Collisional dissociation of heavy mesons in dense QCD matter

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In the framework of the reaction operator approach we calculate and resum the multiple elastic scattering of a fast $q\bar{q}$ system traversing dense nuclear matter. We derive the collisional broadening of the meson’s transverse momentum and the distortion of its intrinsic light cone wave function. The medium-induced dissociation probability of heavy mesons is shown to be sensitive to the opacity of the quark-gluon plasma and the time dependence of its formation and evolution. We solve the system of coupled rate equations that describe the competition between the fragmentation of $c$- and $b$-quarks and the QGP-induced dissociation of the $D$- and $B$-mesons to evaluate the quenching of heavy hadrons in nucleus-nucleus collisions. In contrast to previous results on heavy quark modification, this approach predicts suppression of $B$-mesons comparable to that of $D$-mesons at transverse momenta as low as $p_T \sim 10$ GeV. It allows for an improved description of the large attenuation of non-photonic electrons in central Au+Au reactions at RHIC.

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I. INTRODUCTION

One of the important experimental signatures of the Quark Gluon Plasma (QGP) creation in heavy ion collisions is the detailed suppression pattern of high transverse momenta hadrons [1, 2]. Jet quenching for light quark hadrons in nucleus-nucleus collisions. In contrast to previous results on heavy quark modification, this approach predicts suppression of $B$-mesons comparable to that of $D$-mesons at transverse momenta as low as $p_T \sim 10$ GeV. It allows for an improved description of the large attenuation of non-photonic electrons in central Au+Au reactions at RHIC.

In Eq. (1), $p^+$ is the large light cone momentum of the parton, $z = k^+/p^+$ and $|k| \sim \Lambda_{QCD} \sim 200$ MeV is the deviation from collinearity. We can estimate from the uncertainty principle for the variable conjugate to the non-conserved light cone momentum component $\Delta p^- = (p^-)' - (p^-)$:

$$\Delta y^+ \approx \frac{1}{\Delta p^-} = \frac{2z(1-z)p^+}{k^2 + (1-z)m_h^2 - z(1-z)m_Q^2}.$$ (2)

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With $\Delta y^+ = \tau_{\text{form}} + z_{\text{form}}$, $z_{\text{form}} = \beta_Q \tau_{\text{form}}$ being the formation distance, the formation time reads:

$$\tau_{\text{form}} = \frac{\Delta y^+}{1 + \beta_Q}, \quad \beta_Q = \frac{p_Q}{E_Q}. \quad (3)$$

Clearly, the higher the energy, or $p_T$, and the lighter the hadron - the better the perturbative treatment of partonic energy loss prior to hadronization will be $^{17}$. For a $p_T = 10$ GeV pion at mid-rapidity $\tau_{\text{form}} \approx 25$ fm $\gg L_{QGP}^T$, consistent with the jet quenching assumptions and the robust description of the nuclear modification factor $R_{AA}(p_T)$ versus $\sqrt{s_{NN}}$ and centrality $^{[1,3,4]}$. In contrast, $B$- and $D$-mesons of the same $p_T$ have formation times $\tau_{\text{form}} \approx 0.4$ fm and 1.6 fm, respectively, $\ll L_{QGP}^T$. For the purpose of these estimates we used $\bar{z}(\pi) = 0.7$, $\bar{z}(D) = 0.82$ and $\bar{z}(B) = 0.92$ and the exact hadron masses in Eq. $^{[2]}$. Therefore, at the finite $p_T$ range accessible at RHIC and LHC a conceptually different approach to the description of $D$- and $B$-meson quenching in A+A collisions is required, when compared to light hadrons. In the future in the limit of high transverse momenta, $p_T(D) > 10$ GeV and $p_T(B) > 30$ GeV, improved PQCD description of heavy flavor suppression should seek to also incorporate the traditional partonic energy loss mechanisms.

This Letter summarizes key results from our calculation of the collisional dissociation of heavy mesons in the quark-gluon plasma. In Section II we determine the baseline cross sections for heavy flavor production on the example of the identified $D^0$, $D^+$ and $B^+$ spectra from the Tevatron. The light cone wave functions for $D$- and $B$-mesons are constructed in Section III. In Section IV we use the Gyulassy-Levi-Vitev (GLV) reaction operator approach to derive the collisional broadening of the propagating $q\bar{q}$ system and evaluate the survival and dissociation probabilities of heavy mesons as a function of the cumulative transverse momentum transfer from the medium. In Section V we solve the system of coupled rate equations for the fragmentation into and dissociation of heavy mesons. Phenomenological comparison to RHIC data on the quenching of non-photonic electrons is presented. Our conclusions are given in Section VI.

