Drag of heavy quarks in quark gluon plasma at energies available at the CERN Large Hadron Collider (LHC)

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The drag and diffusion coefficients of charm and bottom quarks propagating through quark gluon plasma (QGP) have been evaluated for conditions relevant to nuclear collisions at the Large Hadron Collider (LHC). The dead cone and Landau-Pomeronchuk-Migdal (LPM) effects on radiative energy loss of heavy quarks have been considered. Both radiative and collisional processes of energy loss are included in the *effective* drag and diffusion coefficients. With these effective transport coefficients, we solve the Fokker-Plank (FP) equation for the heavy quarks executing Brownian motion in the QGP. The solution of the FP equation has been used to evaluate the nuclear suppression factor, R_{AA}, for the nonphotonic single-electron spectra resulting from the semileptonic decays of hadrons containing charm and bottom quarks. The effects of mass on R_{AA} have also been highlighted.

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I. INTRODUCTION

Energy dissipation of heavy quarks in QCD matter is considered one of the most promising probes for the quark gluon plasma (QGP) diagnostics. The energy loss of energetic heavy quarks (Q) while propagating through the QGP medium is manifested in the suppression of heavy-flavored hadrons at high transverse momentum (p_T). The depletion of high-p_T hadrons (D and B mesons) produced in nucleus + nucleus collisions with respect to those produced in proton + proton (pp) collisions has been measured experimentally [1–3] through their semileptonic decays. The two main processes that cause this depletion are (i) elastic collisions and (ii) the bremsstrahlung or radiative loss due to the interaction of the heavy quarks with the quarks, antiquarks, and gluons in the thermal bath created in heavy-ion collisions.

The importance of collisional energy loss in QGP diagnostics was discussed first by Bjorken [4]. The calculations of elastic loss were performed with improved techniques [5,6], and its importance was highlighted subsequently [7,8] in heavy-ion collisions. The collisional energy loss of heavy quarks [9] has gained importance recently in view of the measured nuclear suppression in the p_T spectra of nonphotonic single electrons. Several ingredients, such as inclusions of nonperturbative contributions from the quasihadronic bound state [10], three-body scattering effects [11], the dissociation of heavy mesons because of its interaction with the partons in the thermal medium [12], and employment of running coupling constants and realistic Debye mass [13] have been proposed to improve the description of the experimental data. Wicks *et al.* [14] showed that the inclusion of both elastic and inelastic collisions and the path-length fluctuation reduces the gap between the theoretical and experimental results.

The energy loss of energetic partons by radiation is currently a field of high interest [15–19]. For mass dependence of energy loss due to radiative processes, Dokshitzer and Kharzeev [20] argue that heavy quarks will lose much less energy than light quarks as a result of dead-cone effects [21]. However, Aurenche and Zakharov claim that the radiative process has an anomalous mass dependence [22] because of the finite size of the QGP, which leads to a small difference in energy loss between heavy and light quarks. The mass dependence of the transverse momentum spectrum of the radiated gluons from the heavy quarks is studied in Ref. [23]. The authors found that the medium-induced gluon radiation fills up the dead cone with a reduced magnitude at large gluon energies compared to the radiation from light quarks. For high-energy heavy quarks, the effects of the dead cone, however, are reduced because the magnitude of the angle forbidden for gluon emission behaves as ∼heavy-quark mass/energy [24]. From the study of the mass dependence of the radiative loss, it is shown in Ref. [25] that the very energetic charm (not the bottom) quarks behave like massless partons. Although the authors in Ref. [26] concluded that the suppression of radiative loss for heavy quarks is due to dead-cone effects, it is fair to state that the issue is not settled yet.

The other mechanism that can affect the radiative loss is the Landau-Pomeronchuk-Migdal (LPM) effect [27], which depends on the relative magnitude of two time scales of the system [28]: the formation time (τ_F) and the mean scattering time scale (τ_c) of the emitted gluons. If $\tau_F > \tau_c$, then LPM suppression will be effective. The LPM effect is built into the expression for radiative energy loss of heavy quarks derived in Refs. [23–25,29]. In contrast to those, in the present work, we separately introduce the LPM effects in the energy-loss formula.

