N_c} \left[2\chi_0(|n|, v) \left(i \frac{\partial}{\partial v} + \log \mu^2 \right) + \left(i \frac{\partial}{\partial v} \chi_0(|n|, v) \right) \right] \right\} | v, n \rangle, \quad (16)$$

where, from now on, $\bar{\alpha}_s$ stands for $\bar{\alpha}_s(\mu^2)$, the coupling evaluated at the renormalization point μ in the $\overline{\text{MS}}$ scheme. The first line of Eq. (16) corresponds to the scale invariant sector of the kernel. The function χ_1 for a general conformal spin reads

$$\begin{aligned} \chi_1(n, \gamma) = & \mathcal{S}\chi_0(n, \gamma) + \frac{3}{2}\zeta(3) - \frac{\beta_0}{8N_c}\chi_0^2(n, \gamma) \\ & + \frac{1}{4}\left[\psi''\left(\gamma + \frac{n}{2}\right) + \psi''\left(1 - \gamma + \frac{n}{2}\right) - 2\phi(n, \gamma) - 2\phi(n, 1 - \gamma)\right] \\ & - \frac{\pi^2 \cos(\pi\gamma)}{4 \sin^2(\pi\gamma)(1 - 2\gamma)} \left\{ \left[3 + \left(1 + \frac{n_f}{N_c^3}\right) \frac{2 + 3\gamma(1 - \gamma)}{(3 - 2\gamma)(1 + 2\gamma)} \right] \delta_{n0} \right. \\ & \left. - \left(1 + \frac{n_f}{N_c^3}\right) \frac{\gamma(1 - \gamma)}{2(3 - 2\gamma)(1 + 2\gamma)} \delta_{n2} \right\}. \end{aligned} \quad (17)$$

The definitions $\mathcal{S} = (4 - \pi^2 + 5\beta_0/N_c)/12$, and $\beta_0 = (11N_c - 2n_f)/3$, have been used. The function ϕ can be found in Ref. [13] and reads

$$\begin{aligned} \phi(n, \gamma) = & \sum_{k=0}^{\infty} \frac{(-1)^{(k+1)}}{k + \gamma + \frac{n}{2}} \left(\psi'(k + n + 1) - \psi'(k + 1) \right. \\ & + (-1)^{(k+1)} (\beta'(k + n + 1) + \beta'(k + 1)) \\ & \left. + \frac{\psi(k + 1) - \psi(k + n + 1)}{k + \gamma + \frac{n}{2}} \right), \end{aligned} \quad (18)$$

with

$$4\beta'(\gamma) = \psi'\left(\frac{1 + \gamma}{2}\right) - \psi'\left(\frac{\gamma}{2}\right). \quad (19)$$

In this basis terms with derivatives are associated to the running of the coupling [16], this is the case in the second line of Eq. (16). In particular, the part containing $i\frac{\partial\chi_0}{\partial v}$, which breaks the $v \rightarrow -v$ symmetry, will be shown to give a zero contribution to the cross section when $p_1^2 = p_2^2$. It will also be shown how the term with $i\chi_0\frac{\partial}{\partial v}$ mixes in a non-trivial way the Green's function with the jet vertices.

To represent the cross section in the present formalism the starting point is to project the jet vertices on the basis in Eq. (14):

$$\int d^2\vec{q} \frac{\Phi_{\text{jet}_1}^{(0)}(\vec{q}, p_1^2)}{q^2} \langle \vec{q} | v, n \rangle = \frac{1}{\sqrt{2}} \frac{1}{(\frac{1}{2} - iv)} (p_1^2)^{iv - \frac{1}{2}} \delta_{n,0} \equiv c_1(v) \delta_{n,0}. \quad (20)$$

The $c_2(v)$ projection of $\Phi_{\text{jet}_2}^{(0)}$ on $\langle n, v | \vec{q} \rangle$ is the complex conjugate of (20) with p_1^2 being replaced by p_2^2 . The corresponding inverse relations are

$$\frac{\Phi_{\text{jet}_1}^{(0)}(\vec{q}, p_1^2)}{q^2} = \sum_{n=-\infty}^{\infty} \int_{-\infty}^{\infty} dv c_1(v) \delta_{n,0} \langle n, v | \vec{q} \rangle, \quad (21)$$

$$\frac{\Phi_{\text{jet}_2}^{(0)}(\vec{q}, p_2^2)}{q^2} = \sum_{n=-\infty}^{\infty} \int_{-\infty}^{\infty} dv c_2(v) \delta_{n,0} \langle \vec{q} | v, n \rangle. \quad (22)$$

