Heavy quark energy loss and *D*-mesons in RHIC and LHC energies

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A R T I C L E   I N F O

A B S T R A C T

We obtain the radiative energy loss of a heavy quark in a deconfined medium due to radiation of gluons off them using a recently derived generalised gluon emission spectrum. We find that the heavy flavour loses energy almost in a similar fashion like light quarks through this process. With this, we further analyse the nuclear modification factor for *D*-meson at LHC and RHIC energies. In particular, the obtained result is found to be in close agreement with the most recent data from ALICE collaboration at 2.76 ATeV Pb–Pb collisions. We also discuss the nuclear modification factor due to the collisional energy loss. Furthermore, the result of non-photonic single electron from the decay of both *D*- and *B*-mesons is compared with the RHIC data at 200 AGeV Au–Au collisions, which is also in close agreement.

1. Introduction

The purpose of ongoing relativistic heavy-ion collisions is to understand the properties of nuclear or hadronic matter at extreme conditions. A primary aim lies in the detection of a new state of matter formed in these collisions, the quark–gluon plasma (QGP), where the quarks and gluons are liberated from the nucleons and move freely over an extended region rather than over a limited hadronic volume. Various diagnostic measurements taken at CERN Super Proton Synchrotron (SPS) [1] in the past and at BNL Relativistic Heavy Ion Collider (RHIC) [2–8] in the recent past have provided strong hints for the formation of QGP within the first few fm/*c* of the collisions through the manifestation of hadronic final states. New data from heavy-ion experiments at CERN Large Hadron Collider (LHC) [9–11] have further indicated the formation of such a state of matter.

One of the important features of the plasma produced in heavy-ion collisions is suppressed production of high energy hadrons compared to the case of pp collisions, called jet quenching. The term 'jet quenching', generally, ascribes to the modification of an energetic parton due to its interaction with the coloured medium while passing through it. The basic idea is that the scales of hard (high-$p_\perp$) processes and the medium interactions in the context of heavy-ion collisions, are very distinct in accordance with the uncertainty principle. This provides the fact that the high-$p_\perp$ parton production in *A*–*A* collisions can be computed using perturbative QCD (pQCD), which is quite close to the vacuum rate scaled for binary *N*–*N* collisions in an *A*–*A* collision. The effect of medium is then treated as a final state interaction which is taken into account through the modification of the outgoing parton fragmentation pattern due to parton-medium interactions.

The heavy-ion program at BNL RHIC [8] has clearly revealed that the phenomenon of jet quenching is mainly caused due to the energy loss of the initial hard parton via collisional and radiative processes, prior to hadronisation. The indication for jet quenching in heavy-ion program at CERN LHC has also been [9–11] observed recently. The energy loss encountered by an energetic parton in a QCD medium reveals the dynamical properties of that medium and presently is a field of high interest in view of jet quenching of high energy partons; both light [12–28] and heavy quarks [16,29–51]. Naively, one imagines that the amount of quenching for heavy flavours jet should be smaller than that of light flavours due to the large mass of heavy quarks. However, the single electron data at RHIC [8] exhibit almost a similar suppression for heavy flavoured hadrons compared to that for light hadrons.

The first attempt to estimate the radiative energy loss of heavy flavours in a QGP medium was made in Ref. [32] by using the Gunion–Bertsch (GB) formula [52] of gluon emission for light quark scattering and appropriately modifying the relevant kinematics for heavy quarks. Later the GB-like formula for heavy quarks was reconsidered in Ref. [33] by introducing the mass in the matrix element [34] but only within the small angle approximation. Due to this mass effect, a suppression, known as 'dead cone' effect, in the soft gluon emission off a heavy quark was predicted in comparison to that from a light quark. This resulted in a reduction of heavy quark energy loss induced by the medium [33], which is limited only to the forward direction. However, such a gluon radiation spectrum with a dead cone factor, only applicable

to the forward direction, was also used in the literature [35,37] uniformly for the full range of the emission angle (*i.e.*, both forward and backward directions) of gluon to calculate the heavy quark energy loss in the medium. This can lead to an unphysical result at large angle radiation, as discussed as well as shown in Ref. [53]. Further attempts were also made in the literature [38,42,43,45,46,50,51] to improve the calculation of heavy quark energy loss with various ingredients as well as restrictions. In some cases the energy loss for charm quark was found to be different than the light quark. The subject of heavy quark energy loss is not yet a settled issue and requires more detailed analysis.

In a very recent work [53] the probability of gluon emission off a heavy quark has been generalised by relaxing some of the constraints, *e.g.*, the gluon emission angle and the scaled mass of the heavy quark with its energy, which were imposed in earlier calculations [33,34]. It resulted in a very compact and elegant expression for the gluon radiation spectrum off a heavy quark (*e.g.*, $Q q \to Q q g$) as [53]

$$\frac{dn_g}{d\eta \, dk_\perp^2} = \frac{C_A \alpha_s}{\pi} \frac{1}{k_\perp^2} \mathcal{D}, \tag{1}$$

where the transverse momentum of the emitted massless gluon is related to its energy by $k_\perp = \omega \sin\theta$, and the rapidity, $\eta = -\ln[\tan(\theta/2)]$, is related to the emission angle, and the generalised dead cone is given by

$$\mathcal{D} = \left(1 + \frac{M^2}{s} e^{2\eta}\right)^{-2} = \left(1 + \frac{M^2}{s \tan^2(\frac{\theta}{2})}\right)^{-2}. \tag{2}$$

