Cathode tube effect: Heavy quarks probing the glasma in *p*-Pb collisions

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We study the propagation of charm quarks in the early stage of high energy proton-lead collision, considering the interaction of these quarks with the evolving glasma by means of the Wong equations. Neglecting quantum fluctuations at the initial time, the glasma is made of longitudinal fields, but the dynamics leads to a quick formation of transverse fields; we estimate such a formation time as $\Delta t \approx 0.1$ fm/c, which is of the same order of the formation time of heavy quark pairs $t_{\mathrm{formation}} \approx 1/(2m)$. Limiting ourselves to the simple case of a static longitudinal geometry, we find that heavy quarks are accelerated by the strong transverse color fields in the early stage and this leads to a tilting of the c-quarks spectrum towards higher p_T states. This average acceleration can be understood in terms of drag and diffusion of c-quarks in a hot medium and appears to be similar to the one felt by the electrons ejected by the electron cannon in a cathode tube: we dub this effect as *cathode tube effect*. The tilting of the spectrum affects the nuclear modification factor, R_{pPb}, suppressing this below one at low p_T and making it larger than one at intermediate p_T. We compute $R_{\mathrm{pPb}}(p_T)$ after the evolution of charm quarks in the gluon fields and we find that its shape is in qualitative agreement with the measurements of the same quantity for D-mesons in proton-lead collisions.

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I. INTRODUCTION

The study of the initial condition of the system produced by high energy collisions is a difficult but interesting problem related to the physics of relativistic heavy ion collisions (RHICs), as well as to that of high energy proton-proton (pp) and proton-nucleus (pA) collisions. If the energy of the collision is very large then the two colliding nuclei in the backward light cone can be described within the color-glass-condensate (CGC) effective theory [1–7], in which fast partons are frozen by time dilatation and act as static sources for low momentum gluons: their large occupation number allows for a classical treatment of these fields. The collision of two colored glass sheets, each representing one of the colliding objects in high energy collisions, leads to the formation of strong gluon fields in the forward light cone named as the glasma [8–18]. In the weak coupling regime the glasma consists of longitudinal color-electric and color-magnetic fields; these are again characterized by a large gluon occupation number, $A_\mu^a \simeq 1/g$ with g the QCD coupling, so they can be described by classical field theory namely the Classical Yang-Mills (CYM) theory. Finite coupling brings up quantum fluctuations on the top of the glasma [19–33] that we do not consider in the present article, leaving their inclusion to a forthcoming study. Among the high energy collisions mentioned above, pA are interesting because they allow for both a theoretical and an experimental study of the cold nuclear matter effects (CNME), namely those effects that are not directly related to the formation of the quark-gluon plasma (QGP) and that include shadowing [34] as well as gluon saturation [35–37]; see [38–41] for reviews.

Heavy quarks are excellent probes of the system created in high energy nuclear collisions, both for the preequilibrium part and for the thermalized QGP; see [40–53] and references therein. Their formation time is very small in comparison with the one of light quarks: indeed, this can be estimated as $\tau_{\mathrm{form}} \approx 1/(2m)$ with m the quark mass which gives $\tau_{\mathrm{form}} \leq 0.1$ fm/c for the charm quark. Because heavy quarks are produced immediately after the collision, they can propagate in the evolving glasma fields and probe its evolution. For nucleus-nucleus collisions it is likely that the effect of this propagation is largely washed out by the successive interaction with the bulk QGP; on the other hand, for pA and pp collisions the effect of the interaction with a medium is much smaller because of the smaller lifetime of the latter (if a QGP is created at all); therefore, some effect of the initial propagation in the gluon fields might survive up to the final stage of the evolution after hadronization. Moreover, heavy quark propagation in the background of the gluon field will hardly affect the latter, due to the large mass and the small number of these quarks

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which leads to a negligible color current. Therefore, heavy quarks are ideal probes of the strong gluon fields formed in high energy nuclear collisions.

In this study we focus on c-quarks in high energy p-Pb collisions. The main purpose of our study is to compute consistently the propagation of the heavy quarks in the initial gluon fields, assuming besides that a QGP is not formed, or at least that it is formed at a later stage and its lifetime is quite smaller than the one in Pb-Pb collisions. In particular, we are interested in the nuclear modification factor, R_{pPb}, for D-mesons that has been measured recently [54,55]. In fact, we find that the propagation in gluon fields leads to R_{pPb} that reminds one, at least qualitatively, of the one measured for D-mesons in p-Pb collisions.