The successes of the relativistic hydrodynamical model (see Refs. [30,31] for review) in describing the host of experimental results from Relativistic Heavy Ion Collider (RHIC) [32] indicate that thermalization might have taken place in the system of quarks and gluons formed after the nuclear collisions. The strong final-state interaction of high-energy partons with the QGP, that is, the observed jet quenching [33,34] and the large elliptic flow (v_2) [35,36] in Au + Au collisions at RHIC, indicates the possibility of fast equilibration. On the one hand, the experimental data indicate early thermalization time of ∼0.6 fm/c [37]; on the other hand, the pQCD-based calculations give a thermalization time of ∼2.5 fm/c [38](see also Ref. [39]). The gap between these two time scales suggests that the nonperturbative effects play a crucial role in achieving thermalization. It has also been pointed out that the instabilities [40–43] may derive the system toward faster equilibrium. However, the inclusion of such effects also does not reproduce small thermalization time.

The perturbative QCD (pQCD) calculations indicate that the heavy-quark (Q) thermalization time, τ_i^Q, is larger [38,44] than the light-quarks and gluon thermalization scale, τ_i. Gluons may thermalized even before up and down quarks [45,46]. In the present work, we assume that the QGP is formed at time τ_i. Therefore, the interaction of the nonequilibrated heavy quarks with the equilibrated QGP for the time interval $\tau_i < \tau < \tau_i^Q$ can be treated within the ambit of the Fokker-Plank (FP) equation [47,48]; that is, the heavy quark can be thought of executing Brownian motion [44,45,49–56] in the heat bath of QGP during that interval of time. Therefore, the propagation of heavy quarks through QGP may be treated as the interactions between equilibrium and nonequilibrium degrees of freedom. The FP equation provides an appropriate framework for such studies. The Boltzmann transport equation has recently been applied to study the depletion of high-energy gluons because of its elastic and inelastic interactions with QGP [57].

The article is organized as follows. In the next section, the evolution of the momentum distribution of heavy quarks in QGP is discussed. In Sec. III, we address the issues of radiative energy loss with the dead-cone effect. The nonphotonic electron spectra are discussed in Sec. IV. The initial conditions and space-time evolution are discussed in Sec. V, Sec. VI contains the discussion on the nuclear suppression, and finally Sec. VII is devoted to summary and conclusions.

II. EVOLUTION OF HEAVY-QUARK MOMENTUM DISTRIBUTIONS

The Boltzmann transport equation describing a nonequilibrium statistical system reads as follows:

$$\left[\frac{\partial}{\partial t} + \frac{\mathbf{p}}{E}.\nabla_{\mathbf{x}} + \mathbf{F}.\nabla_{\mathbf{p}}\right] f(x,p,t) = \left[\frac{\partial f}{\partial t}\right]_{\rm col}, \tag{1}$$

where p and E denote momentum and energy, $\nabla_{\mathbf{x}}$ ($\nabla_{\mathbf{p}}$) are spatial (momentum space) gradients, and $f(x,p,t)$ is the phase-space distribution (in the present case, f stands for heavy-quark distribution). The assumption of uniformity in the plasma and absence of any external force leads to

$$\frac{\partial f}{\partial t} = \left[\frac{\partial f}{\partial t}\right]_{\rm col}. \tag{2}$$

The collision term on the right-hand side of the equation can be approximated as (see [50,54] for details)

$$\left[\frac{\partial f}{\partial t}\right]_{\rm col} = \frac{\partial}{\partial p_i}\left[A_i(p)f + \frac{\partial}{\partial p_i}[B_{ij}(p)f]\right], \tag{3}$$

where we have defined the kernels

$$\begin{aligned} A_i &= \int d^3k\,\omega(p,k)\,k_i, \\ B_{ij} &= \int d^3k\,\omega(p,k)\,k_i k_j, \end{aligned} \tag{4}$$

for $|\mathbf{p}| \to 0$, $A_i \to \gamma p_i$, and $B_{ij} \to D\delta_{ij}$, where γ and D stand for drag and diffusion coefficients, respectively. The function $\omega(p,k)$ is given by

$$\omega(p,k) = g\int \frac{d^3q}{(2\pi)^3} f'(q) v\sigma_{p,q\to p-k,q+k}, \tag{5}$$

where f' is the phase-space distribution (in the present case, it stands for light quarks and gluons), v is the relative velocity between the two collision partners, σ denotes the cross section, and g is the statistical degeneracy. The coefficients in the first two terms of the expansion in Eq. (3) are comparable in magnitude because the averaging of k_i involves greater cancellation than the averaging of the quadratic term $k_i k_j$. The higher power of k_i's are smaller [47].