The cross section can then be rewritten as

$$\begin{aligned} \hat{\sigma}(\alpha_s, Y, p_{1,2}^2) &= \frac{\pi^2 \bar{\alpha}_s^2}{2} \sum_{n, n' = -\infty}^{\infty} \int_{-\infty}^{\infty} dv \int_{-\infty}^{\infty} dv' c_1(v) c_2(v') \delta_{n,0} \delta_{n',0} \\ &\times \int \frac{d\omega}{2\pi i} e^{\omega Y} \langle n, v | \hat{f}_\omega | v', n' \rangle. \end{aligned} \quad (23)$$

Making use of the operator representation in Eq. (13), the action of the kernel in Eq. (16) and integration by parts, $\hat{\sigma}$ can be expressed as

$$\begin{aligned} \hat{\sigma}(\alpha_s, Y, p_{1,2}^2) &= \frac{\pi^2 \bar{\alpha}_s^2}{2} \sum_{n = -\infty}^{\infty} \int_{-\infty}^{\infty} dv e^{\bar{\alpha}_s \chi_0(|n|, v) Y} c_1(v) c_2(v) \delta_{n,0} \\ &\times \left\{ 1 + \bar{\alpha}_s^2 Y \left[\chi_1(|n|, v) + \frac{\beta_0}{4N_c} \left(\log(\mu^2) \right. \right. \right. \\ &\left. \left. \left. + \frac{i}{2} \frac{\partial}{\partial v} \log\left(\frac{c_1(v)}{c_2(v)}\right) + \frac{i}{2} \frac{\partial}{\partial v} \right) \chi_0(|n|, v) \right] \right\}. \end{aligned} \quad (24)$$

For the LO jet vertices the logarithmic derivative in Eq. (24) explicitly reads

$$-i \frac{\partial}{\partial v} \log\left(\frac{c_1(v)}{c_2(v)}\right) = \log(p_1^2 p_2^2) + \frac{1}{\frac{1}{4} + v^2}. \quad (25)$$

The angular differential cross section can be calculated considering the following representation of the $c_1 c_2$ product:

$$c_1(v) c_2(v) \delta_{n,0} = \frac{1}{2\sqrt{p_1^2 p_2^2}} \frac{1}{(\frac{1}{4} + v^2)} \left(\frac{p_1^2}{p_2^2}\right)^{iv} \int_{-\pi}^{\pi} \frac{d\phi}{2\pi} e^{in\phi}, \quad (26)$$

with $\phi = \theta_1 - \theta_2 - \pi$. Therefore, in the case where the two resolution momenta are equal, $p_1^2 = p_2^2 \equiv p^2$, the angular differential cross section can be expressed as

$$\begin{aligned} \frac{d\hat{\sigma}(\alpha_s, Y, p^2)}{d\phi} &= \frac{\pi^2 \bar{\alpha}_s^2}{4p^2} \sum_{n = -\infty}^{\infty} \frac{1}{2\pi} e^{in\phi} \int_{-\infty}^{\infty} dv e^{\bar{\alpha}_s \chi_0(|n|, v) Y} \frac{1}{(\frac{1}{4} + v^2)} \\ &\times \left\{ 1 + \bar{\alpha}_s^2 Y \left[\chi_1(|n|, v) - \frac{\beta_0}{8N_c} \chi_0(|n|, v) \right. \right. \\ &\left. \left. \times \left(2 \log\left(\frac{p^2}{\mu^2}\right) + \frac{1}{(\frac{1}{4} + v^2)} \right) \right] \right\}. \end{aligned} \quad (27)$$

The term proportional to $\frac{\partial \chi_0}{\partial v}$ in Eq. (24) gives no contribution after integration as it is an odd function in v . Within NLO accuracy there is freedom to exponentiate the integrand of this result. In this work this is done both for the scale invariant and for the scale dependent terms, in close resemblance with the property of reggeization. Furthermore, renormalization group improved perturbation theory is called for to introduce the replacement

$$\bar{\alpha}_s - \bar{\alpha}_s^2 \frac{\beta_0}{4N_c} \log\left(\frac{p^2}{\mu^2}\right) \rightarrow \bar{\alpha}_s(p^2). \quad (28)$$