Propagation of heavy quarks in the glasma has been studied previously in [53], although within a simplified approach based on a Fokker-Planck equation. Despite studying a simplified situation, the work in [53] is interesting because it shows how the evolution of c-quarks in glasma can be interpreted in terms of drag and diffusion in momentum space, similarly to the evolution in a thermal medium. Here we aim to perform a more complete study of the same problem without relying on the small transferred momentum expansion of [53], as well as including the dynamical evolution of the gluon medium that is missing in [53], in order to quantify the effect of the evolution of c-quarks in the glasma on observables. This is achieved by solving consistently the classical equations of motion of the gluon fields, namely the classical Yang-Mills equations, and of the heavy quarks propagating in the glasma that are the Wong equations. This approach is equivalent to solve the Boltzmann-Vlasov equations for the heavy quarks in a collisionless plasma; as a matter of fact, the Boltzmann-Vlasov equations can be solved by means of the test particles method which amounts to solving the classical equations of motion of the test particles, here represented by the heavy quarks, and these classical equations are just the Wong equations.

The purpose of our study is to estimate the impact of the early stage of p-Pb collisions on R_{pPb}. We mention that the present work should be considered as a preliminary one since we do not include a longitudinal expansion in our calculation; therefore we do not attempt to do a serious comparison with the existing experimental data. While the inclusion of the expansion might reduce the effect on R_{pPb}, we find that the largest part of it comes within ≈ 1 fm/c of evolution; therefore most likely at least part of this effect will remain also in case the longitudinal expansion is taken into account (we will include the longitudinal expansion anyway in a forthcoming paper). Keeping this in mind, whenever we mention that we consider p-Pb collisions at a given energy it means that we have set up the initial color charge distributions on the proton and Pb sides in agreement with what should be done for simulations of realistic collisions, trying to keep both the color charge distributions and the saturation scales as close as possible to what should be done in a complete calculation where expansion is taken into account.

II. GLASMA AND CLASSICAL YANG-MILLS EQUATIONS

In this section we briefly review the glasma and the McLerran-Venugopalan (MV) model [1–3,56]. We remark that in our notation the gauge fields have been rescaled by the QCD coupling $A_\mu \to A_\mu/g$. In the MV model, the static color charge densities ρ_a on the nucleus A are assumed to be random variables that are normally distributed with zero mean and variance specified by the equation

$$\langle \rho_A^a(\boldsymbol{x}_T)\rho_A^b(\boldsymbol{y}_T)\rangle = (g^2\mu_A)^2\varphi_A(\boldsymbol{x}_T)\delta^{ab}\delta^{(2)}(\boldsymbol{x}_T-\boldsymbol{y}_T); \tag{1}$$

here A corresponds to either the proton or the Pb nucleus, a and b denote the adjoint color index. In this work we limit ourselves for simplicity to the case of the $SU(2)$ color group; therefore, $a, b = 1, 2, 3$. In Eq. (1) $g^2\mu_A$ denotes the color charge density and it is of the order of the saturation momentum Q_s [57].

The function $\varphi_A(\boldsymbol{x}_T)$ in Eq. (1) allows for a nonuniform probability distribution of the color charge in the transverse plane. In this article, we study the gluon fields produced in p-Pb collisions. For the case of a the Pb nucleus we assume a uniform probability and take $\varphi(\boldsymbol{x}_T) = 1$. On the other hand, for the proton we use the constituent quark model [58–62]. For each event, we first extract the position of the three valence quarks, $\boldsymbol{x}_i$ with $i = 1, 2, 3$, assuming a Gaussian distribution, namely,

$$\psi(\boldsymbol{x}_T) = e^{-(\boldsymbol{x}_T^2)/(2B_{cq})}; \tag{2}$$

then we build up the probability density

$$\varphi_p(\boldsymbol{x}_T) = \frac{1}{3}\sum_{i=1}^{3} e^{-(\boldsymbol{x}_T^2-\boldsymbol{x}_i^2)/(2B_\rho)}. \tag{3}$$

The two parameters in Eqs. (2) and (3) are $B_{cq} = 3$ GeV and $B_q = 0.3$ GeV. The two parameters above have been fixed by fits of deep inelastic scattering data in [58–62], as well as by requiring that the dipole amplitude obtained within the constituent quark model agrees with that of the spherical proton. We mention that in the literature several sets of parameters are used: we have checked that changing these parameters does not affect substantially our prediction on the nuclear modification factor. In fact, the most important quantity is not the exact distribution of the initial electric fields in the transverse plane but their average magnitude, and this is sensitive to the average value of Q_s, which we fix by the Golec-Biernat-Wusthoff (GBW) fit giving an average Q_s equal to 0.80 GeV at $\sqrt{s} = 5.02$ TeV; in any case, see the discussion below. We remark that this procedure does not correspond to assume that the three valence quarks are the only sources of the large x color charges; indeed, from Eq. (1) it should be obvious to any reader familiar with the MV model that we distribute sea color charges analogously to what is done for the case of a homogeneous $g^2\mu$. In fact, the constituent quark models amount simply to assume that the large x charges from the sea localize around the valence quarks: these act as seeds for the sea charges. The sensitivity on the number of constituent hot spots of color charge has been studied in [61]; in [58–62] the significance of this model in comparison with the simpler Gaussian one is well explained. The Gaussian model of the proton can be used as well in our study, and we will report on this in a future work.

