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CGC beyond eikonal accuracy and its applications in pA collisions

Tolga Altinoluk*
Departamento de Física de Partículas and IGFAE, Universidade de Santiago de Compostela, E-15706 Santiago de Compostela, Galicia-Spain
E-mail: tolga.altinoluk@usc.es

Néstor Armesto
Departamento de Física de Partículas and IGFAE, Universidade de Santiago de Compostela, E-15706 Santiago de Compostela, Galicia-Spain
E-mail: nestor.armesto@usc.es

Guillaume Beuf
Department of Physics, Ben-Gurion University of the Negev, Beer Sheva 84105, Israel
E-mail: beuf@post.bgu.ac.il

Alexis Moscoso
Departamento de Física de Partículas and IGFAE, Universidade de Santiago de Compostela, E-15706 Santiago de Compostela, Galicia-Spain
*Email:*alexis.mrial@gmail.com

We present a method to calculate the corrections to the eikonal approximation in the background field formalism. These corrections associated with the finite width of the target are investigated and gluon propagator in background field is calculated at next-to-next-to-eikonal accuracy. The result is then applied to the single inclusive gluon production cross section at central rapidities and the light-front helicity asymmetry, in pA collisions, in order to analyse these observables beyond the eikonal limit.

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*Speaker.

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1. Introduction

High energy scattering processes are treated usually in eikonal approximation. For the scattering of a dilute projectile on dense target, the high density of the target makes it possible to perform a semi-classical approximation that amounts to replace the target by an intense classical background field. The Color Glass Condensate (CGC, see [1] and references therein) is the effective theory that is used to study such scatterings within the eikonal and semi-classical approximations. In the CGC, one can calculate the observables for such processes in a weak coupling expansion and resum high-energy leading and next-to-leading logarithms. Corrections suppressed by inverse powers of the energy of the collision are systematically neglected within the eikonal limit. We study such power suppressed corrections to the CGC, namely next-to-eikonal contributions due to finite length of the target [2, 3]. A similar analysis in the context of soft gluon resummation for hard scattering amplitudes are considered in [4, 5].

2. Eikonal expansion of the retarded gluon propagator in background field

One of the most important elements to describe the dilute-dense scattering processes at high energy is the retarded gluon propagator in a classical background field. A highly boosted left-moving target can be described by a classical background gluon field $\mathscr{A}_a^-(x^+,\mathbf{x})$. Only the $(-)$ component of this field is enhanced by a Lorentz gamma factor, so the other components are negligible. Moreover, one can also neglect the x^- dependence of this field due to time dilation. In the presence of such a field, it is natural to work in the light-cone gauge $\mathscr{A}_a^+=0$. Since the background field is independent of x^- component, it is convenient to introduce the one dimensional Fourier transform of the retarded gluon propagator

$$G_R^{\mu\nu}(x,y)_{ab}=\int\frac{\mathrm{d}p^+}{2\pi}\,e^{-ik^+(x^--y^-)}\,\frac{1}{2(k^++i\varepsilon)}\,\mathscr{G}^{\mu\nu}_{k^+}(\underline{x};\underline{y})_{ab}\,. \tag{2.1}$$

For our purposes we only need the $(i-)$ component of the retarded gluon propagator which is written in terms of the background scale propagator as

$$\mathscr{G}^{i-}_{k^+}(\underline{x},\underline{y})^{ab}=\frac{i}{k^++i\varepsilon}\partial_{\mathbf{y}^i}\mathscr{G}^{ab}_{k^+}(\underline{x},\underline{y}) \tag{2.2}$$

with the notation $\underline{x}\equiv(x^+,\mathbf{x})$, $\mathbf{x}$ being the transverse component. The scalar propagator $\mathscr{G}^{ab}_{k^+}(\underline{x},\underline{y})$ is the retarded solution of the Green's equation

$$\left[\delta^{ab}\left(i\partial_{x^+}+\frac{\partial^2_{\mathbf{x}}}{2(k^++i\varepsilon)}\right)+g\left(\mathscr{A}^-(\underline{x})\cdot T\right)^{ab}\right]\mathscr{G}^{bc}_{k^+}(\underline{x};\underline{y})=i\,\delta^{ac}\,\delta^{(3)}(\underline{x}-\underline{y})\,. \tag{2.3}$$

