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 Large Hadron Collider (LHC). The dead cone and Landau-Pomeranchuk-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 non-photonic single electron spectra resulting from the semi-leptonic decays of hadrons containing charm and bottom quarks. The effects of mass on \( R_{AA} \) has also been highlighted.

PACS numbers: 12.38.Mh,25.75.-q,24.85.+p,25.75.Nq

I. INTRODUCTION

Energy dissipation of heavy quarks in QCD matter is considered as one of the most promising probe 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 flavoured 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 semi-leptonic decays. The two main processes which 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 were 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 non-photonic single electrons. Several ingredients like inclusions of non-perturbative contributions from the quasi-hadronic bound state [10], 3-body scattering effects [11], the dissociation of heavy mesons due to 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 a field of high current 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 due to dead cone effects [21]. However, Aurenche and Zakharov claim that the radiative process has an anomalous mass dependence [22] due to the finite size of the QGP which leads to small difference in energy loss between a heavy and a light quarks. The mass dependence of the transverse momentum spectrum of the radiated gluons from the heavy quarks is studied in [23]. They 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 a light quarks. For high energy heavy quarks the effects of the dead cone, however, reduces because the magnitude of the angle forbidden for gluon emission behave as \( \sim \) heavy quark mass/energy [24]. From the study of the mass dependence of the radiative loss it is shown in [25] that the very energetic charm (not the bottom) quarks behave like massless partons. Although the authors in [26] concluded that the suppression of radiative loss for heavy quarks is due to dead cone effects but it will be fair to state that the issue is not settled yet.

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

The successes of the relativistic hydrodynamical model (see [30, 31] for review) in describing the host...
of experimental results from Relativistic Heavy Ion Collider (RHIC) indicate that the 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 i.e. the observed jet quenching and the large elliptic flow \( v_2 \) in Au+Au collisions at RHIC indicate the possibility of fast equilibration. On the one hand the experimental data indicate early thermalization time \( \tau \). Therefore, the interaction of the non-equilibrium statistical system reads:

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

where we have defined the kernels

\[
A_i = \int d^3k \omega(p,k)k_i
\]

\[
B_{ij} = \int d^3k \omega(p,k)k_i k_j.
\]

for \( |p| \rightarrow 0 \), \( A_i \rightarrow \gamma p_i \) and \( B_{ij} \rightarrow D \delta_{ij} \) where \( \gamma \) and \( D \) stand for drag and diffusion co-efficients respectively. The function \( \omega(p,k) \) is given by

\[
\omega(p,k) = g \int \frac{d^3q}{(2\pi)^3} f'(q)v\sigma_{p,q\rightarrow p-k,q+k}
\]

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, \( \sigma \) denotes the cross section and \( g \) is the statistical degeneracy. The co-efficients in the first two terms of the expansion in Eq. 8 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.

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

\[
\frac{df}{dt} = \frac{\partial}{\partial p_i} \left[ A_i(p)f + \frac{\partial}{\partial p_i}[B_{ij}(p)f] \right]
\]

The nonlinearity is caused due to the appearance of \( f' \) in \( A_i \) and \( B_{ij} \) through \( w(p,k) \). It arises from the

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

II. EVOLUTION OF HEAVY QUARK MOMENTUM DISTRIBUTIONS

The Boltzmann transport equation describing a non-equilibrium statistical system reads:

\[
\left[ \frac{\partial}{\partial t} + \frac{p}{E} \nabla_x + F \cdot \nabla_p \right] f(x,p,t) = \left[ \frac{\partial f}{\partial t} \right]_{\text{col}}
\]

where \( p \) and \( E \) denote momentum and energy, \( \nabla_x \) (\( \nabla_p \) are spatial (momentum space) gradient 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

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

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

\[
\left[ \frac{\partial f}{\partial t} \right]_{\text{col}} = \frac{\partial}{\partial p_i} \left[ A_i(p)f + \frac{\partial}{\partial p_i}[B_{ij}(p)f] \right]
\]

for \( |p| \rightarrow 0 \), \( A_i \rightarrow \gamma p_i \) and \( B_{ij} \rightarrow D \delta_{ij} \) where \( \gamma \) and \( D \) stand for drag and diffusion co-efficients respectively. The function \( \omega(p,k) \) is given by

\[
\omega(p,k) = g \int \frac{d^3q}{(2\pi)^3} f'(q)v\sigma_{p,q\rightarrow p-k,q+k}
\]

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, \( \sigma \) denotes the cross section and \( g \) is the statistical degeneracy. The co-efficients in the first two terms of the expansion in Eq. 8 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.

