Next-to-Next-to-Leading-Order Logarithmic Corrections at Small Transverse Momentum in Hadronic Collisions

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We study the region of small transverse momenta in $q\bar{q}$- and gg-initiated processes with no colored particle detected in the final state. We present the universal expression of the $\mathcal{O}(\alpha_s^2)$ logarithmically enhanced contributions up to next-to-next-to-leading-order logarithmic accuracy. From there we extract the coefficients that allow the resummation of the large logarithmic contributions. We find that the coefficient known in the literature as $B^{(2)}$ is process dependent, since it receives a hard contamination from the one-loop correction to the leading-order subprocess. We present the general result of $B^{(2)}$ for both quark and gluon channels.

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The process in which a system of not strongly interacting particles of large invariant mass Q^2 [lepton pairs, gauge boson(s), Higgs boson, and so forth] is produced in hadronic collisions is a well studied subject in perturbative QCD [1]. At transverse momenta q_T^2 of order of Q^2 the cross section can be computed by using the standard QCD-improved parton model. When q_T becomes small, the simple perturbative picture is spoiled. This happens because large logarithmic corrections of the form $\log \frac{Q^2}{q_T^2}$ arise due to an incomplete cancellation of soft and collinear singularities between real and virtual contributions. These large logarithmic corrections can be resummed to all orders by using the Collins-Soper-Sterman (CSS) formalism [2].

We consider the class of inclusive hard scattering processes,

$$h_1 h_2 \to A_1 + A_2 \dots A_n + X\,, \qquad (1)$$

where the collision of the hadrons h_1 and h_2 produces a system of not strongly interacting final state particles $A_1 \dots A_n$ carrying total momentum Q and total transverse momentum q_T. According to the CSS formula, and neglecting terms which are finite in the limit $q_T \to 0$, the cross section can be written as (it is assumed that all other dimensionful invariants are of the same order Q^2)

$$\frac{d\sigma}{dq_T^2 dQ^2 d\phi} = \sum_{a,b,c} \int_0^1 dx_1 \int_0^1 dx_2 \int_0^\infty db\, \frac{b}{2} J_0(bq_T) \frac{d\sigma_{c\bar{c}}^{(\mathrm{LO})}}{d\phi} \delta(Q^2 - x_1 x_2 s) \times (f_{a/h_1} \otimes C_{ca})\left(x_1, \frac{b_0^2}{b^2}\right)(f_{b/h_2} \otimes C_{\bar{c}b})\left(x_2, \frac{b_0^2}{b^2}\right) S_c(Q,b)\,, \qquad (2)$$

where $d\phi = dPS(Q \to q_1, q_2, \dots, q_n)$ represents the phase space of the system of noncolored particles, $b_0 = 2e^{-\gamma_e}$, and $\sigma_{c\bar{c}}^{(\mathrm{LO})}$ is the leading-order cross section (i.e., with no final state partons and therefore $q_T = 0$) for the given process ($c, \bar{c}$ can be either $q_f, \bar{q}_{f'}$ or g, g). The function C_{ab} in Eq. (2) is a process-dependent coefficient function, $J_0(bq_T)$ is the Bessel function of the first kind, and $f_{i/h}$ corresponds to the distribution of a parton i in a hadron h. The large logarithmic corrections are exponentiated in the Sudakov form factor

$$S_c(Q,b) = \exp\left\{ -\int_{b_0^2/b^2}^{Q^2} \frac{dq^2}{q^2} \left[A_c(\alpha_s(q^2)) \ln \frac{Q^2}{q^2} + B_c(\alpha_s(q^2)) \right] \right\}. \qquad (3)$$

The functions A_c, B_c, and C_{ab} in Eqs. (2) and (3) have perturbative expansions in α_s,

$$A_c(\alpha_s) = \sum_{n=1}^\infty \left(\frac{\alpha_s}{2\pi}\right)^n A_c^{(n)}, \qquad (4)$$

$$B_c(\alpha_s) = \sum_{n=1}^\infty \left(\frac{\alpha_s}{2\pi}\right)^n B_c^{(n)}, \qquad (5)$$

$$C_{ab}(\alpha_s, z) = \delta_{ab}\delta(1 - z) + \sum_{n=1}^\infty \left(\frac{\alpha_s}{2\pi}\right)^n C_{ab}^{(n)}(z)\,. \qquad (6)$$

In order to obtain the coefficients in Eqs. (4) and (5) at a given order, the differential cross section at small q_T has to be computed at the same order. A comparison with the power expansion in α_s of the resummed result in Eq. (2) allows one to extract the coefficients that control the resummation of the large logarithmic terms.

