Dynamical heavy-quark recombination and the non-photonic single electron puzzle at RHIC

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Abstract. We show that the single, non-photonic electron nuclear modification factor R^e_{AA} is affected by the thermal enhancement of the heavy-baryon to heavy-meson ratio in relativistic heavy-ion collisions with respect to proton-proton collisions. We make use of the dynamical quark recombination model to compute such ratio and show that this produces a sizable suppression factor for the R^e_{AA} at intermediate transverse momenta. We argue that such suppression factor needs to be considered, in addition to the energy loss contribution, in calculations of R^e_{AA}.

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INTRODUCTION

The suppression of single, non-photonic electrons at RHIC [1, 2] is usually attributed to heavy-quark energy losses. However, calculations that successfully describe the nuclear modification factor of charged hadrons fail to describe the single, non-photonic electron nuclear modification factor R^e_{AA} [3, 4, 5]. This has prompted a great deal of effort aimed to better describe the heavy-quark energy loss mechanisms to include not only the radiative part but also the collisional and the medium dynamical properties to compute the radiative piece. As a result, although some improvement in the description of the nuclear modification factor has been gained, it is not yet clear whether the anomalous suppression can be completely attributed to energy losses.

Working along a complementary approach to describe the non-photonic electron yield at RHIC, it has been argued [6, 7] that under the assumption of an enhancement in the heavy-quark baryon to meson ratio, analogous to the case of the proton to pion and the Λ to kaon ratios in Au+Au collisions [8, 9, 10, 11], it is possible to achieve a larger suppression of the nuclear modification factor. The rationale behind the idea is that heavy-quark mesons have a larger branching ratio to heavy-quark baryons, and therefore, when the former are less copiously produced in a heavy-ion environment, the nuclear modification factor decreases, even in the absence of heavy quark energy losses in the plasma.

In a recent work [12], we have quantitatively address this question by making use of a dynamical recombination scenario that accounts for the fact that the probability to form baryons and mesons can depend on a different way on the evolving density during the collision. Here, we present the main points of that work and enhance the discussion. A coalescense model addressing the same goals has been recently presented in Ref. [13].

Recently, following the approach described here, the nuclear modification factor has been estimated [14]. It has also been found that non-negligible contributions from higher-twist processes in large p_T hadron production, may indicate that the recombination mechanism can still be important in that regime [15].

In order to give a qualitative argument that shows how an enhancement in the heavy-quark baryon to meson ratio can suppress the single, non-photonic electron nuclear modification factor, let us look at the p_T integrated R^e_{AA} and to consider that the heavy hadrons are only those containing a single charm,

$$R^{e\ p_{T\,\mathrm{int}}}_{AA} = \frac{1}{\langle n_p \rangle} \frac{N^{\Lambda}_{AA} B^{\Lambda\to e} + N^{D}_{AA} B^{D\to e}}{N^{\Lambda}_{pp} B^{\Lambda\to e} + N^{D}_{pp} B^{D\to e}}, \tag{1}$$

where $\langle n_p \rangle$ is the average number of participants in the collision for a given centrality class, $N^x_{AA\ (pp)}$, refers to the number of x-particles produced in $A+A$ ($p+p$) collisions and $B^{x\to e}$ is the branching ratio for the inclusive decay of x-particles into electrons.

We can bring Eq. (1) into a form that contains the corresponding p_T integrated nuclear modification factor for particles containing charm as:

$$\begin{aligned} R^{e\ p_{T\,\mathrm{int}}}_{AA} &= \frac{1}{\langle n_p \rangle} \left(\frac{N^D_{AA} + N^{\Lambda}_{AA}}{N^D_{pp} + N^{\Lambda}_{pp}} \right) \\ &\times \left(\frac{1 - N^{\Lambda}_{AA}/(N^D_{AA} + N^{\Lambda}_{AA})}{1 - N^{\Lambda}_{pp}/(N^D_{pp} + N^{\Lambda}_{pp})} \right) \left(\frac{1 + CxN^{\Lambda}_{pp}/N^D_{pp}}{1 + xN^{\Lambda}_{pp}/N^D_{pp}} \right), \end{aligned} \tag{2}$$

where

$$C = \frac{N^{\Lambda}_{AA}/N^D_{AA}}{N^{\Lambda}_{pp}/N^D_{pp}}, \qquad\qquad x = \frac{B^{\Lambda\to e}}{B^{D\to e}}. \tag{3}$$