For the proton, $g^2\mu_p\varphi_p(\boldsymbol{x}_T)^{1/2}$ can be understood as an $\boldsymbol{x}_T$-dependent $g^2\mu$ because $\varphi_p(\boldsymbol{x}_T)$ localizes the distribution around the valence quarks; we fix $g^2\mu_p$ for each event assuming that $\langle g^2\mu_p\varphi_p(\boldsymbol{x}_T)^{1/2}\rangle/Q_s = 0.57$, following the result of [57] where the average is defined with $\varphi_p(\boldsymbol{x}_T)$ as a weight function, then estimating Q_s at the relevant energy by using the standard GBW fit [63–65]:

$$Q_s^2 = Q_{s,0}^2\left(\frac{x_0}{x}\right)^{\lambda}, \tag{4}$$

with $\lambda = 0.277$, $Q_0 = 1$ GeV, and $x_0 = 4.1\times10^{-5}$. We remind that whenever we apply this equation to high energy collisions, the relevant value of x for the two colliding objects can be estimated at midrapidity as $\langle p_T\rangle/\sqrt{s}$, where $\langle p_T\rangle$ corresponds to the average p_T of the gluons produced by the collision. For example, at the RHIC energy for $x = 0.01$ we obtain $Q_s = 0.47$ GeV in agreement with the estimate of [66]. At the LHC energy $\sqrt{s} = 5.02$ TeV we find $Q_s = 0.80$ GeV which gives $\langle g^2\mu_p\varphi_p(\boldsymbol{x}_T)^{1/2}\rangle = 1.41$ GeV.

For the Pb nucleus the uncertainty on the Q_s as well as on $g^2\mu$ comes from the different model used to compute Q_s for a large nucleus. Indeed the GBW fit in this case is modified as

$$Q_s^2 = f(A)Q_{s,0}^2\left(\frac{x_0}{x}\right)^{\lambda}, \tag{5}$$

where

$$f(A) = A^{1/3}, \tag{6}$$

within a naive scaling hypothesis, and

$$f(A) = cA^{1/3}\log A, \tag{7}$$

within the IP-Sat model [67]. While other forms of $f(A)$ are possible [68,69], the two above give the higher and lower value of Q_s at the RHIC energy [57]; therefore, we take these two to set the upper and lower estimates of Q_s. Using again $Q_s/g^2\mu = 0.57$, we find $g^2\mu_{\mathrm{Pb}} = 2$ GeV and $g^2\mu_{\mathrm{Pb}} = 3$ GeV at the RHIC energy taking, respectively, the IP-Sat and naive forms; the modified GBW fit then leads to $g^2\mu_{\mathrm{Pb}} = 3.4$ GeV and $g^2\mu_{\mathrm{Pb}} = 5.2$ GeV for the two cases at the LHC energy. We remark that for the Pb nucleus it is enough to take a uniform $\varphi_p(\boldsymbol{x}_T)$ in Eq. (2). Indeed, for large nuclei the saturation scale Q_s, or equivalently $g^2\mu$, is distributed in the transverse plane according to the scaling $Q_s^2 \propto T(\boldsymbol{x}_T)$, where $T(\boldsymbol{x}_T)$ corresponds to the thickness function of the nucleus. This does not resolve the subnuclear scale because it considers the distributions of nucleons on the transverse plane; therefore it is reasonable that on a transverse area as big as that of a nucleon the color charges in Pb distribute with uniform probability. The fact that a uniform $\varphi_p(\boldsymbol{x}_T)$ is enough for the Pb nucleus can be understood qualitatively: the color charge distribution of the Pb nucleus in the transverse plane is obtained integrating over the longitudinal direction, taking into account the contribution of many nucleons distributed along this direction; if every nucleon contributes with its own constituent quarks, when these are projected onto the transverse plane they will appear approximately as distributed uniformly.

The static color sources $\{\rho\}$ generate pure gauge fields outside and on the light cone, which in the forward light cone combine and give the initial glasma fields. In order to determine these fields we firstly solve the Poisson equations for the gauge potentials generated by the color charge distributions of the nuclei A and B, namely,

$$-\partial_\perp^2\Lambda^{(A)}(\boldsymbol{x}_T) = \rho^{(A)}(\boldsymbol{x}_T) \tag{8}$$

(a similar equation holds for the distribution belonging to B). Wilson lines are computed as $V^\dagger(\boldsymbol{x}_T) = e^{i\Lambda^{(A)}(\boldsymbol{x}_T)}$, $W^\dagger(\boldsymbol{x}_T) = e^{i\Lambda^{(B)}(\boldsymbol{x}_T)}$, and the pure gauge fields of the two colliding nuclei are given by $\alpha_i^{(A)} = iV\partial_iV^\dagger$, $\alpha_i^{(B)} = iW\partial_iW^\dagger$. In terms of these fields the solution of the CYM in the forward light cone at initial time, namely the glasma gauge potential, can be written as $A_i = \alpha_i^{(A)} + \alpha_i^{(B)}$ for $i = x, y$ and $A_z = 0$, and the initial longitudinal glasma fields are [8,9]

$$E^z = i\sum_{i=x,y}[\alpha_i^{(B)}, \alpha_i^{(A)}], \tag{9}$$

$$B^z = i([\alpha_x^{(B)}, \alpha_y^{(A)}] + [\alpha_x^{(A)}, \alpha_y^{(B)}]), \tag{10}$$

while the transverse fields are vanishing. It has been suggested that the gauge potentials should be computed by defining the Wilson lines as path-ordered exponentials of multiple layers of color charges in order to describe the propagation of a colored probe through a thick nucleus [57]; we have checked that using multiple layers instead of a single layer of charge does not affect considerably our results, and for the sake of simplicity we report here only the results obtained using one single layer, leaving a more complete report to a forthcoming article.