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

\[
\frac{df}{dt} = \frac{\partial}{\partial p_i} \left[ A_i(p)f + \frac{\partial}{\partial p_i}[B_{ij}(p)f] \right]
\]

The nonlinearity is caused due to 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 which 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 which 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 \( \gamma_i \) 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 [50] for details)

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

During the propagation through the QGP the heavy quarks dissipate energy predominantly by two processes: (i) collisional, e.g. \( gQ \rightarrow qQ, qQ \rightarrow qQ \) and \( qQ \rightarrow qQ \) and (ii) radiative processes, i.e. when the heavy quark emits gluons due to its interaction with the thermal partons in the plasma. Therefore, the drag and diffusion coefficient should include these two processes of energy dissipation.

The elastic collisions of heavy quarks with light quarks \( q \) and gluons \( g \) i.e.: \( gQ \rightarrow qQ, qQ \rightarrow qQ \) and \( qQ \rightarrow qQ \) have been used to evaluate the transport coefficients \( \gamma_{coll} \) and \( D_{coll} \) due to collisional process. At LHC energy one can not ignore the radiative energy loss, therefore, this should also be taken into account through the transport coefficients. The transport coefficient \( \gamma_{coll} \), \( \dot{q} \), which is related to the energy loss \( \dot{E}/dx \) of the propagating partons in the medium, has been used to calculate the shear viscosity to entropy density ratio, \( \eta/s \) [24, 60]. The \( \dot{q} \) is closely related to the diffusion coefficient \( D \) (for detail see [60]). In similar spirit we use \( dE/dx \) to calculate the drag coefficient of the medium and use Einsteins relation, \( D = T M \gamma \) to obtain the diffusion co-efficient 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:

\[
-\frac{dE}{dx} \bigg|_{\text{rad}} = \gamma_{rad} p
\]

where \( \gamma_{rad} \) denotes the drag co-efficient 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, i.e. the collisional process may influence the radiative one, 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 way, we add them up to obtain the effective drag co-efficients, \( \gamma_{eff} = \gamma_{coll} + \gamma_{rad} \) and similarly the effective diffusion coefficient: \( D_{eff} = D_{coll} + D_{rad} \). 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:

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

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

### III. ENERGY DISSIPATION PROCESSES

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

\[
\frac{dn_g}{d\eta d^2k_\perp} = \frac{C_A \alpha_s}{\pi^2} \frac{q_\perp^2}{k_\perp^2 (k_\perp - q_\perp)^2} F^2
\]

where \( k = (k_0, k_\perp, k_3) \) is the four momenta of the emitted gluon and \( q = (q_0, q_\perp, q_3) \) is the four momenta of the exchanged gluon, \( \eta = 1/2 \ln(k_0 + k_3)/(k_0 - k_3) \) is the rapidity and \( C_A = 3 \) 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 [24, 21]:

\[
F = \frac{k_\perp^2}{\omega^2\theta_0^2 + k_\perp^2}
\]

where \( \theta_0 = M/E \).

