Kinematical Correlations for Higgs Boson Plus High P_{T} Jet Production at Hadron Colliders

Peng Sun, C.-P. Yuan, and Feng Yuan

Phys. Rev. Lett. **114**, 202001 — Published 22 May 2015

DOI: 10.1103/PhysRevLett.114.202001

Kinematical correlations for Higgs boson plus High P_T Jet Production at Hadron Colliders

Peng Sun,[1] C.-P. Yuan,[2] and Feng Yuan[1]

[1] *Nuclear Science Division, Lawrence Berkeley National Laboratory, Berkeley, CA 94720, USA*

[2] *Department of Physics and Astronomy, Michigan State University, East Lansing, MI 48824, USA*

Abstract

We investigate the effect of QCD resummation to kinematical correlations in the Higgs boson plus high transverse momentum (P_T) jet events produced at hadron colliders. We show that at the complete one-loop order, the Collins-Soper-Sterman resummation formalism can be applied to derive the Sudakov form factor. We compare the singular behavior of resummation calculation to fixed order prediction in the case that Higgs boson and high P_T jet are produced nearly back-to-back in their transverse momenta, and find a perfect agreement. The phenomenological importance of the resummation effect at the LHC is also demonstrated.

Introduction. Higgs boson discovery at the CERN LHC [1, 2] has stimulated new area of high energy physics research at the colliders, where the precision Higgs physics is at the frontier. This includes Higgs production and decays to investigate the coupling between the Higgs boson and all other particles. The correlation between Higgs and jet production at the LHC will undoubtedly provide important information on the production and further disentangle the electroweak coupling of Higgs boson [3–15]. The goal of this paper is to build a theoretical framework to reduce the uncertainties in the Higgs plus jet production, in particular, in the kinematics of back-to-back azimuthal angular correlation region. In this region, the total transverse momentum of Higgs boson plus jet becomes much smaller than the invariant mass, and the fixed order perturbative calculations suffer from singularities, which will result in large theoretical uncertainties due to factorization/renormalization scale uncertainties [16]. Therefore, we have to perform all order soft gluon resummation to make reliable predictions, and to reduce the theoretical uncertainties.

QCD resummation for this process has its own interest in perturbative QCD. To deal with the divergence in low transverse momentum hard processes, the so-called transverse momentum, or Collins-Soper-Sterman (CSS), resummation is employed [17]. However, the CSS resummation has been mainly applied to the color-neutral particle production, such as inclusive vector boson W/Z and Higgs boson productions. Extension to jet productions in the final state has been much limited. This is not only because of the technique issues associated with the jets in the final state, but also because that the jets carry color and the soft gluon interactions are more complicated than those for color neutral particle production. Nevertheless, there have been progresses in the last few years on the CSS resummation for dijet production in hadronic collisions [18–20]. In this paper, we investigate the CSS resummation for Higgs boson plus one hard jet production,

$$A(P) + B(\bar{P}) \to H + Jet + X \ , \tag{1}$$

where two incoming hadrons carry momenta P and $\bar{P}$, respectively. Because the final state is simpler than that of dijet production, the above process allows us to study the factorization in great detail. Extension to W/Z boson plus jet production shall be straightforward, which are phenomenologically important at the LHC as well.

In the calculations, we apply the effective theory to describe Higgs coupling to gluons in the large top mass limit:

$$\mathcal{L}_{eff} = -\frac{\alpha_s}{12\pi v} F^a_{\mu\nu} F^{a\mu\nu} H, \tag{2}$$

where v is the vacuum expectation value and H the Higgs field, $F^{\mu\nu}$ the gluon field strength tensor and a the color index. Our final resummation formula can be summarized as

$$\frac{d^4\sigma}{dy_h dy_j dP_T^2 d^2q_\perp} = \sum_{ab} \sigma_0 \left[\int \frac{d^2\vec{b}_\perp}{(2\pi)^2} e^{-i\vec{q}_\perp \cdot \vec{b}_\perp} W_{ab\to Hc}(x_1, x_2, b_\perp) + Y_{ab\to Hc} \right] \ , \tag{3}$$