Now, the Mandelstam variable s is given as $s = 2E^2 + 2E \times \sqrt{E^2 - M^2} - M^2$, with E and M, respectively, the energy and mass of the heavy quark. C_A is the Casimir factor for adjoint representation and α_s is the strong coupling constant. In the small angle limit, $\theta \ll \theta_0 (= M/E) \ll 1$, the dead cone in (2) reduces to that in Refs. [33,34] as $(1 + \theta_0^2/\theta^2)^{-2}$ whereas for massless case it becomes unity and (1) reduces to the GB formula [52]. The gluon spectrum for the process, $Q g \to Q gg$, can also be found in Ref. [53]. We also note that the gluon emission spectrum in (1) is obtained in Feynman gauge. The same result is also obtained using light-cone gauge.

In Fig. 1, a Monte Carlo simulation of the above suppression factor (2) (*i.e.*, the scaled gluon emission spectrum off a heavy quark with that of light quark) is displayed. It reveals a forward–backward asymmetry which encompasses the fact that the gluon emission off a heavy quark is as strong as that of light quark at the large angles (backward direction) whereas it is suppressed due to nonzero quark mass at the small angles (forward direction). However, if the energy of the heavy quark is large compared to its mass, the effect of dead cone diminishes, both heavy and light quarks are expected to lose energy almost similarly. This result can have important consequences for a better understanding of heavy flavour energy loss in the context of heavy-ion collisions at RHIC and LHC. In this Letter we intend to use the gluon radiation spectrum in Ref. [53] to obtain the heavy flavour energy loss and attempt to understand the suppression of heavy flavoured hadrons in heavy-ion collisions.

2. Radiative energy loss

Among the interactions that a charged particle undergoes, as it traverses a dense matter, inelastic (*i.e.* radiative) scattering is undoubtedly the most important and interesting one. A number of different energy loss models has also been formulated in the literature (for review see Refs. [27,28]). The basic differences among the different models are the various constraints (*e.g.*, kinematic cuts,
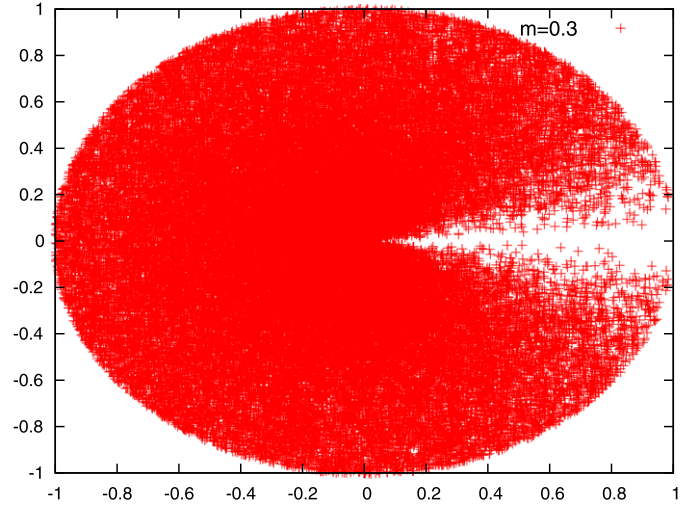
Fig. 1. (Colour online.) A Monte Carlo simulation for the suppression factor in (2) in the full domain of gluon emission angle, θ, off a heavy quark for the scaled mass $m = \frac{M}{\sqrt{s}} = 0.3$. This actually represents a two-dimensional view of the scaled gluon emission probability of a heavy quark with that of a light quark as given in (1). We consider the direction of propagation of a heavy quark is from left to right along the horizontal axis and collide with medium partons at the origin of a circle of unit radius. This simulation has been performed by throwing points at random directions within the full domain of θ but with a probabilistic weight $\mathcal{D}(\theta)$, which would then correspond to a point randomly on the selected θ-line as a '*red plus*' inside the circle of unit radius. The shade with red pluses represents the soft gluon emission zone whereas the conical white zone in the forward direction indicates a dead cone for gluon emission due to the mass effect.

large angle radiation, etc.) implemented to make the calculations manageable. In this section we define the rate of radiative energy loss of a parton with energy E, due to inelastic scatterings with the medium partons in a very canonical way as

$$\frac{dE}{dx} = \frac{\langle \omega \rangle}{\langle \lambda \rangle}, \tag{3}$$

where $\langle \omega \rangle$ and $\langle \lambda \rangle$ are the mean energy of emitted gluons and the mean free path of the traversing quark, respectively.