The scalar Green's equation (2.3) has a form of a Schrodinger equation in $2+1$ dimensions with a matrix potential. Its solution can thus be written as a path integral [6, 7, 8]. For our purposes, we need the discretized form of the path integral, which reads

$$\mathscr{G}^{ab}_{k^+}(\underline{x},\underline{y})=\lim_{N\to\infty}\int\left(\prod_{n=1}^{N-1}d^2\mathbf{z}_n\right)\left[\prod_{n=0}^{N-1}\mathscr{G}_{0,k^+}(z^+_{n+1},\mathbf{z}_{n+1};z^+_n,\mathbf{z}_n)\right]\mathscr{U}^{ab}\left(x^+,y^+;\{\mathbf{z}_n\}\right)\,, \tag{2.4}$$

with the boundary conditions $\mathbf{z}_0 \equiv \mathbf{y}$ and $\mathbf{z}_N \equiv \mathbf{x}$, and

$$z_n^+ = y^+ + \frac{n}{N}(x^+ - y^+)\,. \tag{2.5}$$

The free scalar propagator appearing in the path integral expression, given in equation (2.4) reads

$$\mathcal{G}_{0,k^+}(\underline{x},\underline{y}) = \theta(x^+ - y^+)\,\left(\frac{-ik^+}{2\pi(x^+ - y^+)}\right)\,\exp\left(\frac{ik^+}{2(x^+ - y^+)}(\mathbf{x}-\mathbf{y})^2\right). \tag{2.6}$$

The discretized Wilson line is defined as

$$\mathcal{U}^{ab}(x^+, y^+, \{\mathbf{z}_n\}) = \mathcal{P}_+\left\{\prod_{n=0}^{N-1} \exp\left[ig\frac{(x^+ - y^+)}{N}\left(\mathcal{A}^-(z_n^+, \mathbf{z}_n)\cdot T\right)\right]\right\}^{ab}, \tag{2.7}$$

with $\mathcal{P}_+$ denoting path ordering along the x^+ direction.

In the expressions of the observables related to the dilute-dense scattering at high energy, the background propagator appears typically through its Fourier transform $\int d^2\mathbf{x}\, e^{-i\mathbf{k}\cdot\mathbf{x}}\, \mathcal{G}^{ab}_{k^+}(\underline{x},\underline{y})$. In order to study the corrections to the eikonal limit, we define the medium modification factor $\widetilde{\mathcal{R}}^{ab}_{\underline{k}}(x^+, y^+; \mathbf{x})$ as

$$\int d^2\mathbf{x}\, e^{-i\mathbf{k}\cdot\mathbf{x}}\, \mathcal{G}^{ab}_{k^+}(\underline{x},\underline{y}) = \widetilde{\mathcal{R}}^{ab}_{\underline{k}}(x^+, y^+; \mathbf{y}) \int d^2\mathbf{x}\, e^{-i\mathbf{k}\cdot\mathbf{x}}\, \mathcal{G}_{0,k^+}(\underline{x},\underline{y})\,. \tag{2.8}$$

In order to calculate the expansion of $\widetilde{\mathcal{R}}^{ab}_{\underline{k}}(x^+, y^+; \mathbf{x})$ beyond the eikonal approximation, it is more convenient to consider the large k^+ limit than the limit of the large boost of the target. The two limits of course equivalent due to the invariance of the whole scattering process under longitudinal boosts.