where y_h and y_j are rapidities for the Higgs boson and the jet, P_T for the jet transverse momentum, and $\vec{q}_\perp = \vec{P}_{h\perp} + \vec{P}_T$ for the total transverse momentum of Higgs and the jet. The first term W contains all order resummation and the second term Y comes from the fixed order corrections; σ_0 represents normalization of the differential cross section. In this paper, we will take the dominant $gg \to Hg$ channel as an example to demonstrate how to derive the resummation for W term, which can be written as

$$W_{gg\to Hg}(x_1, x_2, b) \ = \ H_{gg\to Hg}(Q) x_1 f_g(x_1, \mu = b_0/b_\perp) x_2 f_g(x_2, \mu = b_0/b_\perp) e^{-S_{\rm Sud}(Q^2, b_\perp)} \ , \tag{4}$$

at next-to-leading logarithmic (NLL) level, where $Q^2 = s = x_1 x_2 S$ and represents the hard momentum scale, $b_0 = 2e^{-\gamma_E}$, with γ_E being the Euler constant. $f_{a,b}(x,\mu)$ are parton distributions for the incoming partons a and b, and $x_{1,2}$ are momentum fractions of the incoming hadrons carried by the partons. Beyond the NLL, a C function associated with the gluon distribution function will also be included. The Sudakov form factor can be written as

$$S_{\rm Sud}(Q^2, b_\perp) = \int_{b_0^2/b_\perp^2}^{Q^2} \frac{d\mu^2}{\mu^2} \left[\ln\left(\frac{Q^2}{\mu^2}\right) A + B + D_1 \ln\frac{1}{R^2} \right] \ , \tag{5}$$

where R represents the cone size for the jet. Here the parameters A, B, D_1 can be expanded perturbatively in α_s. For $gg \to Hg$ channel, at one-loop order, we have $A = C_A \frac{\alpha_s}{\pi}$, $B = -2C_A\beta_0\frac{\alpha_s}{\pi}$, and $D_1 = C_A\frac{\alpha_s}{2\pi}$. The hard coefficient H can be calculated order by order. From the leading Born diagrams, we have $H^{(0)} = (s^4 + t^4 + u^4 + m_h^8)/(stu)$ [6, 7], where $s = Q^2$, t and u are usual Mandelstam variables for the partonic $2 \to 2$ process.

Comparing to the CSS resummation for inclusive Higgs boson production, we find that the one-loop results for the coefficients A and B agree with those in Ref. [21]. This is because both processes are gluon-gluon initiated processes, and these coefficients come from the gluon splitting contributions. However, because of the final state jet in our case, additional soft gluon radiation will contribute to a resummation factor depending on the jet size, represented by the coefficient D_1 in Eq. (5).

To derive the above resummation, we first calculate $W(b)$ at the complete one-loop order, and show that it can be factorized into the parton distributions and soft and hard factors. The resummation is achieved by solving the associated evolution equations. The asymptotic behavior at low imbalance transverse momentum $q_\perp$ is calculated from the soft and collinear gluon radiations at this order. This asymptotic result will be checked against the full perturbative calculations. Then, we will combine these contribution with those from virtual graphs and collinear jet contributions to derive the one-loop result for $W(b)$.

Asymptotic Behavior at Small-$q_\perp$. The leading order calculations for the process of Eq. (1) comes from the partonic process,

$$g + g \to H + g \ ,$$

which predicts a Delta function at $q_\perp = 0$. At the next-to-leading order (NLO), the real emission diagrams for $g+g \to H+\text{jet}+X$ will contribute to a singular behavior at small-$q_\perp$, in the associate production of Higgs boson and high P_T jet, with additional parton radiation. For the collinear gluon associated with the incoming gluon distribution, they can be easily evaluated, and they are proportional to the gluon-to-gluon splitting kernel at one-loop order. For the soft gluon radiation, we apply the soft gluon approximation in the limit of $q_\perp \ll Q$, and obtain the following expression,