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

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

$$-\frac{dE}{dx} |_{rad} = \Lambda \epsilon = \tau^{-1} \epsilon$$  \hspace{1cm} (12)

where $\epsilon$, the average energy per collision is [63, 64]

$$\epsilon = \langle n_g k_0 \rangle = \int d\eta d^2k_\perp \frac{dn_g}{dp_T} k_0 \Theta(\tau - \tau_F) F^2$$  \hspace{1cm} (13)

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

$$k_\perp^{min} = \mu_D = \sqrt{4\pi\alpha_s T}$$  \hspace{1cm} (14)

The maximum transverse momentum of the emitted gluon is given by:

$$(k_\perp^{max})^2 = \frac{< (s - m^2)^2 >}{4s} = \frac{3ET}{2} - \frac{m^2}{4} + \frac{m^4}{48p_T} \left[ \ln\left( \frac{m^2 + 6ET + 6p_T}{m^2 + 6ET - 6p_T} \right) \right]$$ \hspace{1cm} (15)

Following the procedure of earlier works [50, 63] we evaluate the drag and diffusion coefficients for the elastic processes. Knowing $\gamma_{rad}$ from the radiative processes as described above we obtain the effective drag coefficients and hence effective diffusion coefficient through 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 very small contribution from the former at large $T$. 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 is more than or comparable to that from radiative processes. For the bottom quark we find that the gap between
In the present work the production of heavy quarks (charm and bottom) in hadronic collisions is studied extensively. The production of charm and bottom quarks in pp collisions have been taken from the NLO MNR code\cite{67}. The results from the code may be tested by measuring the distribution of the heavy mesons (containing c and b quarks) in pp collisions at $\sqrt{s_{NN}}=5.5$ TeV. With all these required inputs we solve the FP equation by using the Greens function technique (see\cite{51,65} for details).

**IV. THE NON-PHOTONIC ELECTRON SPECTRA**

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

$$f(z) \propto \frac{1}{z[1 - \frac{z}{1 - \frac{\epsilon_c}{1 - z}^2}]}$$  \hspace{1cm} (16)

for charm quark $\epsilon_c = 0.05$. For bottom quark $\epsilon_b = (M_c/M_b)^2 \epsilon_c$ where $M_c$ ($M_b$) is the charm (bottom) quark mass. The non-photonic single electron spectra originate from the decays of heavy flavoured mesons - e.g. $D \to Xe\nu$ or $B \to Xe\nu$ at mid-rapidity ($y = 0$) can be obtained as follows\cite{69,70}:

$$\frac{dN^e}{ptdp_T} = \int dqt\frac{dN^D}{q_Tdq_T} F(p_T,q_T)$$  \hspace{1cm} (17)

where

$$F(p_T,q_T) = \omega \int \frac{d(\mathbf{p}_T\cdot\mathbf{q}_T)}{2p_Tp_Tq_T} g(\mathbf{p}_T\cdot\mathbf{q}_T/M)$$  \hspace{1cm} (18)

related to the rest frame spectrum for the decay $D \to Xe\nu$ through the following relation\cite{69}:

$$\frac{1}{\Gamma_H} \frac{d\Gamma_H}{dE_e} = \omega g(E_e).$$  \hspace{1cm} (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 p-p 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,

$$R_{AA}(p_T) = \frac{\frac{dN^e}{dp_Tdy}_{Au+Au}}{N_{coll} \times \frac{dN^e}{dp_Tdy}_{p+p}}$$  \hspace{1cm} (21)

called the nuclear suppression factor, will be unity in the absence of any medium. In Eq.\cite{21} $N_{coll}$ stands for the number of nucleon-nucleon interactions in a nucleus+nucleus collision. The experimental data\cite{13} at RHIC energy ($\sqrt{s_{NN}}=200$ GeV) shows substantial suppression ($R_{AA} < 1$) for $p_T \geq 2$ GeV indicating substantial interaction of the plasma particles with charm and bottom quarks from which electrons are originated through the process: $c/b$ (hadronization) $\to D(B)(\text{decay}) \to e+X$. 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 such systems may be studied by solving the hydrodynamic equations:

$$\partial_\mu T^{\mu\nu} = 0$$  \hspace{1cm} (22)

with boost invariance along the longitudinal direction. In the above equation $T^{\mu\nu} = (\epsilon + P)u^\mu u^\nu - g^{\mu\nu}P$, is the energy momentum tensor for ideal fluid, $\epsilon$ is the energy density, $P$ is the pressure and $u^\mu$ 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 co-ordinate dependence 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 [52]).