The dynamical evolution that we study here is given by the CYM equations. In this study we follow [23]; therefore, we refer to that reference for more details. The Hamiltonian density is given by

$$H = \frac{1}{2}\sum_{a,i} E_i^a(x)^2 + \frac{1}{4}\sum_{a,i,j} F_{ij}^a(x)^2, \tag{11}$$

where the magnetic part of the field strength tensor is

$$F_{ij}^a(x) = \partial_i A_j^a(x) - \partial_j A_i^a(x) + \sum_{b,c} f^{abc} A_i^b(x) A_j^c(x); \tag{12}$$

here $f^{abc} = \varepsilon^{abc}$ with $\varepsilon^{123} = +1$. The equations of motion for the fields and conjugate momenta, namely the CYM equations, are

$$\frac{dA_i^a(x)}{dt} = E_i^a(x), \tag{13}$$

$$\frac{dE_i^a(x)}{dt} = \sum_j \partial_j F_{ji}^a(x) + \sum_{b,c,j} f^{abc} A_j^b(x) F_{ji}^c(x). \tag{14}$$

We solve the above equations on a static box in three spatial dimensions as in [23,33].

A. Heavy quarks in the evolving glasma

At the initial time we assume that the momentum distribution of c-quarks is the prompt one obtained within Fixed Order + Next-to-Leading Log (FONLL) QCD, which describes the D-mesons spectra in pp collisions after fragmentation [70–72]:

$$\left.\frac{dN}{d^2 p_T}\right|_{\mathrm{prompt}} = \frac{x_0}{(x_1 + p_T)^{x_2}}; \tag{15}$$

the parameters that we use in the calculations are $x_0 = 6.37 \times 10^8$, $x_1 = 9.0$, and $x_2 = 10.279$. Normalization of the spectrum is not relevant in this article because we are interested in the nuclear modification factor which is a ratio of the final over initial spectrum, and this is unaffected by the overall normalization since the number of heavy quarks is conserved during the evolution; the slope of the spectrum has been calibrated to a collision at 5.02 TeV. Moreover, we assume that the initial longitudinal momentum vanishes (in a longitudinally expanding geometry, this condition can be replaced by the standard Bjorken flow $y = \eta$). Initialization in coordinate space is done as follows: the transverse coordinate distribution is built up by means of the function $\psi(\boldsymbol{x}_T)$ in Eq. (2), because we expect the heavy quarks to be produced in the overlap region of proton and Pb nucleus that coincides with the transverse area of the proton; on the other hand, we use a uniform distribution for the longitudinal coordinate (in a longitudinally expanding geometry, this condition can be replaced by a uniform distribution in spacetime rapidity).

The dynamics of heavy quarks in the evolving glasma is studied by the Wong equations [73,74], that, for a single quark, can be written as

$$\frac{dx_i}{dt} = \frac{p_i}{E}, \tag{16}$$

$$E\frac{dp_i}{dt} = Q_a F_{i\nu}^a p^\nu, \tag{17}$$

$$E\frac{dQ_a}{dt} = -Q_c \varepsilon^{cba} \boldsymbol{A}_b \cdot \boldsymbol{p}, \tag{18}$$

where $i = x, y, z$; here, the first two equations are the familiar Hamilton equations of motion for the coordinate and its conjugate momentum, while the third equation corresponds to the gauge invariant color current conservation. Here $E = \sqrt{\boldsymbol{p}^2 + m^2}$, with $m = 1.5$ GeV corresponding to the charm quark mass. In the third Wong equation Q_a corresponds to the c-quarks color charge; we initialize this by a uniform distribution with support in the range $(-1, +1)$. For each c-quark we produce a $\bar{c}$-quark as well; for this we assume the same initial position of the companion c, opposite momentum and opposite color charge. Solving the Wong equations is equivalent to solving the Boltzmann-Vlasov equations for a collisionless plasma made of heavy quarks, which propagate in the evolving glasma; in fact, the latter equation can be solved by means of the test particle method which amounts to solving the classical equations of motion of the particles in the background of the evolving gluon field. In principle, we should include the heavy quarks color current density on the right hand side of Eq. (14) and compute the backreaction on the gluon fields. However, we neglect this backreaction: this approximation is usually used to study the propagation of heavy probes in a thermal QGP bath and sounds quite reasonable due to the small number of heavy quarks produced by the collision, as well as to their large mass, both of these factors leading eventually to a negligible color current density. On the transverse lattice we do not assume periodic boundary conditions for the heavy quarks: as soon as a heavy quark reaches the boundary of the transverse box we cancel any interaction with the gluon fields and its motion becomes a simple free streaming.