$$\int \frac{d^{D-1}k_g}{(2\pi)^{D-1} 2E_{k_g}} \delta^{(2)}(q_\perp - k_{g\perp}) \left[\frac{p_1\cdot p_2}{p_1\cdot k_g p_2\cdot k_g} + \frac{k_1\cdot p_2}{k_1\cdot k_g p_2\cdot k_g} + \frac{k_1\cdot p_1}{k_1\cdot k_g p_1\cdot k_g} \right] \ , \tag{6}$$

where p_1, p_2 represent the momenta for incoming gluons, k_1 for final state jet, and k_g for the radiated gluon. The above results will lead to soft divergence when Fourier transformed to $b_\perp$ space, for which we will apply the dimension regulation with $D = 4-2\epsilon$. Not all the soft gluon radiation contributes to the finite $q_\perp$. In particular, if the gluon radiation is within the final state jet, its contribution has to be excluded. To evaluate these contributions, we introduce a small offshellness (which is proportional to the cone size R) for k_1 to exclude the gluon radiation inside the jet cone, and further take the narrow jet approximation (NJA) [22, 23], i.e, taking the limit of $R \to 0$. In the NJA, this is equivalent to applying a kinematic cutoff for the radiated gluon.

Adding the soft and collinear gluon radiation together, we derive the asymptotic behavior at small-$q_\perp$,

$$\frac{\alpha_s C_A}{2\pi^2}\frac{1}{q_\perp^2}\int \frac{dx_1' dx_2'}{x_1' x_2'} x_1' g(x_1') x_2' g(x_2') \left[\{\delta(\xi_2-1)\xi_1 \mathcal{P}_{gg}(\xi_1) + (\xi_1 \leftrightarrow \xi_2)\}\right.$$
$$\left. +\delta(\xi_1-1)\delta(\xi_2-1)\left(2\ln\frac{Q^2}{q_\perp^2} - 4\beta_0 + \ln\frac{1}{R^2} + \epsilon\left(\frac{1}{2}\ln^2\left(\frac{1}{R^2}\right) + \frac{\pi^2}{6}\right)\right)\right] , \qquad (7)$$

where $\xi_1 = x_1/x_1'$, $\xi_2 = x_2/x_2'$, $\mathcal{P}_{gg}$ is the gluon splitting kernel and $\beta_0 = (11 - 2N_f/3)/12$, with N_f being the number of effective light quarks. We have kept the $\epsilon = (4-D)/2$ terms, which will contribute when Fourier transforming to $b_\perp$-space. In Eq. (7), the first term comes from collinear gluon radiation. The most important contribution comes from the $\ln(Q^2/q_\perp^2)$ term which is the well-known Sudakov double logarithm in the low transverse momentum limit. Because of the final state jet in the process, we also have a jet size dependent term, similar to dijet production studied in Ref. [20].

It is important to check the above asymptotic behavior against the fixed order calculations in the small transverse momentum limit $q_\perp \ll Q$. In Fig. 1, we plot the comparisons between Eq. (7) and those from the fixed order calculations. We show the $q_\perp$-dependent differential cross section in the low transverse momentum region for the typical kinematics at the LHC. We focus on the $gg \to H + Jet$ production channel, and the jet transverse momentum is in the range between 60 to 100 GeV. Both the Higgs boson and jet are produced in the central rapidity region, with $|y| < 0.5$, and the jet size is set to be $R = 0.5$. The full NLO calculation comes from the MCFM code [24], whereas the asymptotic result from Eq. (7). In the numeric calculations, we have adopted the CT10 PDF set [25]. From this plot, we can clearly see that the asymptotic behavior agrees well with the fixed order calculation in the low $q_\perp$ region. In the right panel of Fig. 1, we plot the comparison as function of the azimuthal angle difference between the Higgs boson and the leading jet $\phi = \phi_H - \phi_j$, which is more relevant for experiment measurement. The back-to-back correlation region corresponds to ϕ around π, where the total transverse momentum $q_\perp \to 0$. Because of this correspondence, again we find that the asymptotic expression agrees well with the fixed order calculations around $\phi = \pi$ and they are divergent. In this region, the QCD resummation is crucial to make reliable predictions, for which we will derive in the following sections.

