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Authors

Xiao, Bo-Wen
Yuan, Feng

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BFKL and Sudakov resummation in higgs boson plus jet production with large rapidity separation

Bo-Wen Xiao^{a,*}, Feng Yuan^{b,*}

^a Key Laboratory of Quark and Lepton Physics (MOE) and Institute of Particle Physics, Central China Normal University, Wuhan 430079, China

^b Nuclear Science Division, Lawrence Berkeley National Laboratory, Berkeley, CA 94720, USA

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ABSTRACT

We investigate the QCD resummation for the Higgs boson plus a high P_T jet production with large rapidity separations in proton–proton collisions at the LHC. The relevant Balitsky–Fadin–Kuraev–Lipatov (BFKL) and Sudakov logs are identified and resummed. In particular, we apply recent developments of the transverse momentum dependent factorization formalism in the impact factors, which provides a systematic framework to incorporate both the BFKL and Sudakov resummations.

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1. Introduction

The production of a Higgs boson in association with a large transverse momentum jet is an important channel at the LHC to investigate the Higgs boson property, in particular, when they are produced with large rapidity separation [1–5]. To explore the full potential to distinguish between different production mechanisms, we need to improve the theoretical computations of this process. There have been great progresses in higher order perturbative calculations in the last few years with next-to-next-to-leading order results available [6–10]. In addition, there exist large logarithms to be resummed to all orders to make reliable theoretical predictions. Because of the large rapidity separation between the two final state particles, an important contribution comes from the so-called Balitsky–Fadin–Kuraev–Lipatov (BFKL) evolution [11], similar to the Mueller–Navelet (MN) dijet production [12]. This physics was first studied in Ref. [13] for Higgs boson plus one or two jets production, where the leading order impact factor for Higgs boson production was derived. Meanwhile, there are Sudakov-type of large logarithms [14,15], which has been shown in Ref. [16] for central rapidity Higgs boson plus jet production. In this paper, we will develop a systematic framework to implement both BFKL and Sudakov resummations for the Higgs boson plus jet production with large rapidity separation at the LHC. This is crucial for phenomenological study to investigate the coupling between the Higgs boson and other particles in the Standard Model.

We focus on the QCD process contributions to the Higgs boson plus jet production,¹

$$p(P_A) + p(P_B) \rightarrow H(y_1, k_{1\perp}) + \text{Jet}(y_2, k_{2\perp}), \quad (1)$$

where the incoming hadrons carry momenta P_A and P_B , two final state particles with rapidities y_1 and y_2 , transverse momenta $k_{1\perp}$ and $k_{2\perp}$, respectively. We take the limit of large rapidity difference $Y = |y_1 - y_2| \sim \frac{1}{\alpha_s} \gg 1$, where as schematically shown in Fig. 1, we can write down the following factorization formula in the momentum space as follows

$$\frac{d^6\sigma(pp \rightarrow H + J)}{dy_1 dy_2 d^2k_{1\perp} d^2k_{2\perp}} = \sum_{b=q,g} \int d^2q_{1\perp} d^2q_{2\perp} V_h(x_1, q_{1\perp}, k_{1\perp}) \times V_b(x_2, q_{2\perp}, k_{2\perp}) f_{BFKL}(q_{1\perp}, q_{2\perp}; Y), \quad (2)$$

where V_b is the impact factor for parton b (quark or gluon), V_h for the Higgs boson, and f_{BFKL} represents the BFKL evolution effects due to gluon radiation in the rapidity interval of Y . This factorization is very much similar to the MN-dijet production process [12], where the dijet are well separated in rapidity. There have been great progresses in theory developments for MN-dijet productions [18–27], and the first detailed experiment measurement

* Corresponding authors.

E-mail addresses: bo.w.xiao@gmail.com (B.-W. Xiao), fyuan@lbl.gov (F. Yuan).

¹ For the electroweak process, such as vector boson fusion contributions, there is no BFKL type of logarithms at higher orders, though the QCD-Sudakov double logarithms still exist [17].