C represents the *enhancement factor* for the ratio of charm baryons to mesons in $A+A$ as compared to $p+p$ collisions and x is the charm baryon to meson relative branching ratios for their corresponding inclusive decays into electrons.

When not integrated over transverse momentum, $1/\langle n_p \rangle \left[(N^D_{AA} + N^{\Lambda}_{AA})/(N^D_{pp} + N^{\Lambda}_{pp}) \right]$ represents the nuclear modification factor for particles with charm. Let us not assume any particular value for this factor and instead concentrate in the other one in Eq. (2), which can be written as

$$T^{e\ p_{T\,\mathrm{int}}}_{AA} = \frac{(1+a)(1+Cxa)}{(1+Ca)(1+xa)}, \tag{4}$$

where $a = N^{\Lambda}_{pp}/N^D_{pp}$. The above quantity is plotted in Fig. 1 as a function of x for different combination of C and a. Notice that the function $T^{e\ p_{T\,\mathrm{int}}}_{AA}$ is less than one when $x < 1$ provided that $Ca > a$.

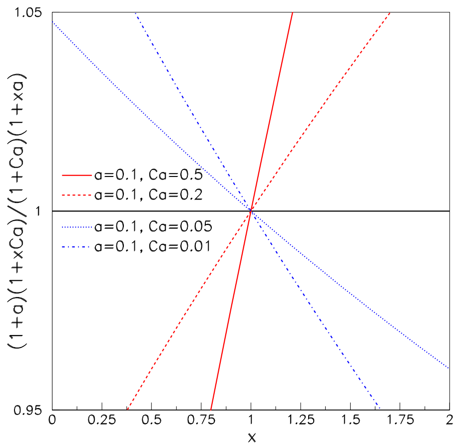

FIGURE 1. (Color online) p_T integrated T_{AA}^{e} as a function of x, the ratio of branching ratios for charmed baryons and mesons to decay inclusively into electrons. Notice that for $x < 1$, $T_{AA}^{e} < 1$ when Ca –the ratio of charm baryons to mesons in A + A– is larger than a –the ratio of charm baryons to mesons in p + p–.

We want to quantitatively address the question of whether the enhancement factor C times a –namely, the heavy-baryon to heavy-meson ratio in Au + Au collisions– can indeed be larger than a –namely, the heavy-baryon to heavy-meson ratio in p + p collisions– and if so, how this affects the behavior of the factor T_{AA}^{e} as a function of p_T.

The rest of the work is organized as follows: After presenting a brief introduction to the dynamical quark recombination model in Sec. II, we proceed in Sec. III to compute the probabilities to form mesons and baryons containing a heavy quark in a relativistic heavy-ion collision environment. In Sec. IV we use these probabilities to write expressions for the meson and baryon transverse momentum distributions. In Sec. V we compute such distributions as well as the baryon to meson ratio. We convolute such ratio with the branching ratios of charmed baryons and mesons to decay into electrons to obtain the p_T unintegrated function T_{AA}^{e} and show that this can be indeed less than 1. Finally we summarize and conclude in Sec. VI.

