NUCLEAR SHADOWING IN HEAVY-FLAVORS HADROPRODUCTION

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Abstract

The light-cone wave function formalism for hadroproduction of heavy-flavors via gluon-gluon fusion is developed. The approach is used for evaluation of nuclear shadowing in heavy-quark production on nuclei at small $x_2 \lesssim 0.1 \cdot A^{-1/3}$. Nuclear attenuation in hadroproduction of heavy-flavors is found to be similar to shadowing effects for heavy-flavor structure functions in deep-inelastic scattering from nuclei. Nevertheless, remaining differences in the theoretical formulation of both processes imply corrections to factorization theorems.

The last years much progress in evaluation of nuclear shadowing in deep inelastic scattering (DIS) at low x has been achieved within the light-cone wave function (LCWF) formalism [1,2] Here we report on the successful extension of the LCWF formalism to nuclear shadowing in heavy quark hadroproduction [3]. This approach for the first time enabled a consistent treatment shadowing effects, which have quite a nontrivial quantum-mechanical origin. We consider $Q\bar{Q}$ production at high energy and moderate values of $x_F = x_1 - x_2$, here $x_{1,2}$ are the light-cone variable for the hadron beam and nuclear target partons producing the heavy $Q\bar{Q}$ pair. Due to the inequality $x_1 x_2 s > 4m_Q^2$ the above region corresponds to $x_2 \ll 1$. At small x_2 the gluon-gluon fusion mechanism dominates in $Q\bar{Q}$ production. Our analysis takes advantage of the typical propagation length of $Q\bar{Q}$ fluctuation of gluons being large at $x_2 \ll 1$:

$$l_{Q\bar{Q}} \sim \frac{2\,\nu_G}{4\,m_Q^2} \approx \frac{1}{M\,x_2}\,,$$

(here $\nu_G = x_b E_{lab}$ denotes the lab frame energy of the projectile gluon and E_{lab} stands for the beam energy). If $l_{Q\bar{Q}}$ is shorter than the average nucleon-nucleon distance in nuclei, $Q\bar{Q}$ production takes place incoherently on all nucleons inside the target and shadowing is absent. However, at $x_2 < x_A \approx 0.1 \cdot A^{-1/3}$ the coherence length will exceed the nuclear size, $l_{Q\bar{Q}} > R_A$, and the formation of $Q\bar{Q}$ pairs takes place coherently on the whole nucleus. In this regime, as in the case of the large coherence length for interaction of $q\bar{q}$ fluctuation of the virtual photon with target nucleus in DIS at small x [1], the shadowing may appear.

At large $x_F \sim 1$ the momentum of the projectile is transfered collectively to a heavy $Q\bar{Q}$ state by several partons in the beam hadron [4]. In this case the heavy quark pair can evidently not be assigned to a single projectile gluon. Consequently, the excitation of heavy-quark components at $x_F \sim 1$ is quite different from the $Q\bar{Q}$ production mechanism at moderate x_F. In this work we focus on moderate x_F only, where the above mechanism is not important. We demonstrate, that despite the infrared divergence of the total cross section of the projectile gluon with the target, the cross section of the $Q\bar{Q}$ production is an infrared stable quantity.

First we will consider the production of heavy $Q\bar{Q}$ pairs in the interaction of a gluon from the projectile with the nucleon target. In the LCWF technique one must carefully distinguish between "bare" and "dressed (physical)" partons. In our specific process the incident gluon will evolve after the interaction with the target into a state which contains a $Q\bar{Q}$ Fock component. If this component is projected onto a final $Q\bar{Q}$ state to calculate the production amplitude of heavy-quark pairs, one has to bear in mind that the physical gluon in the LCWF formalism has a non trivial Fock decomposition which also contains a heavy $Q\bar{Q}$ state. To leading order in the QCD coupling constant, α_S, the incoming gluon G contains bare gluon, two-gluon states, and quark-antiquark components. In the $(z, \vec{r})$ representation the LCWF of the physical incident gluon can be written as (we suppress the virtuality of the incident gluon k_1^2)

$$|G(k_1^2)\rangle = \sqrt{1 - n_{Q\bar{Q}} - n_\xi}|g\rangle + \sum_{z,\vec{r}} \Psi_G(z,\vec{r})|Q\bar{Q}; z, \vec{r}\rangle + \sum_\xi \Psi(\xi)|\xi\rangle\,. \tag{1}$$