Among the set of variables $[k_\perp, \eta, \omega]$ in (1) any two together are sufficient to completely describe an emitted gluon. For convenience we now change the variable duo from $[k_\perp, \eta]$ to $[\omega, \eta]$ as

$$\frac{dn_g}{d\eta \, dk_\perp} \Rightarrow \frac{dn_g}{d\eta \, d\omega}. \tag{4}$$

It is now easy to find mean energy of the emitted soft gluons from the spectrum as

$$\langle \omega \rangle = \left(\int \frac{dn_g}{d\eta \, d\omega} \omega \, d\eta \, d\omega\right) \Big/ \left(\int \frac{dn_g}{d\eta \, d\omega} d\eta \, d\omega\right)$$
$$= \left(\int d\omega \int \mathcal{D} \, d\eta\right) \Big/ \left(\int \frac{1}{\omega} d\omega \int \mathcal{D} \, d\eta\right). \tag{5}$$

Other important quantity in (3) is the mean free path $\langle \lambda \rangle$, which is the average distance covered by the traversing quark between two successive collision, *followed by a soft gluon radiation*. The magnitude of mean free path depends on the characteristics of the system in which the energetic particle is traversing, and it is defined as

$$\langle \lambda \rangle = 1/(\sigma_{2 \to 3} \rho_{\text{qgp}}), \tag{6}$$

where $\sigma_{2 \to 3} \rho_{\text{qgp}} = \rho_q \sigma_{Q q(\bar{q}) \to Q q(\bar{q}) g} + \rho_g \sigma_{Q g \to Q gg}$, $\sigma_{2 \to 3}$ is the cross section of relevant $2 \to 3$ processes and ρ_{qgp} is the density of QGP medium which acts as a background containing target partons, for the high energetic projectile quark. We also note

that the Landau–Pomeranchuk–Migdal (LPM) interference correction may be marginal, which we would estimate below based on the formation time of the emitted gluon along with the kinematical restrictions. Now, we recall the total cross section for $2 \to 3$ processes as given in Ref. [54] as

$$
\begin{aligned}
\sigma_{2 \to 3} &= 2 C_A \alpha_s^3 \int \frac{1}{(q_\perp^2)^2} dq_\perp^2 \int \frac{1}{k_\perp^2} dk_\perp^2 \int \mathcal{D} \, d\eta \\
&= 4 C_A \alpha_s^3 \int \frac{1}{(q_\perp^2)^2} dq_\perp^2 \int \frac{1}{\omega} d\omega \int \mathcal{D} \, d\eta,
\end{aligned}
\tag{7}
$$

where $q_\perp$ is the transverse momentum of the exchanged gluon. Combining (6) and (7) the energy loss in (3) can be written as

$$
\begin{aligned}
\frac{dE}{dx} &= 12 \, \alpha_s^3 \rho_{\mathrm{qgp}} \int\limits_{q_\perp^2 |_{min}}^{q_\perp^2 |_{max}} \frac{1}{(q_\perp^2)^2} dq_\perp^2 \\
&\quad \times \int\limits_{\omega_{min}}^{\omega_{max}} d\omega \, 2 \int\limits_{\eta_{min}}^{\eta_{max}} \mathcal{D} \, d\eta,
\end{aligned}
\tag{8}
$$

where a factor of 2 has been introduced in η integral to cover both upper and lower hemisphere. We note that for $\mathcal{D} = 1$, (8) becomes equivalent to the massless case.

At this point it is important to note that the hierarchy employed in obtaining (1) in Ref. [53] reads as

$$
\sqrt{s}, E \gg \sqrt{|t|} \sim q_\perp \gg \omega > k_\perp \gg m_D,
\tag{9}
$$

where s, u, t are the usual Mandelstam variables and m_D is the Debye screening mass of the thermal gluons. Based on the above hierarchy we obtain the kinematic cuts explicitly on energy–momentum constraints and large angle radiation. The infrared cut-off has been used as

$$
q_\perp^2 |_{min} \simeq \omega_{min}^2 \simeq k_\perp^2 |_{min} \simeq m_D^2 = 4\pi \alpha_s T^2.
\tag{10}
$$

For ultraviolet cut-off on intermediate gluons, we have used [32],

$$
\begin{aligned}
q_\perp^2 |_{max} &= \frac{3}{2} ET - \frac{M^2}{4} \\
&\quad + \frac{M^4}{48 ET \beta_0} \log\left[\frac{M^2 + 6ET(1 + \beta_0)}{M^2 + 6ET(1 - \beta_0)} \right],
\end{aligned}
\tag{11}
$$

where $\beta_0 = (1 - M^2/E^2)^{1/2}$ and T is temperature of thermal background. The ultraviolet cut-off on energy for the emitted soft gluon has been taken as average momentum of the intermediate gluon line as [55]

$$
\omega_{max}^2 \simeq \langle q_\perp^2 \rangle.
\tag{12}
$$

Now, the relation between ω and $k_\perp$, $\omega = k_\perp \cosh \eta$, can be used to obtain bound on η from top, which eventually excludes all collinear singularities for massless case. Finite cut on ω and $k_\perp$ then leads to an inequality,

$$
\cosh \eta > \omega_{max}/k_\perp |_{min},
\tag{13}
$$

from which one can easily obtain the bound on η as

$$
|\eta| < \log\left(\frac{\sqrt{\langle q_\perp^2 \rangle}}{m_D} + \sqrt{\frac{\langle q_\perp^2 \rangle}{m_D^2} - 1} \right).
\tag{14}
$$

We are now in position to discuss the LPM effect which is usually included through a step function $\theta(\tau_i - \tau_f)$ while evaluating the spectrum of the radiated gluon. It basically implies that the

formation time of the gluon, $\tau_f = \langle \omega \rangle / \langle k_\perp^2 \rangle$ must be smaller than the interaction time $\tau_i \sim \Lambda_{\mathrm{QCD}}^{-1} = 0.49/T_C$. This on the other hand imposes a restriction on the phase space of the emitted gluon as $\langle \omega \rangle > 2\Lambda_{\mathrm{QCD}} \approx 4T_C \sim gT \sim \mu_D$, provided $\alpha_s \sim 0.3$, $T_C \sim 170$ MeV and the temperature of the plasma, $T \sim 350$ MeV. Thus, the hierarchy in Eq. (9) excludes the modification of the radiative energy loss due to the LPM interference correction through the infrared regulator, μ_D. Therefore, the present formalism becomes akin to the Bethe–Heitler approximation, in which the scattering centres are well separated and the intensity of the induced radiation from different scatterings is additive.