One-loop calculation of $W(b)$. To calculate the complete one-loop result for $W(b_\perp)$, we have to take into account the following three contributions: (a) virtual graphs contribution to $gg \to Hg$; (b) real gluon contribution associated to the jet; (c), the collinear and soft gluon radiation contribution to finite $q_\perp$. Both (a) and (b) contribute to $\delta^{(2)}(q_\perp)$. All these contributions contain soft divergences, which have to be cancelled out. In the end, we only have collinear divergences associated with the incoming two gluons. To calculate (c), we have to Fourier transform the $q_\perp$-dependent expression of last section into $b_\perp$-space.

Calculations of virtual graphs are available in the literature, and they can be written as [6, 7, 26]

$$\frac{\alpha_s C_A}{2\pi}\left[-\frac{3}{\epsilon^2} + \frac{1}{\epsilon}\left(2\ln\frac{s}{\mu^2} + \frac{tu}{s\mu^2}\right) + \cdots\right] , \qquad (8)$$

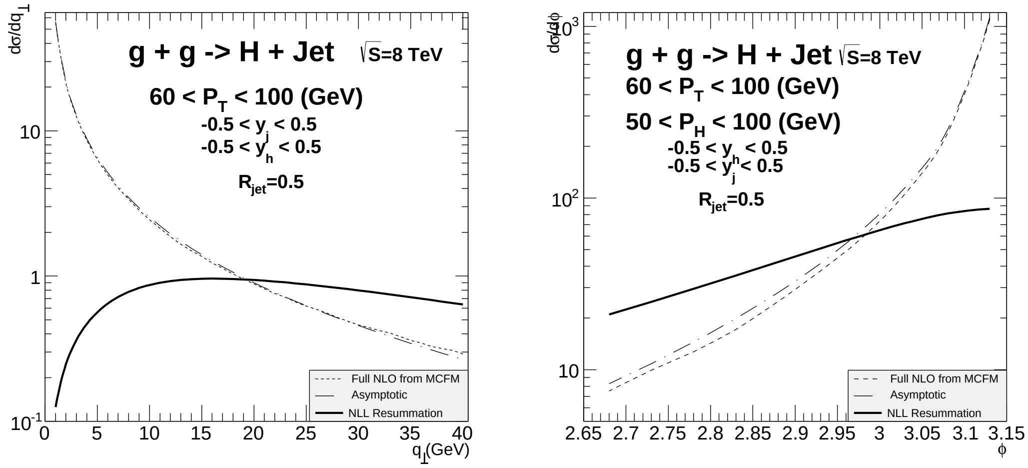

FIG. 1: Higgs boson plus jet production from $gg \to H + jet$ channel at the LHC as functions of total transverse momentum $q_\perp$ (left) and the azimuthal angle difference ϕ (right) of the Higgs boson and the leading jet. The dotted curves represent the result from the MCFM code, whereas the dashed curves from asymptotic result of Eq. (7). For comparison, the resummed cross sections are shown as the solid curves.

where for simplicity, we only kept the divergent terms. Collinear gluon associated with the jet is also easy to carry out, which will depend on the jet algorithm. Following the anti-k_t algorithm, we obtain the following contribution for the gluon jet at the one-loop order [22, 23]:

$$\frac{\alpha_s C_A}{2\pi}\left[\frac{1}{\epsilon^2}+\frac{1}{\epsilon}\left(2\beta_0-\ln\frac{P_T^2R^2}{\mu^2}\right)+\frac{1}{2}\ln^2\left(\frac{P_T^2R^2}{\mu^2}\right)-2\beta_0\ln\frac{P_T^2R^2}{\mu^2}+\frac{67}{9}-\frac{3}{4}\pi^2-\frac{23}{54}N_f\right]\ ,\ (9)$$