With these approximations, the Boltzmann equation reduces to a nonlinear integro-differential equation known as the Landau kinetic equation:

$$\frac{\partial f}{\partial t} = \frac{\partial}{\partial p_i}\left[A_i(p)f + \frac{\partial}{\partial p_i}[B_{ij}(p)f]\right]. \tag{6}$$

The nonlinearity is caused by the appearance of f' in A_i and B_{ij} through $w(p,k)$. It arises from the simple fact that we are studying a collision process that involves two particles: It should, therefore, depend on the states of the two participating particles in the collision process and hence on the product of the two distribution functions. Considerable simplicity may be achieved by replacing the distribution functions of one of the collision partners by their equilibrium Fermi-Dirac or Bose-Einstein distributions (depending on the statistical nature) in the expressions of A_i and B_{ij}. Then Eq. (6) reduces to a linear partial differential equation, usually referred to as the FP equation, describing the interaction of a particle that is out of thermal equilibrium with the particles in a thermal bath of light quarks, antiquarks, and gluons. The quantities A_i and B_{ij} are related to the usual drag and diffusion coefficients, and we denote them by γ_i and D_{ij} respectively (i.e., these quantities can be obtained from the expressions for A_i and B_{ij} by replacing the distribution functions by their thermal counterparts).

The evolution of the heavy-quark momentum distribution (f) while propagating through QGP can be studied by using the FP equation (see Ref. [50] for details):

$$\frac{\partial f}{\partial t} = \frac{\partial}{\partial p_i}\left[\gamma_i(p)f + \frac{\partial}{\partial p_i}[D_{ij}(p)f]\right]. \tag{7}$$

During the propagation through the QGP, the heavy quarks dissipate energy predominantly by two processes: (i) collisional, for example, $gQ \to gQ$, $qQ \to qQ$ and $\bar{q}Q \to \bar{q}Q$, and (ii) radiative processes, that is, when the heavy quark emits gluons as a result of its interaction with the thermal partons in the plasma. Therefore, the drag and diffusion coefficients should include these two processes of energy dissipation.

The elastic collisions of heavy quarks with light quarks (q) and gluons (g), that is, $gQ \to gQ$, $qQ \to qQ$, and $\bar{q}Q \to \bar{q}Q$, have been used to evaluate the transport coefficients ($\gamma_{\rm coll}$ and $D_{\rm coll}$) due to collisional process. At Large Hadron Collider (LHC) energy, one cannot ignore the radiative energy loss; therefore, this should also be taken into account through the transport coefficients. The transport coefficient [58], $\hat{q}$, which is related to the energy loss [59], dE/dx of the propagating partons in the medium, has been used to calculate the shear viscosity-to-entropy density ratio, η/s [26,60]. The $\hat{q}$ is closely related to the diffusion coefficient D (for details, see Ref. [60]). In similar spirit, we use dE/dx to calculate the drag coefficient of the medium and use Einstein's relation, $D = TM\gamma$, to obtain the diffusion coefficient when a heavy quark of mass M is propagating through the medium at temperature T. The action of drag on the heavy quark can be defined through the relation

$$-\left.\frac{dE}{dx}\right|_{\rm rad} = \gamma_{\rm rad}\, p, \tag{8}$$

where $\gamma_{\rm rad}$ denotes the drag coefficient and p is the momentum of the heavy quark. It should be mentioned here that the collisional and the radiative processes are not entirely independent, that is, the collisional process may influence the radiative one, and therefore strictly speaking dE/dx and hence the transport coefficients for radiative and collisional process should not be added to obtain the net energy loss or net value of the drag coefficient. However, in the absence any rigorous method, we add them up to obtain the effective drag coefficients, $\gamma_{\rm eff} = \gamma_{\rm rad} + \gamma_{\rm coll}$, and similarly the effective diffusion coefficient, $D_{\rm eff} = D_{\rm coll} + D_{\rm rad}$. This is a good approximation for the present work because the radiative loss is large compared to the collisional loss at LHC. With these effective transport coefficients, the FP equation reads

$$\frac{\partial f}{\partial t} = \frac{\partial}{\partial p_i}\left[\gamma_{\rm eff}(p) f + \frac{\partial}{\partial p_i}[D_{\rm eff}(p) f]\right], \tag{9}$$

where $\gamma_{\rm eff}$ and $D_{\rm eff}$ contain contributions from both the mechanisms (collisional and radiative). In evaluating the drag coefficient, we have used temperature-dependent strong coupling, α_s, from [61]. The Debye mass, $\sim g(T)T$, is also a temperature-dependent quantity used as a cutoff to shield the infrared divergences arising as a result of the exchange of massless gluons.