DYNAMICAL QUARK RECOMBINATION

It has been shown [17] that the features of the proton to pion ratio can be well described by means of the so called *dynamical quark recombination model* that incorporates how the probability to recombine quarks into mesons and baryons depends on density and temperature. The upshot of the model is that this probability differs for hadrons made up by two and three constituents *with the same mass*, that is to say, the relative population of baryons and mesons can be attributed not only to flow but rather to the dynamical properties of quark clustering in a varying density scenario. A natural question is whether those features remain true for baryons and mesons with one constituent heavy-quark and whether a computed, as opposed to assumed, baryon to meson ratio, can at least partially explain the anomalous single, non-photonic electron suppression at RHIC.

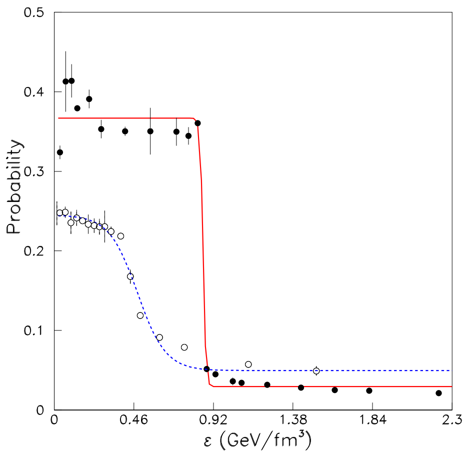

FIGURE 2. (Color online) Probabilities $\mathscr{P}^{B,M}$ to produce charmed baryons and mesons as a function to the energy density ε. Shown are the results of the Monte Carlo simulation for baryons (full circles) and mesons (open circles) together with a fit to these.

The invariant transverse momentum distribution of a given hadron can be written as an integral over the freeze-out, space-time hypersurface Σ of the relativistically invariant phase space particle density $F(x,P)$,

$$E\frac{dN}{d^3P} = g\int_{\Sigma_f} d\Sigma\,\frac{P\cdot u(x)}{(2\pi)^3}F(x,P)\,, \tag{5}$$

where P is the hadron's momentum and $u(x)$ is a future oriented unit four-vector normal to Σ and g is the degeneracy factor for the hadron which takes care of the spin degree of freedom. The function $F(x,P)$ contains the information on the probability that the given hadron is formed.

To allow for a dynamical recombination scenario in a thermal environment, let us assume that the phase space particle density $F(x,P)$ can be factorized into the product of a term containing the thermal occupation number, including the effects of a possible flow velocity, and another term containing the system energy density ε driven probability $\mathscr{P}(\varepsilon)$, for the coalescence of partons into a given hadron. We thus write

$$F(x,P) = e^{-P\cdot v(x)/T}\,\mathscr{P}(\varepsilon)\,, \tag{6}$$

where $v(x)$ is the flow velocity. As we will show, the probability $\mathscr{P}(\varepsilon)$ incorporates in a simple manner the information that the partons that coalesce need to be close in configuration space as well as to have a not so different velocity.

To compute the probability $\mathscr{P}(\varepsilon)$, it has been shown in Ref. [17] (where we refer the reader to for details) that use can be made of the *string flip model* in order to get information about the likelihood of clustering of constituent quarks to form hadrons from an effective quark-quark interaction. In short, the model is a variational quantum Monte Carlo simulation that, taking a set of equal number of all color quarks and antiquarks at a given density, computes the optimal configuration of colorless clusters (baryons or mesons) by minimizing the potential energy of the system. At low densities, the model describes the system of quarks as isolated hadrons while at high densities, this system becomes a free Fermi gas of quarks.

We consider N quarks moving in a three-dimensional box and are described by a variational wave function of the form: $\Psi_\lambda = exp(-\lambda V)\Phi_{FG}$, where λ is the single variational parameter, V is the potential to build either mesons or baryons respectively, and Φ_{FG} is the Fermi-gas wave function given by a product of Slater determinants, which are built up of single-particle wave functions describing a free particle in a box. The square of the variational wave function is the weighting probabilty in the sampling, wich we carry out using metropolis algorithm.