Here $\Psi_G(z, \vec{r})$ is the projection of the physical gluon wave function onto a $Q\bar{Q}$ state (here the light-cone variable z represents the fraction of the gluon momentum carried by the quark and $\vec{r}$ is the transverse separation of the Q-$\bar{Q}$ in the impact parameter space), which up to the color factor and coupling constant coincide with the well known $Q\bar{Q}$ wave function of the virtual photon [1]. $\sum_\xi \Psi(\xi)|\xi\rangle$ represents the light quark-antiquark $q\bar{q}$ and gluon-gluon gg Fock components. The normalization of the bare gluon state in the presence of the $Q\bar{Q}$, $q\bar{q}$ and gg components is given by

$$\langle g|G\rangle = \sqrt{1 - n_{Q\bar{Q}} - n_\xi}\,, \tag{2}$$

where

$$n_{Q\bar{Q}} = \int_0^1 dz \int d^2\vec{r}\, |\Psi_G(z,\vec{r})|^2, \tag{3}$$

determines the weight of $Q\bar{Q}$ states in the physical gluon. Similar, n_ξ denotes the normalization of the light-flavor $q\bar{q}$ and gg components.

After the interaction with the target nucleon at an impact parameter $\vec{b}$, the incident gluon state is transformed into

$$|G;\vec{b}\rangle \Longrightarrow \hat{S}(\vec{b})|G;\vec{b}\rangle\,. \tag{4}$$

Here $\hat{S}(\vec{b})$ is the scattering matrix in the impact parameter representation. The familiar eikonal form [5] of the $\hat{S}$-matrix for the scattering of two partonic systems a and b is

$$\hat{S}(\vec{b}) = \exp\left[-i\sum_{i,j} V(\vec{b}+\vec{b}_i-\vec{b}_j)\hat{T}_i^\alpha \hat{T}_j^\alpha\right], \tag{5}$$

with the one gluon exchange potential (eikonal function)

$$V(\vec{b}) = \frac{\alpha_S}{\pi}\int \frac{d^2\vec{k}}{\vec{k}^2+\mu_g^2}\exp(i\vec{k}\cdot\vec{b})\,.$$

$\hat{T}_{i,j}^\alpha$ are color SU(3) generators acting on the individual partons of a and b at transverse coordinates $\vec{b}_i$ and $\vec{b}_j$ respectively. The effective gluon mass μ_g is introduced as an infrared regulator. We will show below that μ_g can be absorbed into the definition of the target gluon density. To the lowest non trivial order, we have to keep terms up to second order in α_S only. The fact, that the color generators in $\hat{S}$ do not commute is not important to this order in α_S, because our final results will contain the symmetric piece of the product of two generators only. The unitarity of the $\hat{S}$-matrix at every value of $\vec{b}$ relates the probability $P_{Q\bar{Q}}$, to detect a $Q\bar{Q}$ pair in the final state with the probabilities P_G and P_ξ, to find a physical gluon or light-flavor $q\bar{q}$ or gg states:

$$P_G + P_{Q\bar{Q}} + P_\xi = 1\,. \tag{6}$$

P_G can be written as

$$P_G = \frac{1}{8}\sum_{f,i,N'} \langle G^f N'|\hat{S}(\vec{b})|G^i N\rangle \langle G^f N'|\hat{S}(\vec{b})|G^i N\rangle^*\,. \tag{7}$$

where N and N' stand for the initial target nucleon and the final nucleonic three-quark states, "f" and "i" are the color indexes of the final and incident gluon.