In this way the differential distribution can be conveniently rewritten as

$$\frac{d\hat{\sigma}(\alpha_s, Y, p^2)}{d\phi} = \frac{\pi^3 \bar{\alpha}_s^2}{2p^2} \frac{1}{2\pi} \sum_{n=-\infty}^{\infty} e^{in\phi} C_n(Y), \quad (29)$$

with

$$C_n(Y) = \int_{-\infty}^{\infty} \frac{dv}{2\pi} e^{\frac{\bar{\alpha}_s(p^2)Y(\chi_0(|n|,v) + \bar{\alpha}_s(p^2)(\chi_1(|n|,v) - \frac{\beta_0}{8N_c} \frac{\chi_0(|n|,v)}{(\frac{1}{4} + v^2)})}{(\frac{1}{4} + v^2)}}}. \quad (30)$$

Different interesting observables can be constructed with these coefficients and they will be studied in detail in the next section. Due to the large and negative, with respect to the LL terms, size of the NLL corrections it will turn out that the exponentiated form in Eq. (30) is mandatory in order to reach convergent results. This is discussed below.

3. Analysis of convergence and study of observables

To investigate the properties of the different parts of the kernel it is useful to introduce five types of coefficients. The first one is

$$C_n^{\text{LL}}(Y) = \int_{-\infty}^{\infty} \frac{dv}{2\pi} \frac{e^{\bar{\alpha}_s Y \chi_0(|n|,v)}}{(\frac{1}{4} + v^2)}, \quad (31)$$

defined at LL accuracy and previously studied in the literature (see last reference in [3]). When all the NLL terms are exponentiated as in Eq. (30) it will be referred to as C_n^{NLL} , while if the NLL pieces are not exponentiated, as in Eq. (27), it is named $C_n^{\text{Expansion}}$. If in Eq. (30) the χ_1 kernel is removed then the coefficients correspond to the case of LL plus running coupling and it is noted as $C_n^{\text{Running coupling}}$. Finally, the scale invariant contributions can be isolated by setting β_0 to zero in the exponent of Eq. (30) and this will be called $C_n^{\text{Scale invariant}}$.

The coefficient governing the energy dependence of the cross section corresponds to $n = 0$:

$$\hat{\sigma}(\alpha_s, Y, p^2) = \frac{\pi^3 \bar{\alpha}_s^2}{2p^2} C_0(Y). \quad (32)$$

As the NLL corrections are large and negative, when not exponentiated they lead to a non-convergent behavior. This is seen in Fig. 1 where the resolution scale $p = 30$ GeV has been chosen. The values $n_f = 4$ and $\Lambda_{\text{QCD}} = 0.1416$ GeV were taken in $\bar{\alpha}_s(p^2) = 4N_c/(\beta_0 \ln(p^2/\Lambda_{\text{QCD}}^2))$. In this plot it can be seen how the coefficient $C_0^{\text{Expansion}}$ generates an unphysical behaviour in contrast to the exponential rise associated to the other coefficients. Nevertheless it should be noticed that the term proportional to $\bar{\alpha}_s^2 Y$

$$\chi_0 \left(\frac{c_1^{(1)}}{c_1^{(0)}} + \frac{c_2^{(1)}}{c_2^{(0)}} \right) \quad (33)$$

would be added to χ_1 in Eq. (27) if the NLO jet vertex $c^{(0)} + \bar{\alpha}_s c^{(1)}$ was brought into the calculation. The convergence properties of Eq. (27) might well change in that case.

A familiar consequence of introducing the effects of the running of the coupling is that the LL intercept is reduced as can be seen also in Fig. 1. Meanwhile, if the scale invariant sector of

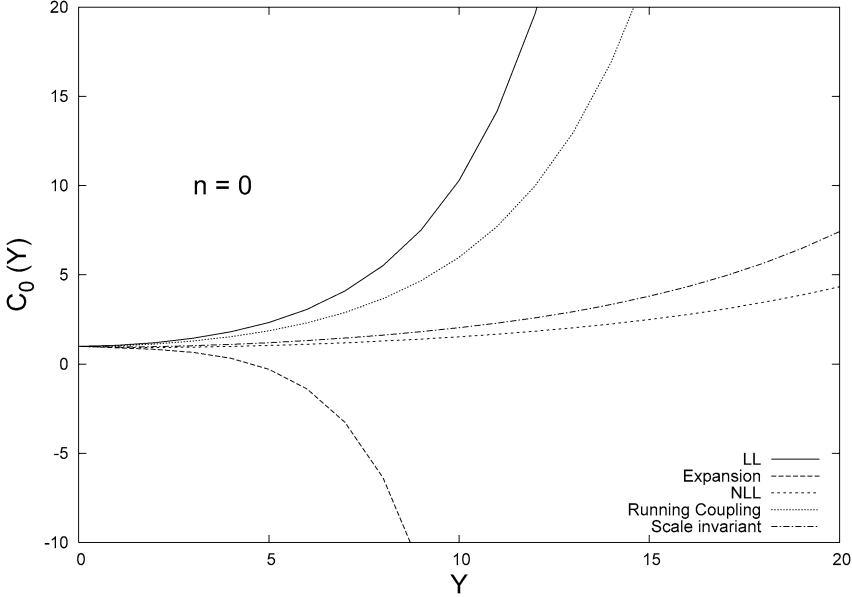


Fig. 1. Evolution in Y of the $C_0(Y)$ coefficient.