III. ENERGY-DISSIPATION PROCESSES

The matrix element for the radiative process (e.g., $Q + q \to Q + q + g$) can be factorized into an elastic process ($Q + q \to Q + q$) and a gluon emission ($Q \to Q + g$). The emitted gluon distribution can be written as [62,63]

$$\frac{dn_g}{d\eta d^2k_\perp} = \frac{C_A\alpha_s}{\pi^2}\frac{q_\perp^2}{k_\perp^2(k_\perp - q_\perp)^2}F^2, \tag{10}$$

where $k = (k_0, k_\perp, k_3)$ is the four-momentum of the emitted gluon and $q = (q_0, q_\perp, q_3)$ is the four-momentum of the exchanged gluon, $\eta = 1/2\, ln(k_0 + k_3)/(k_0 - k_3)$ is the rapidity, $C_A = 3$ is the Casimir invariant of the adjoint representation and $\alpha_s = g^2/4\pi$ is the strong coupling constant.

The effects of quark mass in the gluon radiation is taken into account by multiplying the emitted gluon distribution from massless quarks by F^2, containing the effects of heavy quark mass. F is given by [20,21]

$$F = \frac{k_\perp^2}{\omega^2\theta_0^2 + k_\perp^2}, \tag{11}$$

where $\theta_0 = M/E$.

As the energy loss of a heavy quark is equal to the energy taken away by the radiated gluon, we can estimate the energy loss of the heavy quark by multiplying the interaction rate Λ and the average energy loss per collision ϵ, which is given by the average of the probability of radiating a gluon times the energy of the gluon.

The LPM effects have been taken into account by including a formation-time restriction on the phase space of the emitted gluon in which the formation time, τ_F, must be smaller than the interaction time, $\tau = \Lambda^{-1}$. The radiative energy loss of a heavy quark can be given by

$$-\left.\frac{dE}{dx}\right|_{\rm rad} = \Lambda\epsilon = \tau^{-1}\epsilon, \tag{12}$$

where ϵ, the average energy per collision, is [63,64]

$$\epsilon = \langle n_g k_0\rangle = \int d\eta d^2k_\perp \frac{dn_g}{d\eta d^2k_\perp} k_0 \Theta(\tau - \tau_F) F^2, \tag{13}$$

where $\tau_F = \cosh\eta/k_\perp$. As mentioned before, for the infrared cutoff $k_\perp^{\rm min}$, we choose the Debye screening mass of gluon:

$$k_\perp^{\rm min} = \mu_D = \sqrt{4\pi\alpha_s}T. \tag{14}$$

The maximum transverse momentum of the emitted gluon is given by

$$(k_\perp^{\rm max})^2 = \left\langle\frac{(s-m^2)^2}{4s}\right\rangle = \frac{3ET}{2} - \frac{m^2}{4} + \frac{m^4}{48pT}\ln\left[\frac{m^2 + 6ET + 6pT}{m^2 + 6ET - 6pT}\right]. \tag{15}$$

Following the procedure of earlier works [50,65], we evaluate the drag and diffusion coefficients for the elastic processes. Knowing $\gamma_{\rm rad}$ from the radiative processes as described previously, we obtain the effective drag coefficients and hence effective diffusion coefficient through the Einstein relation. In Figs. 1 and 2, the variation of effective drag and diffusion coefficients with T have been depicted for charm quarks. We observe that the contribution of the radiative loss is large compared to the collisional or elastic one. The difference between the collisional and radiative loss increases with temperature, indicating a very small contribution from

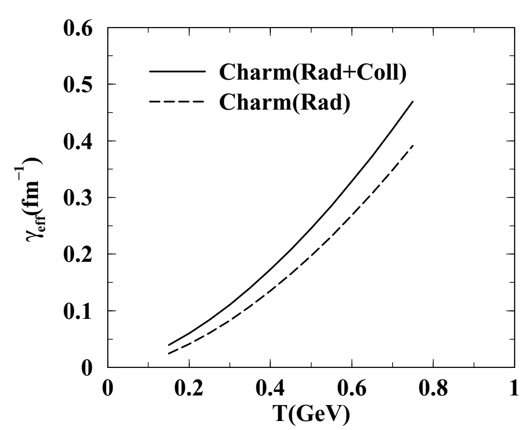

FIG. 1. Variation of effective drag coefficient with temperature for charm quarks.