Now, it is very straightforward to obtain the radiative energy-loss through the inelastic processes, viz., $Q q(\bar{q}) \to Q q(\bar{q}) g$ and $Q g \to Q g g$ [53], for a heavy quark from (8), which reads as

$$
\begin{aligned}
\frac{dE}{dx} &= 24 \alpha_s^3 \left(\rho_q + \frac{9}{4} \rho_g \right) \frac{1}{\mu_g} (1 - \beta_1) \\
&\quad \times \left(\frac{1}{\sqrt{(1 - \beta_1)}} \left[\log(\beta_1)^{-1} \right]^{1/2} - 1 \right) \mathcal{F}(\delta),
\end{aligned}
\tag{15}
$$

where

$$
\begin{aligned}
\mathcal{F}(\delta) &= 2\delta - \frac{1}{2} \log\left(\frac{1 + M^2 e^{2\delta}/s}{1 + M^2 e^{-2\delta}/s} \right) \\
&\quad - \frac{M^2 \cosh \delta/s}{1 + 2M^2 \cosh \delta/s + M^4/s^2},
\end{aligned}
$$

$$
\delta = \frac{1}{2} \log\left[\frac{\log \beta_1^{-1}}{(1 - \beta_1)} \left(1 + \sqrt{1 - \frac{(1 - \beta_1)^{\frac{1}{2}}}{[\log \beta_1^{-1}]^{\frac{1}{2}}}} \right)^2 \right],
$$

$$
s = E^2 (1 + \beta_0)^2, \qquad \beta_1 = \frac{g^2}{\mathcal{C}} \frac{T}{E},
$$

$$
\begin{aligned}
\mathcal{C} &= \frac{3}{2} - \frac{M^2}{4ET} \\
&\quad + \frac{M^4}{48 E^2 T^2 \beta_0} \log\left[\frac{M^2 + 6ET(1 + \beta_0)}{M^2 + 6ET(1 - \beta_0)} \right].
\end{aligned}
\tag{16}
$$

Eq. (15) together with (16) represents radiative energy loss of an energetic quark in a canonical way within the framework of perturbative QCD along with kinematical restrictions for an energetic parton and medium interaction.

3. Heavy quark production in pp collisions

At leading order pQCD, heavy quarks in pp collisions are mainly produced by fusion of gluons ($gg \to Q\overline{Q}$) or light quarks ($q\bar{q} \to Q\overline{Q}$) [56]. The cross-section for the production of heavy quarks from pp collisions at leading order can be expressed as [56,57]:

$$
\begin{aligned}
\frac{d\sigma}{dy_1 \, dy_2 \, dp_T} &= 2x_1 x_2 p_T \sum_{ij} \left[f_i^{(1)}(x_1, Q^2) f_j^{(2)}(x_2, Q^2) \hat{\sigma}_{ij} \right. \\
&\quad \left. + f_j^{(1)}(x_1, Q^2) f_i^{(2)}(x_2, Q^2) \hat{\sigma}_{ij} \right]/(1 + \delta_{ij}),
\end{aligned}
\tag{17}
$$

where i and j are the interacting partons, $f_i^{(1)}$ and $f_j^{(2)}$ are the partonic structure functions and x_1 and x_2 are the fractional momenta of the interacting hadrons carried by the partons i and j. The short range subprocesses for the heavy quark production, $\hat{\sigma} = d\sigma/dt$ are defined as

$$
\frac{d\sigma}{dt} = \frac{1}{16\pi s^2} |\mathcal{M}|^2,
\tag{18}
$$

where $|\mathcal{M}|^2$ for the processes $gg \to Q\overline{Q}$ and $q\bar{q} \to Q\overline{Q}$ can be obtained from Ref. [56]. The running coupling constant α_s at leading order is

$$\alpha_s = \frac{12\pi}{(33 - 2N_f)\ln(Q^2/\Lambda^2)}, \tag{19}$$

where $N_f = 3$ is the number of active flavours and $\Lambda = \Lambda_{\text{QCD}}$. The p_T distribution of production of heavy quarks at leading order supplemented with a K-factor ≈ 2.5 is taken as the baseline for the calculation of the nuclear suppression factor, R_{AA} [44]. Effect of prefactor K is diluted during computation of nuclear modification factor due to its identical effects on both initial and final distributions profiles. Furthermore, the K-factor, if equal for c and b quarks, has not only a diluted effect but can actually be neglected in the ratios. The shadowing effect is considered using EKS98 parameterisation [58] for nucleon structure functions and here we use the CTEQ4M [59] set for nucleon structure function. We use Peterson fragmentation function [60] with parameter $\epsilon_c = 0.06$ and $\epsilon_b = 0.006$ for fragmentation of c quarks into D-mesons and b quarks into B-mesons, respectively.

All the calculations are done assuming the mean intrinsic transverse momentum of the partons to be zero.