the NLL kernel, χ_1 , is also introduced then a further decrease of this rise takes place. The $n = 0$ coefficient is directly related to the normalized cross section

$$\frac{\hat{\sigma}(Y)}{\hat{\sigma}(0)} = \frac{C_0(Y)}{C_0(0)}. \quad (34)$$

The rise in rapidity of this observable is shown in Fig. 2. Clearly the NLL intercept is very much reduced with respect to the LL case. In Fig. 3 large values of Y are considered in order to approach the asymptotic regime. The LL intercept tends to the asymptotic value of $4\bar{\alpha}_s(30) \ln 2 \sim 0.37$ while the NLL result lies around one third of this number.

Independently of the NLL kernel being exponentiated or not, the remaining coefficients with $n \geq 1$ all decrease with Y . This can be seen in the plots of Fig. 4. The consequence of this decrease is that the angular correlations also diminish as the rapidity interval between the jets gets larger. This point can be studied in detail using the mean values

$$\langle \cos(m\phi) \rangle = \frac{C_m(Y)}{C_0(Y)}. \quad (35)$$

$\langle \cos(\phi) \rangle$ is calculated in Fig. 5. The most important consequence of this plot is that the NLL effects dramatically decrease the azimuthal angle decorrelation. This is already the case when only the running of the coupling is introduced but the scale invariant terms make this effect much bigger. This is encouraging from the phenomenological point of view given that the data at the Tevatron typically have lower decorrelation than predicted by LL BFKL or LL with running coupling. It is worth noting that the difference in the prediction for decorrelation between LL and NLL is mostly driven by the $n = 0$ conformal spin. This can be understood looking at the ratio

$$\frac{\langle \cos(\phi) \rangle^{\text{NLL}}}{\langle \cos(\phi) \rangle^{\text{LL}}} = \frac{C_1^{\text{NLL}}(Y) C_0^{\text{LL}}(Y)}{C_0^{\text{NLL}}(Y) C_1^{\text{LL}}(Y)}, \quad (36)$$

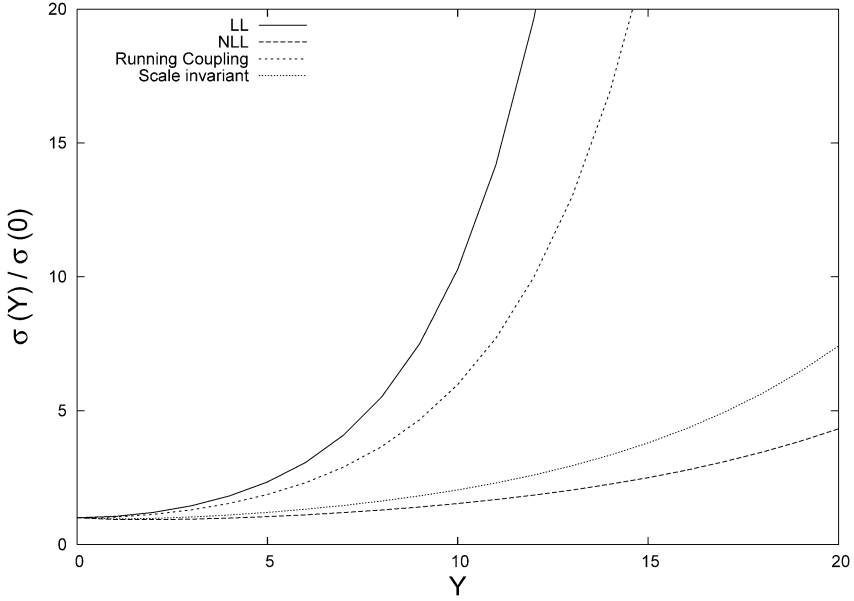


Fig. 2. Evolution of the partonic cross section with the rapidity separation of the dijets.