We can identify the value of the variational parameter λ as being directly proportional to the probability to form a cluster. This fact will be latter exploited to define the density dependent probability $\mathscr{P}(\varepsilon)$ since, as we show below, λ changes from a fixed value at low density (isolated clusters) to zero at high density (Fermi gas).

PROBABILITIES

All the results we present here come from simulations made with 384 particles, 192 quarks and 192 antiquarks, corresponding to having 32 u $(\bar{u})$ plus 32 c $(\bar{c})$ quarks (antiquarks) in the three color charges (anti-charges). The number of quarks corresponds to the second closed shell of a three-dimensional box. The equal amount of light and heavy quarks used in the simulation is not intended to represent the whole system but rather the fraction which will drive the relative recombination. The rest of the light quarks will contribute to determine the thermodynamical properties of the system. To take into account the mass difference between u and c quarks we set $m_c = 10M_u$. We have checked that variations of this particular choice do not affect our relative probabilities.

To determine the variational parameter as a function of density we first select the value of the particle density ρ in the box. Then, we compute the energy of the system as a function of the variational parameter using the Monte Carlo method described in the previous section. The minimum of the energy determines the optimal variational parameter. To get a measure of the probability to form a cluster , we take the variational parameter and divide it by its corresponding value at the lowest density. Notice that since the heavy quarks are not as abundant as the light ones, they do not contribute to the energy density and thus, within the model, this last can be computed by assuming that only light flavors contribute.

In order to find an appropriate measurement of the probability to form baryons and mesons, we multiply this variational parameters by the likelihood to find clusters of baryons made up of two-light, one-heavy quark and mesons made up of one-light, one-heavy quark. This likelihood has to consider the fact that the thermal plasma is mainly made up of light quarks and thus that the number of produced heavy quarks is relatively small. To accomplish this, notice that in a model where the interaction between quarks to form clusters is flavor (as well as color) blind, this likelihood should account only for the combinatorial probabilities.

Consider the case where one starts with a set of n u-quarks, n $\bar{u}$-antiquarks, m c-quarks and m $\bar{c}$-antiquarks, each coming in three colors (we impose that the number of u-quarks be a multiple l of the number of c-quarks, namely, $n = lm$) Using the number of possible colorless (anti)mesons and (anti)baryons that can be formed we can compute the *relative* abundance of baryons with respect to mesons computed under the above assumptions on the number of light and heavy quarks that we start from. Since in the case of mesons we are allowing to consider the case $u\bar{c}$ as well as $\bar{u}c$, we need to include in the counting of the groups of three quarks also the antibaryons. Thus the relative abundance is

$$\frac{c-\text{baryons}+c-\text{antibaryons}}{c-\text{mesons}+c-\text{antimesons}} = \frac{3l}{2(l+1)} \overset{l\to\infty}{\longrightarrow} \frac{3}{2}, \tag{7}$$

since, in the plasma, the number of u-quarks greatly exceeds the number of c-quarks. It can be checked that the assymptotic value 3/2 is rapidly reached, for instance, by taking $l = 30$, the above fraction already becomes 1.475.

Figure 2 shows the probability parameter $\mathscr{P}^{\mathscr{B},\mathscr{M}}(\varepsilon)$ for baryons and mesons,obtained by multiplying the variational parameter with the corresponding fraction of baryon/meson formed at the given energy density. In the case of mesons it corresponds to 1/4 irrespective of the density, while for baryons it has a functional form, since the kind of clusters can be different as density increases. For low densities the ratio of the probabilities becomes 3/2, as expected from the combinatorial described above. Shown in the figure is also a fit to the variational parameters with the functional form

$$f(x) = a_1 + \frac{a_2}{1+exp[(x-x_0)/dx]}. \tag{8}$$

For Baryons $a_1^B = 0.0294$, $a_2^B = 0.3374$, $x_0^B = 0.8604$, $dx^B = 0.0078$, whereas for mesons $a_1^M = 0.0496$, $a_2^M = 0.1953$, $x_0^M = 0.4812$, $dx^M = 0.0813$. We will use this analytical expression to carry out the calculations.