II. $D$- AND $B$-MESON PRODUCTION IN P+P COLLISIONS

Single inclusive hadron production can be calculated to lowest order in the PQCD collinear factorization approach as follows:

$$\frac{d\sigma^H_{NN}}{dy d^2p_T} = K \sum_{abcd} \int d\mu_d d\mu_c \frac{1}{z^2} D_{H/c}(z) \alpha_s^2(\mu_T) |\mathcal{M}_{ab\rightarrow cd}|^2 \times \phi_{a/N}(x_a, \mu_T) \phi_{b/N}(x_b, \mu_T) x_a x_b S^2,$$  \quad (4)

In Eq. $^{(4)}$ $S = (p_A + p_B)^2$ is the squared center of mass energy of the hadronic collision, $x_a = p_a^+/p_A^+$, $x_b = p_b^-/p_B^-$ are the light cone momentum fractions of the incoming partons and $\phi_{a/N}(x_a), \phi_{b/N}(x_b)$ are the parton distribution functions $^{[18]}$. In our calculation the factorization and renormalization scales are set to $\mu^2_T = \mu^2_r = p_T^2 + m_Q^2$. The running of the strong coupling constant is taken to lowest order with $N_f = 4$ active quark flavors and we use $m_c = 1.3$ GeV, $m_b = 4.5$ GeV. Fragmentation functions $D_{H/c}(z)$ of heavy quarks into mesons, derived in $^{[19,20]}$, are used. At lowest order we include $Q + g \rightarrow Q + g, Q + q \rightarrow Q + q$ and $Q + \bar{q} \rightarrow Q + \bar{q}$ processes, using standard heavy quark PDFs, in addition to gluon fusion and light quark annihilation. We find that these processes give a dominant contribution to the single inclusive $D$- and $B$- meson cross section. They ensure faster convergence of the perturbation series and small next-to-leading order $K$-factors. Further details of the perturbative calculation are given in $^{[21]}$.

Figure 1 shows comparison of the PQCD single inclusive spectra, Eq. $^{(4)}$, to Tevatron measurements at $\sqrt{s} = 1.8$ GeV and 1.96 GeV $^{[22,23,24]}$. We observe that a similar description of the $B^+$ cross section can be achieved when compared to $D^0$ and $D^+$ cross sections and note that $b$-quark fragmentation functions are harder ($r = 0.07$) than the $c$-quark fragmentation functions ($r = 0.20$). Here, $r$ is the parameter that determines the hardness of these heavy quark fragmentation
functions \cite{19,20} and \( r_B/r_D \sim m_c/m_b \). We also anticipate that at \( p_T \leq m_b \) large corrections will arise to the LO perturbative calculation of \( B \)-mesons. However, in this region \( \sigma_B/\sigma_{dpT} < \sigma_D/\sigma_{dpT} \) has a smaller contribution to the heavy meson production rate. We use the same choice of scales and \( K = 1.5 \) as the baseline for studying nuclear effects on heavy flavor production at the smaller \( \sqrt{s_{NN}} = 200 \text{ GeV} \) at RHIC and the larger \( \sqrt{s_{NN}} = 5.5 \text{ TeV} \) at LHC. We note that \( K \)-factors cancel in the nuclear modification ratio \( R_{AA}(p_T) \). For baryons, which constitute \( \sim 10\% \) of the total \( c \)- and \( b \)-quark decays, we use softer fragmentation functions with \( r_{baryon} = 1.5 r_{meson} \). In the bottom panel of Fig.\cite{1} results for the electrons, \( 0.5(e^+ + e^-) \), from the semi-leptonic decays of heavy flavor, evaluated in LO PQCD in \( p+p \) collisions at RHIC, are shown. Data is from PHENIX \cite{23}. Harder fragmentation functions, \( r_B = 0.02, r_D = 0.06 \), have also been used for direct comparison to alternative calculations.