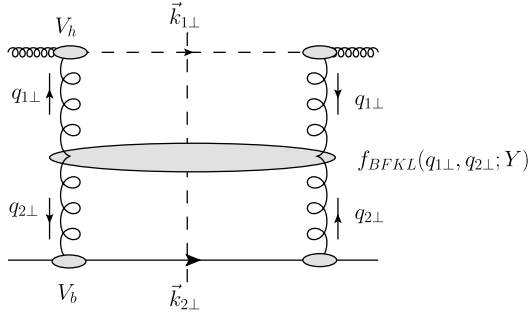


Fig. 1. Schematic factorization for Higgs boson and a hard jet production with large rapidity separation between them at the LHC: f_{BFKL} represents the BFKL evolution with gluon radiation in the rapidity interval between the two final state particles; V_h and V_b for the transverse momentum resummation effects with gluon radiation in the forward rapidity region of the incoming gluon and partons, respectively.

have been performed by the CMS collaboration at the LHC [28]. The experimental results have been interpreted as an evidence for the BFKL dynamics [26]. In our previous publication, we have shown that there exist Sudakov logarithms in MN-dijet productions and these logarithms should be resummed as well [29]. Our results in the following can be applied to MN-dijet processes, and will confirm the factorization formula postulated there. The important difference between the Higgs+Jet process and the MN dijet process is that the Higgs mass can serve as an additional scale which makes the Sudakov resummation a bit more non-trivial. Using the Fourier transform,² it is straightforward to write the above factorization formula in the coordinate space, where the resummation is performed

$$\begin{aligned} & \frac{d^6\sigma(pp \rightarrow H + J)}{dy_1 dy_2 d^2k_{1\perp} d^2k_{2\perp}} \\ &= \sum_b \int \frac{d^2b_{1\perp} d^2b_{2\perp}}{(2\pi)^4} e^{ik_{1\perp} \cdot b_{1\perp} + ik_{2\perp} \cdot b_{2\perp}} \\ & \quad \times \tilde{V}_h(x_1, b_{1\perp}) \tilde{V}_b(x_2, b_{2\perp}) \tilde{f}_{BFKL}(b_{1\perp}, b_{2\perp}; Y). \end{aligned} \quad (3)$$

It is well-known that both the Sudakov resummation and BFKL evolution can be more conveniently carried out in the coordinate space.

In the inclusive production process, the impact factors can be calculated in the collinear factorization approach. However, in the study of the azimuthal angular distribution between the two final state particles, there exist Sudakov double logarithms in the back-to-back correlation kinematics, where, for example, $k_{1\perp}$ is close to $q_{1\perp}$. To resum these large Sudakov type logarithms, we apply the transverse momentum dependent (TMD) factorization [15,30,31] for the impact factors in Eq. (2): V_h is the TMD gluon distribution, and V_b is factorized into the TMD parton distribution and the soft factor associated with the final state jet. The resummation is carried out by solving the relevant evolution equation.

The physical argument for the above factorization is that the higher order gluon radiations can be classified according to the relevant phase space. The most important gluon radiation comes from the large rapidity separation region between the two final state particles, which generates the BFKL evolution effects and can be factorized into the factor f_{BFKL} . In the meantime, the gluon radiations in the forward regions of the incoming quark and gluon are factorized into the TMD parton distributions, with a manifest

rapidity cut-off in their definitions [30]. Therefore, the BFKL and Sudakov contributions are clearly separated out in the gluon radiation phase space and the factorization can be proved accordingly. This will build a systematic framework to implement both BFKL and Sudakov resummation in the process of Eq. (1).

From the resummation point of view, there are two interesting types of logarithms arising from a one-loop calculation for this process, namely, the BFKL type logarithm $\alpha_s Y$ and the Sudakov logarithms. They can be resummed into the factor f_{BFKL} and the impact factors, respectively. As far as the collinear logarithms are concerned, they can be easily dealt with the help of the jet definition and the collinear parton distributions.

The rest of the paper is organized as follows. We take the example of the quark impact factor to demonstrate how the TMD factorization and resummation are applied. Similar results can be obtained for the gluon impact factor. We then calculate the Higgs impact factor, which is factorized into the TMD gluon distribution. Finally, we summarize our results.

