Nuclear suppression at low energy heavy ion collisions

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The effects of non-zero baryonic chemical potential on the drag and diffusion coefficients of heavy quarks propagating through a baryon rich quark gluon plasma have been studied. The nuclear suppression factor, $R_{AA}$ for non-photonic single electron spectra resulting from the semileptonic decays of hadrons containing heavy flavours have been evaluated for low energy collisions. The role of non-zero baryonic chemical potential on $R_{AA}$ has been highlighted.

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I. INTRODUCTION

The nuclear collisions at low energy RHIC run [1, 2] and GSI-FAIR [3] is expected to create a thermal medium with large baryonic chemical potential ($\mu_B$) and moderate temperature ($T$). The heavy flavours namely, charm and bottom quarks may play a crucial role in understanding the properties of such medium because they do not constitute the bulk part of the system and their thermalization time scale is larger than the light quarks and gluons and hence can retain the interaction history very effectively. The perturbative QCD (pQCD) calculations indicate that the heavy quark ($Q$) thermalization time, $\tau_Q$ is larger [4, 5] than the light quarks and gluons thermalization scale $\tau_i$. Gluons may thermalized even before up and down quarks [6, 7]. In the present work we assume that the Quark Gluon Plasma (QGP) is formed at time $\tau_i$. Therefore, the interaction of the non-equilibrated heavy quarks with the equilibrated QGP for the time interval $\tau_i < \tau < \tau_Q$ can be treated within the ambit of the Fokker-Planck (FP) equation [8, 9] i.e. the heavy quark can be thought of executing Brownian motion [4, 6, 10, 17] in the heat bath of QGP during the said interval of time.

As the relaxation time for heavy quarks of mass $M$ at a temperature $T$ are larger than the corresponding quantities for light partons by a factor of $M/T(>1)$ [4] i.e. the light quarks and the gluons get thermalized faster than the heavy quarks, the propagation of heavy quarks through QGP (mainly contains light quarks and gluons) therefore, may be treated as the interactions between equilibrium and non-equilibrium degrees of freedom. The FP equation provide an appropriate frame work for such processes. In case of low energy collisions the radiative energy loss of of heavy quarks will be much smaller than the loss due to elastic processes. Moreover, the thermal production of charm and bottom quarks can be ignored for the range of temperature and baryonic chemical potential under study. Therefore, the FP equation is better applicable in the present situation.

The paper is organized as follows. In the next section a brief description of Fokker Planck equation and the $T$ and $\mu_q$ dependence of the drag and diffusion co-efficients are outlined, the non-photonic electron spectra is discussed in section III, the initial conditions and the space time evolution have been discussed in section IV, section V is devoted to the discussions on nuclear suppression and finally section VI contains the summary and conclusions.

II. THE FOKKER PLANCK EQUATION

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

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

(1)

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

$$\frac{\partial f}{\partial t} = \left[\frac{\partial f}{\partial t}\right]_{\text{col}}$$

(2)

The collision term on the right hand side of the above equation can be approximated as (see [11, 12] for details):

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

(3)

where we have defined the kernels

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

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

(4)
for \( \mathbf{p} \rightarrow 0 \), \( A_i \rightarrow \gamma \rho_i \) 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

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

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, \( \sigma \) denotes the cross section and \( q \) is the statistical degeneracy. The co-efficients 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 \[8\].

With these approximations the Boltzmann equation reduces to a non-linear integro-differential equation known as 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 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 Fokker-Planck equation\[8\] describing the interaction of a particle which is out of thermal equilibrium with the particles in a thermal bath of light quarks, anti-quarks 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 distribution \( f \) is governed by the FP equation \[11\]

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

where \( \gamma_i \) and \( D_{ij} \) are the drag and diffusion coefficients. The elastic collisions of heavy quarks \( Q \) stands for charm and bottom) with thermal light quarks \( q \), anti-quarks \( \bar{q} \) and gluons \( g \) i.e. \( Qq \rightarrow Qq, Q\bar{q} \rightarrow Q\bar{q} \) and \( Qg \rightarrow Qg \) have been used to evaluate the drag and diffusion coefficients as indicated in \[11, 18\]. The thermal distribution of the quarks and the anti-quarks are responsible for the \( T \) and \( \mu_q \) dependence of the drag and diffusion co-efficients. The gluon distribution introduces the \( T \) dependence to these quantities, chemical potential dependence of the drag and diffusion co-efficients originate from the thermal phase space of quarks and anti-quarks.