The total amount of energy dissipated by a heavy quark in the QGP depends on the path length it traverses. Each parton traverse different path length which 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 $(r, \phi)$ in the plasma depends on the number of binary collisions at that point which can be taken as:

$$P(r, \phi) = \frac{2}{\pi R^2}(1 - \frac{r^2}{R^2})^2$$  \hspace{1cm} (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 $(r, \phi)$ in the transverse plane propagate a distance $L = \sqrt{R^2 - r^2\sin^2\phi - rcos\phi}$ in the medium. In the present work we use the following equation for the geometric average of the integral involving drag coefficient $[\int d\tau \gamma(\tau)]$:

$$\Gamma = \frac{\int r dr d\phi P(r, \phi) \int^{L/v}_{0} d\tau \gamma(\tau)}{\int r dr d\phi P(r, \phi)}$$  \hspace{1cm} (24)

where $v$ is the velocity of the propagating partons. Similar averaging has been performed for the diffusion co-efficient. For a static system the temperature dependence of the drag and diffusion co-efficients 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 like $\Gamma$ (Eq. (24)) and hence $R_{AA}$ becomes sensitive to velocity of sound ($c_s$) in the medium. This will be 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 [72]. For the QGP expected to be formed at the LHC we have taken initial temperature, $T_i = 700$ MeV, initial thermalization time $\tau_i = 0.08$ fm/c which reproduces the predicted hadron multiplicity $dN/dy = 2100$ [72] through the relation:

$$T_i^3 \tau_i \approx \frac{2\pi^4}{45\zeta(3)} \frac{1}{4a_{\text{eff}}} \frac{1}{\pi R^2_\Lambda} \frac{dN}{dy}.$$  \hspace{1cm} (25)

where $R_\Lambda$ is the radius of the system, $\zeta(3)$ is the Riemann zeta function and $a_{\text{eff}} = \pi^2 g_{\text{eff}}/90$ where $g_{\text{eff}} (= 2 \times 8 + 7 \times 2 \times 2 \times 3 \times N_F/8)$ is the degeneracy of quarks and gluons in QGP, $N_F$=number of flavours. We have taken the value of the transition temperature, $T_s = 170$ MeV.

The value of $R_{AA}$ is plotted against the $p_T$ of the non-photonic single electron resulting from $D$ decays in Fig. The results show substantial depletion at large $p_T$ indicating large interaction rate of the charm quarks with the thermal medium of partons. The sensitivity of the results on the equation of state is also demonstrated in Fig. 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. It may be mentioned here that $c_s$ increases with temperature. Therefore, due 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 suppression for the bottom quarks are displayed in Fig. We observe quantitatively less suppression compare to charm quarks. The difference between the charm and bottom quarks suppression are affected chiefly by two factors: (i) for different values transport coefficients and (ii) for the
different kind of initial $p_T$ distributions. The bottom quark has less drag coefficients and has harder $p_T$ distributions - both these factors are responsible for the smaller suppression of bottom quark. The present results on $R_{AA}$ may be compared with those obtained in [73] in a different approach.

In Fig 5 we have plotted the ratio: $R_{AA}^{D}/R_{AA}^{B}$ as a function of $p_T$, from where the effect of the mass and the role of the nature (soft or hard) of the initial $p_T$ distributions can be understood (see also [74]).

In Fig. 8 we compare the experimental data obtained by the STAR [1] and PHENIX [2] collaborations for $p_T$ of 0.2 fm/c. The value of initial thermalization time is assumed as 0.2 fm/c. The value of initial thermalization time and initial temperature reproduces the total multiplicity at mid-rapidity, $dN/dy = 1100$. We observe that the data can reasonably be reproduced by taking velocity of sound $c_s = 1/\sqrt{5}$. It should be mentioned here that the inclusion of both radiative and elastic losses in the effective drag enables us to reduce the gap be-

For the theoretical calculations the value of initial and transition temperatures are taken as 400 MeV and 170 MeV respectively. The value of initial thermalization time is assumed as 0.2 fm/c. These values of initial thermalization time and initial temperature reproduces the total multiplicity at mid-rapidity, $dN/dy = 1100$. We observe that the data can reasonably be reproduced by taking velocity of sound $c_s = 1/\sqrt{5}$. It should be mentioned here that the inclusion of both radiative and elastic losses in the effective drag enables us to reduce the gap be-
between the experiment and theory without any enhancement of the pQCD cross section as has been done in our previous work \cite{65}.