We remark that our choice of the formation time, $\tau_{\mathrm{form}} = 1/(2m)$, has to be taken only as a rough estimate: for example, it does not take into account that quarks with a finite value of p_T should be produced earlier because of their larger kinetic energy (the inverse mass should be replaced by their inverse kinetic energy); taking into account that the average p_T of c-quarks is between 1 and 2 GeV, the replacement $\sqrt{p_T^2+m^2} \to m$ in $\tau_{\rm form}$ should be adequate. Moreover, we notice that heavy quarks are always produced in pairs with a vanishing total momentum; besides, we do not have the longitudinal expansion in our calculation. These two factors imply that each time a quark-antiquark pair is formed, it pops out in cells that are at rest so no Lorentz boost of $\tau_{\rm form}$ is needed. In calculations based on relativistic transport the heavy quarks are assumed to do a free streaming between their formation time and the initialization of the quark-gluon plasma phase; see e.g., Ref. [75] and references therein. In our calculation we do not have a free streaming period, and the heavy quarks are formed exactly at their formation time and interact immediately with the gluon background.

III. RESULTS

In Fig. 1, we plot the averaged color-electric fields, measured in lattice units, versus time. Solid lines correspond to the longitudinal fields while dashed lines denote the transverse fields; green and indigo lines correspond to $g^2\mu_{\rm Pb} = 5.2$ GeV and $g^2\mu_{\rm Pb} = 3.4$ respectively. The transverse size of the box is 4 fm, and we have used a transverse lattice with size 91×91 that gives the lattice spacing $\delta x = 0.04$ fm. At the initial time the system is made of purely longitudinal fields, but this configuration is intrinsically unstable and the gluon dynamics leads to the production of transverse fields: within $\Delta t \approx 0.1$ fm/c the bulk is already formed, and at later times the magnitude of the several components of the fields does not change considerably. In our calculations the average values of the field do not change considerably after the initial transient: this is due to the static box geometry used in this calculations. As a matter of fact, we lack the longitudinal expansion which eventually would lead to the decay of the fields. It is well known that when the realistic longitudinal expansion is considered, the transverse fields are still formed quickly, then decay beside the longitudinal ones within a proper time range of $\approx$1 fm/c; therefore the expansion itself should not affect drastically the results presented here. We comment more on this in the Conclusions.

In the upper panel of Fig. 2 we plot the D-meson spectrum, dN/d^2p_T, at the initial time (maroon dashed line) and at $t = 1$ fm/c (green solid line). In the lower panel of the same figure we plot the momentum distribution of c-quarks, dN/dp_T, at the initial time (dashed maroon line), at $t = 0.5$ fm/c (orange dot-dashed line), and at $t = 1$ fm/c (green solid line). We assume $g^2\mu_{\rm Pb} = 5.2$ GeV. In the calculation we have assumed that the formation time of c-quarks is $t_{\rm formation} = 1/(2m_c) \approx 0.06$ fm/c for $m = 1.5$ GeV, but we have checked that lowering this value does not affect considerably the final result. At the

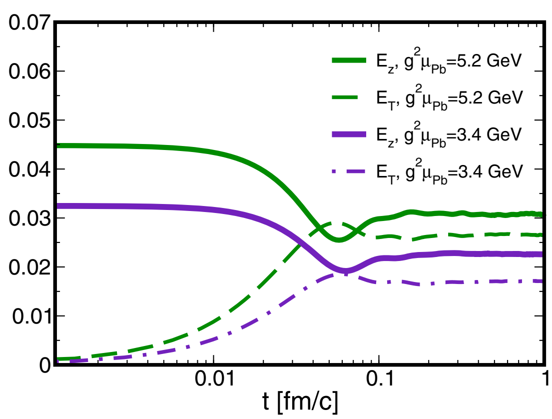

FIG. 1. Averaged color-electric fields for p-Pb collision, measured in lattice units. Solid lines correspond to the longitudinal fields, while dashed lines denote the transverse fields; green and indigo lines correspond to $g^2\mu_{\rm Pb} = 5.2$ GeV and $g^2\mu_{\rm Pb} = 3.4$, respectively. Lattice spacing is $\delta x = 0.04$ fm.