2. Impact factors for the quark and gluon

The partonic scattering of the process described in Eq. (1) comes from quark–gluon and gluon–gluon channels. According to the proposed BFKL factorization, we can separate the calculations into the quark or gluon impact factor and the gluon–Higgs impact factor. Let us take the quark impact factor as an example, which has been studied extensively in the literature [18–27]. In the following, we will investigate how to factorize these results into the TMD parton distributions, soft and hard factors. To simplify the derivation, we take the generic kinematics: the final quark jet has transverse momentum $k_{J\perp}$, the vertical gluon has q_{\perp} . The leading order can be expressed as a Delta function of $\vec{k}_{\perp} = \vec{k}_{J\perp} - \vec{q}_{\perp}$. This translates into a constant in the Fourier transform b_{\perp} -space,

$$\tilde{V}_q(x, b_{\perp})|_{LO} = \tilde{V}_q^{(0)} f_q(x), \quad (4)$$

where $f_q(x)$ represents the quark distribution function and $\tilde{V}_q^{(0)}$ for the leading order factor. At one-loop order, there are virtual and real gluon radiation contributions. The virtual contribution can be written as

$$\Gamma^v = \frac{\alpha_s}{2\pi} \left(\frac{\mu^2}{\bar{q}_{\perp}^2} \right)^{\epsilon} \left\{ C_F \left[-\frac{2}{\epsilon^2} - \frac{3}{\epsilon} - \frac{23}{4} + \frac{3}{2}\pi^2 \right] + \mathcal{K} \right\}, \quad (5)$$

where $\mathcal{K} = C_A \left(\frac{67}{18} - \frac{\pi^2}{6} \right) - \frac{5}{9} N_f$, $C_F = 4/3$ and $N_c = 3$, N_f represents the number of quark flavors. Here we work in the dimensional regulation with $D = 4 - 2\epsilon$ and $\overline{\text{MS}}$ scheme. In the above equation, \bar{q}_{\perp} is the t-channel momentum transfer due to the BFKL factor, and a universal energy dependent term proportional to $C_A \ln(s_0/\bar{q}_{\perp}^2)/\epsilon$ is omitted.³ Together with the similar term from the real gluon radiation, it generates the corresponding BFKL contribution, which can be used to derive the well-known BFKL evolution equation. The detailed procedure can be found in Ref. [29,32]. In the following, we will focus on the QCD dynamics associated with the Sudakov logarithms, and neglect the BFKL part to simplify the derivations.

Real gluon radiation will contribute to a finite transverse momentum k_{\perp} . A schematic diagram is shown in Fig. 2, which has

² We introduce the Fourier transform in b_{\perp} -space, for example, $V_q(x, k_{1\perp}, q_{1\perp}) = \int \frac{d^2b_{1\perp}}{(2\pi)^2} e^{i(k_{1\perp} - q_{1\perp}) \cdot b_{1\perp}} \tilde{V}_q(x, b_{1\perp})$ for the quark Impact Factor.

³ It is very clear that this term corresponds to the BFKL dynamics, since it is proportional to C_A instead of C_F and it depends on the collision energy. It is well-known that the BFKL evolution equation is an energy evolution equation which is proportional to C_A .

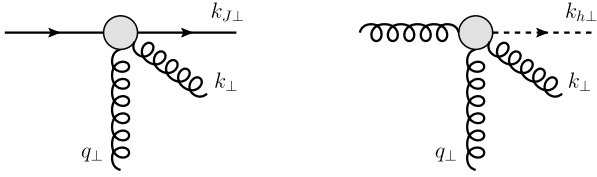


Fig. 2. Typical real gluon radiation diagrams for the quark Impact Factor (left) and Higgs Impact Factor (right) calculations.