### A. The drag and diffusion co-efficients

At low \( \sqrt{s_{\text{NN}}} \) the net baryon density at mid-rapidity is non-zero and its value could be high depending on the value of \( \sqrt{s_{\text{NN}}} \). Therefore, we need to solve the FP equation for non-zero \( \mu_B \). The drag and diffusion coefficients are functions of both the thermodynamical variables: \( \mu_B \) and \( T \).

The energy dependence of the chemical potential has been obtained from the parametrization of the experimental data on hadronic ratios as \[19\] (see also \[20\]),

\[
\mu_B(s_{\text{NN}}) = a(1 + \sqrt{s_{\text{NN}}}/b)^{-1} \tag{8}
\]

where \( a = 0.967 \pm 0.032 \) GeV and \( b = 6.138 \pm 0.399 \) GeV. The parametrization in Eq. \[8\] gives the values of \( \mu_B \) at the freeze-out. The corresponding values at the initial condition are obtained from the baryon number conservation equation. The initial baryonic chemical potential carried by the quarks \( \mu_q = \mu_B/3 \) are shown in table \[1\] for various \( \sqrt{s_{\text{NN}}} \) under consideration.

In the present work we take \( \alpha_s = 0.3 \) because the dependence of the strong coupling on the temperature and baryonic chemical potential is not accurately known yet. The sensitivity of collisional energy loss on the running \( \alpha_s \) is studied in Ref. \[21\] in detail. The variation of the drag coefficients of charm quarks (due to its interactions with quarks and anti-quarks) on the baryonic chemical potential for different \( T \) are displayed in Fig. \[1\]. The drag coefficient for the process: \( Qq \rightarrow Qg \) is \( 8.42 \times 10^{-3} \) \( \text{fm}^{-1} \) \( (1.86 \times 10^{-2} \text{ fm}^{-1}) \) for \( T = 140 \) MeV \( (190 \) MeV) (not displayed in Fig. \[1\]). The \( T \) and \( \mu_q \) dependence of the drag and diffusion co-efficients may be understood as follows. The drag may be defined as the thermal average of the square of the invariant transition amplitude weighted by the momentum transfer for the reactions \( qQ \rightarrow qQ, Q\bar{q} \rightarrow Q\bar{q} \) and \( gQ \rightarrow gQ \). As the temperature of the thermal bath increases the light quarks \( q \) and the gluons move faster and gain the ability to transfer larger momentum during their interaction with the heavy quarks - resulting in the increase of the drag of the heavy quarks propagating through the partonic medium. Since the average momentum of the quarks increases with \( \mu_q \), similar behaviour is expected in the variation of drag with baryonic chemical potential. This
trend is clearly observed in the results displayed in Fig. 4 for charm quark. The drag due to the process $Qq \rightarrow Qq$ is larger than the $Qq \rightarrow Q\bar{q}$ interaction because for non-zero chemical potential, the $Q$ propagating through the medium encounters more $q$ than $\bar{q}$ at a given $\mu_q$. For vanishing chemical potential the contributions from quarks and anti-quarks are same.

In the same way it may be argued that the diffusion coefficient involves the square of the momentum transfer - which should also increase with $T$ and $\mu_q$ as observed in Fig. 2. The diffusion co-efficient for charm quarks due to its interaction with gluons is given by $\sim 1.42 \times 10^{-3}$ GeV$^2$/fm (4.31 $\times 10^{-3}$ GeV$^2$/fm) for $T = 140$ MeV (190 MeV). It may be mentioned here that the drag increases with $T$ when the system behave like a gas. In case of liquid the drag may decrease with temperature (except very few cases) - because a substantial part of the thermal energy goes in making the attraction between the interacting particles weaker - allowing them to move more freely and hence making the drag force lesser. The drag co-efficient of the partonic system with non-perturbative effects may decrease with temperature as shown in Ref. [22] - because in this case the system interacts strongly more like a liquid. The heavy quark momentum diffusion coefficient has been computed [23] at next to leading order within the ambit of hard thermal loop approximations. For $T \sim 400$ MeV the momentum averaged pQCD value (for $\mu_q = 0$) of the diffusion co-efficient obtained in the present work is comparable to the value obtained in [23] in the leading order approximation for the same set of inputs (e.g. strong coupling constant, number of flavours etc). The drag and diffusion coefficients for bottom quarks are displayed in Figs. 3 and 4 respectively, showing qualitatively similar behaviour as charm quarks. The drag co-efficients for bottom quarks due to the process $Qg \rightarrow Qg$ is given by $\sim 3.15 \times 10^{-3}$ fm$^{-1}$ and $6.93 \times 10^{-3}$ fm$^{-1}$ at $T = 140$ MeV and 190 MeV respectively. The corresponding diffusion coefficients are $\sim 1.79 \times 10^{-3}$ GeV$^2$/fm and $5.38 \times 10^{-3}$ GeV$^2$/fm at $T=140$ MeV and 190 MeV respectively.