4. Initial conditions and evolution of the medium

As the heavy quarks are expected to lose most of their energy during the earliest time after the formation of QGP, we can safely neglect the transverse expansion of the plasma while discussing the heavy quark energy loss.

We consider a heavy quark, which is being produced at a point (r, Φ) in a central collision and moves at an angle ϕ with respect to $\hat{r}$ in the transverse plane. If R be the radius of the colliding nuclei, the path length covered by the heavy quark would vary from 0 to $2R$, before it exits the QGP. The distance covered by the heavy quark inside the plasma in a central collision, L, is given by [61]:

$$L(\phi, r) = \sqrt{R^2 - r^2 \sin^2 \phi} - r \cos \phi. \tag{20}$$

We can estimate the average distance travelled by the heavy quarks in the plasma as

$$\langle L \rangle = \frac{\int_0^R r\,dr \int_0^{2\pi} L(\phi, r) T_{AA}(r, b = 0)\,d\phi}{\int_0^R r\,dr \int_0^{2\pi} T_{AA}(r, b = 0)\,d\phi}, \tag{21}$$

where $T_{AA}(r, b = 0)$ is the nuclear overlap function. We estimate $\langle L \rangle$ as 5.78 fm for central Au + Au collisions and 6.14 fm for central Pb + Pb collisions.

The temperature of the plasma at a time τ, assuming a chemically equilibrated plasma can be expressed as [43]

$$T(\tau) = \left(\frac{\pi^2}{1.202} \frac{\rho(\tau)}{(9N_f + 16)} \right)^{\frac{1}{3}}, \tag{22}$$

where the gluon density at time τ is given by [43]:

$$\rho_g(\tau) = \frac{1}{\pi R^2 \tau} \frac{dN_g}{dy}. \tag{23}$$

Here we consider only the gluon density as the heavy quarks lose most of their energy in interaction with gluons. We also add that the gluon multiplicity is taken as 3/2 times the number of charged hadrons and the initial temperature is obtained using (22), assuming an initial time.
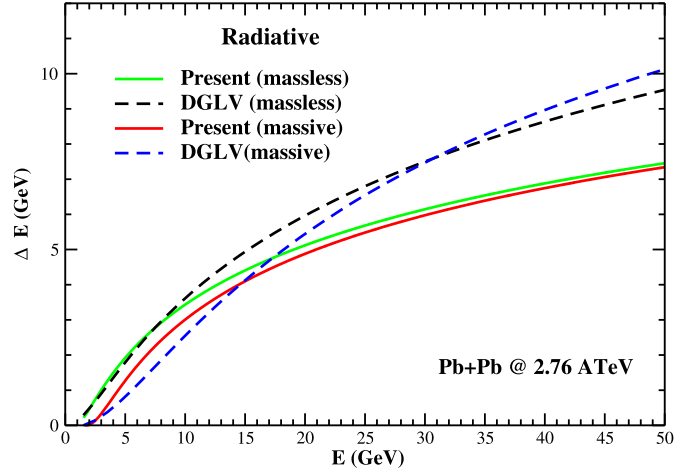
Fig. 2. (Colour online.) Comparison of average energy loss for light quark and charm quark with mass 1.5 GeV in a deconfined quark matter produced in Pb–Pb collision at 2.76 ATeV in the present and DGLV [43] formalisms. For both cases the characteristics of the deconfined matter are treated in the same footing, i.e., the strong coupling $\alpha_s = 0.3$ and the average path length, $\langle L \rangle \approx 6.14$ fm, traversed by an energetic quark in a deconfined medium produced in such collisions.

We take $(\frac{dN_g}{dy}) \approx 1125$ for Au + Au collisions at 200 AGeV [62], ≈ 2855 for Pb + Pb collisions at 2.76 ATeV [10] and ≈ 4050 for Pb + Pb collisions at 5.5 ATeV [63]. We assume that the heavy quark having rapidity in the central region moves along the fluid of identical rapidity. This kind of approximation has been used earlier in literature [64–66].

We calculate the initial temperature of QGP formation T_0 at 200 ATeV as 400 MeV, at 2.76 ATeV as 525 MeV and at 5.5 ATeV as 590 MeV, assuming the initial time of QGP formation as $\tau_0 = 0.2$ fm/c. The critical temperature T_c for the existence of QGP is taken as ≈ 170 MeV. The time, by which the plasma will reach the critical temperature, τ_c is found to be ≈ 2.627 fm/c at 200 AGeV, 5.9038 fm/c at 2.76 ATeV and 8.375 fm/c at 5.5 ATeV, assuming Bjorken's cooling law, $T^3\tau = \text{constant}$.
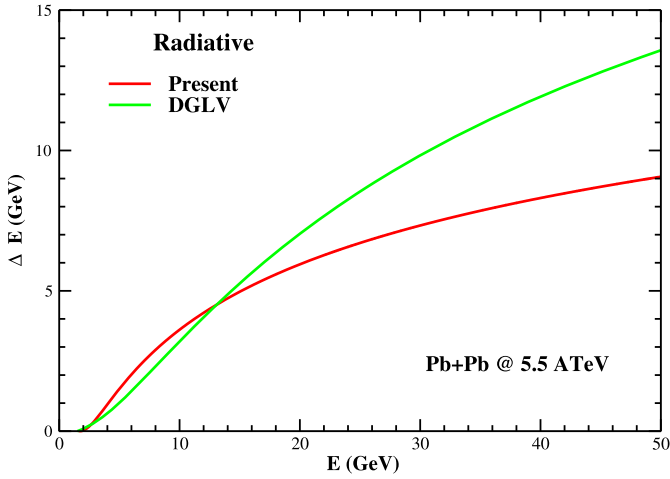
The average path length of the heavy quark inside the plasma is calculated as follows. The velocity v_T of a heavy quark can be expressed as $p_\perp/m_T$, where m_T is the transverse mass. Thus, the heavy quark would cross the plasma in a time $\tau_L = \langle L \rangle / v_T$. Now, if $\tau_c \geqslant \tau_L$, the heavy quark would remain inside the QGP during the entire period, τ_0 to τ_L. But if $\tau_c < \tau_L$, it would remain inside QGP only while covering the distance $v_T \times \tau_c$. Thus, we further approximate the expanding and cooling plasma with one at a temperature of T at $\tau = \langle L \rangle_{\text{eff}}/2$, where $\langle L \rangle_{\text{eff}} = \min[\langle L \rangle, v_T \times \tau_c]$ (see Ref. [43]).