where the divergent and logarithmic terms are independent of jet algorithm, and the rest of the finite terms depend on the algorithm [27]. We note that the above result includes contributions from final state gluon splitting into a gluon pair or a quark-antiquark pair at the NLO, via the initial state gluon-gluon fusion processes. At the NLO, we also need to renormalize the effective ggH coupling, cf. Eq. (2), which yields the following contribution:

$$\frac{\alpha_s C_A}{2\pi}\left(\frac{Q^2}{\mu^2}\right)^{-\epsilon}\left(-\frac{3}{\epsilon}\right)2\beta_0\ . \tag{10}$$

Here, we have set the renormalization scale as Q^2 to simplify the final expression.

The Fourier transformation of the $q_\perp$-dependent expression, cf. Eq. (7), to the $b_\perp$-space, cf. $W(b_\perp)$, contains both double pole ($1/\epsilon^2$) and single pole ($1/\epsilon$) contributions in the dimensional regularization scheme. Among them, the soft divergence is cancelled out by those from Eqs.. (8) and (9). The collinear divergence in terms of $(1/\epsilon)\ln(1/R)$ associated with the final state jet is also cancelled out between the jet contributions from Eq. (9) and the Fourier transformation of Eq. (7) in $b_\perp$-space. In addition, the finite $\ln^2(P_T^2R^2/\mu^2)$ terms are cancelled out after summing over Eqs. (7), (8) and (9). The above cancellations provide important cross checks for our derivations. Finally, there are only divergences coming from the collinear divergences of the gluon distributions. After renormalizing the gluon distributions from the incoming hadrons, we obtain the finite contribution at one-loop order,

$$W^{(1)}(b) = H^{(0)}\frac{\alpha_s C_A}{2\pi}\left\{\ln\frac{b_0^2}{b^2\bar{\mu}^2}\left[\delta(\xi_2-1)\xi_1\mathcal{P}_{gg}(\xi_1)+(\xi_1\leftrightarrow\xi_2)\right]+\delta(\xi_1-1)\delta(\xi_2-1)\right.$$
$$\left.\times\left[-\left(\ln\frac{Q^2b_\perp^2}{b_0^2}\right)^2+\left(4\beta_0-\ln\frac{1}{R^2}\right)\ln\frac{Q^2b_\perp^2}{b_0^2}\right]\right\}+H^{(1)}\delta(\xi_1-1)\delta(\xi_2-1)\ , \quad (11)$$

where a common integral of the parton distributions as that in Eq. (7) is implicit but not shown. In the above result, the leading and sub-leading logarithmic terms are evident, and the remaining hard coefficient $H^{(1)}$ is

$$\begin{aligned} H^{(1)} \;=\; & H^{(0)}\frac{\alpha_s C_A}{2\pi}\left[\ln^2\left(\frac{Q^2}{P_T^2}\right)+2\beta_0\ln\frac{Q^2}{P_T^2R^2}+\ln\frac{1}{R^2}\ln\frac{Q^2}{P_T^2}-2\ln\frac{-t}{s}\ln\frac{-u}{s}\right.\\ & +\ln^2\left(\frac{\tilde{t}}{m_h^2}\right)-\ln^2\left(\frac{\tilde{t}}{-t}\right)+\ln^2\left(\frac{\tilde{u}}{m_h^2}\right)-\ln^2\left(\frac{\tilde{u}}{-u}\right)+2\mathrm{Li}_2\left(1-\frac{m_h^2}{Q^2}\right)\\ & \left.+2\mathrm{Li}_2\left(\frac{t}{m_h^2}\right)+2\mathrm{Li}_2\left(\frac{u}{m_h^2}\right)+\frac{67}{9}+\frac{\pi^2}{2}-\frac{23}{54}N_f\right]+\delta H^{(1)}\ , \end{aligned} \quad (12)$$

where $\tilde{t}=m_h^2-t$, $\tilde{u}=m_h^2-u$, and $\delta H^{(1)}$ represents terms not proportional to $H^{(0)}$ and can be found in Refs. [6, 7]. The above will enter into final resummation result as one-loop correction to Eq. (4).