So far, we have discussed the suppression of the non-photonic electron produced in nuclear collisions due to the propagation of the heavy quark in the partonic medium in the pre-hadronization era. However, the suppression of the D mesons in the post-hadronization 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 post-hadronization era is found to be small \cite{72}, 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 loss 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 NLO MNR \cite{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 semi-leptonic 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 non-photonic 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.

Acknowledgment: We are grateful to Matteo Cacciari for providing us the heavy quarks transverse momentum distribution and also for useful discussions. This work is supported by DAE-BRNS project Sanction No. 2005/21/5-BRNS/2455.

---

[1] B. I. Abeleb et al. (STAR Collaboration), Phys. Rev. Lett. 98, 192301 (2007).
[2] A. Adare et al. (PHENIX Collaboration), Phys. Rev. Lett. 98, 172301 (2007).
[3] S. S. Adler et al. (PHENIX Collaboration), Phys. Rev. Lett. 96, 032301 (2006).
[4] J. D. Bjorken, Fermilab preprint 82/59-THY (1982) unpublished.
[5] M. H. Thoma and M. Gyulassy, Nucl. Phys. B 351, 491 (1991).
[6] A. Peshier, Phys. Rev. Lett. 97, 212301 (2006).
[7] A. K. Dutt-Mazumder, J. Alam, P. Roy and B. Sinha, Phys. Rev. D 71, 094016 (2005).
[8] M. G. Mustafa and M. H. Thoma, Acta Phys. Hung. A 22, 93 (2005).
[9] E. Braaten and M. H. Thoma, Phys. Rev. D 44, 2625 (1991).
[10] H. van Hees, M. Mannarelli, V. Greco and R. Rapp, Phys. Rev. Lett. 100, 192301 (2008).
[11] C. M. Ko and W. Liu, Nucl. Phys. A 783, 23c (2007).
[12] A. Adil and I. Vitev, Phys. Rev. Lett. B 649, 139 (2007).
[13] P. B. Gossiaux and J. Aichelin, Phys. Rev. C 78, 014904 (2008).
[14] S. Wicks, W. Horowitz, M. Djordjevic and M. Gyulassy, Nucl. Phys. A 784, 426 (2007).
[15] M. Gyulassy, P. Levai and I. Vitev, Nucl. Phys. B 571, 197 (2000); M. Gyulassy, P. Levai and I. Vitev, Phys. Rev. Lett. 85, 5535 (2000); M. Gyulassy and X.-N. Wang, Nucl. Phys. B 420, 583 (1994).
[16] H. Zhang, J. F. Owens, E. Wang and X. N. Wang, Phys. Rev. Lett. 98, 212301 (2007).
[17] R. Baier, Y. L. Dokshitzer, S. Peigne and D. Schiff, Phys. Lett. B 345, 277 (1995); R. Baier, Y. L. Dokshitzer, A. H. Mueller and D. Schiff, Nucl. Phys. B 531, 403 (1998).
[18] C. A. Salgado and U. A. Wiedemann, Phys. Rev. Lett. 89, 092303 (2002).
[19] P. Jacobs and X. N. Wang, Prog. Part. Nucl. Phys. 54, 443 (2005); R. Baier, D. Schiff, B. G. Zakharov, Ann. Rev. Nucl. Part. Sci. 50, 37 (2000).
[20] Y. L. Dokshitzer and D. E. Kharzeev, Phys. Lett. B, 519, 199 (2001).
[21] R. K. Ellis, W. J. Stirling and B. R. Webber, QCD and Collider Physics, Cambridge University Press, Cambridge, 1996.
[22] B. G. Zakharov, JETP Lett. 86, 444(2007); P. Aurenche and B. G. Zakharov, JETP Lett. 90, 237 (2009).
[23] N. Armesto, C. A. Salgado and U. A. Wiedemann, Phys. Rev. D 69, 114003 (2004).
[24] B. W. Zhang, E. Wang and X.-N. Wang, Phys. Rev. Lett. 93, 072301 (2004).
[25] N. Armesto, A. Dainese, C. A. Salgado and U. A. Wiedemann, Phys. Rev. D 71, 054027 (2005).
[26] Roy A. Lacey et al., Phys. Rev. Lett. 103, 142302(2009)
[27] M. Gyulassy and X.-N. Wang, Nucl. Phys. B 420, 583 (1994).
[28] S. Klein, Rev of Modern Phys, 71, 1501(1999)
[29] M. Djordjevic and M. Gyulassy, Nucl. Phys. A 733, 265 (2004).
[30] P. Huovinen and P. V. Ruuskanen, Ann. Rev. Nucl. Part. Sci. 56, 163 (2006).
[31] D. A. Teaney. \texttt{arXiv:0905.2433} [nucl-th].