III. THE NON-PHOTONIC ELECTRON SPECTRA

After obtaining the drag and diffusion coefficients we need the initial heavy quark momentum distributions for solving the FP equation. For low collision energy rigorous QCD based calculations for heavy flavour production is not available (for higher $\sqrt{s_{NN}} = 200$ GeV see Ref. [24] for rigorous QCD calculations). In the present work this is obtained from pQCD calculation [25, 26] for the processes: $gg \rightarrow Q\bar{Q}$ and $q \bar{q} \rightarrow Q\bar{Q}$. Here we intend to deal with the nuclear suppression factor, $R_{AA}$, which involves the ratio of the momentum distribution functions. Therefore, the final results may not be too sensitive to the initial distributions because of some cancellations that may take place in the ratio.

With the initial condition mentioned above the FP equation has been solved for the heavy quarks. We convolute the solution with the fragmenta-

FIG. 1: Variation of the drag coefficient of charm quark due to its interactions with light quarks and anti-quarks as a function of $\mu_q$ for different temperatures.

FIG. 2: Variation of the diffusion coefficient of charm quark due to its interactions with light quarks and anti-quarks as a function of $\mu_q$ for different temperatures.
use Peterson function \([27]\) given by:

\[
f(z) \propto \frac{1}{[z + (\frac{1}{2} - \frac{\epsilon_c}{2z})^2]}\]  

(9)

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.

We evaluate the electron spectra from the decays of heavy mesons originating from the fragmentation of the heavy quarks propagating through the QGP medium formed in heavy ion collisions. In the same way the electron spectrum from the p-p collisions can be obtained from the charm and bottom quark distribution which goes as initial conditions to the solution of FP equation. The ratio of these two quantities gives the nuclear suppression, \(R_{AA}\) as:

\[
R_{AA}(p_T) = \frac{dN_{ee}^{Au+Au}}{d^5p_Tdy} \times \frac{N_{coll}}{d^5p_Tdy} \]

(14)
called the nuclear suppression factor, will be unity in the absence of any medium. In the above equation \(N_{coll}\) denotes the number of nucleon nucleon collisions in Au+Au interaction. However, the experimental data \([31,32]\) 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 which are originated through the process: \(c(b)\) (hadronization) \(\rightarrow D(B)\) (decay) \(\rightarrow e + X\). 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.

IV. THE INITIAL CONDITIONS FOR THE SPACE-TIME EVOLUTION

The nuclear suppression for heavy quarks depend on the parameters like initial temperature \((T_i)\), thermalization time \((\tau_i)\), equation of state (EOS) and the transition temperature \((T_c)\).

We assume that the system reaches equilibration at a time \(\tau_i\) after the collision at temperature \(T_i\) which are related to the produced hadronic (predominantly mesons) multiplicity through the following relation:

\[
T_i^3 \tau_i \approx \frac{2\pi^4}{45\zeta(3)} \frac{1}{4a_{eff}} \frac{1}{\pi R_A^2} \frac{dN}{dy}.
\]

(15)
where $R_A$ is the radius of the system, $\zeta(3)$ is the Riemann zeta function and $a_{eff} = \pi^2 g_{eff}/90$ where $g_{eff} = 2 \times 8 + 7 \times 2 \times 3 \times N_F/8$ is the degeneracy of quarks and gluons in QGP, $N_F$ = number of flavours.

The value of the multiplicities for various $\sqrt{s_{NN}}$ have been calculated from the Eq. below [33]:

$$\frac{dN}{dy} = \frac{dn_{pp}}{dy} \left[ (1 - x) \frac{N_{part}}{2} + x < N_{coll} > \right]$$

(16)

$N_{coll}$ is the number of collisions and contribute $x$ fraction to the multiplicity $dn_{pp}/dy$ measured in $pp$ collision. The number of participants, $N_{part}$ contributes a fraction $(1 - x)$ to $dn_{pp}/dy$, which is given by

$$\frac{dn_{pp}}{dy} = 2.5 - 0.25ln(s) + 0.023ln^2(s)$$

(17)

The values of $N_{part}$ and $N_{coll}$ are estimated for $(0-5\%)$ centralities by using Glauber Model. The value of $x$ depends very weakly on $\sqrt{s_{NN}}$ [34], in the present work we have taken $x = 0.1$ for all the energies.