In this Letter we study the behavior of cross sections at small transverse momenta at second order in α_s in both the quark and the gluon channels. We find that the analytic form of the logarithmically enhanced contributions can be computed perturbatively in a universal manner by using the recent knowledge on the infrared behavior of tree-level [3,4] and one-loop [5] QCD amplitudes. In this way, we are able to extract the coefficients $A_c^{(1)}$, $B_c^{(1)}$, $C_{ab}^{(1)}$, $A_c^{(2)}$, and $B_c^{(2)}$ for *any* $q\bar{q}$ or gg initiated process in the class

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(1). Details on our calculation will be given elsewhere [6]. Here we present and discuss only our main results.

By following Ref. [7] we multiply the differential cross section, calculated at parton level, by q_T^2 and take moments with respect to $z = Q^2/s$ defining the dimensionless quantity:

$$\Sigma(N) = \int dz\, z^N \frac{q_T^2 Q^2}{d\sigma_0/d\phi} \frac{d\sigma}{dq_T^2\, dQ^2\, d\phi}. \quad (7)$$

In the quark channel, for the sake of simplicity and in order to compare our result for $\Sigma(N)$ to the one originally obtained for Drell-Yan in Ref. [7], we restrict our attention to the *nonsinglet* contribution to the cross section defined by

$$\sigma^{\mathrm{NS}} = \sum_{ff'} (\sigma_{q_f \bar{q}_{f'}} - \sigma_{q_f q_{f'}}). \quad (8)$$

To have $q_T \neq 0$ at least one gluon must be emitted; thus $\Sigma(N)$ has the expansion

$$\Sigma(N) = \frac{\alpha_s}{2\pi} \Sigma^{(1)}(N) + \left(\frac{\alpha_s}{2\pi}\right)^2 \Sigma^{(2)}(N) + \cdots. \quad (9)$$

In the following we will systematically neglect in $\Sigma(N)$ all contributions that vanish as $q_T \to 0$.

In order to compute the small q_T behavior of $\Sigma(N)$ our strategy is as follows. The singular behavior at small q_T is dictated by the infrared (soft and collinear) structure of the relevant QCD matrix elements. At $\mathcal{O}(\alpha_s)$ this structure has been known for long time [3]. Recently, the universal functions that control the soft and collinear singularities of tree-level and one-loop QCD amplitudes at $\mathcal{O}(\alpha_s^2)$ have been computed [4,5].

By using this knowledge, and exploiting the simple kinematics of the leading-order subprocess, we were able to construct *improved* factorization formulas that allow us to control *all* infrared singular regions avoiding any problem of double counting [6]. We have used these improved formulas to approximate the relevant matrix elements and compute the small q_T behavior of $\Sigma(N)$ in a completely universal manner.

The calculation at $\mathcal{O}(\alpha_s)$ is straightforward, and we recover the well-known results,

$$\Sigma^{(1)}_{q\bar{q}}(N) = 2C_F \log\frac{Q^2}{q_T^2} - 3C_F + 2\gamma^{(1)}_{qq}(N) \quad (10)$$

and

$$\Sigma^{(1)}_{gg}(N) = 2C_A \log\frac{Q^2}{q_T^2} - 2\beta_0 + 2\gamma^{(1)}_{gg}(N). \quad (11)$$

Here $C_F = \frac{N_c^2-1}{2N_c}$, $C_A = N_c$, and $T_R = 1/2$ are the $\mathrm{SU}(N_c)$ QCD color factors; $\beta_0 = \frac{11}{6}C_A - \frac{2}{3}n_f T_R$ and $\gamma^{(1)}_{qq}(N)$, $\gamma^{(1)}_{gg}(N)$ are the quark and gluon one-loop anomalous dimensions, respectively. From Eqs. (10) and (11) one obtains

$$A^{(1)}_a = 2C_a, \qquad B^{(1)}_a = -2\gamma_a, \qquad a = q, g, \quad (12)$$

where C_a and γ_a are the coefficients of the leading $(1-z)^{-1}$ singularity and $\delta(1-z)$ term in the one-loop Altarelli-Parisi kernels P_{aa}, respectively,

$$C_q = C_F, \qquad C_g = C_A,$$
$$\gamma_q = \frac{3}{2} C_F, \qquad \gamma_g = \beta_0. \quad (13)$$

