Low Transverse Momentum Heavy Quark Pair Production to Probe Gluon Tomography

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Abstract

We derive the transverse momentum dependent (TMD) factorization for heavy quark pair production in deep inelastic scattering, where the total transverse momentum is much smaller than the invariant mass of the pair. The factorization is demonstrated at one-loop order, in both Ji-Ma-Yuan and Collins-11 schemes for the TMD definitions, and the hard factors are calculated accordingly. Our result provides a solid theoretical foundation for the phenomenological investigations of the gluon TMDs in this process, and can be extended to other similar hard processes, including dijet (di-hadron) production in DIS.
Nucleon tomography in terms of various quark and gluon distribution functions has attracted great attention from both experiment and theory sides in recent years. Among these concepts, the transverse momentum dependent (TMD) parton distributions offer a clear picture of internal parton structure of the nucleon, and most importantly, they can be measured in hard QCD scattering processes. Together with the generalized parton distributions (GPDs), they reveal the nucleon tomography in most complete fashion, i.e., the three-dimension imaging of partons in nucleon. Theory and experiment developments toward this tomography has been one of focused topics in the planned electron-ion collider (EIC). The TMD quark distributions can be systematically investigated in the semi-inclusive hadron production in deep inelastic scattering processes (SIDIS), and a great deal have been learned from previous experiments, and will be mostly covered in the future experiments, such as JLab 12 GeV upgrade and the EIC. TMD quark distributions can also be studied in the Drell-Yan lepton pair production in pp collisions. On the other hand, the gluon TMDs are not easy to access. This is because photon does not couple to gluon at leading power in the hard scattering processes. It has been suggested to probe the gluon TMDs through two-particle productions in DIS or pp collision processes, such as heavy quark pair production in DIS, direct two photon production in pp collisions. The associated QCD dynamics for two-particle production has attracted intensive theoretical studies in recent years, including the factorization property for the hard processes and the universality of the parton distributions. It was generally believed that a TMD factorization shall apply to two-particle production in DIS processes. However, it has never been explicitly written down in a factorization form of TMD distributions. In the following, we will, for the first time, derive the TMD factorization formula for low transverse momentum heavy quark pair production. This result shall provide a solid theoretical foundation for phenomenological studies of gluon tomography in hard process, and open a new window for QCD studies for various other processes as well, such as dijet (di-hadron) production in DIS.

We focus on the heavy quark pair production in DIS process,

\[ \gamma^* + p \rightarrow c\bar{c}[M_{c\bar{c}}, p_\perp] + X, \]

where the transverse momentum of the pair \( p_\perp \) is much smaller than the invariant mass \( M_{c\bar{c}} \), and we keep the virtuality of the photon the same order as \( M_{c\bar{c}} \). In the following calculations, we denote incoming photon momentum as \( q \), \( P_A \) for nucleon (along \( +\hat{z} \) direction), \( k_1 \) and \( k_2 \) for the heavy quark and antiquark, respectively. We further introduce two dimensionless vectors \( n_c = k_1/m_c \) and \( n_{\bar{c}} = k_2/m_c \) to represent the directions of two final state particle: \( n_c^2 = n_{\bar{c}}^2 = 1 \). In addition, two light-like vectors are adopted: \( n = (1^-, 0^+, 0^\perp) \) and \( \bar{n} = (0^-, 1^+, 0^\perp) \), where \( \pm \) of a momentum is defined as \( P^\pm = (P^0 \pm P^z)/\sqrt{2} \). Therefore, \( P_A \) is dominated by its plus component. In this kinematics, the differential cross section is sensitive to the transverse momentum dependence of the gluon in the nucleon, and we can formulate a TMD factorization. Because the final state carries color, the naive TMD factorization has to be modified. From our following calculations, we find that an additional soft factor shall be included in the factorization formula. Therefore, the differential cross