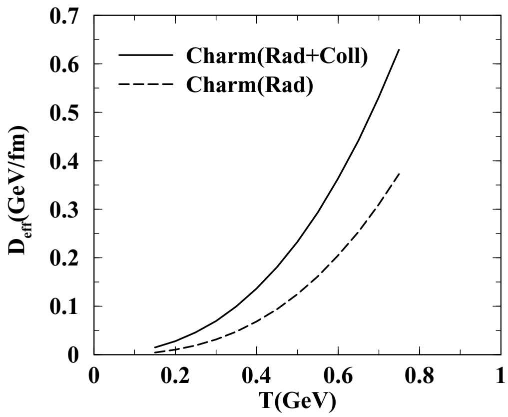

FIG. 2. Variation of effective diffusion coefficient with temperature for charm quarks.

the former at large T. A similar difference is reflected in the diffusion coefficients as we have used Einstein's relation to obtain it from the drag coefficients. We observe that at low T and p_T, the contributions from collisional processes are more than or comparable to those from radiative processes. For the bottom quark, we find that the gap between the drag coefficients with radiative and elastic processes is smaller (compared to charm quarks) at lower temperature domain. Quantitatively, the value of drag is smaller for bottom quarks than for charm quarks because of their larger relaxation time. However, the qualitative behavior is similar to charm quarks as shown in Fig. 3. The diffusion coefficient of the bottom quark is large (Fig. 4) compared to the charm quark because of the large mass of the former introduced through Einstein's relation.

On obtaining the effective drag and diffusion coefficients, next we need to know the initial heavy-quark momentum distributions to solve the FP equation. The production of charm and bottom quarks in hadronic collisions has been studied extensively [66]. In the present work, the p_T distribution of charm and bottom quarks in pp collisions have been taken from the NLC MNR code [67]. The results from the code may be tested by measuring the production cross sections of heavy mesons (containing c and b quarks) in pp collisions at $\sqrt{s_{\rm NN}} = 5.5$ TeV.

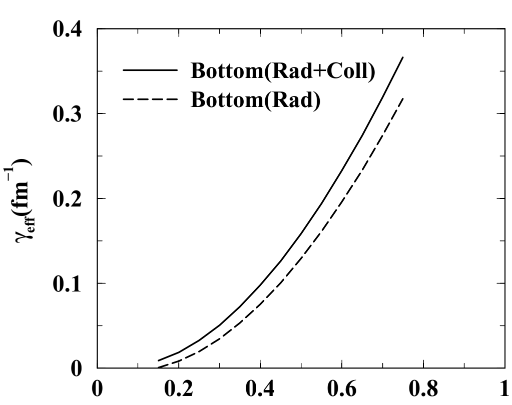

FIG. 3. Same as Fig. 3 for bottom quarks.

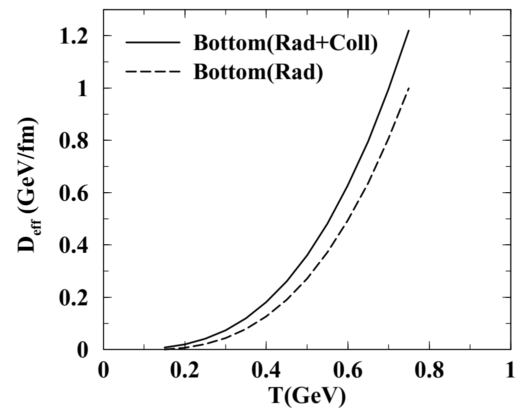

FIG. 4. Same as Fig. 2 for bottom quarks.

With all these required inputs, we solve the FP equation by using the Green function technique (see [51,65] for details).

IV. THE NONPHOTONIC ELECTRON SPECTRA

The FP equation has been solved for the heavy quarks with the initial condition mentioned previously. We convolute the solution with the fragmentation functions of the heavy quarks to obtain the p_T distribution of the heavy mesons (B and D) ($dN^{D,B}/q_T dq_T$). For heavy-quark fragmentation, we use the Peterson function [68] given by

$$f(z) \propto \frac{1}{z\left\{1 - \frac{1}{z} - [\epsilon_c/(1-z)]\right\}^2} \tag{16}$$

for charm quark $\epsilon_c = 0.05$. For bottom quarks, $\epsilon_b = (M_c/M_b)^2 \epsilon_c$, where M_c (M_b) is the charm (bottom) quark mass. The nonphotonic single-electron spectra originate from the decays of heavy-flavored mesons; for example, $D \to Xe\nu$ or $B \to Xe\nu$ at midrapidity ($y = 0$) can be obtained as follows [69,70]:

$$\frac{dN^e}{p_T dp_T} = \int dq_T \frac{dN^D}{q_T dq_T} F(p_T, q_T), \tag{17}$$