At this order it is possible to obtain also the coefficient $C^{(1)}_{ab}$ by considering the q_T integrated distribution and including the renormalized virtual correction to the LO amplitude $c\bar{c} \to A_1 + A_2 \ldots A_n$, summed over spins and colors, which, at $\mathcal{O}(\epsilon^0)$, can be written as

$$\mathcal{M}^{(0)\dagger}_{c\bar{c}}(\phi)\mathcal{M}^{(1)}_{c\bar{c}}(\phi) + \text{c.c.} = \frac{\alpha_s}{2\pi}\left(\frac{4\pi\mu^2}{Q^2}\right)^{\epsilon} \frac{\Gamma(1-\epsilon)}{\Gamma(1-2\epsilon)} \left(-\frac{2C_c}{\epsilon^2} - \frac{2\gamma_c}{\epsilon} + \mathcal{A}_c(\phi)\right) |\mathcal{M}^{(0)}_{c\bar{c}}(\phi)|^2. \quad (14)$$

(All our results are obtained using the factorization and renormalization prescriptions of the $\overline{MS}$ scheme and within the framework of conventional dimensional regularization.) In Eq. (14) the structure of the poles in $\epsilon = (4-d)/2$ is universal [8] and fixed by the flavor of the incoming partons. The *finite* part $\mathcal{A}$ (which can depend on the kinematics of the final state noncolored particles) depends instead on the particular process in the class (1) we want to consider. In the case of Drell-Yan we have [9]

$$\mathcal{A}^{\mathrm{DY}}_q = C_F\left(-8 + \frac{2}{3}\pi^2\right), \quad (15)$$

whereas for Higgs production in the $m_{\mathrm{top}} \to \infty$ limit the finite contribution is [10]

$$\mathcal{A}^H_g = 5C_A + \frac{2}{3}C_A\pi^2 - 3C_F \equiv 11 + 2\pi^2. \quad (16)$$

By using the information in Eq. (14) we obtain, for $C^{(1)}_{ab}$,

$$C^{(1)}_{ab}(z) = -\hat{P}^{\epsilon}_{ab}(z) + \delta_{ab}\delta(1-z) \times \left(C_a\frac{\pi^2}{6} + \frac{1}{2}\mathcal{A}_a(\phi)\right), \quad (17)$$

where $\hat{P}^{\epsilon}_{ab}(z)$ is the $\mathcal{O}(\epsilon)$ term in the Altarelli-Parisi $\hat{P}_{ab}(z,\epsilon)$ splitting kernel, given by

$$\begin{aligned} \hat{P}^{\epsilon}_{qq}(z) &= -C_F(1-z), \\ \hat{P}^{\epsilon}_{gq}(z) &= -C_F z, \\ \hat{P}^{\epsilon}_{qg}(z) &= -2T_R z(1-z), \\ \hat{P}^{\epsilon}_{gg}(z) &= 0. \end{aligned} \quad (18)$$

At order α_s the coefficients $A^{(1)}_a$ and $B^{(1)}_a$ are fully determined by the *universal* Altarelli-Parisi splitting functions. The function $C^{(1)}_{ab}$ depends instead on the process through the one-loop corrections to the LO matrix element. The general expression in Eq. (17) reproduces correctly the coefficient $C_{ab}^{(1)}$ for Drell-Yan [7], Higgs production in the $m_{\rm top} \to \infty$ limit [11], and $\gamma\gamma$ production [12]. (The coefficient $C_{qq}^{(1)}$ for ZZ production in Ref. [13] is not correct.)

At second order in α_s, two different contributions to $\Sigma^{(2)}(N)$ have to be considered: the real correction corresponding to the emission of one extra parton (i.e., two gluons or a $q\bar{q}$ pair) with respect to the $\mathcal{O}(\alpha_s)$ contribution, and its corresponding virtual correction.

The double-real emission contribution is the most difficult to compute. One has to integrate over the phase space of the two unresolved final state partons keeping q_T fixed and finally perform the z integration in Eq. (7). We find that, likewise $\Sigma^{(1)}(N)$, this contribution to $\Sigma^{(2)}(N)$ is *process independent;* i.e., it does not depend on the particular process in the class (1) we want to consider.