been calculated in the literature. Its contribution can be expressed as,

$$\frac{\alpha_s}{2\pi^2} \left[\frac{1+z^2}{1-z} - \epsilon(1-z) \right] \left\{ C_F \frac{(1-z)^2 q_\perp^2}{k_\perp^2 (k_\perp - (1-z)q_\perp)^2} \right\}, \quad (6)$$

where $(1-z)$ is the momentum fraction of the incoming quark carried by the radiated gluon with transverse momentum k_\perp . Clearly, there are two important contributions from the singularities in the above equation: (1) collinear gluon radiation associated with the incoming quark when $k_\perp \rightarrow 0$; (2) soft gluon radiation associated with the final state jet when $k_\perp \sim (1-z)q_\perp$. We take the leading power contribution in the limit of $k_\perp \ll q_\perp$, where soft gluon radiation with $z \rightarrow 1$ plays an important role. By applying the plus function prescription to separate out the collinear gluon radiation from the incoming quark, we are left with the following term,

$$\frac{\alpha_s}{2\pi^2} C_F \frac{1}{k_\perp^2} \delta(1-z) \int \frac{d\alpha}{\alpha} (1 + (1-\alpha)^2) \frac{\alpha^2 q_\perp^2}{(k_\perp - \alpha q_\perp)^2}, \quad (7)$$

which contains the soft divergence at $k_\perp \rightarrow 0$ and the collinear divergence associated with the final state jet at $k_\perp - \alpha q_\perp \rightarrow 0$. Following the same procedure described in Refs. [33–35], we apply the anti- k_t jet algorithm and the narrow jet approximation [36,37] which lead to

$$\frac{\alpha_s}{2\pi^2} C_F \frac{1}{k_\perp^2} \delta(1-z) \left[\ln \frac{q_\perp^2}{k_\perp^2} + \ln \frac{1}{R^2} + \epsilon \left(\frac{1}{2} \ln^2 \frac{1}{R^2} + \frac{\pi^2}{6} \right) \right], \quad (8)$$

where R represents the jet size. When Fourier transformed into b_\perp -space with respect to k_\perp , the above result will contain a soft divergence in terms of $1/\epsilon^2$, which will be cancelled out by the virtual contribution in Eq. (5). Adding them together, we find the one-loop result for \tilde{V}_q as,

$$\begin{aligned} \tilde{V}_q(x, b_\perp)|_{\text{NLO}} &= \tilde{V}_q^{(0)} \int \frac{dx'}{x'} f_q(x') \frac{\alpha_s}{2\pi} \left\{ C_F \mathcal{P}_{qq}(z) \left(-\frac{1}{\epsilon} - \ln \frac{q_\perp^2 b_\perp^2}{c_0^2} \right) \right. \\ &\quad - (1-z)C_F + \delta(1-z) \left[C_F \left(-\frac{1}{2} \ln^2 \left(\frac{q_\perp^2 b_\perp^2}{c_0^2} \right) \right) \right. \\ &\quad \left. \left. + \left(\frac{3}{2} - \ln \frac{1}{R^2} \right) \ln \frac{q_\perp^2 b_\perp^2}{c_0^2} \right) + \mathcal{K} + \Delta I_q \right] \right\}, \quad (9) \end{aligned}$$

where $x' = x/z$ and $\tilde{V}_q^{(0)}$ represents the leading order normalization as mentioned above, $c_0 = 2e^{-\gamma_E}$, $\mathcal{P}_{qq}(z)$ is the quark-quark splitting kernel and $\Delta I_q = C_F \left[\frac{3}{2} \ln \frac{1}{R^2} + \frac{3}{4} + \frac{2}{3} \pi^2 \right]$. In reaching the above expression, we have also included the jet contribution [35]. Clearly, there are Sudakov double and single logarithms. The above result can be factorized into the TMD quark distribution and the soft factor associated with the final state jet. Here we follow the Collins 2011 scheme for the definition of TMDs, which are defined with soft factor subtraction [30] as follows

$$f_q^{(sub.)}(x, b_\perp, \mu_F, \zeta_c) = f_q^{unsub.}(x, b_\perp) \sqrt{\frac{S_2^{\bar{n},v}(b_\perp)}{S_2^{n,\bar{n}}(b_\perp) S_2^{n,v}(b_\perp)}}, \quad (10)$$