[32] I. Arsene \textit{et al.} (BRAHMS Collaboration), Nucl. Phys. A \textbf{757}, 1 (2005); B. B. Back \textit{et al.} (PHOBOS Collaboration), Nucl. Phys. A \textbf{757}, 28 (2005); J. Adams \textit{et al.} (STAR Collaboration), Nucl. Phys. A \textbf{757}, 102 (2005); K. Adcox \textit{et al.} (PHENIX Collaboration), Nucl. Phys. A \textbf{757}, 184 (2005).

[33] S. S. Adler \textit{et al.} (PHENIX Collaboration), Phys. Rev. Lett. \textbf{96}, 202301 (2006).

[34] J. Adams \textit{et al.} (STAR Collaboration), Phys. Rev. Lett. \textbf{91}, 072304 (2003).

[35] S. S. Adler \textit{et al.} (PHENIX Collaboration), Phys. Rev. Lett. \textbf{91}, 182301 (2003).

[36] K. H. Ackemann \textit{et al.} (STAR Collaboration), Phys. Rev. Lett. \textbf{86}, 402 (2003).

[37] P. Arnold, J. Lenaghan, G. D. Moore and L. G. Phys. Rev. Lett. \textbf{94}, 072302 (2005).

[38] R. Baier, A. H. Mueller, D. Schiff and D. T. Son, Phys. Lett. B \textbf{539}, 46 (2002).

[39] P. Romatschke and R. Venugopalan, Phys. Rev. Lett. \textbf{96}, 062302 (2006).

[40] S. Mrowczynski, Phys. Lett. B \textbf{314} (1993) 118; Phys. Rev. C \textbf{49}, 2191 (1994); Phys. Lett. B \textbf{393}, 26 (1997).

[41] P. Romatschke and M. Strickland, Phys. Rev. D \textbf{68}, 036004 (2003).

[42] P. Arnold, J. Lenaghan and G. D. Moore, J. High Energy Phys \textbf{08}, 002 (2003).

[43] P. Arnold, G. D. Moore and L. G. Yaffe, J. High Energy Phys. \textbf{01}, 030 (2003).

[44] G. D. Moore and D. Teaney, Phys. Rev. C \textbf{71}, 064904 (2005).

[45] J. Alam, S. Raha and B. Sinha, Phys. Rev. Lett. \textbf{73}, 1895 (1994).

[46] E. Shuryak, Phys. Rev. Lett. \textbf{68}, 3270 (1992).

[47] E. M. Lifshitz and L. P. Pitaevskii, Physical Kinetics, Butterworth-Hienemann, Oxford 1981.

[48] R. Balescu, Equilibrium and Non-Equilibrium Statistical Mechanics (Wiley, New York, 1975).

[49] S. Chakraborty and D. Syam, Lett. Nuovo Cim. \textbf{41}, 381 (1984).

[50] B. Svobitsky, Phys. Rev. D \textbf{37}, 2484 (1988).