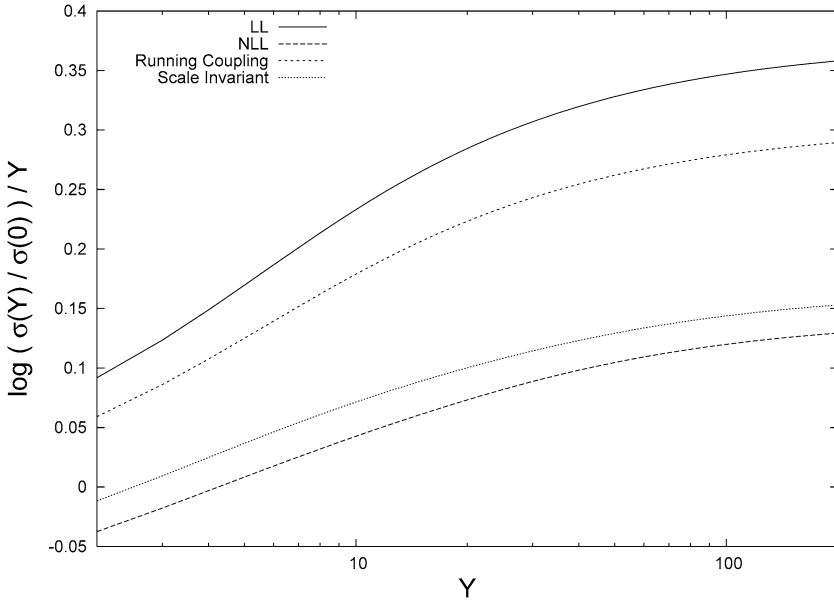


Fig. 3. Intercepts for the cross section as a function of rapidity.

and noticing that

$$1.2 > \frac{C_1^{\text{NLL}}(Y)}{C_1^{\text{LL}}(Y)} > 1. \tag{37}$$

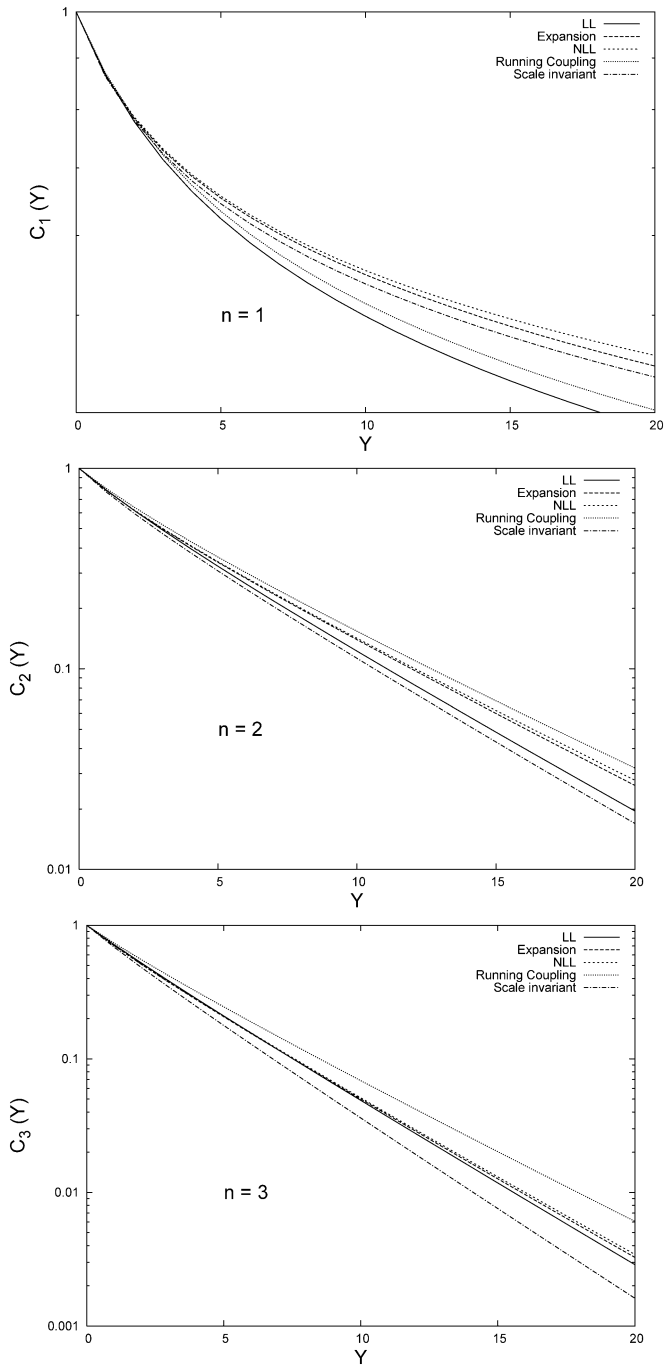


Fig. 4. Evolution in Y of the $C_n(Y)$ coefficients for $n = 1, 2, 3$.