where b_\perp is the Fourier conjugate variable respect to the transverse momentum k_\perp , μ_F the factorization scale and $\zeta_c^2 = x^2(2v \cdot P)^2/v^2 = 2(xP^+)^2 e^{-2y_n}$ with y_n the rapidity cut-off in the Collins-2011 scheme. The second factor corresponds to the soft factor subtraction with n and \bar{n} as the light-front vectors $n = (1^-, 0^+, 0_\perp)$, $\bar{n} = (0^-, 1^+, 0_\perp)$, whereas v is an off-light-front four-vector $v = (v^-, v^+, 0_\perp)$ with $v^- \gg v^+$. The un-subtracted TMD reads as

$$\begin{aligned} f_q^{unsub.}(x, k_\perp) &= \frac{1}{2} \int \frac{d\xi^- d^2 \xi_\perp}{(2\pi)^3} e^{-ix\xi^- P^+ + i\vec{\xi}_\perp \cdot \vec{k}_\perp} \\ &\quad \times \langle PS | \bar{\psi}(\xi) \mathcal{L}_n^\dagger(\xi) \gamma^+ \mathcal{L}_n(0) \psi(0) | PS \rangle, \quad (11) \end{aligned}$$

with the gauge link defined as

$$\mathcal{L}_n(\xi) \equiv \exp \left(-ig \int_0^{-\infty} d\lambda v \cdot A(\lambda n + \xi) \right).$$

The light-cone singularity in the un-subtracted TMDs is cancelled out by the soft factor as in Eq. (10) with S^{v_1, v_2} defined as

$$S_2^{v_1, v_2}(b_\perp) = \langle 0 | \mathcal{L}_{v_2}^\dagger(b_\perp) \mathcal{L}_{v_1}^\dagger(b_\perp) \mathcal{L}_{v_1}(0) \mathcal{L}_{v_2}(0) | 0 \rangle. \quad (12)$$

Following the similar idea, we introduce a subtracted soft factor associated with the final state jet,

$$S_J(b_\perp, \mu_F) = \sqrt{\frac{S_{n, n_1}(b_\perp) S_{n_1, \bar{n}}(b_\perp)}{S_{n, \bar{n}}(b_\perp)}}, \quad (13)$$

where n_1 represents the jet direction. One-loop calculation leads to the following result,

$$S_J^{(1)} = \frac{\alpha_s}{2\pi} C_F \left[\ln \frac{1}{R^2} \ln \frac{b_\perp^2 \mu_F^2}{c_0^2} + \frac{1}{2} \ln^2 \left(\frac{1}{R^2} \right) + \frac{\pi^2}{6} \right], \quad (14)$$

again with narrow jet approximation, from which we obtain the anomalous dimension $\gamma^{(s)} = \frac{\alpha_s}{2\pi} C_F \ln(1/R^2)$. Together with the result for the quark distribution from Ref. [30,38], the following TMD factorization can be verified at one-loop order,

$$\tilde{V}_q(x, b_\perp) = f_q^{(sub.)}(x, b_\perp, \mu_F, \zeta_c) S_J(b_\perp, \mu_F) H(q_\perp, \mu_F). \quad (15)$$

Furthermore, in order to eliminate the large logarithms in the hard factor $H^{(1)}$, we have to choose the appropriate scales as $\mu_F^2 = \zeta_c^2 = \bar{q}_\perp^2$. This corresponds to the factorization that the TMD quark distribution only contains contribution from the gluon radiation in the forward region of the incoming quark. The gluon radiation in the central region (rapidity interval between the two final state particles) belongs to the BFKL evolution. Finally, following the Collins-Soper-Sterman (CSS) resummation approach [15], we obtain the all order result as follows

$$\begin{aligned} \tilde{V}_q(x, b_\perp) &= \tilde{V}_q^{(0)} e^{-S_q(q_\perp, b_\perp)} C \otimes f_q(x, \bar{\mu} = c_0/b) \\ &\quad \times \left[1 + \frac{\alpha_s}{2\pi} (\mathcal{K} + \Delta I_q) \right], \quad (16) \end{aligned}$$