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This can be generalized to real photon case as well, where a similar TMD factorization can be derived.
section can be written as

\[ \frac{d\sigma(\gamma^* p \rightarrow c\bar{c} + X)}{d^2p_{\perp}dM_{c\bar{c}}^2d\cos\theta} = \sigma_0 \int \frac{d^2b_{\perp}}{(2\pi)^2} e^{ip_{\perp} \cdot b_{\perp}} W_{c\bar{c}}(x, b_{\perp}) , \]  

(2)

where \( \sigma_0 \) represents the leading Born diagrams contribution,

\[ \sigma_0 = \frac{\pi^2M_{c\bar{c}}^2N_cC_F\alpha_s^2\alpha^2}{(N_c^2-1)(S+Q^2)Q^4} \left\{ \frac{t_1}{u_1} + \frac{u_1}{t_1} + \frac{4m_c^2Q^2}{t_1u_1}(1 - \frac{m_c^2}{t_1u_1}) \right\} \right. 
- \left. \frac{2Q^2}{t_1u_1}(Q^2 + \frac{2t_1u_1 - Q^4}{t_1u_1}m_c^2) - \frac{2Q^4}{t_1u_1} , \]  

(3)

here \( Q^2 = -q^2 \) with \( q \) the momentum of the virtual photon, \( \beta = \sqrt{1-4m_c^2/M_{c\bar{c}}^2} \), and the variable \( \tilde{Q} \) is defined as \( \tilde{Q}^2 = M_{c\bar{c}}^2 + Q^2 \). Then the longitudinal-momentum fraction for the incident gluon from nucleon \( x \) can be written as \( x = \tilde{Q}^2/(S + Q^2) \) with the photon-hadron center-of-mass energy \( S = (q + P_A)^2 \). In the above expression, \( \theta \) is the scattering angle between final heavy quark and nucleon in the photon-hadron center-of-mass frame. \( t_1 = (k_1 - xP_A)^2 - m_c^2 \), and \( u_1 = (k_2 - xP_A)^2 - m_c^2 \). In the impact parameter \( b_{\perp} \)-space, the TMD factorization for \( W_{c\bar{c}} \) can be written as

\[ W_{c\bar{c}}(x, b_{\perp}) = H(\tilde{Q}, \mu)xg(x, b_{\perp}, \tilde{Q}, \mu)S(b_{\perp}, \mu) , \]  

(4)

where the hard factor \( H \), the soft factor \( S \), and the TMD gluon distribution all depend on the factorization scale \( \mu \). Schematically, this factorization can be viewed as in Fig. 1, where the photon scatters off gluon to produce the heavy quark pair (hard factor), and the transverse momentum of the final state comes from the gluon distribution (lower part) and the soft gluon radiation from the final state. Compared to the SIDIS, we will find that in the above factorization we have the soft function \( S \) instead of the TMD fragmentation function. Naive TMD definitions for the gluon distribution and soft function \( S \) contain the light-cone singularities, and a regulation introduces the scheme dependence. In the following, we will show calculations in two schemes: Ji-Ma-Yuan scheme \[10\] and Collins-11 scheme \[11\]. We first derive the soft factor in this process, and calculate the associated hard factors and the gluon TMDs. In the end, we will derive the resummation formalism. We will find that the final results do not depend on the schemes.
II. SOFT GLUON RADIATION

A key point to demonstrate the QCD factorization is to show that the leading power gluon radiation can be included into the various factors in the factorization formula. In this analysis, a power counting method is crucial to achieve the final factorization result. For example, the gluon radiation associated with the incoming gluon contributes to the collinear and soft gluon part, which can be absorbed into the gluon distribution and soft factor in the factorization formula. This part of contribution is similar to those in the SIDIS and low transverse momentum Drell-Yan lepton pair production in pp collisions. Now, we turn to the final state radiation. Because of heavy quark mass, soft radiation from the quark pair contributes to the leading power of the differential cross section. The resummation of all order soft gluon radiation associated with heavy quark moving in $n_c$-direction can be summarized into a Wilson line in that direction. This has also been applied to formulate the threshold resummation in heavy quark pair production [12, 13]. By applying this technique, we can summarize the gluon radiation from the initial state gluon and the final state heavy quark pair as three kinds of Wilson lines. For final-state quark and anti-quark, we have