III. LIGHT CONE WAVE FUNCTIONS

In order to calculate the effect of the QGP medium on a meson traveling through, we need a model for the meson state in terms of its partonic degrees of freedom \((x_i,k_i)\). Here \( x_i = k_i^+/P^+ \), \( \sum_{i=1}^{n} x_i = 1 \) are the light cone momentum fractions and \( k_i \), \( \sum_{i=1}^{n} k_i = K \), are the internal parton transverse momenta. To determine the typical momenta \((k_i^2)\), we examine the numerical results of potential model calculations of \( D \)- and \( B \)-meson mass spectra and decay widths \cite{24}. Solving the Dirac equation for the relativistic light quark in the potential of the heavy quark \( V = -\xi/r + br \), one can achieve a very good description of the lowest-lying and excited heavy meson states. The radial wave function of the light quark \( r(r) = r^2|\psi(|r|)|^2 \) is found to have its maximum at \( a_0 = 2 - 3 \text{ GeV}^{-1} \) \cite{25}. Fourier transforming a Gaussian distribution, which features the same maximum of the radial wave function \( r(r) \), we determine the momentum distribution: \( \psi(r) \sim e^{r^2/(2a_0^2)} \rightarrow \psi(k) \sim e^{k^2 a_0^2/2}. \) Evaluation of the mean transverse momentum squared is straightforward:

\[
\langle k^2 \rangle = N_k^2 2\pi \int_{-1}^{1} d\cos(\theta) \cos^2(\theta) \times \int_{0}^{\infty} k^2 dk e^{k^2 a_0^2} = \frac{1}{2a_0^2}, \tag{5}
\]

where \( N_k^2 \) ensures the proper normalization of the distribution in momentum space. Taking \( a_0 = 2.5 \text{ GeV}^{-1} \) for the light parton we obtain \( \langle k^2 \rangle = 0.08 \text{ GeV}^2 \).

To determine the longitudinal momentum fractions of the quarks in a boosted heavy meson we express the state in terms of its multi-parton Fock components as follows \cite{28,29}:

\[
|\Psi_M, P^+, P \rangle = \sum_{n \geq 2} \int \prod_{i=1}^{n} \frac{dx_i d^2 k_i}{(2\pi)^3 \sqrt{2x_i}} \delta \left( \sum_{i=1}^{n} x_i - 1 \right) \times \delta^2 \left( \sum_{i=1}^{n} k_i \right) \psi(k_i, x_i) |n; k_i + x_i P, x_i P^+ \rangle, \tag{6}
\]

where the light cone wave functions \( \psi(k_i, x_i) \) of the partons in the meson are universal. Note that these do not depend on the external meson momenta, \( P^+, P \). We have chosen the following normalization for the single parton states:

\[
\langle x' P^+, k' ; 1 | 1; k, x P^+ \rangle = (2\pi)^3 2x \delta(x - x') \delta^2(k - k'). \tag{7}
\]

In our calculations we take the lowest lying \( n = 2 \) Fock component. A heavy meson moving in the positive light cone direction will then be described by

\[
|\psi(K, \Delta k; x, m_1, m_2) |^2 = \text{Norm}^2 e^{-\frac{\Delta k^2 + (4m^2 - 1) x + 4m^2 x^2}{4x(1-x)\Lambda^2}} \times \delta^2(K), \tag{8}
\]

where the overall large light cone momentum has been integrated out. In Eq. \cite{26} \( K = k_1 + k_2 \) and \( \Delta k = k_1 - k_2 \). We assume here that \( x \) is the momentum fraction carried by the heavy quark \( Q \) and use \( m_u = m_d = 0.005 \text{ GeV} \), \( m_c = 1.3 \text{ GeV} \) and \( m_b = 4.5 \text{ GeV} \). The light cone wave function satisfies the requirement that the maximum in the longitudinal momentum density distribution is achieved when the constituent partons of the meson are at the same rapidity or, equivalently, \( m_{T_1}/x_1 = m_{T_2}/x_2 \) \cite{26,29}. Fixing the light quark momentum fraction squared \( \langle (\Delta k^2/2) \rangle \) to the value obtained in Eq. \cite{27}, we can determine \( \Lambda = 0.735 \text{ GeV} \) and \( \Lambda = 1.055 \text{ GeV} \) for \( D \)-mesons and \( B \)-mesons, respectively.

Figure\cite{2} show the momentum density distribution of the heavy quark inside the heavy mesons from Eq. \cite{26} plotted versus the light cone momentum fraction \( x \) and its half relative transverse momentum \( \Delta k/2 = |\Delta k/2| \).
The following normalization, $\int d^2 \Delta k dx \ |\psi(x, \Delta k)|^2 = 1$, was used. We note the momentum distribution is fairly narrow in $\Delta k/2$ and momentum transfers from the medium with $\mu \geq 1$ GeV [1, 3, 4] at the initial stages of the evolution of the QGP density may easily dissociate the heavy mesons if they tend to form early $\ll L_T^{QGP}$, see Eqs. (11) and (12). Integrating over $\Delta k$ we obtain the parton distribution function of the heavy quark inside the meson:

$$\phi_{QM}(x) = \int d^2 \Delta k d^2 K |\psi(K, \Delta k; x, m_1, m_2)|^2 = \text{Norm}^2 4\pi x(1-x)\Lambda^2 e^{-\frac{(x(1-x)+m_2^2)}{2\Lambda^2}}. \quad (9)$$

We note that the heavy c- and b-quark distributions, Eq. (11), closely resemble in shape the fragmentation functions $D_{H/Q}(z)$ [13, 20] and peak toward larger values of $x$ with increasing heavy quark mass.