5. Results and discussion

In Fig. 2 a comparison of average radiative energy loss of an energetic quark traversing in a deconfined quark matter produced in Pb–Pb collision at 2.76 ATeV in the present calculation with Djordjevic, Gyulassy, Levai and Vitev (DGLV) formalism in Refs. [22,43]. As can be seen both light and heavy quarks in the present formalism, within the gluon emission spectrum of $\mathcal{O}(\alpha_s)$ and $\mathcal{O}(1/k_\perp^2)$ as given in (1), lose energy in a similar fashion for $E \geqslant 10$ GeV since the effect of mass is small compared to the energy. However, it is slightly less than that of a light quark for $E \leqslant 10$ GeV, due to the dead cone suppression at small angles. In addition the results from the present calculation differ from that of DGLV [22,43] one. These differences arise mainly because of the proper kinematic cuts for gluon emission as well as the method used to obtain energy loss. The various cuts in the present as well as in DGLV formalism are in
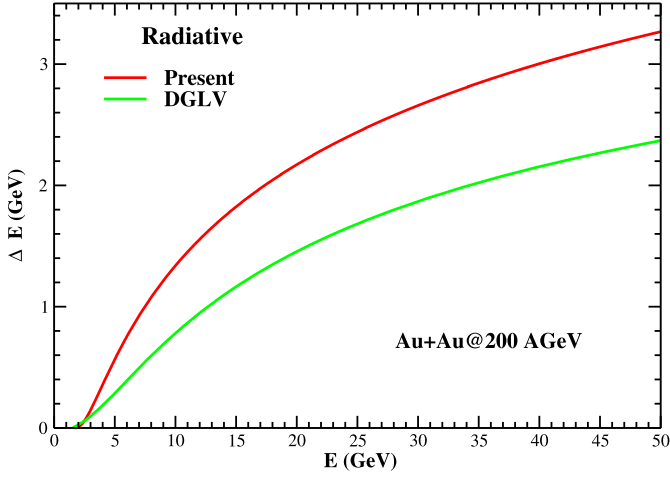
Fig. 3. (Colour online.) Same as Fig. 2 but only for charm quark in Au–Au collision at 200 AGeV with $\langle L \rangle = 5.78$ fm.
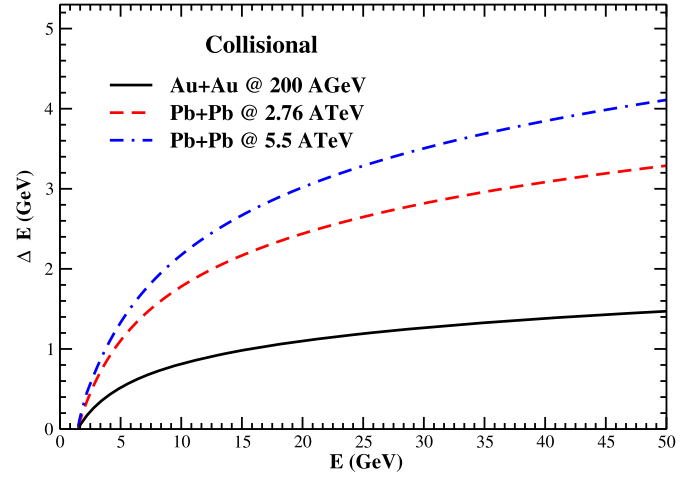

Fig. 5. (Colour online.) Collisional energy loss of charm quark [31] in Pb–Pb collision at 2.76 ATeV and 5.5 ATeV at LHC, and 200 AGeV at RHIC energies.

Fig. 4. (Colour online.) Same as Fig. 3 in Pb–Pb collision at 5.5 ATeV with $\langle L \rangle =$ 6.14 fm.