TMD Factorization and Resummation. In order to carry out the resummation, we factorize the above one-loop results into the TMD parton distributions, and soft and hard factors, following the CSS procedure [17]. In $b_\perp$-space, this factorization can be written as

$$W\left(Q,b_\perp\right)=x_1\,f_g(x_1,b_\perp,\zeta,\mu_F,\rho)x_2f_g(x_2,b_\perp,\bar{\zeta},\mu_F,\rho)H^{TMD}_{gg\to Hg}(Q,\mu_F,\rho)S_{gg\to Hg}(b_\perp,\mu_F,\rho)\ ,$$

where we have followed Ji-Ma-Yuan scheme to define the TMD gluon distribution f_g [27]. In this scheme, an off-light-cone vector v ($\bar{v}$) is introduced to regulate the light-cone singularity, $\zeta^2=(2v\cdot P)^2/v^2$ (and $\bar{\zeta}^2=(2\bar{v}\cdot\bar{P})^2/\bar{v}^2$). The dependence on $\rho=(2v\cdot\bar{v})^2/v^2\bar{v}^2$ and the factorization scale μ_F cancel out among different factors. The TMD gluon distribution is the same as that describes the low transverse momentum Higgs boson production in hadronic collisions [27]. Therefore, we can use the results obtained there to carry out the resummations associated with the incoming gluon distributions. In particular, an evolution equation can be derived for the TMD distributions respect to ζ,

$$\frac{\partial}{\partial\ln\zeta}f_g(x,b_\perp,\zeta)=\left(K(b_\perp,\mu)+G(\zeta,\mu)\right)\times f_g(x,b_\perp,\zeta)\ , \quad (13)$$

where K and G are soft and hard parts in the evolution kernel. They obey the renormalization group equation with the so-called cusp anomalous dimension γ_K. The solution of the above evolution follows the CSS formalism, which resums the large logarithms of $\ln(\zeta^2b_\perp^2)$ [17, 27]. Final resummation results are obtained by setting $\zeta^2=\bar{\zeta}^2=\rho Q^2$. Additional resummation effects come from the soft factor, which is defined as

$$S_{gg\to Hg} \;=\; f_{abc}f_{a'b'c'}\langle 0|\mathcal{L}^\dagger_{vad}(b_\perp)\mathcal{L}_{\bar{v}be}(b_\perp)\mathcal{L}^\dagger_{ncf}(b_\perp)\mathcal{L}_{nc'f}(0)\mathcal{L}^\dagger_{\bar{v}b'e}(0)\mathcal{L}_{va'd}(0)|0\rangle\ . \quad (14)$$

This includes the soft gluon interactions between the final state jet defined in the n-direction (along the jet) and the incoming patrons defined by the two vectors v and $\bar{v}$. A renormalization group equation for the soft factor can be calculated, and the anomalous dimension is found to be $\gamma^{(s)} = \frac{\alpha_s C_A}{2\pi}\left(\ln \rho^2 + \ln \frac{1}{R^2}\right)$. Because of soft gluon interactions between the final state jet and incoming partons, the anomalous dimension depends on the jet size, which results into the D_1 coefficient in the resummation form factor Eq. (5).

By combining the evolutions of the TMD gluon distributions and the soft factor, we arrive at the final resummation results in Eq. (4). The coefficients in Eqs. (4) and (5) at one-loop order can be obtained from the results in previous section. As we mentioned above, the TMD gluon distributions in Higgs plus jet production are the same as those for inclusive Higgs production. Therefore, the A coefficients in Eq. (4) follow that in Higgs resummation, which solely come from the evolutions of the gluon distributions [17].