IV. COLLISIONAL DISSOCIATION OF HEAVY MESONS

The GLV reaction operator formalism [30, 31] was developed for calculating the induced radiative energy loss of hard quarks or gluons when they pass through a dense medium. In this approach, the multi-parton dynamics is described by a series expansion in $\chi = \int_0^{L_T^{QGP}} \sigma_{el}(z) \rho(z) dz = L_T^{QGP}/(\lambda)$, the mean number of interactions that a fast projectile undergoes along its trajectory. Each interaction is represented by a reaction operator that summarizes the unitarized basic scattering between the propagating system and the medium. The summation to all orders in opacity, $\chi$, is achieved by a recursion of the reaction operator. For the case of collisional interactions of individual fast partons in dense QCD matter, their diffusion in transverse momentum space has been derived to leading power [32] and leading power corrections [32].

The transverse momentum transfer $q_n$ at position $n$ to a fast parton when it scatters on the soft constituents of the medium is distributed according to the normalized differential cross section

$$\frac{1}{\sigma_{el}(n)} \frac{d\sigma_{el}(n)}{d^2 q_n} = \frac{\mu^2(n)}{\pi(q_n^2 + \mu^2(n))^2}. \quad (10)$$

In Eq. (10), $\mu(n) = gT(n)$ is the thermally generated Debye screening mass, $\sigma_{el}(n) \approx C_\mu 2\pi\alpha_s^2/\mu^2(n)$ with $C_\mu = 9/4, 1, 4/9$ for gg, gq, qq, respectively, and the mean free path $\Lambda(n) = 1/\sigma_{el}(n)\rho(n)$. Two momentum transfers are necessary to build one power of the elastic scattering cross section in Eq. (10), allowing for three $t = \infty$ on-shell cuts in the forward scattering Feynman diagrams, see Fig. 3. Momentum flow in the legs of the direct- and virtual-interaction terms is constrained as follows:

$$\text{Dir.} \sim \delta^2(q_n - q_n'), \quad \text{Vir.} \sim -\frac{1}{2} \delta^2(q_n + q_n'). \quad (11)$$

For further details on the derivation of this formalism, see [30, 31, 32, 33].

The diagrams that are relevant to a single in-medium interaction of the quark-antiquark system are shown in Fig. 3. We work in terms of the momenta $K$ and $\Delta k$, defined in the previous Section. If the momentum distribution of the $q\bar{q}$ system that has undergone $n$ scatterings is $\propto |M_n^*(K, \Delta k)M_n(K, \Delta k)|$, it can be related to the momentum density prior to the last collision as follows:

$$M_{n-1}^*(K, \Delta k) \left[ e^{-q_n \cdot \nabla K} e^{-q_n \cdot \nabla \Delta k} \otimes \left( e^{-q_n \cdot \nabla \Delta k} e^{-q_n \cdot \nabla ^*} + e^{-q_n \cdot \nabla \Delta k} e^{+q_n \cdot \nabla ^*} \right) - 1 \otimes \left( 2 \left( 1 + e^{-2q_n \cdot \nabla \Delta k} e^{+0_n \cdot \nabla ^*} \right) + e^{-0_n \cdot \nabla ^*} - 2e^{-2q_n \cdot \nabla ^*} K_{n-1}(K, \Delta k) \right) \right]. \quad (12)$$

Here $e^{-q_n \cdot \nabla K, \Delta k}$ are the momentum shift operators [31]. In Eq. (12) we assume that the overlap between the amplitude and its conjugate varies slowly with $q_n$ and symmetrize the momentum shifts in the last two virtual terms around $q_n = 0$. This allows us to write at the level of the momentum distributions (squared amplitudes):

$$|M_n(K, \Delta k)|^2 \propto \left[ 2 \left( e^{-q_n \cdot \nabla K} - 1 \right) \cosh \left( e^{-q_n \cdot \nabla \Delta k} \right) + 2 \left( e^{-q_n \cdot \nabla ^*} - 1 \right) \right] |M_{n-1}(K, \Delta k)|^2. \quad (13)$$