where \otimes represents the convolution in x and $f_q(x, \bar{\mu})$ the integrated quark distribution. Following the so-called ‘‘TMD’’ scheme [39,40] in CSS resummation, the hard and soft factors at the appropriate scale lead to the coefficients at α_s , represented by \mathcal{K} and ΔI_q . The Sudakov factor can be written as

$$S_q(q_\perp, b_\perp) = \int_{c_0^2/b_\perp^2}^{q_\perp^2} \frac{d\mu^2}{\mu^2} \left[A_q \ln \frac{q_\perp^2}{\mu^2} + B_q + D_q \ln \frac{1}{R^2} \right], \quad (17)$$

with $A_q = \sum_i A_q^{(i)}$, $A_q^{(1)} = D_q^{(1)} = \frac{\alpha_s}{2\pi} C_F$, $B_q^{(1)} = -\frac{3}{2} A_q^{(1)}$, and the C coefficient function is $C^{(1)} = \frac{\alpha_s}{2\pi} C_F (1-x)$.

Similar calculations can be performed for the gluon impact factor,

$$\begin{aligned} \tilde{V}_g(x, b_\perp) &= \tilde{V}_g^{(0)} e^{-S_g(q_\perp, b_\perp)} C \otimes f_g(x, \bar{\mu} = c_0/b) \\ &\times \left[1 + \frac{\alpha_s}{2\pi} (\mathcal{K} + \Delta I_g) \right], \end{aligned} \quad (18)$$

with one-loop results as $A_g^{(1)} = D_g^{(1)} = \frac{\alpha_s}{2\pi} C_A$, $B_g^{(1)} = -2\beta_0 A^{(1)}$, and $\Delta I_g = C_A \left(2\beta_0 \ln \frac{1}{R^2} - \frac{\pi^2}{6} \right) - \frac{N_f}{6}$ with $\beta_0 = \frac{11}{12} - \frac{N_f}{18}$ and N_f being the number of flavors. The C coefficient vanishes at one-loop order. In the BFKL factorization, the quark and gluon impact factors are universal, which means that they are same as those in MN-dijet processes. Indeed, we can apply the above impact factors and obtain the consistent results as those in Ref. [29].

3. Impact factor for the higgs boson

The computation procedure of the quark and gluon Impact Factors can be applied to the Higgs impact factor as well. The leading order impact factor has been computed in Ref. [13,41]. It is again a Delta function of $\vec{k}_\perp = \vec{k}_{h\perp} - \vec{q}_\perp$, where $k_{h\perp}$ and q_\perp are transverse momenta of the final state Higgs boson and the vertical gluon, respectively. In the following we will focus on how the factorization works and derive the associated impact factor at one-loop order. To simplify the derivation, we will apply the effective theory approach for the Higgs boson production in the heavy top quark limit [42,43]. We leave the finite top quark mass corrections [41] for a future study, where we expect that the factorization will still be valid, although the impact factor will be modified accordingly.

At one-loop order, the virtual graph contribution in the gluon-to-Higgs boson impact factor can be deduced from that in Higgs boson plus jet production by taking the limit of the large rapidity separation between the final state particles [44,45],

$$\Gamma^\nu = \frac{\alpha_s}{2\pi} \left(\frac{\mu^2}{\bar{q}_\perp^2} \right)^\epsilon \left\{ N_c \left[-\frac{1}{\epsilon^2} + \frac{1}{\epsilon} \left(\ln \frac{\bar{m}^2}{\bar{q}_\perp^2} - 2\beta_0 \right) \right] + \Delta I_h + \mathcal{K} \right\} \quad (19)$$

where $\bar{m}^2 = m_h^2 + \bar{q}_\perp^2$ with Higgs mass m_h ,

$$\Delta I_h = C_A \left[\pi^2 + 2\text{Li}_2(x_q) + \ln(x_q) \ln \frac{(1+x_q)^2}{x_q} \right]$$

and x_q is defined as $x_q = \bar{q}_\perp^2/m_h^2$. Again, we have subtracted the universal energy dependent term related to the BFKL evolution. The contribution from the real gluon radiation shown in Fig. 2 can be summarized as