where

$$F(p_T, q_T) = \omega \int \frac{d(\mathbf{p}_T.\mathbf{q}_T)}{2p_T \mathbf{p}_T.\mathbf{q}_T} g(\mathbf{p}_T.\mathbf{q}_T/M), \tag{18}$$

where M is the mass of the heavy mesons (D or B), $\omega = 96(1 - 8m^2 + 8m^6 - m^8 - 12m^4 ln m^2)^{-1} M^{-6}$ $(m = M_X/M)$, and $g(E_e)$ is given by

$$g(E_e) = \frac{E_e^2\left(M^2 - M_X^2 - 2ME_e\right)^2}{(M - 2E_e)}, \tag{19}$$

related to the rest-frame spectrum for the decay $D \to Xe\nu$ through the following relation [69]:

$$\frac{1}{\Gamma_H}\frac{d\Gamma_H}{dE_e} = \omega g(E_e). \tag{20}$$

We evaluate the electron spectra from the decays of heavy mesons originating from the fragmentation of the heavy quarks propagating through the QGP formed in heavy-ion collisions. Similarly, the electron spectrum from the pp collisions can be obtained from the charm and bottom quark distribution, which goes as the initial conditions to the solution of FP equation. The ratio of these two quantities, R_{AA}, then gives

$$R_{AA}(p_T) = \frac{(dN^e/d^2p_Tdy)^{\mathrm{Au+Au}}}{N_{\mathrm{coll}} \times (dN^e/d^2p_Tdy)^{\mathrm{p+p}}}, \tag{21}$$

called the nuclear suppression factor, which will be unity in the absence of any medium. In Eq. (21) N_{coll} stands for the number of nucleon-nucleon interactions in a nucleus + nucleus collision. The experimental data [1–3] at RHIC energy ($\sqrt{s_{\mathrm{NN}}} = 200$ GeV) shows substantial suppression ($R_{AA} < 1$) for $p_T \geqslant 2$ GeV, indicating substantial interaction of the plasma particles with charm and bottom quarks from which electrons are originated through the process $c(b)$ (hadronization) $\longrightarrow$ $D(B)$(decay) $\longrightarrow$ $e + X$. The loss of energy of high-momentum heavy quarks propagating through the medium created in Au + Au collisions causes a depletion of high-p_T electrons.

V. SPACE-TIME EVOLUTION

The system formed in nuclear collisions at relativistic energies evolves dynamically from the initial to the final state. The time evolution of such systems may be studied by solving the hydrodynamic equations

$$\partial_\mu T^{\mu\nu} = 0, \tag{22}$$

with boost invariance along the longitudinal direction [71]. In this equation, $T^{\mu\nu} = (\epsilon + P)u^\mu u^\nu - g^{\mu\nu}P$ is the energy-momentum tensor for an ideal fluid, ϵ is the energy density, P is the pressure, and u^μ is the hydrodynamic four-velocity. It is expected that the central-rapidity region of the system formed after nuclear collisions at LHC energy is almost net baryon free. Therefore, the equation governing the conservation of net baryon number need not be considered here. The radial coordinate dependencies of T have been parametrized as in Ref. [52]. Some comments on the effects of the radial flow are in order here. The radial expansion will increase the size of the system and hence decrease the density of the medium. Therefore, with radial flow, the heavy quark will traverse a larger path length in a medium of reduced density. These two oppositely competing phenomena may have negligible net effects on the nuclear suppression (see also Ref. [52]).

The total amount of energy dissipated by a heavy quark in the QGP depends on the path length it traverses. Each parton traverse a different path length that depends on the geometry of the system and on the point where its is created. The probability that a parton is produced at a point (r, ϕ) in the plasma depends on the number of binary collisions at that point, which can be taken as

$$P(r, \phi) = \frac{2}{\pi R^2}\left(1 - \frac{r^2}{R^2}\right)\theta(R - r), \tag{23}$$

where R is the nuclear radius. It should be mentioned here that the expression in Eq. (23) is an approximation for the collisions with zero impact parameter. A parton created at (r, ϕ) in the transverse plane propagates a distance $L = \sqrt{R^2 - r^2\sin^2\phi} - r\cos\phi$ in the medium. In the present work, we use the following equation for the geometric average of the integral involving drag coefficient $[\int d\tau\gamma(\tau)]$:

$$\Gamma = \frac{\int rdrd\phi P(r,\phi)\int^{L/v} d\tau\gamma(\tau)}{\int rdrd\phi P(r,\phi)}, \tag{24}$$

where v is the velocity of the propagating partons. Similar averaging has been performed for the diffusion coefficient. For a static system, the temperature dependence of the drag and diffusion coefficients of the heavy quarks enter via the thermal distributions of light quarks and gluons through which it is propagating. However, in the present scenario, the variation of temperature with time is governed by the equation of state or velocity of sound of the thermalized system undergoing hydrodynamic expansion. In such a scenario, the quantities such as Γ [Eq. (24)] and hence R_{AA} becomes sensitive to velocity of sound (c_s) in the medium. This is shown in the next section.