In Eq. (13) we also used the symmetry of the momentum transfer distribution, Eq. (11), relative to the transformation $q_n \rightarrow -q_n$. The basic step in Eq. (13) allows us to resume the interactions to all orders in opacity and relate the final momentum distribution of the evolved $q\bar{q}$
system to the one in Eq. (8):
\[ |\psi_f(K, \Delta k)|^2 = \sum_{n=0}^{\infty} \frac{2^\nu \nu!}{n!} \prod_{i=1}^{n} d^2q_i \frac{1}{\sigma_{el}} d^2q_i \]
\[ \times \left[ \left( e^{-q_n \nabla_k} - 1 \right) \cos(-q_n \cdot \nabla \Delta k) \right. \]
\[ + \left. \left( e^{-q_n \nabla \Delta k} - 1 \right) \right] |\psi_0(K, \Delta k)|^2 \].

(14)

An approximate closed form for Eq. (13) can be obtained by Fourier transforming to the impact parameter space \((B, b)\) conjugate to \((K, \Delta k)\):
\[ |\tilde{\psi}_f(B, b)|^2 = |\tilde{\psi}_0(B, b)|^2 \exp \left[ 2\chi \int d^2q \frac{1}{\sigma_{el}} d^2q \left( e^{iq \cdot B} - 1 \right) \cos(q \cdot b) + (e^{q \cdot b} - 1) \right] \].

(15)

A closed form for the integral in Eq. (19) does not exist even for simple forms of the differential elastic scattering cross section. The cosine term couples the broadening of the total momentum \(K\) of the quark-antiquark system to the distortion of the light cone wave function in \(\Delta k\).

We can still calculate the final momentum distribution by considering only the leading effect of the coupling:
\[ \cos(q \cdot B) = 1 - (q \cdot B)^2/2 + \ldots \].

Next, we perform the integrals over the Yukawa potential. With \(q = |q|\),
\[ \int d^2q \frac{\mu^2}{\pi(\mu^2 + q^2)^2} e^{iq \cdot \phi} \cos(\phi) = b \mu K_1(b \mu) \].

(16)

The \(b \rightarrow B\) contribution without interference yields a result similar to Eq. (19) and the second term in the expansion of the cosine gives
\[ -\frac{\mu^2 B^2}{2\pi} \int_0^\pi d^2q \frac{q^2}{(\mu^2 + q^2)^2} \left( e^{iB \cos(\phi)} - 1 \right) \cos^2(\phi) = \frac{3\mu^2 B^2}{32} \left( 1 + 2 \log(B^2 \mu^2) \right) \].

(17)

In Eq. (19) we have taken \(B\) and \(b\) in the same direction and chosen 1/\(B\) as the upper limit of the \(q\) integral to obtain an estimate for the upper limit of the interference contribution. The key to evaluating the average K- and \(\Delta k\)-broadening of the partons is the small \(b \mu\) expansion in Eqs. (19) and (17):
\[ b \mu K_1(b \mu) = 1 - \frac{b^2 \mu^2}{2} \left[ \ln \left( \frac{2e^{-\gamma_E}}{b \mu} \right) + \frac{1}{2} \right] + O(b^4 \mu^4) \].

(18)

Note that the leading correction that arises from Eq. (17), \(3\mu^2 B^2/16\), is small. Keeping terms \(\propto \chi \mu^2 \xi\), where \(\xi = \ln(2e^{-\gamma_E}/(\mu b)) + 1/2 - 3/16 \geq 0(1)\), and incorporating only the leading effect of the coupling term in Eq. (17) we find:
\[ |\tilde{\psi}_f(B, b)|^2 = |\tilde{\psi}_0(B, b)|^2 e^{-b^2(\chi \mu^2 \xi)} e^{-B^2(\chi \mu^2 \xi)}. \]

(19)

We treat \(\xi\) as approximately constant when compared to the power behavior of \(b^2\) and \(B^2\) when Fourier transforming Eq. (19) back to momentum space:
\[ |\psi_f(K, \Delta k)|^2 = \int d^2B d^2\Delta k e^{-B \cdot \Delta k} e^{-B \cdot B} K e^{-B^2(\chi \mu^2 \xi)} \times e^{-b^2(\chi \mu^2 \xi)} \left( \text{Norm}^2 4\pi x(1-x) \Lambda^2 e^{-b^2(x(1-x)\Lambda^2)} \right) \times e^{-\frac{K^2}{4\pi \chi \mu^2 \xi}} \left[ \text{Norm}^2 \frac{x(1-x) \Lambda^2}{\chi \mu^2 \xi + x(1-x) \Lambda^2} \right] \times e^{-\frac{\Delta k^2}{4(\chi \mu^2 \xi + x(1-x) \Lambda^2)}} e^{-\frac{m^2(1-x) + m^2 x}{x(1-x) \Lambda^2}} ] \].