The time evolution of the temperature and the baryon density ($n_B$) have been obtained by solving the following equations:

$$\partial_\tau T^{\mu\nu} = 0, \quad \partial_\mu n_B^\mu = 0$$

(18)

in (1+1) dimension with boost invariance along the longitudinal direction [35]. In the above equation $T^{\mu\nu} = (\epsilon + P)u^\mu u^\nu - g^{\mu\nu}P$, is the energy momentum tensor and $n_B^\mu = n_B u^\mu$ is the baryonic flux, where $\epsilon$ is the energy density, $P$ is the pressure, and $u^\mu$ is the hydrodynamic four velocity. The radial co-ordinate dependence of $T$ and $n_B$ have been parametrized as in Ref. [13]. The velocity of sound for the QGP phase is taken as $c_s = 1/\sqrt{3}$. 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 effects may have negligible effects on the nuclear suppression (see also [13]). Moreover, at lower collision energies (as in the present case) the amount of radial flow will not be as substantial as in $\sqrt{s_{NN}} = 200$ GeV.

The total amount of energy dissipated by a heavy quark in the QGP depends on the path length it traverses. Each parton traverses different path length which depends on the geometry of the system and on the point where its is produced. The probability that a parton is created at a point $(r, \phi)$ in the plasma depends on the number of binary collisions at that point which can be taken as [13]:

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

(19)

where $R$ is the nuclear radius. A parton created at $(r, \phi)$ in the transverse plane propagate 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 averaging of the drag coefficient:

$$\Gamma = \int r dr d\phi P(r, \phi) \int^L_0 d\tau \gamma(\tau)$$

(20)

where $v$ is the velocity of the propagating partons. Similar averaging has been performed for the diffusion co-efficient. For a static system the $T$ and $\mu_q$ dependence of the drag and diffusion co-efficients of the heavy quarks enter via the thermal distributions of light quarks, anti-quarks and gluons through which it is propagating. However, in the present scenario the variation of the temperature and the

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**TABLE I**: The values center of mass energy , $dN/dy$, initial temperature ($T_i$) and quark chemical potential - used in the present calculations.

| $\sqrt{(s_{NN}) (GeV)}$ | $dN/dy$ | $T_i (MeV)$ | $\mu_q (MeV)$ |
|--------------------------|---------|-------------|---------------|
| 39                       | 617     | 240         | 62            |
| 27                       | 592     | 199         | 70            |
| 17.3                     | 574     | 198         | 100           |
| 7.7                      | 561     | 197         | 165           |
baryon density with time are governed by the equation of state or the velocity of sound of the thermalized system undergoing hydrodynamic expansion. In such a scenario the quantities like $\Gamma$ (Eq. 20) and hence $R_{AA}$ becomes sensitive to velocity of sound in the medium.

![Graph showing nuclear suppression factor, $R_{AA}$ as function of $p_T$ for various $\sqrt{s_{NN}}$. Inset: the nuclear suppression factor due to the interaction of $D$ meson in a thermal medium of pions and nucleons.](image)

**FIG. 6: Nuclear suppression factor, $R_{AA}$ as function of $p_T$ for various $\sqrt{s_{NN}}$. Inset: the nuclear suppression factor due to the interaction of $D$ meson in a thermal medium of pions and nucleons.**

V. THE NUCLEAR SUPPRESSION

To demonstrate the effect of non-zero baryonic chemical potential we evaluate $R_{AA}$ for $\mu_q = 200$ MeV and $\mu_q = 0$ for a given $T_0 = 200$ MeV. The results are displayed in Fig. 5- representing the combined effects of temperature and baryon density on the viscous drag and diffusion. The viscous drag on the heavy quarks due to its interaction with quarks is larger than that of its interactions with the anti-quarks (Fig. 1). Resulting in larger suppression in the former case than the later. The net suppression of the electron spectra from the Au+Au collisions compared to p+p collisions is effected by quarks, anti-quarks and gluons. The results for net suppressions are displayed for $\mu_q = 200$ MeV (dashed line) and $\mu_q = 0$ (with asterisk). The experimental detection of the non-zero baryonic effects will shed light on the net baryon density (and hence baryon stopping) in the central rapidity region. However, whether the effects of non-zero baryonic chemical potential is detectable or not will depend on the overall experimental performance.