VI. THE NUCLEAR SUPPRESSION

The p_T dependence of R_{AA} is sensitive to the nature of the initial (prior to the interaction with the medium) distribution of heavy quarks [67]. For the QGP expected to be formed at the LHC, we have taken initial temperature $T_i = 700$ MeV and initial thermalization time $\tau_i = 0.08$ fm/c, which reproduces the predicted hadron multiplicity $dN/dy = 2100$ [72] through the relation

$$T_i^3\tau_i \approx \frac{2\pi^4}{45\zeta(3)}\frac{1}{4a_{\mathrm{eff}}}\frac{1}{\pi R_A^2}\frac{dN}{dy}, \tag{25}$$

where R_A is the radius of the system, $\zeta(3)$ is the Riemann ζ function, and $a_{\mathrm{eff}} = \pi^2 g_{\mathrm{eff}}/90$, where g_{eff} ($=2 \times 8 + 7 \times 2 \times 2 \times 3 \times N_F/8$) is the degeneracy of quarks and gluons in QGP and N_F is the number of flavors. We have taken the value of the transition temperature to be $T_c = 170$ MeV.

The value of R_{AA} is plotted against the p_T of the nonphotonic single electron resulting from D decays in Fig. 5. The results show substantial depletion at large p_T, indicating a large interaction rate of the charm quarks with the thermal medium of partons. The sensitivity of the results on the equation

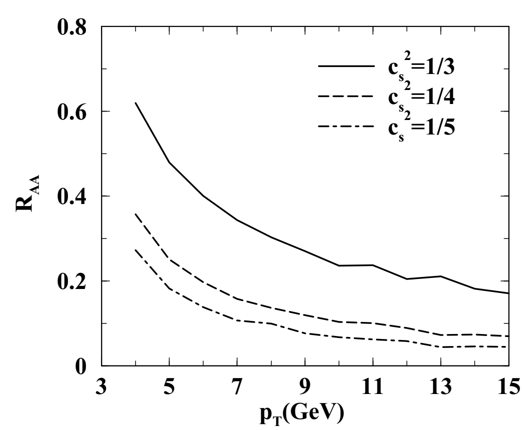

FIG. 5. Nuclear suppression factor R_{AA} as a function of p_T for various equation of state for nonphotonic single electrons resulting from D-mesons decay.

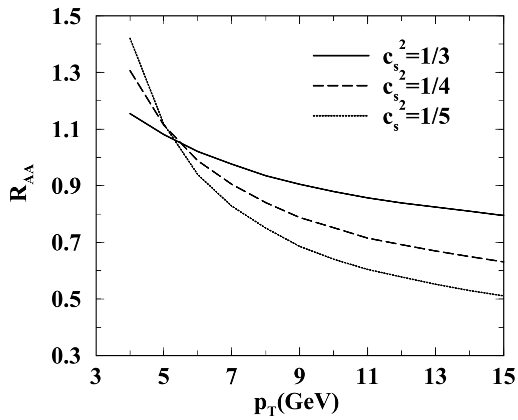

FIG. 6. Same as Fig. 5 for B mesons.

of state is also demonstrated in Fig. 5. A softer equation of state (lower value of c_s) makes the expansion of the plasma slower, enabling the propagating heavy quarks to spend more time interacting in the medium and hence lose more energy before exiting from the plasma, which results in less particle production at high p_T. This is clearly demonstrated in Fig. 5. It may be mentioned here that c_s increases with temperature. Therefore, because of the higher initial temperature of the QGP formed at LHC, the value of c_s may be larger than that of QGP formed at RHIC energies. Keeping this in mind, we predict the nuclear suppression factors for three values of c_s: $1/\sqrt{3}$ (maximum possible), $1/\sqrt{4}$, and $1/\sqrt{5}$ (Fig. 5).

The nuclear suppressions for the bottom quarks are displayed in Fig. 6. We observe quantitatively less suppression as compared to charm quarks. The difference between the charm quarks' and bottom quarks' suppression is affected chiefly by two factors: (i) different values of transport coefficients and (ii) different kinds of initial p_T distributions. The bottom quark has less drag coefficients and harder p_T distributions: Both these factors are responsible for the smaller suppression of the bottom quark. The present results on $R_{\rm AA}$ may be compared with those obtained in Ref. [73] in a different approach.