$$L_{n_c}(\xi) = P \exp \left( -ig \int_0^{+\infty} d\lambda n_c \cdot A(\lambda n_c + \xi) \right),$$

$$L_{n_{\bar{c}}}(\xi) = P \exp \left( ig \int_0^{+\infty} d\lambda n_{\bar{c}} \cdot A(\lambda n_{\bar{c}} + \xi) \right),$$

which are in the fundamental representation, $A^{\mu} = A_c^{\mu} T^c$. While for the soft gluon radiation from the incoming gluon, we have

$$L_{\bar{v}}(\xi) = P \exp \left( ig \int_0^{+\infty} d\lambda \bar{v} \cdot A(\lambda \bar{v} + \xi) \right),$$

which is in the adjoint representation, $A^{\mu} = -if_{abc} A_b^{\mu}$. The vector $\bar{v}$ is along the momentum $P_A$. The soft function in the factorization formula Eq. (4) contains contributions from the above three Wilson lines. For the process of Eq. (1), there is only one color base combination between them, and the soft function can be written as,

$$\bar{S}(b_\perp, \mu, \rho) = \frac{\int_0^\pi \frac{(\sin \phi)^{-2\epsilon}}{a_1} d\phi \langle 0 | L^\dagger_{\bar{v}ca}(b_\perp) \text{Tr} \left[ L^\dagger_{n_c}(b_\perp) T^a L^\dagger_{n_{\bar{c}}}(b_\perp) L_{n_{\bar{c}}}(0) T^a L_{n_c}(0) \right] L_{\bar{v}ac}(0) | 0 \rangle }{\text{Tr}[T^d T^d]},$$

with $a_1 = \frac{\sqrt{\pi(2-2\epsilon)}}{\Gamma(1-\epsilon)}$ in the $(4-2\epsilon)$ dimension. In the above definition, we have integrated out the azimuthal angle $\phi$ between $n_{c,\perp}$ and $p_{\perp}$. Intuitively, $\bar{v}$ shall be chosen along the momentum direction of the nucleon $P_A$ with $\bar{v} = \bar{n}$. However, this definition contains a light-cone singularity. Regulating this singularity introduces the scheme-dependence. In our calculations, we follow two different schemes: Ji-Ma-Yuan [10] and Collins-11 [11]. We will show that the factorization works for both schemes, and the hard factors can be calculated accordingly.

\[\text{There is no light-cone singularity associated with } L_{n_c,n_{\bar{c}}} \text{ because of } n_{c,\bar{c}}^2 \neq 0.\]
In the Ji-Ma-Yuan scheme, the light-cone gauge link is chosen slightly off-light-cone: $\bar{v}^2 \neq 0$ with $\bar{v}^+ \gg \bar{v}^-$. In Fig. 2, we plot the real gluon radiation contribution to the soft factor. Similarly, we have virtual diagrams. Adding them together, we find that, at one-loop order,

$$
\overline{S}^{(1)}_{JMY}(b_\perp, \mu, \rho) = \frac{\alpha_s}{2\pi} \left\{ C_A \ln \frac{c_0^2}{b_1^2 \mu^2} \left( B_{final} + \ln \rho^2 + \ln \frac{\tilde{Q}^2}{\zeta^2} - 1 \right) + C_{final} \right\},
$$