As it is discussed in details in [2, 3], the large k^+ expansion of the medium modification factor is performed in two steps. First, the k^+ expansion is calculated for fixed $\mathbf{k}/k^+$(semi-classical expansion) and then the result is re-expanded for small $\mathbf{k}/k^+$(small angle expansion). At the end, after performing both the semi-classical and the small angle expansion, the medium modification factor $\widetilde{\mathcal{R}}^{ab}_{\underline{k}}(x^+, y^+; \mathbf{x})$ reads

$$\begin{aligned}\widetilde{\mathcal{R}}_{\underline{k}}(x^+, y^+; \mathbf{y}) &= \mathcal{U}(x^+, y^+; \mathbf{y}) + \frac{(x^+ - y^+)\mathbf{k}^i}{k^+}\,\mathcal{U}^i_{[0,1]}(x^+, y^+; \mathbf{y}) + i\frac{(x^+ - y^+)}{2k^+}\,\mathcal{U}_{[1,0]}(x^+, y^+; \mathbf{y}) \\ &+ \frac{(x^+ - y^+)^2\mathbf{k}^i\mathbf{k}^j}{(k^+)^2}\,\mathcal{U}^{ij}_{[0,2]}(x^+, y^+; \mathbf{y}) + i\frac{(x^+ - y^+)^2\mathbf{k}^i}{2(k^+)^2}\,\mathcal{U}^i_{[1,1]}(x^+, y^+; \mathbf{y}) - \frac{(x^+ - y^+)^2}{4(k^+)^2}\,\mathcal{U}_{[2,0]}(x^+, y^+; \mathbf{y})\,.\end{aligned} \tag{2.9}$$

Here $\mathcal{U}(x^+, y^+; \mathbf{y})$ is the usual Wilson line operator and $\mathcal{U}_{[i,j]}(x^+, y^+; \mathbf{y})$ are the *decorated Wilson line* operators that appear beyond the strict eikonal limit. For the details of the calculation and the explicit expressions of the decorated Wilson lines, we refer to the reader [3].

3. Single inclusive gluon production in pA collisions

In the CGC formalism, a highly boosted left-moving nucleus is usually described by a classical gluon shockwave $\mathcal{A}^\mu_a(x) = \delta^{\mu -}\,\delta(x^+)\,\mathcal{A}^-_a(\mathbf{x})$ in the light-cone gauge $A^+_a = 0$. That field has indeed a vanishing longitudinal width and no x^- dependence in the limit of infinite boost.

Consider instead a background field

$$\mathscr{A}_a^{\mu}(x) = \delta^{\mu -}\, \mathscr{A}_a^{-}(x^+, \mathbf{x}) \tag{3.1}$$

with a finite support along the x^+ direction, from $x^+ = 0$ to $x^+ = L^+$. In the case of a large nucleus, this should be the dominant finite-boost correction with respect to the usual gluon shockwave. On the other hand, a highly boosted right-moving proton, considered as dilute, is described by a classical color current

$$j_a^{\mu}(x) = \delta^{\mu -}\, j_a^{+}(x) \tag{3.2}$$

with zero width along x^-: $j_a^+(x) \propto \delta(x^-)$. Let us consider a proton-nucleus collision with a particular impact parameter $\mathbf{B}$ and choose the center of the nucleus as the reference point for the transverse plane, so that a generic point $\mathbf{x}$ in the transverse plane is at a distance $|\mathbf{x}-\mathbf{B}|$ from the center of the proton and at a distance $|\mathbf{x}|$ from the center of the nucleus.Then, the color current $j_a^+(x)$ can be written as $j_a^+(x) = \delta(x^-)\, \mathscr{U}^{ab}(x^+, -\infty, \mathbf{x})\rho^b(\mathbf{x}-\mathbf{B})$ where ρ^b is the transverse density of color charges inside the proton before it reaches the nucleus, and $\mathscr{U}^{ab}(x^+, -\infty, \mathbf{x})$ is the Wilson line implementing the color precession of these color charges in the background field $\mathscr{A}_a^-(x^+, \mathbf{x})$ of the nucleus. One can define the Fourier transform of the color charge density

$$\rho^a(\mathbf{y}-\mathbf{B}) = \int \frac{d^2\mathbf{q}}{(2\pi)^2} e^{i\mathbf{q}\cdot(\mathbf{y}-\mathbf{B})} \tilde{\rho}(\mathbf{q}) \tag{3.3}$$