In Fig. 7, we have plotted the ratio $R^D_{\rm AA}/R^B_{\rm AA}$ as a function of p_T, from where the effect of the mass and the role of the nature

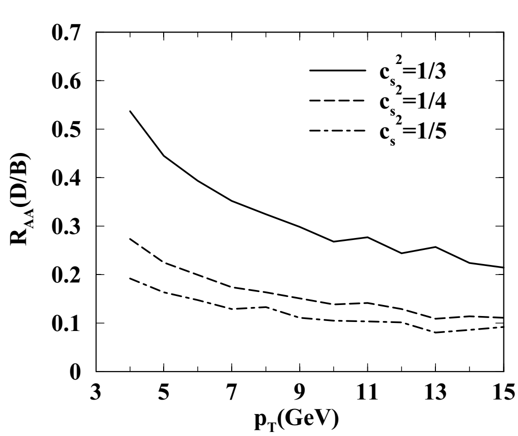

FIG. 7. Variation of the ratio of nuclear suppression factor $R_{\rm AA}$ for charm to bottom quarks as a function of p_T.

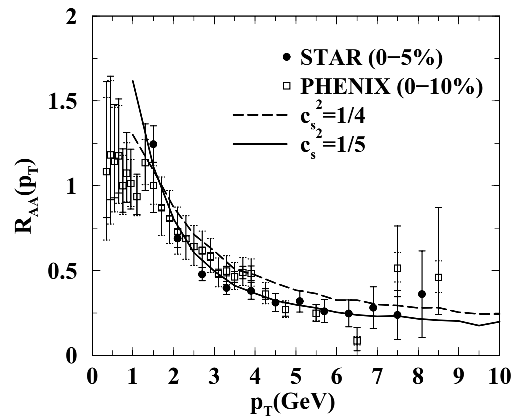

FIG. 8. Comparison of $R_{\rm AA}$ obtained in the present work with the experimental data obtained by STAR and PHENIX collaboration for $\sqrt{s_{\rm NN}} = 200$ GeV. The experimental data of STAR and PHENIX collaborations are taken from Refs. [1] and [2] respectively.

(soft or hard) of the initial p_T distributions can be understood (see also Ref. [74]).

In Fig. 8, we compare the experimental data obtained by the STAR [1] and PHENIX [2] collaborations for Au + Au collisions at $\sqrt{s_{\rm NN}} = 200$ GeV with theoretical results obtained in the present work. For the theoretical calculations, the values of initial and transition temperatures are taken as 400 MeV and 170 MeV respectively. The value of initial thermalization time is assumed to be 0.2 fm/c. These values of initial thermalization time and initial temperature reproduce the total multiplicity at midrapidity, $dN/dy = 1100$. We observe that the data can reasonably be reproduced by taking velocity of sound c_s as $1/\sqrt{5}$. It should be mentioned here that the inclusion of both radiative and elastic losses in the effective drag enables us to reduce the gap between the experiment and theory without any enhancement of the pQCD cross section, as has been done in our previous work [65].

So far, we have discussed the suppression of the nonphotonic electron produced in nuclear collisions resulting from the propagation of the heavy quark in the partonic medium in the prehadronization era. However, the suppression of the D mesons in the posthadronization era (when both the temperature and density are lower than the partonic phase) should in principle be also taken into account. The suppression of the D mesons in the posthadronization era is found to be small [75], indicating the fact that the hadronic medium (of pions and nucleons) is unable to drag the D mesons strongly.

VII. SUMMARY AND CONCLUSIONS

We have evaluated the drag and diffusion coefficients containing both the elastic and radiative losses for charm and bottom quarks. We found that the radiative loss is dominant over its collisional counterpart. In the radiative process, dead-cone and LPM effects are taken into account. With these transport coefficients and initial charm and bottom p_T distributions from the NLC MNR [67] code, we have solved the FP equation. The solution of FP equation has been used to predict nuclear suppression factors to be measured through the semileptonic decays of heavy mesons (D and B) for LHC conditions. We find that the suppression is quite large, indicating that the heavy quarks undergo substantial interactions in the QGP medium. The ratio of the suppression for D and B quarks has also been evaluated to understand the effects of mass on the suppression. The same formalism has been applied to study the experimental data on nonphotonic single-electron spectra measured by STAR and PHENIX collaborations at the highest RHIC energy. The data is well reproduced without any enhancement of the pQCD cross section.

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