As an example, in Fig. 1, we show the resummation results, as compared to the fixed order calculations. When Fourier transforming the $b_\perp$-expression to obtain the transverse momentum distribution, we follow the b_* prescription of CSS resummation [17], and apply the non-perturbative form factors following the parameterizations in Refs. [28]. The final result is not sensitive to the choice of the non-perturbative form factor. In the numeric calculations, we have used parameters $A^{(1,2)}$, $B^{(1)}$, $D_1^{(1)}$ and $H^{(0,1)}$ in the resummation formula Eq. (4). All these coefficients are obtained from our one-loop calculation, except that of $A^{(2)}$, for which we take from the resummation for inclusive Higgs production $gg \to H$ [21]. We would like to emphasize that $H^{(1)}$ correction is of order 1 in the kinematics shown in Fig. 1, which highlights the importance of next-to-leading corrections. From these plots, we can clearly see that the resummations are important in the kinematic region where the fixed order calculations have singular behavior.

In the plots of Fig. 1, we include the NLO perturbative calculations (or, LO in $q_\perp$ distribution) in the comparisons. It would be desirable to compare to the NNLO calculations which are unfortunately not yet available. We notice that the inclusive cross section for Higgs plus one jet production has recently been calculated at NNLO in Ref. [10]. We hope that in the near future, these results can be made available to compute the differential cross sections in $q_\perp$ distribution or azimuthal angular distribution, from which we can further improve the theoretical predictions in Fig. 1. The existing NLO calculations for Higgs boson plus two jets production [24] could be added, after proper phase space integral and subtracting the double counting contribution, to improve the prediction for large $q_\perp$ or $\Delta\phi$ away from π region. It is however beyond the scope of this paper.

The preliminary experimental data on Higgs boson plus jet production have demonstrated the powerful reach for Higgs physics at the LHC [13–15]. These data have been mainly compared to parton shower Monte Carlo programs. The fixed order QCD calculations are divergent in the back-to-back azimuthal correlation region around $\Delta\phi \approx \pi$ (see Fig. 1), where all order resummation calculation is essential for making reliable theory predictions. The combination of the resummation technique derived in this paper and the fixed order calculations will provide high precision theory description of the associated production of Higgs boson and high energy jet events for the distributions shown in Fig. 1 and various other kinematical observables. This is of crucial importance for Higgs physics study at the LHC. We will carry out a detailed phenomenological studies along this line in the future.

Conclusions. In Summary, we have derived all order soft gluon resummation for Higgs boson plus high energy jet production. The expansion of our resummation formula agrees well with the fixed order calculations in the low transverse momentum region of Higgs boson and the jet, where we showed that resummation effects have to be included to have a reliable prediction.

Our derivations are based on a complete one-loop perturbative calculation. The results have been cross checked in various respects. It demonstrates that the final resummation formalism is consistent in the framework of CSS resummation. These results will provide important guidelines for future developments in electroweak boson plus jet production processes at the LHC. Extension to Higgs boson (or electroweak boson) plus two jets production shall be followed as well, which is a potential channel to investigate the unique production mechanism for Higgs boson at the collider.

In our calculations, we have applied the narrow jet approximation. This enables us to derive explicitly the one-loop analytical results and show the complete cancellation of the infrared divergences. This is important for demonstrating the factorization in the TMD formalism. The consistency in our derivation shall encourage future developments of a general set-up for the final state jet without the NJA, which is undoubtedly a more challenging calculation. Nevertheless, we expect the leading results will remain the same, which originate from soft gluons radiated out of the initial state partons and the final state jet.

This material is based upon work supported by the U.S. Department of Energy, Office of Science, Office of Nuclear Physics, under contract number DE-AC02-05CH11231, and by the U.S. National Science Foundation under Grant No. PHY-0855561 and PHY-1417326.

