Competition of Heavy Quark Radiative and Collisional Energy Loss in Deconfined Matter

P.B. Gossiaux, J. Aichelin, T. Gousset and V. Guiho
SUBATECH, Université de Nantes, EMN, IN2P3/CNRS
4 rue Alfred Kastler, 44307 Nantes cedex 3, France
E-mail: pol.gossiaux@subatech.in2p3.fr

Abstract. We extend our recently advanced model on collisional energy loss of heavy quarks in a quark gluon plasma (QGP) by including radiative energy loss. We discuss the approach and present first preliminary results, including a comparison of the role of both types of energy loss for experimental data. We draw the conclusion that the present nuclear modification factor $R_{AA}$ data on non-photonic single electrons does not permit to “select” between this two types of energy-loss mechanisms.

1. Introduction

One of the main objectives of the studies of ultrarelativistic heavy ion collisions in heavy ion collisions is the exploration of the properties of a quark gluon plasma (QGP), which is most probably formed in these collision. To achieve this objective is all but easy because this plasma, if created at all, will live only for a very short time before it suffers a phase transition (or more precisely a cross over) towards the hadronic phase. After this transition the hadrons continue to interact weakening the possible information on the plasma phase.

Indeed, the multiplicity of light particles, those which contain only up, down and strange quarks, is well reproduced in statistical models indicating that these multiplicities are established at temperatures close to the transition temperature between the plasma and the hadronic phase and carry therefore little or no information on the plasma interior. The multiplicity of the observed resonances presents evidence that, after the multiplicity is established, hadronic interactions modify considerably the momenta of the hadrons. In view of the fact that many of the involved elementary cross section are little known or unknown it is difficult if not impossible to asses the spectra of the hadrons at the moment of their formation, a prerequisite for an analysis of the plasma properties with help of these particles.

The observation that the bulk part of the observed particles is not sensitive to the QGP properties limits severely the possibilities to study the interior of the QGP and it is all but easy to identify those observables which may serve for this purpose. Those which have been advanced include the high $p_T$ hadrons which originate from jets which
do not come to equilibrium with the plasma as well as the $p_T$ and $v_2$ distribution of heavy mesons which contain either a $c$ or a $b$ quark.

Heavy quarks are produced in hard binary initial collisions between the incoming nucleons. Their production cross sections are known from pp collisions and their spectra can be calculated (see f.i. [1]) in the framework of perturbative Quantum ChromoDynamics (pQCD), up to a systematic error band due to the choice of various scales. In heavy ion collisions, these spectra can be taken bona fide as the initial transverse momentum distribution of heavy quarks. Comparing this distribution with that measured allows to define the nuclear modification factor $R_{AA} = (d\sigma_{AA}/d^2p_T) / (N_c d\sigma_{pp}/d^2p_T)$, where $N_c$ is the number of initial binary collisions between projectile and target. The deviation of $R_{AA}$ from unity essentially reflects the interaction of the heavy quark with the plasma because the hadronic cross sections of heavy mesons are small. Present data shows that “high” $p_T$ heavy quarks ($p_T \gtrsim m$) do not come to thermal equilibrium with the QGP; therefore $R_{AA}$ contains the information on the interaction of heavy quarks while they traverses the plasma. In addition, the distribution of heavy quarks at the moment of their creation is isotropic in azimuthal direction, therefore the elliptic flow $v_2 = \langle \cos 2(\phi - \phi_R) \rangle$, where $\phi (\phi_R)$ is the azimuthal angle of the emitted particle (reaction plane) is 0. The observed finite $v_2$ value of the observed heavy meson can only originate from interactions between light QGP constituents and the heavy quarks. The simultaneous description of $R_{AA}$ and $v_2$ – presently the only observables for which data exist – gives then the hope to better understand the interaction of those heavy quarks with the QGP and thus its profound nature. Unfortunately the experimental results depend not only on the elementary interaction but also on the description of the expansion of the QGP. Therefore the ultimate aim is to control the expansion by results on the light meson sector. This has not been achieved yet and therefore it is difficult to asses the influence of the expansion on the observables.

Coming to facts, the $R_{AA}$ of $\approx 0.2-0.3$ value observed for large $p_T$ ($p_T \sim 4-8$ GeV) non-photonic single-electrons (n.p.s.e.) [2, 3], of the same order of the pionic $R_{AA}$, is associated with 2 related puzzles, hereafter referred as “single electron puzzle” (s.e.p.):

- **level-1 s.e.p.**: One cannot understand such a small value within the framework of pQCD;
- **level-2 s.e.p.**: Even if one introduces some degree of freedom in the pQCD inspired models, i.e. some free parameter as f.i. $\hat{q}$ or $dN_g/dy$, it is not possible to reproduce both $R_{AA}$ of n.p.s.e. and pions within the same global framework.