BARYON TO MESON RATIO

In order to quantify how the different probabilities to produce sets of three quarks as compared to sets of two quarks affect the particle's yields as the energy density changes during hadronization, we need to resort to a model for the space-time evolution of the collision. We take Bjorken's scenario which incorporates the fact that initially, expansion is longitudinal, that is, along the beam direction which we take as the $\hat{z}$ axis and include transverse flow as a small effect on top of the longitudinal expansion. In this scenario, the relation between the temperature T and the 1+1 proper-time τ is given by

$$T = T_0 \left(\frac{\tau_0}{\tau}\right)^{v_s^2}, \tag{9}$$

where $\tau = \sqrt{t^2 - z^2}$. Equation (9) assumes that the speed of sound v_s changes slowly with temperature. For simplicity we take v_s as a constant equal to the ideal gas limit $v_s^2 = 1/3$. We also consider that hadronization takes place on hypersurfaces Σ characterized by a constant value of τ and therefore $d\Sigma = \tau\rho\ d\rho\ d\phi\ d\eta$, where η is the spatial rapidity and ρ, ϕ are the polar transverse coordinates. Thus, the transverse spectrum for a hadron species H is given as the average over the hadronization interval of the right hand-side of Eq. (5), namely

$$E\frac{dN^H}{d^3P} = \frac{g}{\Delta\tau}\int_{\tau_0}^{\tau_f} d\tau \int_{\Sigma} d\Sigma\, \frac{P\cdot u(x)}{(2\pi)^3} F^H(x,P), \tag{10}$$

where $\Delta\tau = \tau_f - \tau_0$.

To find the relation between the energy density ε –that the probability $\mathscr{P}$ depends upon– and T, we resort to lattice simulations. For the case of two flavors (since the heavy quark does not thermalize), a fair representation of the data [16] is given by the analytic expression

$$\varepsilon/T^4 = a\left[1+\tanh\left(\frac{T-T_c}{bT_c}\right)\right], \tag{11}$$

with $a = 4.82$ and $b = 0.132$. We take $T_c = 175$ MeV.

Considering the situation of central collisions, we assume that there is no dependence of the particle yield on the transverse polar coordinates. Furthermore we consider that the space-time and momentum rapidities are completely correlated. Under these assumptions it is possible to write the hadron transverse momentum distribution as

$$\frac{dN}{p_T\,dp_T\,dy} = g\frac{m_T\Delta y}{4\pi}\frac{\rho_{\text{nucl}}^2}{\Delta\tau}\int_{\tau_0}^{\tau_f} \tau d\tau \mathscr{P}(\tau) I_0(p_T \sinh\eta_T/T)e^{-\cosh\eta_T/T}. \tag{12}$$

where ρ_{nucl} is the radius of the colliding nuclei and I_0 is the Bessel function I of order zero, y is the $1+1$ momentum rapidity.

Armed with the expression, we now proceeed to apply the analysis to the computation of the charmed meson and baryon distributions.

RESULTS

Figure 3 shows examples of the transverse momentum distributions for mesons and baryons obtained from Eq. (12). We set the mases of the charmed baryons and mesons as $m^B = 2.29$GeV (corresponding to Λ_c) and $m^B = 1.87$GeV (corresponding to D). We take the initial hadronization time as $\tau_0 = 1$fm, at an initial temperature $T_0 = 200$MeV and the final hadronization temperature as $T_f = 100$MeV, corresponding, according to Eq. (9), to a final time $\tau_0 = 8$fm. Shown are the cases with $v_T = 0$ and $v_T = 0.4$. Notice that a finite transverse expansion velocity produces a broadening of the distributions, as expected.