(8)

where $c_0 = 2e^{-\gamma_E}$, $\rho^2 = (2v \cdot \bar{v})^2/v^2\bar{v}^2$, $\zeta^2 = x^2(2v \cdot P_A)^2/v^2$, and $v$ is another non-light-like vector $v^2 \neq 0$ with $v^- \gg v^+$. B and C functions are defined as

$$
B_{final} = \frac{1}{2N_c} \frac{1 + \beta^2}{\beta} \ln \frac{1 - \beta}{1 + \beta} - \frac{2C_F}{N_c} + \ln \frac{t_1u_1}{\tilde{Q}^2 m_c^2},
$$

$$
C_{final} = \frac{1}{2N_c} f_{c\bar{c}} + 2C_F \ln \frac{t_1u_1}{M_{c\bar{c}}^2 m_c^2} - C_A \text{Li}_2(1 - \frac{t_1u_1}{M_{c\bar{c}}^2 m_c^2}).
$$

(9)

The large $N_c$ suppressed term $f_{c\bar{c}}$ can be further decomposed into

$$
f_{c\bar{c}} = -\frac{(\beta^2 + 1)}{2\beta} f_{c\bar{c}}^a + \frac{2(\beta^2 + 1)(1 - \beta \cos \theta)}{1 - \beta^2} f_{c\bar{c}}^b,
$$

(10)

where the factor $f_{c\bar{c}}^a$ can be written as

$$
f_{c\bar{c}}^a = \left( \ln \frac{b_1}{b_4} \right)^2 - \left( \ln \frac{b_3}{b_2} \right)^2 + 2\ln \frac{b_3b_4}{b_1b_2} \ln \frac{b_1 \cot \frac{\theta}{2}}{b_2} + 2 \left( \text{Li}_2(\frac{b_2}{b_4}) + \text{Li}_2(\frac{b_4}{b_1}) - \text{Li}_2(\frac{b_2}{b_3}) - \text{Li}_2(\frac{b_3}{b_1}) \right),
$$

(11)

with $b_1 = 1 + \beta$, $b_2 = 1 - \beta$, $b_3 = 1 + \beta \cos \theta$, $b_4 = 1 - \beta \cos \theta$. The analytic expression of $f_{c\bar{c}}^b$ is complicated to present in the paper, so we will reserve the integral

$$
f_{c\bar{c}}^b = \int_0^1 \frac{\ln(ay^2 + 1)}{-ay^2 + y(a + c - 1) + 1},
$$

(12)
with \( a = \beta^2 \sin^2 \theta / (b_1 b_2) \) and \( c = b_2^2 / (b_1 b_2) \).

Similarly, the non-light-like vector \( v \) defined above was introduced to regulate the light-cone singularity in the TMD gluon distribution in Ji-Ma-Yuan scheme, for which we have \(^{[14]}\)

\[
xg^{\text{unsub.}}(x, k_\perp, \mu, \zeta, \rho) = \int \frac{d\xi d^2\xi_\perp}{P^2(2\pi)^3} e^{-ixP^+\xi^- + i\vec{k}_\perp \cdot \vec{\xi}_\perp}
\times \left\langle P | F^+_{\mu}(\xi^- , \xi_\perp) L^{(1)}_{\text{vab}}(\xi^- , \xi_\perp) L_{\text{vbc}}(0, 0, \mu) F^+(0)|P \right\rangle ,
\]

where the associated gauge link is in adjoint representation. The above gluon distribution contains not only collinear gluon contribution, but also the soft gluon contribution, which is defined as

\[
S^{a,\bar{v}}(b_\perp) = \langle 0 | L^\dagger_{\text{vb}a}(b_\perp) L^\dagger_{\text{v} a}(b_\perp) L_{\text{vb}c}(0) L_{\text{v} c}(0)|0 \rangle / (N_c^2 - 1). \tag{14}
\]

Therefore the gluon distribution after subtraction is defined as

\[
xg_{\text{JMY}}(x, b_\perp, \mu, \zeta, \rho) = xg^{\text{unsub.}}(x, b_\perp, \mu, \zeta, \rho) / S^{a,\bar{v}}(b_\perp), \tag{15}
\]

which will enter into the factorization formula in Eq. (4).