[51] H. van Hees, R. Rapp, Phys. Rev. C \textbf{71}, 034907 (2005).

[52] S. Turbide, C. Gale, S. Jeon and G. D. Moore, Phys. Rev. C \textbf{72}, 014906 (2005).

[53] J. Bjorkerak and R. Venugopalan, Phys. Rev. C \textbf{63}, 024909 (2001).

[54] P. Roy, J. Alam, S. Sarkar, B. Sinha and S. Raha, Nucl. Phys. A \textbf{624}, 687 (1997).

[55] M G. Mustafa and M. H. Thoma, Acta Phys. Hung. A \textbf{22}, 93 (2005).

[56] P. Roy, A. K. Dutt-Mazumder and J. Alam, Phys. Rev. C \textbf{73}, 044911 (2006).

[57] O. Fochler, Z. Xu and C. Greiner, \texttt{arXiv:1003.4380} [hep-ph].

[58] R. Baier, Yu.L. Dokshitzer, A.H.Mueller, S. Peigne, D. Schiff, Nucl. Phys. B, \textbf{484}, 265(1997).

[59] R. Baier, Nucl. Phys. A. \textbf{715}, 209(2003).

[60] A. Majumder, B. Mueller and X.N. Wang, Phys. Rev. Lett.\textbf{99}, 192301(2007).

[61] O. Kaczmarek and F. Zantow, Phys. Rev. D, \textbf{71}, 114510(2005).

[62] J.F. Gunion, G.Bertsch, Phys. Rev. D, \textbf{25}, 746(1982).

[63] X. W. Chang et al, Chin. Phys. Lett. \textbf{22}, 72 (2005).

[64] M. G. Mustafa, D. Pal, D. K. Srivastava and M. H. Thoma, Phys. Lett. B, \textbf{428},234(1998).

[65] S$\bar{\text{S}}$ K Das, J. Alam and P. Mohanty, Phys. Rev. C \textbf{80}, 054916 (2009).

[66] M. Cacciari, S. Frixione, M.L. Mangano, P. Nason and G. Ridolfi, J. High Ener. Phys. \textbf{0407}, 033 (2004); M. Cacciari and P. Nason, Phys. Rev. Lett. \textbf{89}, 122003 (2002); M. Cacciari and P. Nason, J. High Ener. Phys. \textbf{0309}, 006 (2003); M. Cacciari, P. Nason and R. Vogt, Phys. Rev. Lett. \textbf{95}, 122001 (2005).

[67] M. L. Mangano, P. Nason and G. Ridolfi, Nucl. Phys. B \textbf{373}, 295 (1992).

[68] C. Peterson \textit{et al.}, Phys. Rev. D \textbf{27}, 105 (1983).

[69] M. Gronau, C. H. Llewellyn Smith, T. F. Walsh, S. Wolfram and T. C. Yang, Nucl. Phys. B \textbf{123}, 47 (1977).

[70] A. Ali, Z. Phys. C \textbf{1}, 25 (1979).

[71] J. D. Bjorken, Phys. Rev. D \textbf{27} , 140(1983).

[72] N. Armesto, N. Borghini, S. Jeon, U.A. Wiedemann (ed), J. Phys. G: Nucl. Part. Phys. \textbf{35}, 054001 (2008).

[73] M. Djordjevic, M. Gyulassy and S. Wicks, Phys. Rev. Lett. \textbf{94}, 112301 (2005).

[74] N. Armesto, M. Cacciari, A. Dainese, C. A. Salgado and U. A. Wiedemann, J. Phys. G \textbf{32}, 5421 (2006); N. Armesto, M. Cacciari, A. Dainese, C. A. Salgado and U. A. Wiedemann, J. Phys. G \textbf{35}, 054001 (2008); W. A. Horowitz, J. Phys. G \textbf{35}, 054001 (2008); I. Vitev, J. Phys. G \textbf{35}, 054001 (2008).

[75] S. K Das, J. Alam, P. Mohanty and B. Sinha, Phys. Rev. C \textbf{81}, 044912 (2010).