and the gluon-nucleus reduced amplitude, $\overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q})$, as

$$\mathscr{M}_{\lambda}^{a}(\underline{k}, \mathbf{B}) = \int \frac{d^2\mathbf{q}}{(2\pi)^2}\, e^{-i\mathbf{q}\cdot\mathbf{B}}\; \overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q})\;\; \tilde{\rho}^b(\mathbf{q})\,, \tag{3.4}$$

where the gluon production amplitude is $\mathscr{M}_{\lambda}^{a}(\underline{k}, \mathbf{B})$ is given as

$$\mathscr{M}_{\lambda}^{a}(\underline{k}, \mathbf{B}) = \varepsilon^{i*}(2k^+) \lim_{x^+\to\infty} \int d^2\mathbf{x} \int dx^- e^{ik\cdot x} \int d^4y\; G_R^{i-}(x,y)_{ab}\; j_b^+(y). \tag{3.5}$$

The cross section for single inclusive gluon production can be written as

$$k^+ \frac{d\sigma}{dk^+\, d^2\mathbf{k}} = \int \frac{d^2\mathbf{q}}{(2\pi)^2}\; \varphi_p(\mathbf{q})\; \frac{\mathbf{q}^2}{4}\; \frac{1}{N_c^2-1} \sum_{\lambda\, \text{phys.}} \left\langle \overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q})^{\dagger}\, \overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q}) \right\rangle_A . \tag{3.6}$$

where $\varphi_p(\mathbf{q})$ is the unintegrated gluon distribution. The second observable that we consider is the light-front helicity asymmetry of the produced gluon. The calculation of this asymmetry is almost identical to single inclusive gluon production cross section, except that instead of summing over the helicity $\lambda = \pm 1$, one takes the difference between $\lambda = +1$ and $\lambda = -1$ contributions, i.e.

$$k^+ \frac{d\sigma^+}{dk^+ d^2\mathbf{k}} - k^+ \frac{d\sigma^-}{dk^+ d^2\mathbf{k}} = \int \frac{d^2\mathbf{q}}{(2\pi)^2} \varphi(\mathbf{q}) \frac{\mathbf{q}^2}{4} \frac{1}{N_c^2-1} \sum_{\lambda \text{phys}} \lambda \left\langle \overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q})^{\dagger}\, \overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}, \mathbf{q}) \right\rangle_A . \tag{3.7}$$

By using the expanded expression of the medium modification factor one can write the reduced gluon-nucleus amplitude at next-to-next-to-eikonal accuracy as

$$\overline{\mathscr{M}}_{\lambda}^{ab}(\underline{k}; \mathbf{q}) = i\,\varepsilon_{\lambda}^{i*} \int d^2\mathbf{y}\, e^{i\mathbf{y}\cdot(\mathbf{q}-\mathbf{k})} \bigg\{ 2\left(\frac{\mathbf{k}^i}{\mathbf{k}^2} - \frac{\mathbf{q}^i}{\mathbf{q}^2}\right) \mathscr{U}(L^+, 0; \mathbf{y})$$

$$+\left(\frac{L^+}{k^+}\right)\left[\left(\delta^{ij}-2\,\frac{\mathbf{q}^i\mathbf{k}^j}{\mathbf{q}^2}\right)\mathscr{U}^j_{[0,1]}(L^+,0;\mathbf{y})-i\frac{\mathbf{q}^i}{\mathbf{q}^2}\,\mathscr{U}_{[1,0]}(L^+,0;\mathbf{y})\right]$$

$$+\left(\frac{L^+}{k^+}\right)^2\left[-2\frac{\mathbf{q}^i}{\mathbf{q}^2}\mathbf{k}^j\mathbf{k}^l\,\mathscr{U}^{jl}_{[0,2]}(L^+,0;\mathbf{y})-i\frac{\mathbf{q}^i\mathbf{k}^j}{\mathbf{q}^2}\,\mathscr{U}^j_{[1,1]}(L^+,0;\mathbf{y})+\frac{1}{2}\frac{\mathbf{q}^i}{\mathbf{q}^2}\,\mathscr{U}_{[2,0]}(L^+,0;\mathbf{y})\right.$$