Using the explicit form of the $\hat{S}$-matrix and applying closure for the final nucleonic states N', in order α_S^2 we can represent (7) as elastic scattering $\hat{S}$-matrix element for an interaction with nucleon of a color-singlet two gluon state $(GG)_1$

$$P_G = \langle (GG)_1 N|\hat{S}(\vec{b})|(GG)_1 N\rangle\,. \tag{8}$$

Making use of the Fock state decomposition (1) of the dressed gluon, we obtain

$$\begin{aligned} P_G &= (1 - n_{Q\bar{Q}} - n_\xi)\langle (gg)_1 N|\hat{S}(\vec{b})|(gg)_1 N\rangle \\ &+ 2\int dz d^2\vec{r}\,|\Psi_G(z,\vec{r})|^2 \langle (Q\bar{Q}g)_1 N|\hat{S}(\vec{b})|(Q\bar{Q}g)_1 N\rangle \\ &+ 2\int d\xi |\Psi(\xi)|^2 \langle (\xi g)_1 N|\hat{S}(\vec{b})|(\xi g)_1 N\rangle\,. \end{aligned} \tag{9}$$

To the considered order in perturbative QCD the contributions of different flavors in (6) and (9) do not interfere. We therefore suppress $q\bar{q}$ and gg terms in the following. Since in the color-singlet $(gg)_1$ state in Eq. (9) both gluons enter at the same impact parameter we have

$$\langle (gg)_1 N|\hat{S}(\vec{b})|(gg)_1 N\rangle = 1\,. \tag{10}$$

From Eqs. (9), (10) and the explicit form (3) for $n_{Q\bar{Q}}$, we obtain

$$P_{Q\bar{Q}} = 2\int dz d^2\vec{r}\,|\Psi_G(z,\vec{r})|^2 \langle (Q\bar{Q}g)_1 N|\left[1-\hat{S}(\vec{b})\right]|(Q\bar{Q}g)_1 N\rangle. \tag{11}$$

After integrating over the impact parameter, we finally obtain the cross section for the inclusive production of heavy-flavor quark-antiquark pairs in gluon-nucleon collisions:

$$\begin{aligned}\sigma(GN \to Q\bar{Q}\,X) &= 2\int dz d^2\vec{r}\,|\Psi_G(z,\vec{r})|^2 \int d^2\vec{b}\,\langle (Q\bar{Q}g)_1 N|\left[1-\hat{S}(\vec{b})\right]|(Q\bar{Q}g)_1 N\rangle \\ &= \int dz d^2\vec{r}\,|\Psi_G(z,\vec{r})|^2 \sigma_3(\vec{r}_1,\vec{r}_2) = \langle \sigma_{3N}\rangle_{Q\bar{Q}g}\,,\end{aligned} \tag{12}$$

where

$$\begin{aligned}\sigma_{3N} = \sigma_3(\vec{r}_1,\vec{r}_2) &= 2\int d^2\vec{b}\langle (Q\bar{Q}g)_1 N|\left[1-\hat{S}(\vec{b})\right]|(Q\bar{Q}g)_1 N\rangle \\ &= \frac{9}{8}\left[\sigma(r_1)+\sigma(r_2)-\frac{1}{9}\sigma(r)\right]\,.\end{aligned} \tag{13}$$

As indicated σ_{3N} is the total interaction cross section for the scattering of a color-singlet three-parton state $(Q\bar{Q}g)_1$ from a nucleon. The g-Q and g-$\bar{Q}$ separation in the impact parameter space are labeled $\vec{r}_1$ and $\vec{r}_2$ respectively. In Born approximation $\sigma(r)$ does not depend on x_2. The x_2-dependence of the dipole cross section in the limit of small x_2 generated by the higher Fock components can be evaluated through the generalized BFKL equation [6].

After integration over the gluon virtuality we obtain the cross section for hadron nucleon collisions, which to leading order in $log(4m_Q^2/\mu_g^2)$ is given by

$$\frac{d\sigma_{Q\bar{Q}}(h,N)}{dx_1} \approx g_1(x_1,\mu^2=4m_Q^2)\,\sigma(GN\to Q\bar{Q}\,X;x_2,k_1^2=0)\,. \tag{14}$$

Notice, that we reproduce the dependence on the gluon density of the beam as in the parton model ansatz. The parton model dependence on the gluon density of the target can be reproduced if one takes into account the behavior of the dipole cross section at small r

$$\sigma(r) \to \sigma(x_2,r) = \frac{\pi^2}{3}r^2\,\alpha_S(r)\left[x_2\,g(x_2,k_2^2\approx\frac{\mathcal{A}}{r^2})\right]\,, \tag{15}$$

where $\mathcal{A}\approx 10$ [7].