To substantiate level-1 s.e.p., we notice that the elastic cross section and hence the collisional energy loss introduced in early approaches [4] has to be multiplied [5] by an artificial $K$ factor of the order of 10 to match the experimental data. These

‡ Although one observes similar fact for light-quark leading hadrons, it is not clear whether the dominant basic microscopic processes are of similar nature for these two probes, due to the heavy quark finite-mass and the flavor exchange mechanism.
early calculations, however, used ad hoc assumptions on the (fixed) coupling constant $\alpha_s$ and on the infrared regulator $\mu$. This last drawback was partly cured in [6], where an HTL regularization is used for the collisional energy loss, but where the spatial diffusion coefficient $D_s$ one should privilege to reproduce the experimental data is still found much smaller than its pQCD value§. Of course, collisional energy loss is only one source of the energy loss. For light particle the radiative energy loss is even more important. Its importance for heavy quarks has been addressed in numerous publications ([7, 8, 9, 10, 11, 12, 13] and references therein, to cite but a few) and is however still debated. In [10, 11] and in previous works of the authors, the heavy quark quenching was predicted to be undeniably larger than later observed in experiment. This work was pursued in [13], where both radiative and collisional processes are implemented, as well as path length fluctuations (not considered before by the authors)∥. With the parameters chosen – fixed $\alpha_s = 0.3$ and $dN_g/dy = 1000$ – and both additional mechanisms [13] could maintain the good agreement found for the pion $R_{AA}$ [14] in the GLV approach with pure radiative energy loss and fixed path length, while the n.p.s.e. $R_{AA}$ was reduced w.r.t. [10, 11] but still found to exceed the experimental values...a conclusion that supports the so-called level-2 s.e.p.. In [9], Armesto et al. have extended the Wiedemann-Salgado path-integral formalism for radiative energy loss [15] to the case of finite-mass parton. Later on, they applied their formalism [12] to study the n.p.s.e. quenching in the case of AA collisions, under the same assumptions used in a previous work dedicated to the light hadrons [16] (from which they have extracted $\hat{q} \in [4, 14]$ GeV$^2$/fm). Although Armesto et al. argue that “claims of inconsistency between theory and experiment are not supported”, they obtain a n.p.s.e. $R_{AA} \in [0.4, 0.5]$ for $p_T = 5–8$ GeV/$c$, i.e. not enough quenching from their theoretical “prediction” (once calibrated on light hadron quenching).

These puzzles have casted doubt, whether pQCD is the right framework to describe these interactions and quenching at $p_T = 5–20$ GeV/$c$, and alternative approaches were proposed, based on the advocated existence of non-perturbative degrees of freedom in the QGP [17] or on the AdS/CFT conjecture [18, 19]. In a more conservative way, Peshier [20] has argued that the collisional energy loss could be vastly increased if a running coupling is taken to evaluate the elastic cross section. All alternative approaches have the common property to hold larger stopping power for heavy quarks than pQCD; the last proposal has however the appealing feature that collisional energy loss presents a weaker dependence on the parton mass (than the radiative one), what seems to be favored by the level-2 s.e.p.. Inspired by [20], we have recently advanced an approach for the collisional energy loss of heavy quarks in the QGP [21, 22, 23] in which a) $\mu$ has been fixed by the demand that more realistic calculations using the hard thermal loop approach give the same energy loss as the Born type pQCD calculation and b) the coupling constant

§ One should moreover notice that no $b$-quark contribution to n.p.s.e. was considered in [5, 6], which is known to increase the $R_{AA}$.
∥ but however with a rather simple implementation of medium evolution (no radial expansion, medium considered at a fixed Bjorken time)
is running and constrained by the sum rule advanced by Dokshitzer [24]. Both these improvements increased the cross section especially for small momentum transfers and reduced therefore the necessary $K$ (used in our previous work [5]) factor to 2.

Faced with the theoretical uncertainties affecting $R_{AA}$, a current has developed [25] which suggests that the “unknown” parameters that basically encode the interaction frequency with the medium (like $\hat{q}$ and $dN_g/dy$) should be constrained using experimental data. It was even argued by Lacey et al. [26] that a global analysis of pions and n.p.s.e. $R_{AA}$ was “consistent with jet quenching dominated by radiative energy loss for both heavy and light partons”. However, scenario’s relying on a finite contribution of collisional energy loss were not considered in [26]. One of the goals of this contribution is to complement the work of Lacey et al. by considering models containing both types of energy loss and the “competition” between them in order to reproduce the data.