(20)

In providing a physics interpretation to Eq. (20), we note that the first part of our result represent the broadening in the momentum distribution of the meson itself, \((K^2) = 4\chi \mu^2 \xi\). It is twice the size of the broadening for an individual parton. The second part, critically important for this work, represents the distortion in the intrinsic momentum distribution of the quarks, which leads to the meson decay.

We first integrate out the distribution in \(K\). The small acoplanarity that may arise in heavy meson or non-photonic electron triggered correlations [21, 34] is neglected in this work. Since our wave functions are real, they are given by the square root of the \(\Delta k\) and \(x\) momentum distributions, Eqs. (8) and (20). In our calculation the initial-state \(\psi_0(\Delta k, x) = \psi_{\Lambda}(\Delta k, x)\) represents the inclusive D- or B-mesons, respectively. The final-state \(\psi_f(\Delta k, x) = a \psi_f(\Delta k, x) + (1-a)\psi_{\bar{q}}(\Delta k, x)\) denotes a superposition of the meson and a dissociated \(q\bar{q}\) pair. The survival probability is given by \(a^2\) if the light cone momentum distributions are normalized to unity. We readily obtain:
\[ P_s(\chi \mu^2 \xi) = \left| \int d^2\Delta k dx \psi_f(\Delta k, x) \psi_0(\Delta k, x) \right|^2 \]
\[ = \left| \int d^2\Delta k dx \text{Norm}^2 \left[ \frac{x(1-x) \Lambda^2}{\chi \mu^2 \xi + x(1-x) \Lambda^2} \right] \times e^{-\frac{\Delta k^2}{4(\chi \mu^2 \xi + x(1-x) \Lambda^2)}} e^{-\frac{m^2(1-x) + m^2 x}{x(1-x) \Lambda^2}} \right|^2 \]
\[ = \left| \int dx \text{Norm}^2 4\pi x(1-x) \Lambda^2 e^{-\frac{m^2(1-x) + m^2 x}{x(1-x) \Lambda^2}} \frac{2 \sqrt{x(1-x) \Lambda^2} \sqrt{\chi \mu^2 \xi + x(1-x) \Lambda^2}}{\sqrt{x(1-x) \Lambda^2} + \sqrt{\chi \mu^2 \xi + x(1-x) \Lambda^2}} \right|^2 \].

(21)

Eq. (21) is one of the main theoretical results derived in this Letter. Although the integral over the heavy quark light cone momentum fraction \(x\) has to be taken numerically, direct comparison of the integrand to Eq. (9) shows that the meson survival probability \(P_s(\chi \mu^2 \xi) \leq 1\).
Equality is reached when $\chi \mu^2 \xi = 0$, i.e. in the absence of interactions in the medium. We finally note that typical values of $\xi$ can be estimated from the requirement $b \mu \ll 1$: for $b \mu = 0.25 - 0.1$ we find $\xi \approx 2 - 3$, respectively.

V. HEAVY MESON SUPPRESSION PHENOMENOLOGY

We are now ready to calculate the suppression of heavy hadrons, $H(c)$ and $H(b)$, from collisional interactions in the QGP. The fragmentation time, averaged over the final-state mesons and baryons is calculated as follows:

$$\frac{1}{\langle \tau_{\text{form}}(p_T, t) \rangle} = \left[ \sum_i \int_0^1 dz D_{H_i/Q}(z) \times \tau_{\text{form}}(z, p_T, m_Q, t) \right]^{-1}.$$  \hspace{1cm} (22)

In Eq. (22) the fragmentation functions are normalized to the fragmentation fractions for $Q \to H(Q)$, which in the QCD factorization approach are assumed to be universal, $\int_0^1 dz D_{H_i/Q}(z) = f_{H_i}(Q)$, $\sum_i f_{H_i}(Q) = 1$. The dissociation time, on the other hand, is externally driven by the dynamics of the evolving bulk medium and calculated by taking the logarithmic derivative of the dissociation probability,

$$P_d(p_T, m_Q, t) = 1 - P_s(p_T, m_Q, t),$$  \hspace{1cm} (23)

derived in Eq. (21), as follows:

$$\frac{1}{\langle \tau_{\text{diss}}(p_T, t) \rangle} = \frac{\partial}{\partial t} \ln P_d(p_T, m_Q, t).$$  \hspace{1cm} (24)