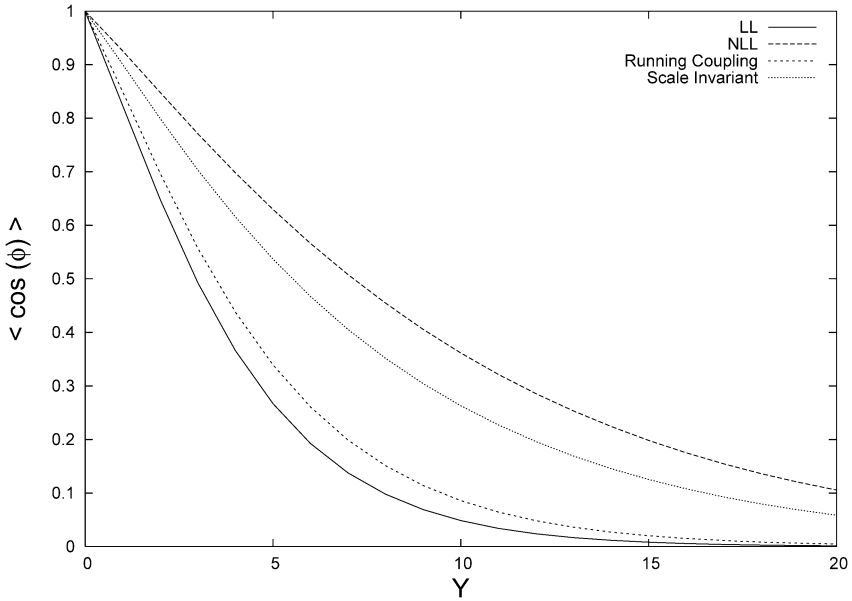


Fig. 5. Dijet azimuthal angle decorrelation as a function of their separation in rapidity.

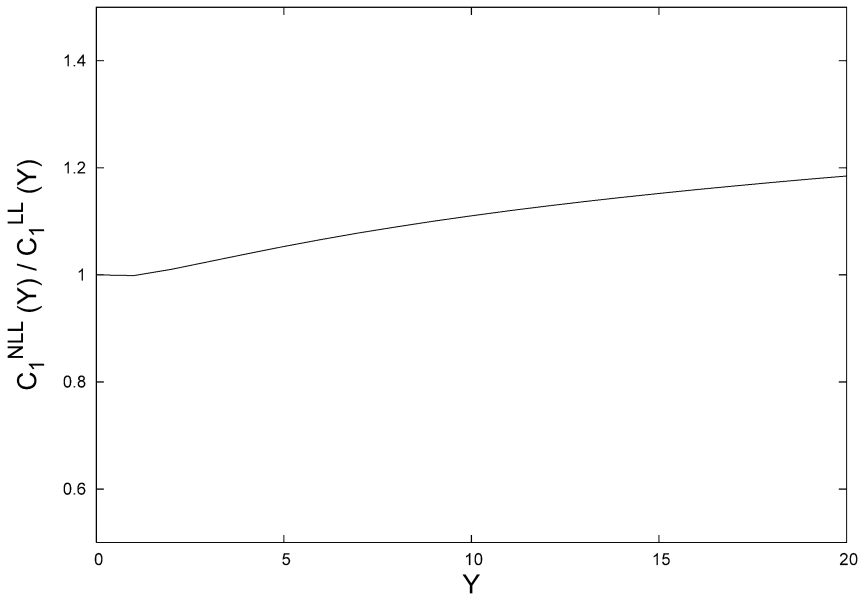


Fig. 6. Comparative ratio between the NLL and LL coefficients for $n = 1$ conformal spin.