$$\left.\left.+\frac{i}{4}\left(\mathbf{k}^2\delta^{ij}-2\mathbf{k}^i\mathbf{k}^j\right)\mathscr{U}^j_{(\mathrm{A})}(L^+,0;\mathbf{y})+\frac{\mathbf{k}^j}{4}\,\mathscr{U}^{ij}_{(\mathrm{B})}(L^+,0;\mathbf{y})+\frac{i}{4}\,\mathscr{U}^i_{(\mathrm{C})}(L^+,0;\mathbf{y})\right]\right\}^{ab}. \tag{3.8}$$

After squaring the reduced gluon-nucleus amplitude without summing over the helicity λ, one can realise that the transverse momentum structure of the strict eikonal and next-to-next-to-eikonal terms are symmetric under the exchange of $i \leftrightarrow j$ where i and j are the indices of the transverse components of the physical polarisation vectors in the amplitude and complex conjugate amplitude. On the other hand, the transverse momentum structure of the next-to-eikonal terms is anti-symmetric under the same exchange.

In the calculation of the single inclusive gluon production cross section one should sum over the gluon polarisations, which leads to

$$\sum_\lambda \varepsilon^{i*}_\lambda \varepsilon^j_\lambda = \delta^{ij} \tag{3.9}$$

Since the transverse momentum structure of the next-to-eikonal terms are anti-symmetric under the exchange of $i \leftrightarrow j$, this contribution to the single inclusive gluon production cross section vanishes. On the other hand, the light-front helicity asymmetry is calculated by taking the difference between the $\lambda = +1$ and $\lambda = -1$ contributions. In this case, one uses

$$\sum_\lambda \lambda\,\varepsilon^{i*}_\lambda \varepsilon^j_\lambda = i\mathscr{E}^{ij} \tag{3.10}$$

where $\mathscr{E}^{ij}$ is the anti-symmetric matrix with $\mathscr{E}^{12} = +1$. Thus, the light-front helicity asymmetry gets the contribution only from the next-to-eikonal terms. The final expression for the single inclusive gluon cross section reads at next-to-next-to-eikonal accuracy

$$k^+\frac{d\sigma}{dk^+d^2\mathbf{k}} = \int\frac{d^2\mathbf{q}}{(2\pi)^2}\varphi(\mathbf{q})\frac{\mathbf{q}^2}{4}\int d^2\mathbf{r}e^{i\mathbf{r}\cdot(\mathbf{q}-\mathbf{k})}\left\{4C^i(\mathbf{k},\mathbf{q})C^j(\mathbf{k},\mathbf{q})O(\mathbf{r})+\left(\frac{L^+}{k^+}\right)^2\left[2\frac{\mathbf{q}^i}{\mathbf{q}^2}C^i(\mathbf{k},\mathbf{q})\right.\right.$$

$$\times\left[-2\mathbf{k}^l\mathbf{k}^m O^{l,m}_{[0,2]}(\mathbf{r})-2i\mathbf{k}^l O^l_{[1,1]}(\mathbf{r})+O_{[2,0]}(\mathbf{r})\right]+C^i(\mathbf{k},\mathbf{q})\left[i(\mathbf{k}^2\delta^{il}-2\mathbf{k}^i\mathbf{k}^l)O^l_{(\mathrm{A})}(\mathbf{r})+\mathbf{k}^m O^{im}_{(\mathrm{B})}(\mathbf{r})\right.$$

$$\left.\left.\left.+iO^i_{(\mathrm{C})}(\mathbf{r})\right]+\tilde{C}^{li}(\mathbf{k},\mathbf{q})\left[\tilde{C}^{li}(\mathbf{k},\mathbf{q})O^{l;m}_{[0,1];[0,1]}(\mathbf{r})+2i\frac{\mathbf{q}^i}{\mathbf{q}^2}O^l_{[0,1];[1,0]}(\mathbf{r})\right]+\frac{1}{\mathbf{q}^2}O_{[1,0];[1,0]}(\mathbf{r})\right\} \tag{3.11}$$