We use the same initial soft gluon rapidity density $dN^g/dy$ as in the calculation of the $\pi^0$ quenching [4] in central Au+Au and Cu+Cu collisions at RHIC and central Pb+Pb collisions at LHC and the cumulative momentum transfer is given by

$$\chi \mu^2 \xi = \beta Q \frac{\mu^2}{\lambda_0} \xi \ln \frac{t}{t_0}.$$  \hspace{1cm} (25)

In Eq. (25) $t_0 \leq t$ and $t_0 = 0.6 \text{ fm}$, consistent with QGP formation times used in hydrodynamic simulations of bulk observables [53], and $\beta_Q = dz/dt$. Let us denote by

$$f^Q(p_T, t) = \frac{d\sigma^Q(t)}{dyd^2p_T}, f^Q(p_T, t = 0) = \frac{d\sigma_{\text{PQCD}}^Q}{dyd^2p_T},$$  \hspace{1cm} (26)

$$f^H(p_T, t) = \frac{d\sigma^H(t)}{dyd^2p_T}, f^H(p_T, t = 0) = 0,$$  \hspace{1cm} (27)

the double differential cross sections for the heavy quarks and hadrons (mesons+baryons). Initial conditions are also specified above, in particular the heavy quark distribution is given by the perturbative $c$- and $b$-quark jet cross section. The fragmentation fraction of $b$-quarks into $B_c$-mesons is very small. In our work it is neglected and the rate equations that describe the competition between $b$- and $c$-quark fragmentation and $D$- and $B$-meson dissociation decouple for different heavy quark flavors. Including the loss and gain terms we obtain:

$$\partial_t f^Q(p_T, t) = -\frac{1}{\langle \tau_{\text{form}}(p_T, t) \rangle} f^Q(p_T, t) + \frac{1}{\langle \tau_{\text{diss}}(p_T, z, t) \rangle} \int_0^1 dx \frac{1}{x^2} \phi_{Q/H}(x) f^H(p_T/x, x),$$ \hspace{1cm} (28)

$$\partial_t f^H(p_T, t) = -\frac{1}{\langle \tau_{\text{diss}}(p_T, t) \rangle} f^H(p_T, t) + \frac{1}{\langle \tau_{\text{form}}(p_T, z, t) \rangle} \int_0^1 dz \frac{1}{z} D_{H/Q}(z) f^Q(p_T/z, t).$$ \hspace{1cm} (29)

In Eqs. (28) and (29) $z$ and $x$ are typical fragmentation and dissociation momentum fractions and we have checked that in the absence of a medium, $\tau_{\text{diss}}(p_T, t) \to \infty$, we recover the PQCD spectrum of heavy hadrons from vacuum jet fragmentation.

We solve the rate equations, Eqs. (28) and (29), numerically using the initial conditions, Eqs. (20) and (21), and obtain the final spectra of heavy hadrons at $t \gg \tau_{\text{QGP}}$, $\tau_{\text{form}}$ and $\tau_{\text{diss}}$. The nuclear modification factor, $R_{AA}(p_T)$, which arises from the collisional dissociation of $D$- and $B$-mesons, is shown in Fig. 4 for physical situations expected to be prevalent at RHIC and LHC. In
spite of the inherently different physics mechanisms that drive the suppression of light and heavy hadrons, comparison to the quenching of the $\pi^0$ 4 can be done if the same model for the QGP properties, such as the gluon rapidity density $dN_g/dy$, is used. The theoretical uncertainty comes from varying the parameter $\xi$ in its natural range and $\mu_0$ and $\lambda_0$ are set by the overall gluon multiplicity and the QGP formation time, $\tau_0$, using thermodynamic arguments. The top panels show results for central Au+Au and Cu+Cu collisions at the maximum RHIC energy \( \sqrt{s_{NN}} = 200 \text{ GeV} \). The bottom panels present $R_{AA}(p_T)$ for two different $dN/dy$ extrapolations in central Pb+Pb collisions at LHC.