[1] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **716**, 1 (2012).
[2] S. Chatrchyan *et al.* [CMS Collaboration], Phys. Lett. B **716**, 30 (2012).
[3] S. Dawson, Nucl. Phys. B **359**, 283 (1991).
[4] A. Djouadi, M. Spira and P. M. Zerwas, Phys. Lett. B **264**, 440 (1991).
[5] D. de Florian, M. Grazzini and Z. Kunszt, Phys. Rev. Lett. **82**, 5209 (1999).
[6] V. Ravindran, J. Smith and W. L. Van Neerven, Nucl. Phys. B **634**, 247 (2002).
[7] C. J. Glosser and C. R. Schmidt, JHEP **0212**, 016 (2002).
[8] J. M. Campbell, R. K. Ellis and G. Zanderighi, JHEP **0610**, 028 (2006).
[9] J. M. Campbell, R. K. Ellis, R. Frederix, P. Nason, C. Oleari and C. Williams, JHEP **1207**, 092 (2012) [arXiv:1202.5475 [hep-ph]].
[10] R. Boughezal, F. Caola, K. Melnikov, F. Petriello and M. Schulze, JHEP **1306**, 072 (2013).
[11] S. Gangal and F. J. Tackmann, Phys. Rev. D **87**, no. 9, 093008 (2013).
[12] S. Dittmaier, S. Dittmaier, C. Mariotti, G. Passarino, R. Tanaka, S. Alekhin, J. Alwall and E. A. Bagnaschi *et al.*, arXiv:1201.3084 [hep-ph]; S. Heinemeyer *et al.* [LHC Higgs Cross Section Working Group Collaboration], arXiv:1307.1347 [hep-ph].
[13] G. Aad *et al.* [ATLAS Collaboration], JHEP **1409**, 112 (2014) [arXiv:1407.4222 [hep-ex]].
[14] G. Aad *et al.* [ATLAS Collaboration], Phys. Lett. B **738**, 234 (2014) [arXiv:1408.3226 [hep-ex]].
[15] G. Aad *et al.* [ATLAS Collaboration], arXiv:1408.7084 [hep-ex].
[16] I. W. Stewart and F. J. Tackmann, Phys. Rev. D **85**, 034011 (2012) [arXiv:1107.2117 [hep-ph]].
[17] J. C. Collins, D. E. Soper and G. F. Sterman, Nucl. Phys. B **250**, 199 (1985).
[18] A. Banfi, M. Dasgupta and Y. Delenda, Phys. Lett. B **665**, 86 (2008); A. Banfi and M. Dasgupta, JHEP **0401**, 027 (2004).
[19] A. H. Mueller, B. -W. Xiao and F. Yuan, Phys. Rev. D **88**, 114010 (2013).
[20] P. Sun, C. -P. Yuan and F. Yuan, arXiv:1405.1105 [hep-ph].
[21] D. de Florian and M. Grazzini, Phys. Rev. Lett. **85**, 4678 (2000).
[22] B. Jager, M. Stratmann and W. Vogelsang, Phys. Rev. D **70**, 034010 (2004).

[23] A. Mukherjee and W. Vogelsang, Phys. Rev. D **86**, 094009 (2012).

[24] J. M. Campbell and R. K. Ellis, Nucl. Phys. Proc. Suppl. **205-206**, 10 (2010); MCFM home page, http://mcfm.fnal.gov.

[25] J. Gao, M. Guzzi, J. Huston, H. -L. Lai, Z. Li, P. Nadolsky, J. Pumplin and D. Stump *et al.*, Phys. Rev. D **89**, 033009 (2014).

[26] C. R. Schmidt, Phys. Lett. B **413**, 391 (1997).

[27] X. Ji, J. P. Ma and F. Yuan, Phys. Rev. D **71**, 034005 (2005); JHEP **0507**, 020 (2005).

[28] F. Landry, R. Brock, P. M. Nadolsky and C. P. Yuan, Phys. Rev. D **67**, 073016 (2003); P. Sun, C. -P. Yuan and F. Yuan, Phys. Rev. D **88**, 054008 (2013); P. Sun, J. Isaacson, C.-P. Yuan and F. Yuan, arXiv:1406.3073 [hep-ph].