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Forward hadron production in pA collisions beyond leading order

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Abstract

In this talk, we report our recent progress on pinning down the cause of negativity in the NLO single inclusive hadron production in pA collisions at forward rapidity.

Keywords: Perturbative calculations, Factorization, Color transparency

1. Introduction

Parton densities in hadrons are strongly enhanced when the hadrons are accelerated at high energy colliders. However, due to the merging of soft partons, the parton density stops increasing when the energy is large enough to reach the saturation regime. Parton density distributions at high energy are well described by the color glass condensate (CGC) effective theory [1]. To study saturation effects, the ideal way is to study processes where a dilute projectile interacts with a dense target. This motivates the intensive study of forward hadron production in proton-nucleus (pA) scatterings at high energies.

At high energy, the cross section for this process can be factorized into the convolution of a perturbatively-calculable hard part, non-perturbative coefficients describing the parton densities in the incoming hadrons and the hadronization of final-state partons into hadrons. The unintegrated gluon distribution in the dense target is related to the dipole correlator $S(\mathbf{r} = \mathbf{x} - \mathbf{y}) = \left\langle \frac{1}{N_c} \mathrm{Tr}\, U(\mathbf{x}) U^\dagger(\mathbf{y}) \right\rangle$ in momentum space. Its rapidity evolution is described by the well-known Balitsky-Kovchegov (BK) equation [2, 3]. The leading-order (LO) cross section [4] for single inclusive forward hadron production in pA collisions has been derived, and numerical implementations [5, 6, 7, 8, 9] of the LO cross section are found to be consistent with experimental measurements, however with a rather large overall normalization factor. It is therefore instructive to understand how these results would change at higher orders.

The cross section for forward hadron production in pA collisions was calculated at next-to-leading order (NLO) [10, 11] and the first numerical implementation of these expressions showed that the cross section for the production of hadrons at large transverse momenta [12] is negative. Several proposals [13, 14, 15] have suggested to implement the explicit kinematic constraint or 'Ioffe time' cutoff to solve this issue, which however, did not remove the unphysical negativity completely.

In this proceeding, we analyse the cause of the negativity at large transverse momentum at NLO, and try to fix this issue. For the sake of simplicity, we will only consider the $q \to q$ channel in the following as it has a similar behavior as the full NLO cross section. We will use the simple Golec-Biernat and Wüsthoff (GBW) [16] parametrization for the dipole correlator, and the correction from NLO BK equation will be discussed at the very end.

2. Formalism

The NLO cross section for single inclusive hadron production at forward rapidity can be read from Ref. [11] which will be referred to as 'CXY'. By leaving

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out an overall integration over the impact parameter **b**, the differential multiplicity for the quark channel reads

$$\frac{\mathrm{d}N^{pA\to hX}}{\mathrm{d}^2\mathbf{p}\mathrm{d}y_h} = \int_\tau^1 \frac{\mathrm{d}z}{z^2} D_{h/q}(z) x_p q(x_p) \frac{\mathcal{S}^{(0)}(k_\perp)}{(2\pi)^2} \tag{1}$$
$$+\frac{\alpha_s}{2\pi^2}\int \frac{\mathrm{d}z}{z^2} D_{h/q}(z) \int_{\tau/z}^1 \mathrm{d}\xi \frac{1+\xi^2}{1-\xi}\frac{x_p}{\xi} q\left(\frac{x_p}{\xi}\right)$$
$$\times\left\{C_F \mathcal{I}(k_\perp,\xi) + \frac{N_c}{2}\mathcal{J}(k_\perp,\xi)\right\}$$
$$-\frac{\alpha_s}{2\pi^2}\int \frac{\mathrm{d}z}{z^2} D_{h/q}(z) \int_0^1 \mathrm{d}\xi \frac{1+\xi^2}{1-\xi} x_p q\left(x_p\right)$$
$$\times\left\{C_F \mathcal{I}_v(k_\perp,\xi) + \frac{N_c}{2}\mathcal{J}_v(k_\perp,\xi)\right\},$$

where

$$\mathcal{I}(k_\perp,\xi) = \int \frac{\mathrm{d}^2\mathbf{q}}{(2\pi)^2}\mathcal{S}(q_\perp)\left[\frac{\mathbf{k}-\mathbf{q}}{(\mathbf{k}-\mathbf{q})^2} - \frac{\mathbf{k}-\xi\mathbf{q}}{(\mathbf{k}-\xi\mathbf{q})^2}\right]^2, \tag{2}$$

$$\mathcal{J}(k_\perp,\xi) = \int \frac{\mathrm{d}^2\mathbf{q}}{(2\pi)^2}\mathcal{S}(q_\perp)\Big[\frac{2(\mathbf{k}-\xi\mathbf{q})\cdot(\mathbf{k}-\mathbf{q})}{(\mathbf{k}-\xi\mathbf{q})^2(\mathbf{k}-\mathbf{q})^2} \tag{3}$$
$$-\int \frac{\mathrm{d}^2\mathbf{l}}{(2\pi)^2}\frac{2(\mathbf{k}-\xi\mathbf{q})\cdot(\mathbf{k}-\mathbf{l})}{(\mathbf{k}-\xi\mathbf{q})^2(\mathbf{k}-\mathbf{l})^2}\mathcal{S}(l_\perp)\Big],$$

$$\mathcal{I}_v(k_\perp,\xi) = \mathcal{S}(k_\perp)\int \frac{\mathrm{d}^2\mathbf{q}}{(2\pi)^2}\left[\frac{\mathbf{k}-\mathbf{q}}{(\mathbf{k}-\mathbf{q})^2} - \frac{\xi\mathbf{k}-\mathbf{q}}{(\xi\mathbf{k}-\mathbf{q})^2}\right]^2, \tag{4}$$