To study this question, we have incorporated a radiative energy-loss mechanism in our model [21, 22, 23], not as sophisticated as those cited earlier but which – to our belief – incorporates the dominant aspects in the case of heavy quarks. We will describe the model in section 2 and we report and discuss our first results in section 3.

2. The Model

Extending our approach to radiative energy loss, we will first focus on gluon emission in single high energy processes, as described in QCD by the 5 matrix elements depicted in fig. 1. The commutation relation $T^g T^h = T^h T^g - i f_{ghc} T^c$ allows to regroup the 5 matrix elements.
elements into 3 combinations, each of them being independently gauge invariant:

\[ iM_{h,q,l,q}^{QED} = C_{I/Q}i(M_{1/3} + M_{2/4}) \quad \text{and} \quad iM_{V}^{QCD} = C_{V}i(M_{1} + M_{3} + M_{5}), \quad (1) \]

where h.q. (l.q.) marks the emission of the gluon from the heavy (light) quark line. \( C_{I}, \ C_{Q}, \ \text{and} \ C_{V} \) are the associated color algebra matrix elements. The matrix elements labeled as “QED” are the bremsstrahlung diagrams already observed in Quantum Electrodynamics (QED), whereas that labeled “QCD” is the diagram specific to QCD. The QCD diagram is the main object of interest here, as it can be shown to dominate the energy loss of heavy quark.

We evaluate the matrix elements in the so-called “scalar QCD” approximation, which shows to be appropriate at small or moderate gluon fractional momentum \( x \). The matrix elements in scalar QCD (see ref.[27]) are given by

\[
\begin{align*}
 iM_{1}^{SQCD} & = C_{I}(ig)^{3}(p_{b} + p_{3})^{\mu} \frac{\not{p}_{3} - \not{p}_{b}}{(p_{3} - p_{b})^{2}} D_{\mu\nu}^{5}[p_{3} - p_{b}] \left[ (p_{a} + p_{1} - k)^{\nu}(2p_{a} - k) \cdot \epsilon - (p_{a} - k)^{\nu} \epsilon \right] \\
 iM_{5}^{SQCD} & = C_{V}(ig)^{3} D_{\mu\nu}^{5}(p_{3} - p_{b}) D_{\nu\sigma}^{\prime}[p_{1} - p_{a}] \left[ g_{\mu\nu}(p_{a} - p_{1} + p_{3} - p_{b}) + g_{\nu\sigma}(p_{1} - p_{a} - k) + g_{\sigma\mu}(p_{b} - p_{3} + k) \right] \epsilon^{\sigma} \times \frac{(p_{3} + p_{b})^{\mu}(p_{a} + p_{1})^{\nu}}{(p_{3} - p_{b})^{2}(p_{1} - p_{a})^{2}}. (2)
\end{align*}
\]

\( M_{3} \) is obtained by replacing \( p_{a} \rightarrow p_{b} \) and \( p_{1} \rightarrow p_{3} \) in \( M_{1} \). Using light-cone gauge and keeping only the leading term in \( \sqrt{s} \) we see that the square of the matrix element factorizes

\[ |M|^{2} = |M_{el}(s, t)|^{2} P_{g}(m, t, \vec{k}_{T}, x) \quad (3) \]

with \( |M_{el}(s, t)|^{2} = g^{4}\alpha_{s}^{2} \frac{m^{2}}{\tilde{T}^{2}} \) being the matrix element squared for the elastic cross section in a fixed \( \alpha_{s} \) Coulomb-like interaction. \( P_{g}(m, t, s, \vec{k}_{T}) \) describes the distribution function of the produced gluons. To discuss the physics we adopt the following light-cone coordinates \( \{p^{+}, p^{-}, \vec{p} \} \), with scalar product defined as \( p_{a}p_{b} = \frac{p_{a}^{2}p_{b}^{2} + p_{a}^{+}p_{b}^{+}}{2} - p_{a}p_{b} \):

\[
p_{a} = \{ \sqrt{t - m^{2}}, \frac{m^{2}}{\sqrt{s - m^{2}}} \}, \quad p_{b} = \{ 0, \sqrt{s - m^{2}}, 0 \} \quad (4)
\]

for the entrance channel and

\[ k = \{ k^{+}, 0, \vec{k}_{T} \}, \quad p_{1} = p_{a} + q - k = \{ p_{1}^{+}, \frac{m_{T}^{2}}{p_{1}^{+}}, \vec{k}_{T} - \vec{q}_{T} \}, \quad p_{3} = p_{b} - q = \{ \frac{q^{2}}{p_{3}^{2}}, p_{3}^{-}, -\vec{q}_{T} \} \quad (5)\]