Figure 4 shows the charmed baryon to meson ratio obtained from the ratio of the above transverse momentum distributions. Shown is the range for this ratio when varying the transverse expansion velocity v_T from 0 to 0.4. Notice that for a finite v_T, this ratio goes above 1 for $p_T 3.5$GeV.

We now proceed to compute the p_T unintegrated function T^e_{AA}. For this purpose, we take that the posible charmed mesons decaying inclusively into electrons or positrons

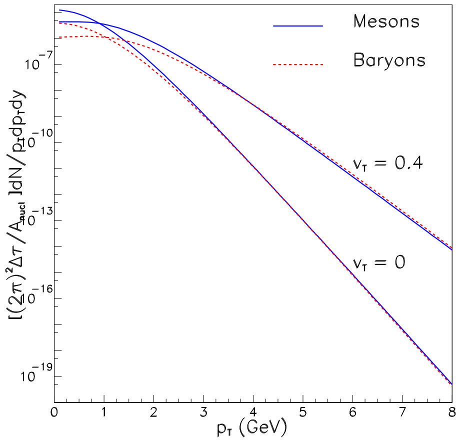

FIGURE 3. (Color online) Charmed baryon and meson transverse momentum distributions. The parameters used in the calculation are $m^B = 2.29$ GeV, $m^M = 1.87$ GeV, $\tau_0 = 1$ fm, $T_0 = 200$ MeV, $T_f = 100$ MeV, corresponding to a final time $\tau_f = 8$ fm. Shown are the cases with $v_T = 0$ and $v_T = 0.4$.

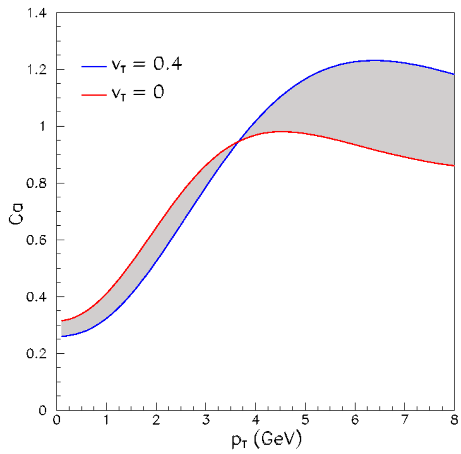

FIGURE 4. (Color online) Charmed baryon to meson ratio, Ca, as a function of transverse momentum. The parameters used in the calculation are $m^B = 2.29$ GeV, $m^M = 1.87$ GeV, $\tau_0 = 1$ fm, $T_0 = 200$ MeV, $T_f = 100$ MeV, corresponding to a final time $\tau_f = 8$ fm. Shown is a range when varying the transverse expansion velocity v_T from 0 (upper curve at low momenta) to 0.4 (lower curve at low momenta).

are $D^{\pm}(B^{D^{\pm}\to e\pm} = 16.0\%)$, D^0, $D^0(B^{D^0,D^0\to e\pm} = 6.53\%)$, $D_s^{\pm}(B^{D_s^{\pm}\to e\pm} = 8\%)$ and that the possible charmed baryons decaying inclusively into electrons or positrons are Λ_c, $\Lambda_c(B^{\Lambda_c\Lambda_c\to e\pm} = 4.5\%)$. Thus, we get $x = 0.14$. We also approximate the masses of all the charmed mesons considered to be equal to the mass of the $D^{\pm}$ mesons.

From Eq. (4) we see that, without integrating over p_T, the dependence of the transverse momentum comes from $a = (dN^{\Lambda}_{pp}/dp_T)/(dN^{D}_{pp}/dp_T)$ and the product $Ca = (dN^{\Lambda}_{AA}/dp_T)/(dN^{D}_{AA}/dp_T)$. The integrated ratio a^{int} has been computed in Ref. [7] using