$$\frac{\alpha_s}{2\pi^2} \frac{1}{k_\perp^2} C_A \left\{ \mathcal{P}_{gg}(z) + \delta(1-z) \left[\ln \frac{\bar{m}^2}{k_\perp^2} - 2\beta_0 \right] \right\}, \quad (20)$$

where the terms associated with the BFKL evolution have been subtracted from the real gluon contribution, see the discussions in the following section. Adding the above two terms together, we obtain the following result in the b_\perp -space,

$$\begin{aligned} \tilde{V}_h(x, b_\perp)|_{\text{NLO}} &= \tilde{V}_h^{(0)} \int \frac{dx'}{x'} f_g(x') \left\{ \frac{\alpha_s}{2\pi} [C_A \mathcal{P}_{gg}(z) + \delta(1-z) (\Delta I_h + \mathcal{K})] \right. \\ &+ \frac{\alpha_s}{2\pi} C_A \delta(1-z) \left[-\frac{1}{2} \ln^2 \left(\frac{q_\perp^2 b_\perp^2}{c_0^2} \right) \right. \\ &\left. \left. + \left(\ln \frac{\bar{m}^2}{q_\perp^2} - 2\beta_0 \right) \ln \frac{c_0^2}{q_\perp^2 b_\perp^2} \right] \right\}, \end{aligned} \quad (21)$$

where $\tilde{V}_h^{(0)}$ represents the leading order factor [13], $z = x/x'$, $f_g(x')$ is the gluon distribution and $\mathcal{P}_{gg}(z)$ for the gluon-gluon splitting kernel. Again, the above result can be factorized into the TMD gluon distribution,

$$\tilde{V}_h = x f_g(x, b_\perp, \mu_F, \zeta_c) H(q_\perp, \mu_F), \quad (22)$$

for which we will choose the factorization scale $\mu_F^2 = \bar{q}_\perp^2$ and $\zeta_c^2 = \bar{m}^2$ to eliminate the large logarithms in the hard factor. All order resummation is achieved by solving the energy evolution equation for the TMD gluon distribution,

$$\begin{aligned} \tilde{V}_h(x, b_\perp) &= \tilde{V}_h^{(0)} e^{-S_h(q_\perp, b_\perp)} C \otimes f_g(x, \bar{\mu} = c_0/b) \\ &\times \left[1 + \frac{\alpha_s}{2\pi} (\mathcal{K} + \Delta I_h) \right], \end{aligned} \quad (23)$$

where the Sudakov factor can be written as

$$S_h(q_\perp, b_\perp) = \int_{c_0^2/b_\perp^2}^{q_\perp^2} \frac{d\mu^2}{\mu^2} \left[A_h \ln \frac{\bar{m}^2}{\mu^2} + B_h \right]. \quad (24)$$

We find that $A_h = A_g$, $B_h = B_g$, which is because they come from the same TMD gluon distribution, and the C coefficient function vanishes at one-loop order.

4. BFKL evolution

In this section, we will examine the associated BFKL evolution equations for the above processes. In particular, we want to compare the associated contributions between the MN-dijet scattering and Higgs plus jet production with large rapidity separation. The goal is to show that the BFKL evolution has been consistently taken into account with the impact factor calculations.

In the quark impact factor calculation, we have the following term associated with the BFKL gluon radiation,

$$\frac{\alpha_s}{2\pi^2} C_A \int \frac{dz}{z} \frac{2(1-z) k_\perp \cdot (k_\perp - z q_\perp)}{k_\perp^2 (k_\perp - z q_\perp)^2}. \quad (25)$$

Here z integral is limited by $z > k_\perp^2/s_\Lambda$, where $s_\Lambda = 2p_1^- p_A^+$ represents the invariant mass cut-off with p_1 for the incoming quark momentum. This cut-off will be combined with the other particle in the final state to obtain the boost invariant evolution for the BFKL gluon radiation. There is no final state jet divergence, because the $z \sim k_\perp/q_\perp$ is regulated by the numerator. Further calculations can be performed by averaging the azimuthal angle between k_\perp and q_\perp , from which we find the integrand vanishes in the region of $z > k_\perp/q_\perp$. Therefore, the final result will be

$$\frac{\alpha_s}{2\pi^2} C_A \frac{1}{k_\perp^2} \ln \frac{s_\Lambda^2}{q_\perp^2 k_\perp^2}. \quad (26)$$