Fig. 6. (Colour online.) Nuclear modification factor R_{AA} for D-mesons with both collisional and radiative energy loss in Pb–Pb collision at 2.76 ATeV. The data are from Ref. [67] but only the systematic error bars are shown here.

close proximity except the gluon emission in DGLV is constrained only to the forward emission angles [27], $\theta \leqslant \pi/2$, whereas in the present calculation the full range of θ is taken care off through the variable η as shown in (13) and (14).
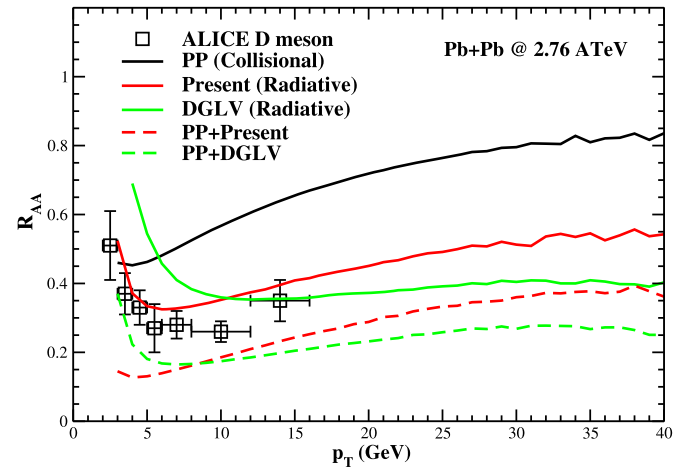
In Figs. 3 and 4 we have displayed average energy loss of a charm quark in a deconfined quark matter, respectively, at 200 AGeV Au–Au collision at RHIC and 5.5 ATeV Pb–Pb collision at LHC. We find that at RHIC energies the average energy loss of a charm quark in our formalism is higher than that of the DGLV formalism for the considered energy range, $(0 < E < 50)$ GeV, of the charm quark. On the other hand Fig. 4 is qualitatively similar to Fig. 2 in terms of comparison of two formalism for heavy quark. As seen the average energy loss of charm quark is larger in the present formalism only in the domain, $(0 < E < 15)$ GeV, of the charm quark and beyond which it is less compared to the DGLV formalism. The difference, in fact, increases as energy of the quark increases.

In Fig. 5 we display a comparison of collisional energy loss of charm quark as calculated by Peigne and Peshier (PP) in Ref. [31] for RHIC and LHC energies. As seen the collisional energy loss increases with the increase in centre of mass energy of the colliding ions.

In Fig. 6 the nuclear suppression factor, R_{AA}, for D-meson is displayed considering both radiative and collisional energy loss and

compared with the ALICE data [67] at 2.76 ATeV. As can be seen the differences in radiative energy loss between the present and DGLV formalism discussed in Fig. 2 for 2.76 ATeV in Pb–Pb collisions is clearly reflected in Fig. 6. For the present calculation it is manifested in gradual increase of R_{AA} of D-meson [67] for transverse momentum, $p_\perp > 5$ GeV whereas in DGLV case it remains almost constant. The suppression factor obtained in the present formalism with radiative energy loss is in close agreement with the most recent data from ALICE collaboration at 2.76 ATeV [67]. On the other hand the inclusion of the collision contribution is found to suppress R_{AA} further in both cases. As found the data suggest that the collisional contribution may be small. Nonetheless, more data in the high-$p_\perp$ domain are necessary to know the actual trend of the energy loss of charm quark and will finally constrain the various energy loss and jet quenching model in the literature. We also expect a similar rise in light hadrons for high $p_\perp$ since both light and heavy quarks lose energy in a similar fashion as shown in Fig. 2. However, we note that the ALICE data on R_{AA} for inclusive charge hadrons [68] at 2.76 ATeV in Pb–Pb collision has also shown a similar increasing trend as $p_\perp$ increases. It is natural to believe that such data are completely dominated by the contribution from light hadrons. For completeness, we also display R_{AA} for LHC energy at 5.5 ATeV in Fig. 7.
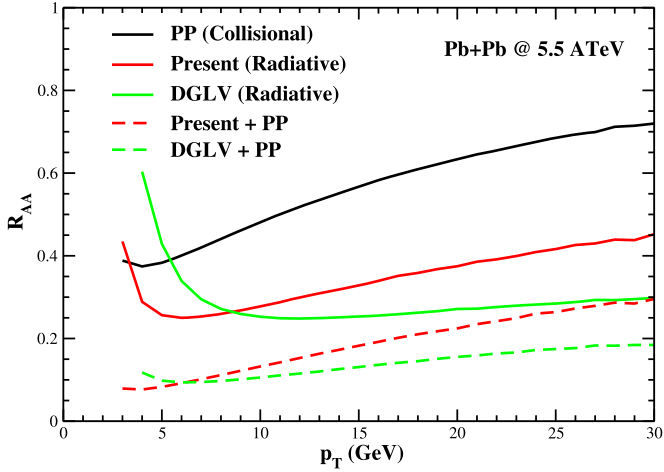
Fig. 7. (Colour online.) Nuclear modification factor, R_{AA}, for D-mesons in Pb–Pb collision at 5.5 ATeV.
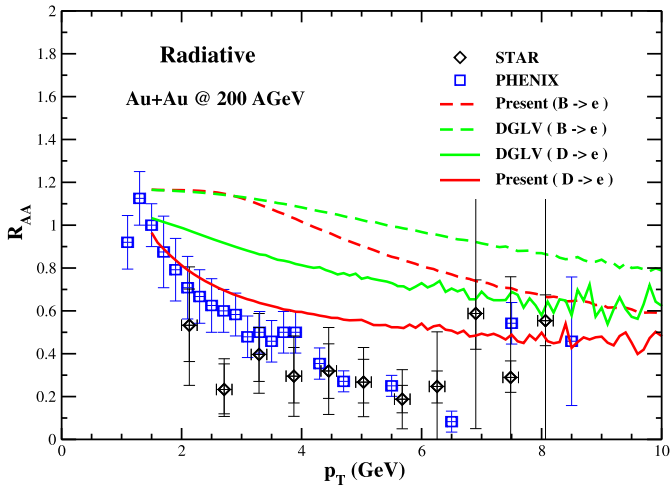

Fig. 8. (Colour online.) R_{AA} with only radiative energy loss for non-photonic single electron from the decay of individual D- and B-mesons in Au–Au collision at 200 AGeV. The data are from Ref. [8]. Both systematic and statistical error bars are shown for STAR data whereas only systematic error bars are displayed for PHENIX data.