This ratio is calculated in Fig. 6. This point is a consequence of the good convergence in terms of asymptotic intercepts of the NLL BFKL calculation for conformal spins larger than zero. In particular the $n = 1$ case is special in that the property of zero intercept at LL, $\chi_0(1, 1/2) = 0$, is

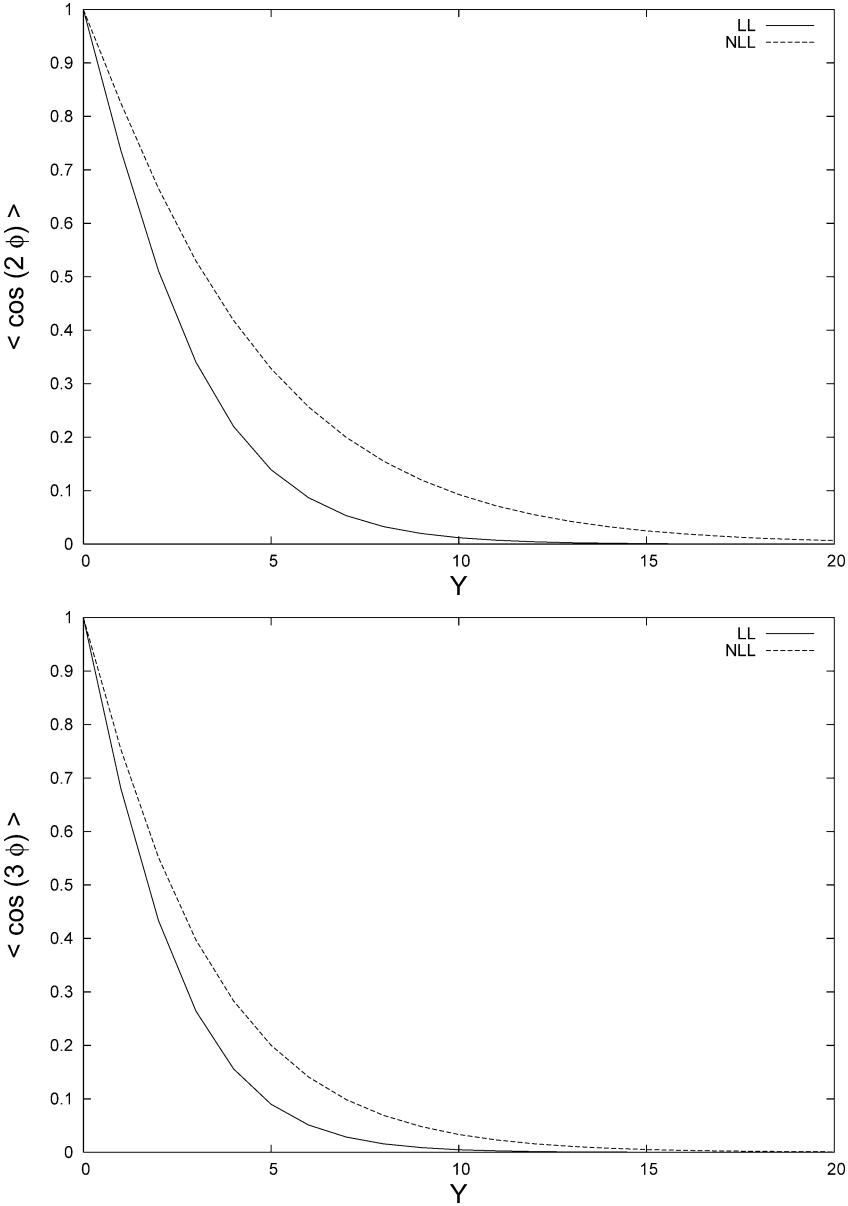


Fig. 7. Dijet azimuthal angle decorrelation as a function of their separation in rapidity.

preserved under radiative corrections since

$$\chi_1\left(1, \frac{1}{2}\right) = S\chi_0\left(1, \frac{1}{2}\right) + \frac{3}{2}\zeta(3) - \frac{\beta_0}{8N_c}\chi_0^2\left(1, \frac{1}{2}\right) + \frac{\psi''(1)}{2} - \phi\left(1, \frac{1}{2}\right) \quad (38)$$

is also zero. For completeness the $m = 2, 3$ cases for $\langle \cos(m\phi) \rangle$ are shown in Fig. 7. These distributions are relevant because they prove the structure of the higher conformal spins. The

trend is the same as previously discussed: the correlation increases as higher order corrections in the small x resummation are included.

4. Conclusions

An analytic procedure has been presented to calculate the effect of higher order corrections in the description of Mueller–Navelet jets where two jets with moderately high and similar transverse momentum are produced at a large relative rapidity separation in hadron–hadron collisions. This is a promising observable to study small x physics at the large hadron collider at CERN given its large energy range. The focus of the analysis has been on those effects with direct origin in the NLO BFKL kernel, while the jet vertices have been considered at LO accuracy. It has been shown how the growth with energy of the cross section is reduced when going from a LL to a NLL approximation, and how the azimuthal angle decorrelations largely decrease due to the higher order effects. The present study has been performed at partonic level while the implementation of a full analysis, including parton distribution functions, NLO jet vertices and the investigation of collinearly improved kernels, will be published elsewhere.