The virtual contribution is simpler to compute, and we find it to be *process dependent*. More importantly, its process dependence is fully determined by the function $\mathcal{A}$ appearing in the one-loop correction to the LO subprocess [see Eq. (14)].

In the following, for the sake of simplicity, we present the total results for $\Sigma^{(2)}(N)$ corresponding to the choice of the factorization and renormalization scales fixed to Q^2. Since we are interested in extracting the coefficients $A_{q,g}^{(2)}$ and $B_{q,g}^{(2)}$, as in the $\mathcal{O}(\alpha_s)$ case we concentrate on the *diagonal* $q\bar{q}$ and gg contributions to $\Sigma^{(2)}(N)$.

In the quark (nonsinglet) channel we obtain

$$\begin{aligned}
\Sigma_{q\bar{q}}^{(2)}(N) = {} & \log^3\frac{Q^2}{q_T^2}[-2C_F^2] + \log^2\frac{Q^2}{q_T^2}[9C_F^2 + 2C_F\beta_0 - 6C_F\gamma_{qq}^{(1)}(N)] \\
& + \log\frac{Q^2}{q_T^2}\bigg[C_F^2\left(\frac{2}{3}\pi^2 - 7\right) + C_FC_A\left(\frac{35}{18} - \frac{\pi^2}{3}\right) - \frac{2}{9}C_Fn_fT_R + 2C_F\mathcal{A}_q(\phi) + (2\beta_0 + 12C_F)\gamma_{qq}^{(1)}(N) \\
& \qquad - 4[\gamma_{qq}^{(1)}(N)]^2 + 4C_F^2\left(\frac{1}{(N+1)(N+2)} - \frac{1}{2}\right)\bigg] \\
& + \bigg[C_F^2\left(-\frac{15}{4} - 4\zeta(3)\right) + C_FC_A\left(-\frac{13}{4} - \frac{11}{18}\pi^2 + 6\zeta(3)\right) - 3C_F\mathcal{A}_q(\phi) + C_Fn_fT_R\left(1 + \frac{2}{9}\pi^2\right) \\
& \quad + 2\gamma_{(-)}^{(2)}(N) + 2C_F\gamma_{qq}^{(1)}(N)\left(\frac{\pi^2}{3} + 2\frac{1}{(N+1)(N+2)}\right) + 2\gamma_{qq}^{(1)}(N)\mathcal{A}_q(\phi) \\
& \quad - 2C_F(\beta_0 + 3C_F)\left(\frac{1}{(N+1)(N+2)} - \frac{1}{2}\right)\bigg], \qquad (19)
\end{aligned}$$

whereas for the gluon channel the result is

$$\begin{aligned}
\Sigma_{gg}^{(2)}(N) = {} & \log^3\frac{Q^2}{q_T^2}[-2C_A^2] + \log^2\frac{Q^2}{q_T^2}[8C_A\beta_0 - 6C_A\gamma_{gg}^{(1)}(N)] \\
& + \log\frac{Q^2}{q_T^2}\bigg[C_A^2\left(\frac{67}{9} + \frac{\pi^2}{3}\right) - \frac{20}{9}C_An_fT_R + 2C_A\mathcal{A}_g(\phi) + 2\beta_0[\gamma_{gg}^{(1)}(N) - \beta_0] \\
& \qquad - 4[\gamma_{gg}^{(1)}(N) - \beta_0]^2 - 4n_f\gamma_{gq}^{(1)}(N)\gamma_{qg}^{(1)}(N)\bigg] \\
& + \bigg[C_A^2\left(-\frac{16}{3} + 2\zeta(3)\right) + 2C_Fn_fT_R + \frac{8}{3}C_An_fT_R - 2\beta_0\left(\mathcal{A}_g(\phi) + C_A\frac{\pi^2}{6}\right) \\
& \quad + 2\gamma_{gg}^{(2)}(N) + 2\gamma_{gg}^{(1)}(N)\left(\mathcal{A}_g(\phi) + C_A\frac{\pi^2}{3}\right) + 4C_Fn_f\gamma_{qg}^{(1)}(N)\frac{1}{(N+2)}\bigg]. \qquad (20)
\end{aligned}$$

In Eq. (19), $\gamma_{(-)}^{(2)}(N)$ is the nonsinglet spacelike two-loop anomalous dimension [14]; in Eq. (20), $\gamma_{gg}^{(2)}(N)$ is the singlet spacelike two-loop anomalous dimension [15]; $\zeta(n)$ is the Riemann ζ function [$\zeta(3) = 1.202\ldots$], and the function $\mathcal{A}_a(\phi)$ is defined through Eq. (14). The coefficients $\frac{1}{(N+1)(N+2)}$ and $\frac{1}{(N+2)}$ have origin on the N moments of $-\hat{P}_{qq}^{\epsilon}(z)$ and $-\hat{P}_{gq}^{\epsilon}(z)$, respectively.