Fig. 9. (Colour online.) R_{AA} with collisional and radiative energy-loss for non-photonic single electron from the combined decay of both D- and B-mesons in Au–Au collision at 200 AGeV. The data are from Ref. [8]. Both systematic and statistical error bars are shown for STAR data whereas only systematic error bars are displayed for PHENIX data.

Fig. 10. (Colour online.) Same as Fig. 9 in Pb–Pb collision at 2.76 ATeV.

In Fig. 8 the nuclear suppression factors for individual decay of D- and B-mesons to non-photonic single electron are displayed considering only the radiative energy loss for RHIC energy at 200 AGeV. As expected the contribution from the B decay is small compared to that of D decay. In Fig. 9 the total contribution of single electron from D and B decay is shown considering both radiative and collisional energy loss. It is found that the contributions of the collisional energy loss are important at RHIC energy. We also compare our results with that of DGLV. In Fig. 10, we give prediction for single electron result for LHC energy at 2.76 ATeV.

6. Summary

We obtain the radiative energy loss of a heavy quark akin to the Bethe–Heitler approximation by considering the most generalised gluon emission multiplicity expression derived very recently. This suggests that both energetic heavy and light quarks lose energy due to gluon emission almost similarly and the mass plays a role only when the energy of the quark is of the order of it. The hierarchy used for simplifying the matrix element as well as
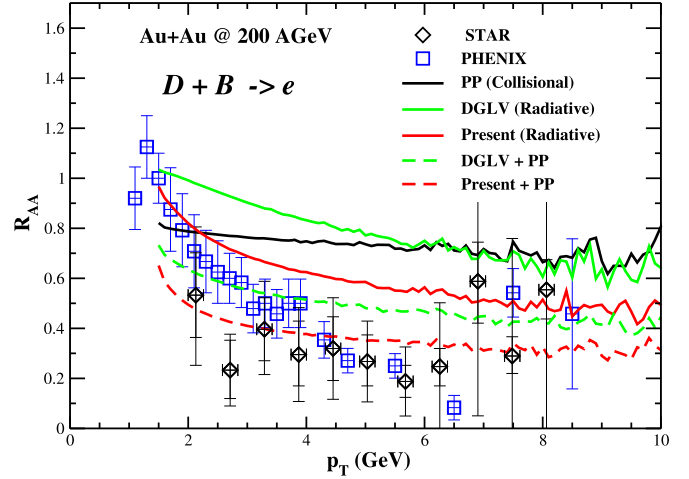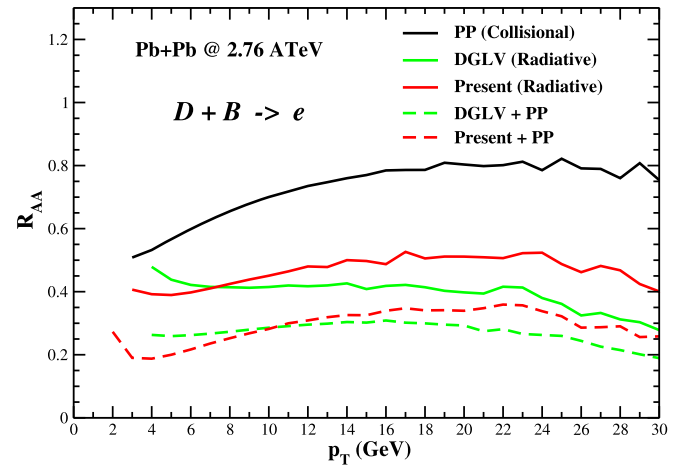
for obtaining the gluon radiation spectrum imposes a restriction on the phase space of the emitted gluon in which the formation time is estimated to be less than the interaction time. This suggests that the LPM interference correction may be marginal. Further, we compare our results with the DGLV formalism and it is found to differ significantly. To compute the nuclear suppression factor for D-meson we consider both radiative and collision energy loss along with longitudinal expansion of the medium. The nuclear modification factor for D-meson with radiative energy loss obtained in the present formalism has an increasing trend at high $p_\perp$ and found to agree closely with the very recent data from ALICE collaboration at 2.76 ATeV. When the collisional counter part is added independently, the further suppression is obtained in the nuclear modification factor. On the other hand the non-photonic single electron data at 200 AGeV RHIC energy requires contributions from collisional energy loss as well from B decay. However, it is necessary to obtain both radiative and collisional energy loss from the same formalism to minimise the various uncertainties, which is indeed a difficult task. Moreover, data at high-$p_\perp$ region with improved statistics are required to remove prejudice on different energy loss and jet quenching models.

Acknowledgements

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