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References

- [1] L.N. Lipatov, *Sov. J. Nucl. Phys.* 23 (1976) 338;
V.S. Fadin, E.A. Kuraev, L.N. Lipatov, *Phys. Lett. B* 60 (1975) 50;
V.S. Fadin, E.A. Kuraev, L.N. Lipatov, *Sov. Phys. JETP* 44 (1976) 443;
V.S. Fadin, E.A. Kuraev, L.N. Lipatov, *Sov. Phys. JETP* 45 (1977) 199;
I.I. Balitsky, L.N. Lipatov, *Sov. J. Nucl. Phys.* 28 (1978) 822;
I.I. Balitsky, L.N. Lipatov, *JETP Lett.* 30 (1979) 355.
- [2] A.H. Mueller, H. Navelet, *Nucl. Phys. B* 282 (1987) 727.
- [3] V. Del Duca, C.R. Schmidt, *Phys. Rev. D* 49 (1994) 4510;
V. Del Duca, C.R. Schmidt, *Phys. Rev. D* 51 (1995) 2150;
W.J. Stirling, *Nucl. Phys. B* 423 (1994) 56.
- [4] S. Abachi, et al., D0 Collaboration, *Phys. Rev. Lett.* 77 (1996) 595;
B. Abbott, et al., D0 Collaboration, *Phys. Rev. Lett.* 84 (2000) 5722;
B. Abbott, et al., D0 Collaboration, *Phys. Rev. Lett.* 80 (1998) 666;
F. Abe, et al., CDF Collaboration, *Phys. Rev. Lett.* 77 (1996) 5336;
F. Abe, et al., CDF Collaboration, *Phys. Rev. Lett.* 78 (1997) 4307, Erratum.
- [5] L.H. Orr, W.J. Stirling, *Phys. Rev. D* 56 (1997) 5875;
L.H. Orr, W.J. Stirling, *Phys. Lett. B* 429 (1998) 135;
L.H. Orr, W.J. Stirling, *Phys. Lett. B* 436 (1998) 372;
J.R. Andersen, V. Del Duca, S. Frixione, C.R. Schmidt, W.J. Stirling, *JHEP* 0102 (2001) 007;
J. Kwiecinski, A.D. Martin, L. Motyka, J. Outhwaite, *Phys. Lett. B* 514 (2001) 355.
- [6] V.S. Fadin, L.N. Lipatov, *Phys. Lett. B* 429 (1998) 127;
G. Camici, M. Ciafaloni, *Phys. Lett. B* 430 (1998) 349.
- [7] J.R. Andersen, A. Sabio Vera, *Nucl. Phys. B* 679 (2004) 345;
J.R. Andersen, A. Sabio Vera, *Nucl. Phys. B* 699 (2004) 90.
- [8] J.R. Andersen, A. Sabio Vera, *Phys. Lett. B* 567 (2003) 116.
- [9] A. Sabio Vera, hep-ph/0307046;

- A. Sabio Vera, hep-ph/0408008;
A. Sabio Vera, hep-ph/0510046;
J.R. Andersen, hep-ph/0406241;
J.R. Andersen, AIP Conf. Proc. 792 (2005) 726;
M. Dobbs, et al., hep-ph/0403100;
M. Dittmar, et al., hep-ph/0511119.
- [10] J. Bartels, D. Colferai, G.P. Vacca, Eur. Phys. J. C 24 (2002) 83;
J. Bartels, D. Colferai, G.P. Vacca, Eur. Phys. J. C 29 (2003) 235.
- [11] G.P. Salam, JHEP 9807 (1998) 019;
C.R. Schmidt, Phys. Rev. D 60 (1999) 074003;
J.R. Forshaw, D.A. Ross, A. Sabio Vera, Phys. Lett. B 455 (1999) 273;
A. Sabio Vera, Nucl. Phys. B 722 (2005) 65.
- [12] D.Y. Ivanov, A. Papa, Nucl. Phys. B 732 (2006) 183.
- [13] A.V. Kotikov, L.N. Lipatov, Nucl. Phys. B 582 (2000) 19.
- [14] J.R. Andersen, A. Sabio Vera, JHEP 0501 (2005) 045.
- [15] L.N. Lipatov, Sov. Phys. JETP 63 (1986) 904, Zh. Eksp. Teor. Fiz. 90 (1986) 1536.
- [16] J.R. Forshaw, D.A. Ross, A. Sabio Vera, Phys. Lett. B 498 (2001) 149.