In the case of Mueller-Navelet dijet productions, we can perform the same calculation for the other impact factor and introduce $\bar{s}_\Lambda = 2p_2^+ p_B^-$ with the following expression,

$$\frac{\alpha_s}{2\pi^2} C_A \frac{1}{k_\perp^2} \ln \frac{\bar{s}_\Lambda^2}{q_\perp^2 k_\perp^2}. \quad (27)$$

By taking into account the kinematic relation for the massless particles in the final state of the MN-dijet production, we find that $s_\Lambda \bar{s}_\Lambda = s s_y = s_y^2$ where s_y defined by the jet transverse momentum q_\perp and the rapidity difference between the two final state particles $s_y = q_\perp^2 e^{\Delta Y}$. Therefore, the BFKL evolution can be simplified as

$$\frac{\alpha_s}{2\pi^2} C_A \frac{2}{k_\perp^2} \ln \frac{s_y^2}{q_\perp^2 k_\perp^2}. \quad (28)$$

The above is the universal BFKL evolution contribution, which only depends on the transverse momentum and the rapidity between the two final state particles. We expect the same BFKL contribution from the Higgs plus jet production as well. This provides an important cross check for the above calculations.

From the details of the gluon-Higgs impact factor calculation, we find that there is only the following term contributing to the BFKL evolution,

$$\frac{\alpha_s}{2\pi^2} C_A \frac{2}{k_\perp^2} \ln \frac{s_\Lambda}{k_\perp^2}, \quad (29)$$

and all other power suppressed terms drop out from the calculations in Ref. [16]. It is interesting to note that the above term can be separated into two terms,

$$\frac{\alpha_s}{2\pi^2} C_A \frac{2}{k_\perp^2} \ln \frac{s_\Lambda}{k_\perp^2} = \frac{\alpha_s}{2\pi^2} C_A \frac{1}{k_\perp^2} \left[\ln \frac{\tilde{m}^2}{k_\perp^2} + \ln \frac{s_\Lambda^2}{\tilde{m}^2 k_\perp^2} \right], \quad (30)$$

where the first term contributes to the Sudakov logs (as in Eqs. (20) and (24)), and the second term gives the BFKL evolution after combined with the BFKL term from the quark impact factor calculation. The latter is achieved by taking into account the following identity from the kinematics of Higgs boson plus jet production,

$$s_\Lambda \bar{s}_\Lambda^2 = s^2 s_y^2 = s_y^4 \tilde{m}^2 / q_\perp^2, \quad (31)$$

in the limit of large rapidity separation (ΔY) between the Higgs boson and the produced jet, where again s_y is defined as $s_y = q_\perp^2 e^{\Delta Y}$.

5. Summary and discussions

The final resummation results for the BFKL and Sudakov resummation effects in the Higgs boson plus jet production with large rapidity separation are obtained by substituting the results in Eqs. (16), (18), (23) into Eq. (3). An important cross check has been performed by comparing to the derivation in Ref. [16] with only Sudakov resummation, and we find the complete agreement.

The factorization method developed in this paper can have great impact in LHC physics. A potential application is to study in Higgs plus two jets production where the final state three particles are well separated in rapidity. This channel is an important place to study the vector boson fusion contribution in Higgs boson production at the LHC, where we need to understand the QCD resummation contributions accurately.

Theoretically, both BFKL and Sudakov resummations are the important corner stones in the perturbative QCD applications to high energy hadronic collisions. Recently, there have been strong interests [33,46–51] to combine these two resummations consistently in the hard scattering processes at various collider experiments. A detailed comparison of different approaches deserves a future

study, in particular, between the TMD factorization formalism of [33,46] and those of direct computation in Ref. [48]. Our results in this paper is a step further toward a systematic framework to deal with both physics. We anticipate more applications in the future, in particular, for multi-jets events at the LHC, such as three-jet or four-jet productions [52–54].

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