The N-dependent part of the results in Eqs. (19) and (20) agrees with the one obtained from the second order expansion of Eq. (2) (see, e.g., Ref. [16]). By comparing also the N-independent part we obtain for $A^{(2)}$

$$A_a^{(2)} = KA_a^{(1)}, \qquad K = C_A\left(\frac{67}{18} - \frac{\pi^2}{6}\right) - n_fT_R\frac{10}{9}, \qquad (21)$$

in agreement with the results of Refs. [17,18]. Moreover, we find that $B^{(2)}$ can be expressed as

$$B_a^{(2)} = -2\delta P_{aa}^{(2)} + \beta_0\left(\frac{2}{3}C_a\pi^2 + \mathcal{A}_a(\phi)\right), \qquad a = q, g, \qquad (22)$$

where $\delta P_{aa}^{(2)}$ are the coefficients of the $\delta(1-z)$ term in the two-loop splitting functions $P_{aa}^{(2)}(z)$ [14,15], and are given by

$$\begin{aligned}\delta P_{qq}^{(2)} &= C_F^2\left(\frac{3}{8} - \frac{\pi^2}{2} + 6\zeta(3)\right) \\ &\quad + C_F C_A\left(\frac{17}{24} + \frac{11\pi^2}{18} - 3\zeta(3)\right) \\ &\quad - C_F n_f T_R\left(\frac{1}{6} + \frac{2\pi^2}{9}\right), \\ \delta P_{gg}^{(2)} &= C_A^2\left(\frac{8}{3} + 3\zeta(3)\right) - C_F n_f T_R - \frac{4}{3} C_A n_f T_R\,.\end{aligned} \tag{23}$$

From Eq. (22) we see that $B^{(2)}$, besides the $-2\delta P_{aa}^{(2)}$ term which matches the expectation from the $\mathcal{O}(\alpha_s)$ result, receives a *process-dependent* contribution controlled by the one-loop correction to the LO amplitude [see Eq. (14)]. We conclude that the Sudakov form factor in Eq. (3) is actually process dependent beyond next-to-leading-order logarithmic accuracy. The interpretation of this result will be given elsewhere [19].

However, by using the general expression in Eq. (22) it is possible to obtain $B^{(2)}$ for a given process just by computing the one-loop correction to the LO amplitude for that process. For the case of Drell-Yan, by using Eq. (15), our result for $\Sigma_{q\bar{q}}^{(2)}(N)$ agrees with the one of Ref. [7], and we confirm

$$\begin{aligned}B_q^{(2)DY} &= C_F^2\left(\pi^2 - \frac{3}{4} - 12\zeta(3)\right) \\ &\quad + C_F C_A\left(\frac{11}{9}\pi^2 - \frac{193}{12} + 6\zeta(3)\right) \\ &\quad + C_F n_f T_R\left(\frac{17}{3} - \frac{4}{9}\pi^2\right).\end{aligned} \tag{24}$$

In the case of Higgs production in the $m_{\mathrm{top}} \to \infty$ limit, by using Eq. (16) we find

$$\begin{aligned}B_g^{(2)H} &= C_A^2\left(\frac{23}{6} + \frac{22}{9}\pi^2 - 6\zeta(3)\right) + 4 C_F n_f T_R \\ &\quad - C_A n_f T_R\left(\frac{2}{3} + \frac{8}{9}\pi^2\right) - \frac{11}{2} C_F C_A\,.\end{aligned} \tag{25}$$

In particular, this result allows one to improve the present accuracy of the matching between resummed predictions [20] and fixed-order calculations [21].

To summarize, we have studied the logarithmically enhanced contributions at small transverse momentum in hadronic collisions at second order in perturbative QCD. The calculation was performed in a process-independent manner, allowing us to show that the Sudakov form factor is actually process dependent beyond next-to-leading-order logarithmic accuracy. We have provided a general expression for the coefficient $B^{(2)}$ for both quark and gluon initiated processes.

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