The results for $R_{AA}$ are shown in Fig. 6 for various $\sqrt{s_{NN}}$ with inputs from table I. We observe that at large $p_T$ the suppression is similar for all energies under consideration. This is because the collisions at high $\sqrt{s_{NN}}$ are associated with large temperature but small baryon density at mid-rapidity- which is compensated by large baryon density and small temperature at low $\sqrt{s_{NN}}$ collisions. Low $p_T$ particles predominantly originate from low temperature and low density part of the evolution where drag is less and so is the nuclear suppression.

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. We have estimated $R_{AA}$ for $D$ mesons due to $t_s$ interaction with pions and nucleons and found that it has a value value closer to unity, indicating the fact that the hadronic medium (of pions and nucleons) is unable to drag the $D$ mesons. Therefore, the measured depletion in $R_{AA}$ for the non-photonic electron will indicate the presence of partonic medium and the amount of depletion may the used to characterize the thermal medium.

It has been shown in [38] that a large enhancement of the pQCD cross section is required for the reproduction of experimental data on elliptic flow at RHIC energies. In our earlier work [18] we have evaluated the $R_{AA}$ for non-photonic single electron spectra resulting from the semileptonic decays of hadrons containing heavy flavours and observed that the data from RHIC collisions at $\sqrt{s_{NN}} = 200$ GeV are well reproduced by enhancing the pQCD cross sections by a factor 2 and with an equation of state $P = \epsilon/4$. In the same spirit we evaluate $R_{AA}$ with twice enhanced pQCD cross section and keeping all other quantities unaltered (Fig. 7). The results in Fig. 7 show stronger suppression as compared to the results displayed in Fig. 6 but it is similar in all the energies under consideration.

VI. SUMMARY AND CONCLUSIONS

We have studied the effects of baryonic chemical potential and temperature on the drag and diffusion coefficients of heavy quarks moving in a thermalized system of quarks and gluons. We have observed that both the drag and diffusion coefficients increase with temperature and chemical potential. When we have enhanced the pQCD cross section for the interaction of the heavy quarks with the thermal system by a factor of two - the resulting suppressions in $R_{AA}$ are between $20\% - 30\%$ for $\sqrt{s_{NN}} = 39 - 7.7$ GeV. The radiative energy loss [39] (see [14] for
FIG. 7: Same as Fig. 6 with enhancement of cross section by a factor of 2.

...a review) of heavy quarks is suppressed due to dead cone effects and has been neglected in the present work. Moreover, at low collisions energies the collisional loss is dominant over its radiative counterpart (see [15, 16, 17] for details). It may be mentioned here that the theoretical formalism, the FP equation is applicable better for heavy quarks than light quarks and gluons (because of their frequent productions and annihilations). However, the production of charm and bottom quarks are smaller at low energy collisions making the measurements of non-photonic single electron spectra and hence $R_{AA}$ for heavy quarks difficult. The detection of the non-zero baryonic chemical potential effects observed in the present work through the nuclear suppression factor will help in determining the net baryon density (and hence baryon stopping) in the mid-rapidity region. However, whether such effects is detectable experimentally or not will depend on the overall experimental performance.

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I. Krivoruchenko, Phys. Rev. C 73, 035204 (2006).

[37] J. Haidenbauer, G. Krein, U. -G. Meissner and A. Sibirtsev, Eur. Phys. J. A 33, 107 (2007).

[38] D. Molnar and M. Gyulassy, Nucl. Phys. A 697, 495 (2002).

[39] A. Dainese, C. Loizides and G. Paic, Eur. Phys. J. C 38 (2005) 461; C. Loizides, ibid. 49, 339 (2007).

[40] 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).

[41] H. Zhang, J. F. Owens, E. Wang and X. N. Wang, Phys. Rev. Lett. 98, 212301 (2007).

[42] R. Baier, Y. L. Dokshitzer, S. Peigne and D. Schiff, Phys. Lett. B 345, 277 (1995); R. Baier, Y. L. Dokshitzer, A. H. Mueller, S. Peigne and D. Schiff, Nucl. Phys. B 531, 403 (1998);

[43] C. A. Salgado and U. A. Wiedemann, Phys. Rev. Lett. 89, 092303 (2002).

[44] 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).

[45] H. van Hees, M. Mannarelli, V. Greco and R. Rapp, Phys. Rev. Lett. 100, 192301 (2008).

[46] C. M. Ko and W. Liu, Nucl. Phys. A 783, 23c (2007).

[47] A. K. Dutt-Mazumder, J. Alam, P. Roy and B. Sinha, Phys. Rev. D 71, 094016 (2005).
$R_{AA}^{1/2} = 39$ GeV
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