The above evaluation of $\sigma(GN\to Q\bar{Q}X)$ is based on the diagonalization of the parton $\hat{S}$-matrix in the $(z,\vec{r})$-representation at high energy, when the lifetime of the $Q\bar{Q}$ state becomes large as compared to the nucleon radius. Evidently, the same situation takes place for nucleus target at $l_{Q\bar{Q}} \gtrsim R_A$. Consequently to obtain the $Q\bar{Q}$ production cross section for nuclear targets it is enough to substitute σ_{3A} for σ_{3N} in Eq. (12). The cross section σ_{3A} for the scattering of a three-parton, color singlet, $(Q\bar{Q}g)_1$ state from a nucleus with mass number A can be written through σ_{3N} in the conventional Glauber form

$$\sigma_{3A} = 2\int d^2\vec{b}\left\{1-\left[1-\frac{1}{2A}\,\sigma_{3N}T(\vec{b})\right]^A\right\} \approx 2\int d^2\vec{b}\left\{1-\exp\left[-\frac{1}{2}\,\sigma_{3N}T(\vec{b})\right]\right\}\,. \tag{16}$$

Here $\vec{b}$ is the impact parameter of the $(Q\bar{Q}g)$-nucleus scattering process, $T(\vec{b})$ stands for the optical thickness of the nucleus

$$T(\vec{b}) = \int_{-\infty}^{+\infty} dz\, n_A(\vec{b}, z)\,,$$

with the nuclear density $n_A(\vec{b}, z)$ normalized to $\int d^3\vec{r}\, n_A(\vec{r}) = A$. Thus, we arrive at the following expression for the cross section for nucleus target

$$\sigma(GA \to Q\bar{Q}\,X; x_2, k_1^2) = \\ 2\int dz d^2\vec{r}\,|\Psi_G(z,r)|^2 \int d^2\vec{b}\left\{1 - \exp\left[-\frac{1}{2}\,\sigma_{3N}T(\vec{b})\right]\right\} = \langle\sigma_{3A}\rangle_{Q\bar{Q}g}\,. \quad (17)$$

Nuclear shadowing is usually discussed in terms of the nuclear transparency

$$T_A = \frac{\sigma_{Q\bar{Q}}(h, A)}{A\,\sigma_{Q\bar{Q}}(h, N)}\,.$$

In the impulse approximation $T_A = 1$. The driving contribution to nuclear attenuation results from the coherent interaction of the three-parton $(Q\bar{Q}g)$ state with two nucleons inside the target:

$$T_A = 1 - \frac{1}{4}\frac{\langle\sigma_{3N}^2\rangle_{Q\bar{Q}g}}{\langle\sigma_{3N}\rangle_{Q\bar{Q}g}}\int d^2\vec{b}\,T^2(\vec{b}) \approx 1 - \frac{3A}{16\pi R_{ch}^2}\frac{\langle\sigma_{3N}^2\rangle_{Q\bar{Q}g}}{\langle\sigma_{3N}\rangle_{Q\bar{Q}g}}\,, \quad (18)$$

where $R_{ch} \approx 1.1 fm A^{1/3}$ is mean square nucleus charge radius. For open-charm production in the accessible region of small $x_2 \sim 0.01$ evaluation of nuclear transparency through Eq. (18) yields

$$1 - T_A \sim 6 \cdot 10^{-3} \cdot A^{1/3}\,. \quad (19)$$

Eq. (18) is counterpart of the formula for nuclear transparency in DIS, in which σ_{3N} is replaced by the dipole cross section σ_{2N} for color singlet $Q\bar{Q}$ pair. Thus, we have shown that in heavy-flavor hadroproduction the $Q\bar{Q}$ Fock states in the projectile gluon, as in the case of $q\bar{q}$ Fock states of the virtual photon in DIS, leads to shadowing. However, due to the difference between σ_{3N} and σ_{2N} the nuclear shadowing in these processes differs. Consequently it cannot be ascribed to nuclear modifications of gluon structure functions, which defies the conventional parton model factorization.