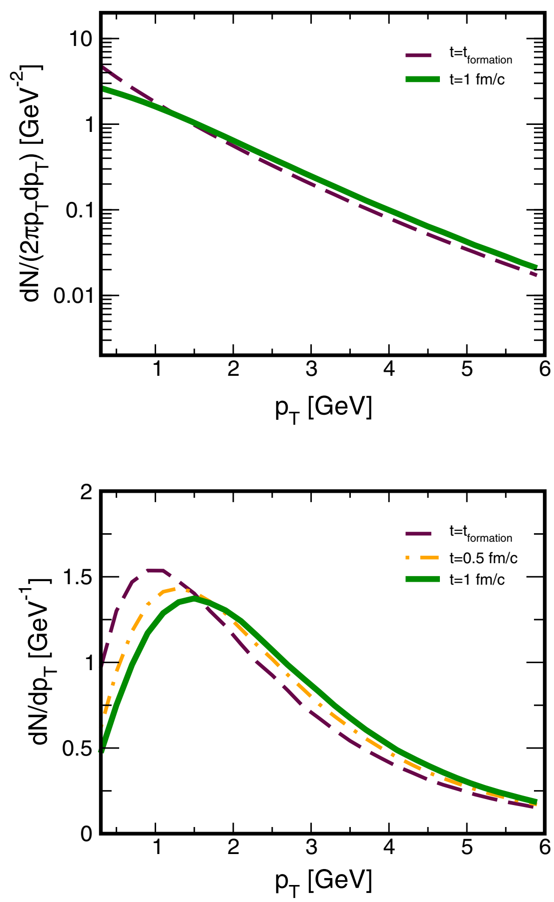

FIG. 2. In the upper panel, we plot the D-meson spectrum, dN/d^2p_T, at the initial time (maroon dashed line) and at $t = 1$ fm/c (green solid line). In the lower panel, we plot the momentum distribution of c-quarks, dN/d^2p_T, at the initial time (dashed maroon line), at $t = 0.5$ fm/c (orange dot-dashed line), and at $t = 1$ fm/c (green solid line). We take $g^2\mu_{\rm Pb} = 5.2$ GeV.

end of the evolution we adopt a standard fragmentation for the charm quark to D-meson [76], with

$$f(z) \propto \frac{1}{z(1 - \frac{1}{z} - \frac{\epsilon_c}{1-z})^2}, \tag{19}$$

where $z = p_D/p_c$ is the momentum fraction of the D-meson fragmented from the charm quark, and ϵ_c is a free parameter to fix the shape of the fragmentation function in order to reproduce the D-meson production in pp collisions [75], namely, $\epsilon_c = 0.06$. In the lower panel of Fig. 2, we plot the c-quark distribution dN_c/dp_t at the initial time (maroon dashed line), at $t = 0.5$ fm/c (orange dot-dashed line), and at $t = 1$ fm/c (green solid line). We notice that the main effect of the interaction of the heavy quarks with the gluon field is to empty the low p_T states of the c-quarks and fill the states with higher values of p_T; this effect looks similar to the acceleration that electric charges would feel in the background of a transverse field.

In order to understand better the interaction of the c-quarks with the evolving glasma fields, we prepare initializations in which we put all the c-quarks in a very thin p_T bin to obtain a δ-like distribution; the evolution of this distribution is studied again by means of the Wong equations. This is done in order to better understand the interaction of the glasma with different p_T modes. The results of this are shown in Fig. 3, in which we plot the distribution function dN_c/dp_T at the initial time (solid black lines), at $t = 0.5$ fm/c (green dashed lines), and at $t = 1$ fm/c (solid red lines) for several values of the initial p_T. We notice that in all the cases examined here the interaction with the glasma fields leads to the spreading of dN/dp_T, which is very similar to the standard diffusion in momentum space encountered in a Brownian motion. In addition to this, for low p_T we find that diffusion is flanked by a drag towards higher values of p_T: this results in an average acceleration of the c-quarks, and it is similar to

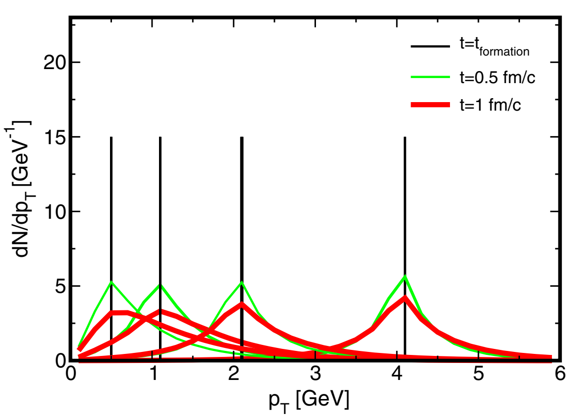

FIG. 3. Evolution of δ-distribution functions of c-quarks in the glasma fields. Black solid lines correspond to the initializations, green dashed lines to $t = 0.5$ fm/c, and red solid lines to $t = 1$ fm/c. We take $g^2\mu_{\rm Pb} = 5.2$ GeV.

what we would expect putting low-p_T quarks in a hot medium. A more quantitative comparison of the evolution of heavy quarks in glasma and in a hot plasma will be the subject of a forthcoming article.