To demonstrate the TMD factorization at one-loop order, we calculate the differential cross section on an on-shell gluon target. In the perturbative calculations, we take the leading power contribution in the limit \( \bar{Q} \gg q_\perp \). It can be shown, in this limit, diagram by diagram that they can be factorized into the TMD gluon distribution, soft factor and the hard factor as in the factorization formula. In particular, the one-loop result for the TMD gluon distribution has been calculated in Ref. \(^{[14]}\) in Ji-Ma-Yuan scheme. By subtracting the gluon distribution, we find that the hard factor can be written as

\[
H^{(1)}_{\text{JMY}}(\mu, \rho) = \frac{\alpha_s}{\pi} C_A \left\{ \left( \beta_0 - \frac{B_{f,\text{final}}}{2} \right) + \ln \frac{\rho}{2} - \frac{3}{4} \ln \frac{\bar{Q}^2}{\mu^2} + \frac{\ln^2 \rho}{4} - \frac{3}{4} \ln \rho + \frac{\pi^2}{6} + \frac{7}{4} + B^{V,g_\gamma}_{f} \right\},
\]

where we have taken \( \zeta^2 = \rho \bar{Q}^2 \) for convenience, and \( B^{V,g_\gamma}_{f} \) comes from the finite contribution of the virtual diagrams.

### III. COLLINS-11 SCHEME

Subtraction of the light-cone singularity is essential to establish the TMD factorization. Collins introduced a subtraction scheme where the parton distribution and soft factor do not contain light-cone singularity \(^{[11]}\). According to this new scheme, the TMD gluon is defined as

\[
xg_{\text{JCC}}(x, b_\perp, \mu, \zeta_c) = xg^{\text{unsub.}}(x, b_\perp) \sqrt{\frac{S^{a,v}(b_\perp)}{S^{a,v}(b_\perp) S^{a,v}(b_\perp)}} , \tag{17}
\]

where \( \zeta_c = x^2(2v \cdot P_A)^2 / v^2 = 2(xP_A^+)^2 e^{-2y_n} \) with \( y_n \) the rapidity cut-off in Collins-11 scheme. Calculating up to one-loop order, we have

\[
xg^{(1)}_{\text{JCC}}(x, b_\perp, \mu, \zeta_c) = \frac{\alpha_s}{2\pi} C_A \left\{ 2 \left( \frac{1}{\epsilon} + \ln \frac{\zeta_c^2}{b_\perp^2 \mu^2} \right) \mathcal{P}_{g\to g}(x)
+ \delta(1 - x) \left\{ 2\beta_0 \ln \frac{b_\perp^2 \mu^2}{c_0^2} + \frac{1}{2} \left( \ln \frac{\zeta_c^2}{\mu^2} \right)^2 - \frac{1}{2} \left( \ln \frac{\zeta_c^2 b_\perp^2}{c_0^2} \right)^2 \right\} \right\}, \tag{18}
\]
where the gluon splitting kernel $P_{gg}(x) = \frac{x}{(1-x)_+} + \frac{x}{x} + x(1-x) + \delta(x-1)\beta_0$ with $\beta_0 = (11N_c - 2n_f)/(12N_c)$.

The soft factor can be defined similarly,

$$S_{\text{JCC}}(b_\perp, \mu) = \frac{\alpha_s}{2\pi} \left\{ C_A \ln \frac{\bar{Q}_0^2}{b_\perp^2\mu^2} (B_{\text{final}} + \ln \frac{\bar{Q}_0^2}{c_0^2}) + C_{\text{final}} \right\} ,$$

at one-loop order. From the above results, we immediately obtain the hard factor as

$$H_{\text{JCC}}^{(1)}(\mu) = \frac{\alpha_s}{\pi} \frac{C_A}{\pi} \left\{ (\beta_0 - B_{\text{final}}^2/2) \ln \frac{\bar{Q}_0^2}{\mu^2} - \frac{1}{4} \ln^2 \frac{\bar{Q}}{\mu} - \frac{\pi^2}{12} + B_{\gamma} V_f \right\} ,$$

where $\zeta_c$ has been chosen as $\bar{Q}$ in the above calculation.