and the light-front helicity asymmetry reads

$$k^+\frac{d\sigma^+}{dk^+d^2\mathbf{k}}-k^+\frac{d\sigma^-}{dk^+d^2\mathbf{k}} = \int\frac{d^2\mathbf{q}}{(2\pi)^2}\varphi(\mathbf{q})\mathbf{q}^2\int d^2\mathbf{r}e^{i\mathbf{r}\cdot(\mathbf{q}-\mathbf{k})}\left(\frac{L^+}{k^+}\right)\mathscr{E}^{ij}$$

$$\times C^j(\mathbf{k},\mathbf{q})\left[i\tilde{C}^{li}(\mathbf{k},\mathbf{q})O^l_{[0,1]}(\mathbf{r})+\frac{\mathbf{q}^i}{\mathbf{q}^2}O_{[1,0]}(\mathbf{r})\right] \tag{3.12}$$

where $C^i(\mathbf{k},\mathbf{q}) = \left(\frac{\mathbf{k}^i}{\mathbf{k}^2}-\frac{\mathbf{q}^i}{\mathbf{q}^2}\right)$, $\tilde{C}^{ij}(\mathbf{k},\mathbf{q}) = \left(\delta^{ij}-2\mathbf{k}^i\frac{\mathbf{q}^j}{\mathbf{q}^2}\right)$ and the dipole operators are defined as

$$O(\mathbf{r}) = \int d^2\mathbf{b}\frac{1}{N_c^2-1}\left\langle \mathrm{tr}\left[\mathscr{U}\left(\mathbf{b}+\frac{\mathbf{r}}{2}\right)\mathscr{U}^\dagger\left(\mathbf{b}-\frac{\mathbf{r}}{2}\right)\right]\right\rangle_A, \tag{3.13}$$

$$O^{i\cdots j}_{[\alpha,\beta]}(\mathbf{r}) = \int d^2\mathbf{b}\frac{1}{N_c^2-1}\left\langle \mathrm{tr}\left[\mathscr{U}^{i\cdots j}_{[\alpha,\beta]}\left(\mathbf{b}+\frac{\mathbf{r}}{2}\right)\mathscr{U}^{\dagger}\left(\mathbf{b}-\frac{\mathbf{r}}{2}\right)\right]\right\rangle_A , \tag{3.14}$$

$$O^{i\cdots j;l\cdots m}_{[\alpha,\beta];[\gamma,\delta]}(\mathbf{r}) = \int d^2\mathbf{b}\frac{1}{N_c^2-1}\left\langle \mathrm{tr}\left[\mathscr{U}^{i\cdots j}_{[\alpha,\beta]}\left(\mathbf{b}+\frac{\mathbf{r}}{2}\right)\mathscr{U}^{l\cdots m\dagger}_{[\gamma,\delta]}\left(\mathbf{b}-\frac{\mathbf{r}}{2}\right)\right]\right\rangle_A . \tag{3.15}$$

4. Conclusions

We have presented a systematic eikonal expansion of the retarded gluon propagator in a background field. The corrections to the strict eikonal limit that are the due the finite width of the target are calculated at next-to-next-to-eikonal accuracy. The eikonal expansion performed at the level of the gluon background propagator is then applied to high energy dilute-dense scattering processes within the CGC framework. Two different observables have been analysed, in pA collisions at midrapidity, within this framework: the single inclusive gluon production cross section and the light-front helicity asymmetry of produced gluons. For the single inclusive gluon cross section, it has been shown that next-to-eikonal terms vanish and the first non-vanishing corrections to the strict eikonal limit that appear at next-to-next-to-eikonal order have been calculated. On the other hand, for the light-front helicity asymmetry, it has been shown that both the strict eikonal and the next-to-next-to-eikonal terms vanish and the leading contribution to this observable turns out to be the next-to-eikonal terms.

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