From the point of view of the conventional Regge phenomenology Eq. (18) is related to the contribution of the valence intermediate state $(Q\bar{Q}g)$ in the double Pomeron exchange in interaction of $Q\bar{Q}g$ system with nucleus. The typical mass squared of the valence $Q\bar{Q}g$ state $\sim 4m_Q^2$. Besides, this low mass valence intermediate states the high mass triple Pomeron intermediate states with the mass spectrum $\propto 1/M^2$ can contribute to shadowing. In terms of the higher Fock components of the projectile to leading order in α_S the triple Pomeron shadowing correction is connected with $Q\bar{Q}gg$ states. In leading log-x_2 approximation we can consider this four-body system as originating from the valence $Q\bar{Q}g$ states after radiation of a soft gluon. The transverse size of the $Q\bar{Q}gg$ system is $\sim \max(1/m_Q, 1/\mu_g)$ The analysis of low-x HERA data on F_2 within the generalized BFKL equation yields for the effective gluon mass the value $\mu_g \approx 0.75$ GeV, which is considerably less than heavy quark masses. Thus, namely the transverse distribution of the soft gluons controls the triple Pomeron shadowing correction. The same situation takes place for DIS. In the both cases the interaction of the higher Fock-states with the target can be treated as an interaction of an octet-octet color dipole, because the transverse distance between the soft gluon and the valence ($Q\bar{Q}g$ for $Q\bar{Q}$ hadroproduction and $Q\bar{Q}$ for DIS) system is considerably greater than the size of the valence system. One can show [3] that to leading order in $\log(m_Q^2/\mu_g^2)$ the probability distributions for radiation of soft gluon by the color singlet $Q\bar{Q}g$ and $Q\bar{Q}$ states are proportional to $\langle\sigma_{3N}\rangle_{Q\bar{Q}g}$ and $\langle\sigma_{2N}\rangle_{Q\bar{Q}}$, respectively. As a result the triple Pomeron shadowing corrections to nuclear transparency for these two processes must be close to each other. The analysis of the triple Pomeron diagram for DIS [8] yields the shadowing correction which agrees well with the standard Regge formula

$$1 - T_A \approx \frac{3A}{R_{ch}^2} A_{3\mathbb{P}} \log\left(\frac{x_A}{x}\right) \sim 1.5 \cdot 10^{-2} A^{1/3} \log\left(\frac{x_A}{x}\right). \qquad (20)$$

with the triple Pomeron coupling $A_{3\mathbb{P}} \approx 0.16\,\mathrm{GeV}^{-2}$ as determined in the photoproduction experiment [9]. Eq. (20) can be used for a qualitative estimate of the triple Pomeron shadowing for $Q\bar{Q}$ hadroproduction. For $x_2 \sim 0.01$ the inclusion of the triple Pomeron contribution increases the shadowing from the valence contribution (19) by the factor ~ 2. Thus, for this region of x_2 we get $1 - T_A \sim 12 \cdot 10^{-3} \cdot A^{1/3}$. Approximated by the ansatz $T_A \approx A^{\alpha-1}$, we get $1 - \alpha \sim 15 \cdot 10^{-3}$, in agreement with recent experimental results on the open-charm production at moderate x_F [10,11]. It is worth calling that in DIS due to the large valence component related to light quarks the triple Pomeron diagram becomes important only at $x \lesssim 0.001$. The detailed evaluation of the shadowing related to the $Q\bar{Q}gg$ state in heavy quark hadroproduction will be presented elsewhere.

In summary, we demonstrated that at moderate x_F nuclear shadowing for the $Q\bar{Q}$ hadroproduction through gluon-gluon fusion can be considered as shadowing for diffractive excitation of the projectile gluon into $Q\bar{Q}$ state. The result is infrared stable and the interaction with comovers is not important. The shadowing cannot be described in terms of modification of the target nucleus gluon density, which defies the factorization theorems.

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