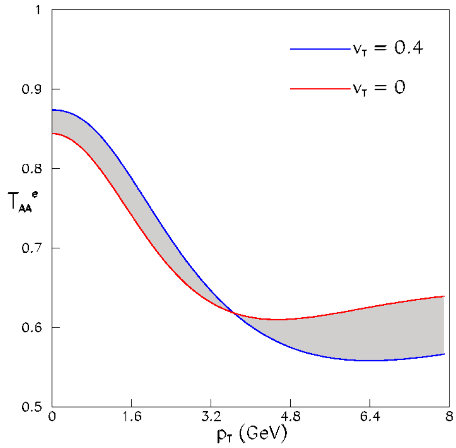

FIGURE 5. (Color online) Suppression factor T_{AA}^e as a function of transverse momentum. The parameters used in the calculation are $m^B = 2.29$ GeV, $m^M = 1.87$ GeV, $\tau_0 = 1$ fm, $T_0 = 200$ MeV, $T_f = 100$ MeV, corresponding to a final time $\tau_f = 8$ fm, $x = 0.14$, a=0.073. Shown is a range for the transverse expansion velocity form $v_T = 0$ (upper curve at low p_T) and $v_T = 0.4$ (lower curve at low p_T.

a Pythia simulation, with the result $a^{int} = 0.073$. We have also performed a simulation using Pythia at NLO with 100,000 events and have found that with such statistics, the ratio of charmed baryons to charmed mesons in p + p collissions at $\sqrt{s_{NN}} = 200$ GeV is flat up to $p_T \approx 5$ GeV and consistent with the value reported in Ref. [7]. Therefore, for simplicity we take a as a constant equal to the above quoted number. Thus

$$T_{AA}^e \frac{(1+a^{int}}{(1+xa^{int}} \frac{1+x(dN_{AA}^{\Lambda}/dp_T)/(dN_{AA}^D/dp_T)}{1+(dN_{pp}^{\Lambda}/dp_T)/(dN_{pp}^D/dp_T)}. \tag{13}$$

Figure 5 shows T_{AA}^e as a function of p_T. We have used a range of values for the transverse expansion velocity between $v_T = 0$ and $v_T = 0.4$. We see that for the chosen evolution parameters, T_{AA}^e is indeed smaller than 1 and thus it contributes to the suppression of the single non-photonic electron nuclear modification factor R_{AA}^e.

CONCLUSIONS

In this work we have shown that the anomalous suppression of the single non-photonic electron nuclear modification factor R_{AA}^e can be partially understood by realizing that this quantity is affected by an enhancement in the charmed baryon to meson ratio at intermediate p_T in Au + Au collisions. This enhancement is due to the fact that in this region, thermal recombination becomes the dominant mechanism for hadron production. We have made used of the DQRM to calculate this ratio and have shown that for moderate and even for vanishing transverse expansion velocities, it indeed can be larger than the charmed baryon to meson ratio in p + p collisions. This enhancement in turn produces that the T_{AA}^e is below 1 and thus contributes to the suppression factor introduced by considering energy losses due to the propagation of heavy flavors in the plasma.

It is worth to keep in mind some important features concerning the results of the present calculation: First, notice that we have not included the momentum shift introduced by energy losses when computing the transverse distributions of charmed mesons and baryons. This is so because for R^e_{AA}, energy losses should be included in the prefactor of the function T^e_{AA}. In this sense, in order to avoid a double counting of the effect, the ratio that goes into the calculation of this last function is the raw ratio.

Second, it is expected that at some value of p_T , fragmentation becomes the dominant mechanism for hadron production and therefore that the charmed baryon to meson ratio decreases above that p_T value, given that fragmentation produces more mesons than baryons. Third, we have considered finite values of transverse flow for charmed mesons and baryons even thought it might be questionable that heavy flavors also flow as light flavors do. Nevertheless, there seems to be some experimental support for heavy quark flow [18]. In this sense, the flow strength range we have considered is only for moderate values. Notice however that even in the absence of flow the suppression factor keeps being less than 1. Some of these issues will be the subject of a future work to appear elsewhere.

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