Our results agree with those of [53]; in the present work, however, we find a very tiny drag of high-p_T quarks towards lower values of p_T. In fact, we have run additional simulations starting with $p_T = 10$ GeV, but the evolution of the distribution function that we have found is very similar to the one at 4 GeV in Fig. 3; we have additionally checked that the qualitative result is unaffected by changing the gluon energy density. We have done this by tuning the $g^2\mu_{\rm Pb}$ to some lower value [the energy density in the MV model scales approximately as $(g^2\mu)^4$ at fixed lattice spacing], but the time evolution of the distribution function is unaffected by this and the lower gluon density results merely in a slower momentum diffusion. We think that the lack of drag for high p_T is due to the lack of the gluon radiation in the equations of motion. While we plan to add this term in a future work, we comment here that we do not expect gluon radiation to play a dominant role at low p_T; as a matter of fact, it is known that gluon radiation becomes important for high p_T, so while this mechanism will affect the energy loss of high-p_T charm and beauty quarks, it should not affect considerably the results that we have already found at low p_T.

The drag and diffusion of the c-quarks in momentum space has an effect on the nuclear modification factor of D-mesons, defined as

$$R_{\rm pPb} = \frac{(dN/d^2p_T)_{\rm evolved}}{(dN/d^2p_T)_{\rm prompt}}, \tag{20}$$

where the prompt spectrum is given by Eq. (15) after fragmentation and $(dN/d^2p_T)_{\rm evolved}$ corresponds to the spectrum obtained by fragmentation of the c-quark spectrum after the evolution in the glasma fields. In Fig. 4 we plot the nuclear modification factor for the D-mesons that we obtain within our calculation. The result is shown for two values of $g^2\mu_{\rm Pb}$ for the Pb nucleus at $\sqrt{s} = 5.02$ TeV, namely, $g^2\mu_{\rm Pb} = 3.4$ GeV (dashed blue line) and $g^2\mu_{\rm Pb} = 5.2$ GeV (solid green line), as discussed in the previous section. Experimental data correspond to the backward rapidity region (namely to the proton side) obtained by the LHCb collaboration [55]. We remark that although we show experimental data here, we do not aim to a precise fit of these by our calculation because we miss the longitudinal expansion: data are shown only to quantify the order of magnitude of our result, while a closer comparison with data will be the subject of a forthcoming study. We have chosen to show these data rather than the averaged published by the ALICE collaboration because those are an average of the forward and backward rapidity region, and in this case the CNME are very important and should be included in our initial state. We have checked however that

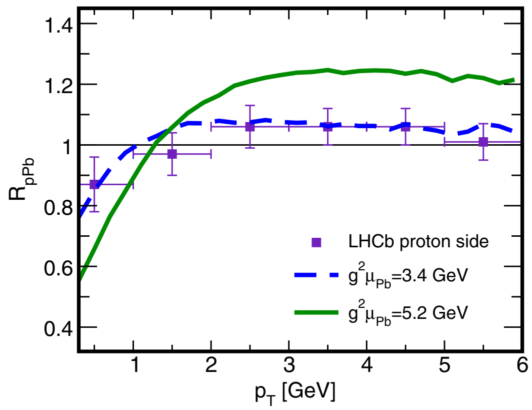

FIG. 4. Nuclear suppression factor versus p_T. We plot the results for two values of $g^2\mu$ for the Pb nucleus, namely, $g^2\mu = 3.4$ GeV (orange dot-dashed line) and $g^2\mu = 5.2$ GeV (orange dashed line). The solid green line corresponds to the R_{pPb} for the D-mesons, obtained assuming a standard fragmentation scenario for the c-quarks. Simulations have been stopped at $t = 1$ fm/c. Data correspond to the backward rapidity side (namely the proton side) of the LHCb collaboration [55].

including these effects in the initial state does not affect the drag and diffusion of c-quarks in the evolving glasma (results will be reported elsewhere).

Figure 4 is the main result of the present article: it shows that R_{pPb} can get a substantial deviation from one because of the interaction of the c-quarks with the evolving gluon fields in the glasma in the very early stage of a high energy p-Pb collision. As explained above, this result is due to the diffusion of heavy quarks in momentum space accompanied by a drag of the low p_T quarks towards higher momenta. The net effect that we find is very different from what is usually discussed in the heavy quark community, namely energy loss. In fact, our results suggest that in the very early stage heavy quarks can gain energy rather than lose it, because they are formed almost immediately after the collision and probe the strong gluon fields of the glasma, while energy loss will be substantial only in the presence of a medium, namely, of the quark-gluon plasma that forms in a later stage. Most likely, this energy gain can be understood even in simpler terms considering that low and intermediate p_T heavy quarks are injected at the formation time into a system with a very large energy density; therefore it appears natural that during their propagation they get energy rather than lose it.

This effect is interesting not only for its straightforward application to heavy quarks; as a matter of fact, since it comes from the propagation in the strong gluon fields of the evolving glasma, the c-quarks probe these fields. The fact that the qualitative shape of our R_{pPb} resembles that measured in experiments might suggest that at least part of the measured R_{pPb} comes from the propagation of the c-quarks in the glasma, and might be considered as the signature of the glasma itself. In this regard a more quantitative statement will be put in a forthcoming article when the longitudinal expansion will be included, and the amount of this effect will be compared to CNME.