One important feature of this approach is the sensitivity to the build-up of the QGP density. If $\tau_0 << \tau_{\text{form}}$ and $\rho(t)$ decreases rapidly, hadrons would not have formed at times when the dissociation mechanism is most efficient. This leads to reduced suppression, which can be seen in the high-$p_T$ behavior of the $D$-meson $R_{AA}$ and understood when one recalls that $\tau_{\text{form}} \propto p_T$. Since the distortion of the light cone wave functions reflects the accumulated squared transverse momentum transfer, the reduced suppression at $p_T < m_Q$ arises from the velocity factor $\beta_Q$ in Eq. (24). We observe in Fig. 4 that initial formation time $\tau_0 = 0.6$ fm, consistent with the one used in hydrodynamic simulations to describe bulk QGP observables 33, the moderate- and high-$p_T$ suppression of heavy hadron production is comparable to that of lighter hadrons. We have checked that for QGP formation time $\tau_0 = 0.2$ fm the $p_T > 5$ GeV heavy quark quenching is reduced by a factor of $\sim 1.5$. Finally, one should note that dissociation/fragmentation both emulate energy loss by shifting the quarks/hadrons to lower transverse momenta. For example, $B$- and $D$-meson attenuation is sensitive to the partonic slope. For this reason the suppression at LHC is found to be comparable or slightly smaller than the one calculated at RHIC in spite of the larger QGP densities and temperatures.

Contrary to calculations that emphasize radiative and collisional heavy quark energy loss 3, 10, 14, 15, QGP-induced dissociation predicts $B$-meson suppression comparable to or larger than that of $D$-mesons at transverse momenta as low as $p_T \sim 10$ GeV, see Fig. 4. This is due to the significantly smaller formation times for $H(b)$ relative to $H(c)$. Thus, each fragmentation/dissociation cycle proceeds at a much faster rate for $b$-quarks/$B$-mesons. This is an example where the large mass facilitates the hadron suppression mechanism. The observable effect of this faster rate in the quenching of the final hadron distributions is amplified by the significance of the early $t \sim \tau_0$ hot and dense stage in the dynamical evolution of the QGP. Our simulations show that the rates in Eqs. (28) and (29) may play a more important role than the effective fractional energy losses $\epsilon = \Delta E/E \approx 1 - \tilde{\epsilon}$ and $\epsilon \approx 1 - \tilde{\epsilon}$, which are larger for $c$-quarks/$D$-mesons.

Inclusive non-photon electron data for the most central Au+Au collisions at $\sqrt{s_{NN}} = 200 \text{ GeV}$ from PHENIX 5 and STAR 7 are compared to the current calculation in Fig. 5. We used the PYTHIA event generator 36 to simulate the $(e^+ + e^-)/2$ spectra coming from the primary decays of the $D$- and $B$-mesons in our baseline $p+p$ and QGP-modified $A+A$ results. We find that the single electron suppression can be as large as a factor of four and approximates well the quenching extracted from available data 3, 4, 5, 6. We emphasize that such agreement between theory and experiment is not achieved at the cost of neglecting the contribution of the $B$-mesons to the non-photon electron spectra.

VI. CONCLUSIONS

It has been recently suggested that collisional dissociation of heavy quarkonia in the quark-gluon plasma 37 may be a possible explanation for the suppression of their production rate in nucleus-nucleus collisions. It is thus surprising that until now a similar physics mechanism has not been considered for open heavy flavor. In this Letter, we investigated the perturbative QCD dynamics of open charm and beauty production and, in the framework of the reaction operator (GLV) approach 30, 31, 32, 33, extended to composite $q\bar{q}$ systems, derived the medium-induced dissociation probability for heavy $D$- and $B$-mesons traversing dense nuclear matter. We showed that the effective energy loss, which arises from the sequential fragmentation and dissociation of heavy quarks and mesons, is sensitive to the interplay between the formation times of the hadrons and the QGP and the detailed expansion dynamics of hot nuclear matter. The proposed new attenuation mechanism, which stems from the short formation times of $D$- and $B$-mesons and underlies the suppression of the inclusive non-photon ele-
tions in nucleus-nucleus collisions at RHIC, was found to be compatible with the measured large, factor of four to five, quenching for heavy flavor.

Previous studies, based on radiative and collisional parton energy loss \cite{10, 11, 14} and heavy quark diffusion \cite{15} under-predict the suppression of non-photonic electrons in central Au+Au collisions due to the small $b$-quark quenching. A natural consequence of the approach, presented in this Letter as a viable alternative to existing calculations, is that $B$-mesons are attenuated as much as $D$-mesons at transverse momenta as low as $p_T \sim 10$ GeV. While we anticipate that a comparable description of the attenuation of the non-photonic electrons may be achieved when partonic energy loss is combined with quark-resonance interactions near the QCD phase transition in Langevin transport simulations, the hierarchy $R_{AA}(H(c)) \ll R_{AA}(H(b))$ will not be changed in the accessible transverse momentum range \cite{38}. We conclude that robust experimental determination of the dominant mechanism for in-medium modification of open heavy flavor would require direct and separate measurements of the $B$- and $D$-meson $R_{AA}$ distributions versus $p_T$ and centrality in collisions of heavy nuclei.

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