**IV. RESUMMATION**

The large logarithms in the fixed order perturbative calculations as we have shown above can be resummed by applying the Collins-Soper-Sterman resummation \cite{15}. In particular, in this case, we can derive the energy evolution equation for the TMD gluon distribution, and the renormalization group equation for the soft and hard factors. By solving these equations, we resum the large logarithms. The final expression for $W(b_\perp, \bar{Q})$ can be written as

$$W(x, b_\perp, \bar{Q}^2) = g(x, b_\perp, \bar{Q}_0, \bar{Q}_0) S(b_\perp, \bar{Q}_0) H(\bar{Q}, \bar{Q}) e^{-S_{\text{Sud}}(\bar{Q}, \bar{Q}_0)} ,$$

where $\bar{Q}_0$ is chosen such that the intrinsic TMD gluon distribution at the input scale $\bar{Q}_0$. All the large logarithms is included in the Sudakov form factor,

$$S_{\text{Sud}} = -\int_{\bar{Q}_0}^{\bar{Q}} \frac{d\mu}{\mu} \left( \ln \frac{\bar{Q}}{\mu} \gamma_K(\mu) - \gamma_S(\mu, 1) + \frac{\alpha_s C_A}{\pi} (1 - 2\beta_0 - \ln \frac{\bar{Q}_0^2 b_\perp^2}{c_0^2}) \right) ,$$

where

$$\gamma_K(\mu) = \frac{2\alpha_s(\mu) C_A}{\pi} ,$$

$$\gamma_S(\mu, \rho) = -\frac{\alpha_s(\mu) C_A}{\pi} (B_{\text{final}} + \ln \rho - 1) .$$

We notice that the $\rho$-dependence cancels out in the above Sudakov form factor, as well as in the final expression in the resummed form. In particular, in the Sudakov factor $S_{\text{Sud}}$, $\gamma_S$ takes value at $\rho = 1$ after the resummation. We would like to emphasize that the final results agree with each other between the Ji-Ma-Yuan and Collins-11 schemes of the TMD definitions.
The above factorization results can be carried out for the spin-dependent observables in this process as well, in particular, for the single transverse spin asymmetry from the gluon Sivers function, which will have a similar factorization formula with the soft function defined above. The associated single spin asymmetry can be written as \( S_\perp \times p_\perp \) where \( S_\perp \) is the transverse spin vector and \( p_\perp \) defined above as the total transverse momentum of the heavy quark pair. This tells that the heavy quark pair production in DIS at the planned EIC will provide important information on the gluon Sivers function \([1]\).

In addition, our results can also be extended to the linearly polarized gluon distribution contribution in the above process \([3]\). However, because this term is proportional to \( \cos 2\phi \) where \( \phi \) is the azimuthal angle between \( k_{1\perp} \) and \( p_\perp \), the soft function \( S \) of Eq. (7) will have to be modified to include explicit dependence on \( \phi \). Nevertheless, a TMD factorization can be formulated for this case as well.

V. CONCLUSION

In this paper, by studying one-loop correction to heavy quark pair production at low transverse momentum, we have derived the associated TMD factorization formalism in DIS process. The light-cone singularity regulation in the TMD and soft factor has been performed in both Ji-Ma-Yuan and Collins-11 schemes, and the hard factors were calculated accordingly. The final results have been shown to be independent of the regulation schemes. The factorization result of this paper should provide an important step further to investigate the gluon tomography in nucleon in hard processes. A number of extensions shall follow, including QCD factorization studies for di-jet (di-hadron) production in DIS, heavy quark pair and di-jet production in \( pp \) collisions. In particular, a recent calculation of \( t\bar{t} \) pair production in \( pp \) collision has been analyzed in the soft-collinear-effective theory \([16]\). A detailed comparison with the current calculation showed that they are consistent with each other. We will address these issues in future publications.

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