We dub the effect summarized in Figs. 2 and 4 as the *cathode tube effect*. The reason for this name is easy to understand. As a matter of fact, the cathode tubes are devices in old televisions, in which an electron cannon ejects electrons and these are accelerated and deflected by electric field before they hit a fluorescent screen. *Mutatis mutandis*, the same effect, takes place in the early stage of high energy p-Pb collisions; indeed, here (color-)electric fields accelerate the prompt c-quarks that are injected into the bulk by the inelastic collisions among the protons on the one hand and the nucleons in Pb on the other hand (using this analogy, the electron cannon is here replaced by the p and Pb projectiles).

IV. CONCLUSIONS AND OUTLOOK

We have studied consistently the propagation of c-quarks in the evolving strong gluon fields allegedly produced in high energy p-Pb collisions. As the initial condition we have taken the standard glasma with longitudinal color-electric and color-magnetic fields, adapted in order to take into account the finite size of the system; for the initialization of the c-quarks we have considered the standard FONLL perturbative production tuned in order to reproduce the D-mesons spectrum in proton-proton collisions. We have set up the saturation scale for both the proton and the Pb nucleus in order to reproduce the expected one at $\sqrt{s} = 5.02$ TeV; for this reason, even if we do not include the longitudinal expansion in the calculation, we discuss the gluon fields produced in p-Pb collisions at this energy.

We have computed the nuclear modification factor, R_{pPb}, for these collisions; the result is summarized in Fig. 4. Although we do not aim to reproduce the experimental data because of the lack of the longitudinal expansion, we have found that the qualitative shape of our R_{pPb} resembles that measured by the LHCb collaboration on the proton side. Leaving a detailed comparison with experimental data to a future project, for the time being, we emphasize that the propagation of c-quarks in the evolving glasma has only been partly studied within a small transferred momentum approximation and assuming a static gluonic medium [53], so this article aims to start to fill this gap and paves the way for more complete studies. The lack of longitudinal expansion in our calculations might affect quantitatively the effect on the modification factor since the dilution of the energy density at later times will lead to the lowering of the average magnitude of the fields, an effect that is not considered in the present work. However, it is well known that even in the case of the expansion the transverse fields form quickly, and their magnitude remains substantial for a proper time range of ≈ 1 fm/c. Therefore, on the basis of the existing results we expect that the effect on the nuclear modification factor will be present also in the case of the expanding geometry. We also mention that initial fluctuations can be added [33], and these bring transverse fields already at the initial time; these should affect the propagation of heavy quarks and might balance in part the effect of the expansion.

We remark that we have not assumed the formation of a hot medium, namely the QGP, in this calculation. Indeed, although there is a lot of evidence that the QGP is formed in Pb-Pb collisions, such a strong evidence is missing at the moment for p-Pb collisions. We will consider more closely this problem in the future, by coupling our evolution of the c-quark spectrum to relativistic transport and to Langevin dynamics, in order to estimate quantitatively the effect of a hot medium on R_{pPb}.

We have preliminarily studied the effect of the propagation of the c-quarks in the evolving glasma in the case of Pb-Pb collisions at the LHC energy. In this case we have checked that a propagation for approximately 0.3 fm/c, which is a standard initialization time for QGP in relativistic transport and hydro simulations, is enough to obtain a substantial effect. Although the R_{PbPb} in this case cannot be compared directly with the experimental data due to the much longer propagation in the hot QGP, the effect of the early propagation in the gluon fields should not be ignored. Again, we will couple our results to relativistic transport [46,77] in the near future in order to quantify how the tilting of the c-quarks spectrum produced in the preequilibrium phase affects the late stage dynamics of heavy quarks in Pb-Pb collisions.

We have not included for simplicity the effects of CNME on the prompt spectrum in our calculations. Because of the lack of these effects in the present calculation, most likely the cathode tube effect is more relevant for the proton side of the p-Pb collision in which shadowing and/or gluon saturation should not give a substantial contribution. In fact, experimental data show that suppression of R_{pPb} is more pronounced on the Pb side, where both shadowing and gluon saturation are expected to give substantial deviations from the perturbative QCD prompt production of c-quarks. In addition to these, we have not included here the quantum fluctuations on the top of the glasma: these fluctuations add a transverse electric field at the initial time therefore they will enhance the cathode tube effect.

In the future it will be interesting to address the question whether the cathode tube can contribute to the collective flows. In fact, the cathode tube is the result of the interaction of the heavy quarks with the evolving background intense gluon field: if an initial anisotropy in the transverse plane coordinates exists at the initial time (this is easy to obtain due to event-by-event fluctuations in the color charge distribution in the proton) and if the gluon system expands along the directions of the highest pressure gradients, then it might be possible that a hydrodynamic flow develops in the transverse plane as well as some amount of v_2 for the heavy quarks (the same would be true for v_3, v_4, and so on). We will consider all these important effects in forthcoming studies.

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