$$\mathcal{J}_v(k_\perp,\xi) = \mathcal{S}(k_\perp)\Big[\int \frac{\mathrm{d}^2\mathbf{q}}{(2\pi)^2}\frac{2(\xi\mathbf{k}-\mathbf{q})\cdot(\mathbf{k}-\mathbf{q})}{(\xi\mathbf{k}-\mathbf{q})^2(\mathbf{k}-\mathbf{q})^2} \tag{5}$$
$$-\int \frac{\mathrm{d}^2\mathbf{q}}{(2\pi)^2}\frac{\mathrm{d}^2\mathbf{l}}{(2\pi)^2}\frac{2(\xi\mathbf{k}-\mathbf{q})\cdot(\mathbf{l}-\mathbf{q})}{(\xi\mathbf{k}-\mathbf{q})^2(\mathbf{l}-\mathbf{q})^2}\mathcal{S}(l_\perp)\Big]$$

with $\mathbf{p} = z\mathbf{k}$, $x_p = p_\perp e^{y_h}/(z\sqrt{s})$, $\tau = zx_p$, $x_g = p_\perp/(z\sqrt{s})e^{-y_h}$, $p_\perp = |\mathbf{p}|$, $q_\perp = |\mathbf{q}|$, $k_\perp = |\mathbf{k}|$, and $l_\perp = |\mathbf{l}|$. The color dipole in momentum space is $\mathcal{S}(k_\perp) = \mathcal{S}(k_\perp,\mathbf{b}) = \int \mathrm{d}^2\mathbf{r} e^{-i\mathbf{k}\cdot\mathbf{r}} S(\mathbf{r})$, while the superscript (0) stands for the tree level unrenormalized color dipole. In these expressions, ξ is the longitudinal momentum fraction of the incoming quark carried by the outgoing one after the radiation of a gluon with longitudinal momentum fraction $1-\xi$.

Apparently, there are two types of divergences in the NLO multiplicity: the collinear divergence and the rapidity divergence, occurring in terms proportional to C_F and N_c respectively. The collinear divergence is regulated by dimensional regularization and absorbed into the DGLAP evolution of PDFs $q(x)$ and FFs $D_{h/q}(z)$. The rapidity divergence in the N_c−terms arises when $\xi \to 1$, i.e., emitting a soft gluon collinear to the target. Naturally, this divergence should be put into the evolution of the dense target. In CXY, the rapidity divergence is regulated by introducing the renormalized dipole

$$\mathcal{S}(k_\perp) = \mathcal{S}^{(0)}(k_\perp)$$
$$+ 2\alpha_s N_c \int_0^1 \frac{\mathrm{d}\xi}{1-\xi}\left[\mathcal{J}(k_\perp,1) - \mathcal{J}_v(k_\perp,1)\right], \tag{6}$$

which in coordinate space is the integral form of BK equation. This procedure successfully absorbs the rapidity divergence into the evolution of the target, however it leads to the negativity of the NLO cross section at large transverse momenta. An oversubtraction in the rapidity regularization is found to be the cause of this negativity. This is easily demonstrated by the behaviour of the N_c−terms at large transverse momenta: $\mathcal{J}(k_\perp,\xi) - \mathcal{J}_v(k_\perp,\xi) \sim \frac{\xi}{k_\perp^4}$, a positive and linearly increasing function of ξ. To reduce the oversubtraction, one can, instead of Eq. (6), introduce a factorization scale ξ_f between 0 and 1 in the rapidity regularization,

$$\mathcal{S}(k_\perp) = \mathcal{S}^{(0)}(k_\perp)$$
$$+ 2\alpha_s N_c \int_{\xi_f}^1 \frac{\mathrm{d}\xi}{1-\xi}\left[\mathcal{J}(k_\perp,1) - \mathcal{J}_v(k_\perp,1)\right]. \tag{7}$$

This reduces to the CXY subtraction for $\xi_f = 0$, while less positive contributions will be included in the renormalized dipole for $\xi_f > 0$ at large $k_\perp$.

3. Results

Having now the general definition of the renormalized dipole with ξ_f, we will demonstrate its effect on the cross section by varying the regularization scale. For simplicity, we will use the GBW model [16] for the dipole correlator, which is

$$S(\mathbf{r}) = e^{-\mathbf{r}^2 Q_s^2/4}, \quad \mathcal{S}(k_\perp) = \frac{4\pi}{Q_s^2} e^{-k_\perp^2/Q_s^2}, \tag{8}$$

with the saturation scale $Q_s^2 = cA^{1/3}Q_{s0}^2\left(\frac{x_0}{x}\right)^\lambda$, and A being the atomic number of the target nucleus, $c = 0.56$, $Q_{s0} = 1$ GeV, $x_0 = 3.04\times10^{-4}$ and $\lambda = 0.288$. The simplified result for the NLO cross section in the GBW model is expressed in [17] for finite N_c rather than in the large N_c limit taken in CXY. This allows for a clear separation of the collinear divergence in the C_F−terms and the rapidity divergence in the N_c−terms. We use the DSS [18] and MSTW 2008 [19] NLO parametrizations for the FFs $D_{h/q}(z)$ and quark PDFs $q(x)$ respectively. We will evaluate the multiplicity at RHIC energy such that $\sqrt{s} = 200$ GeV, $\alpha_s = 0.2$, $\mu^2 = 10$ GeV^2 and $y_h = 3.2$.

The results with fixed values of ξ_f are shown in Fig. 1. The vertical line in the figure indicates the point where $p_\perp \approx Q_s(x_g)$. On the upper panel, we show the multiplicity as a function of transverse momentum for different values of ξ_f. The LO result is shown as the black