for the exit channel, with \( m_{T}^{2} = m^{2} + (\vec{k}_{T} - \vec{q}_{T})^{2}, \quad p_{3}^{-} \approx \frac{\vec{p}_{3}^{-} - \vec{p}_{3}^{+} - x \vec{p}_{a}^{+} - \frac{q^{2}}{p_{3}^{2}}}{m_{T}^{2}} \)

where \( \bar{x} \triangleq 1 - x \). In this coordinate system, the invariant momentum transfer is \( t = q^{2} \approx q_{T}^{2} \) while \( x \) is defined as \( x = k^{+}/p_{a}^{+} \) and represents the relative longitudinal momentum fraction of the gluon with respect to the incoming heavy quark: \( |M_{SQCD}|^{2} \)

has moreover a very simple form:

\[
|M_{SQCD}^{SQCD}|^{2} = g^{2} D_{QCD}^{4x^{2}}|M_{el}|^{2} \left( \frac{\vec{k}_{T}^{2}}{k_{T}^{2} + x^{2}m^{2}} - \frac{\vec{k}_{T} - \vec{q}_{T}}{(\vec{q}_{T} - \vec{k}_{T})^{2} + x^{2}m^{2}} \right)^{2} \quad (6)
\]

with the color factor \( D_{QCD} = C_{A} \ast C_{el}^{qq} = \frac{2}{3} \). The first term in the bracket describes the emission from the incoming heavy quark line, the second term the emission from
the gluon line. This shows that in light cone gauge and in this coordinate system in leading order of $\sqrt{s}$ the matrix element for the emission from the light quark do not contribute. In the case of massless quarks we recover the matrix elements of Gunion and Bertsch (GB) of [28]. From this factorized form, one deduces a similar expression for the radiative differential SQCD cross section when $s \gg m^2$:

$$
\frac{d\sigma^{\text{Q} \rightarrow \text{Q} \phi}}{d^2q_T} \approx \frac{d\sigma_{\text{el}}^{\text{Q} \rightarrow \text{Q} \phi}}{d^2q_T} P^{\text{SQCD}}(m, \vec{q}, k_T, x) \Theta \left( (m_T + q_T)^2 + \frac{x^2 k_T^2}{(m_T + q_T)^2} - x s \right),
$$

where

$$
\frac{d\sigma_{\text{el}}^{\text{Q} \rightarrow \text{Q} \phi}}{d^2q_T} = \frac{2}{9} \times \frac{4\alpha_s^2}{t^2},
$$

and where the Heaviside function traduces the phase-space boundary which translates into $k_T \leq k_{\text{t,max}}(\phi(\vec{q}, k_T))$, with $k_{\text{t,max}}^2 \approx (1 - x) s - (\sqrt{m^2 + q_T^2} + q_T)^2$ at small and moderate $x$. Apart from a small region located close to the phase space boundary, $P^{\text{SQCD}}$ admits the simple form:

$$
P^{\text{SQCD}}(m, \vec{q}, k_T, x) \approx C_A \alpha_s \langle \frac{\vec{k}_T}{k_T^2 + x^2 m^2} - \frac{\vec{k}_T - \vec{q}}{(\vec{k}_T - \vec{q})^2 + x^2 m^2} \rangle^2,
$$

that recovers the GB result at small $x$ when $m$ is taken to 0:

$$
P^{\text{SQCD}}(m = 0, \vec{q}, k_T, x \ll 1) = \frac{C_A \alpha_s \langle q_T^2 k_T^2 (\vec{k}_T - \vec{q})^2 \rangle}{\pi^2}.
$$

We now proceed along the following strategy: in our Monte Carlo numerical framework MC@nHQ (results presented in section 3), we have sampled the probability distribution associated to (7) using (9), the collisional cross section described in [21] instead of (8), and a rigorous implementation of the phase-space constrain. In the rest of this section, we will aim at establishing some synthetic relations useful for the physical interpretation; for this purpose, we will neglect the phase-space constrain in (7). The integral over $d\phi(\vec{k}_T, \vec{q})$ can thus be done analytically pretty easily by rewriting

$$
A = \frac{q_T^2 + 2x^2 m^2}{e((k_T - q_T)^2 + x^2 m^2)} - \frac{x^2 m^2}{(k_T - q_T)^2 + x^2 m^2},
$$

with $e \triangleq k_T^2 + x^2 m^2$. One gets

$$
P^{\text{SQCD}}(m, q_T, k_T, x) \triangleq \int d\phi P^{\text{SQCD}}(m, \vec{q}, k_T, x) \frac{2N_c \alpha_s}{\pi} (1 - x)
$$

$$
= \left[ \frac{q_T^2 + 2x^2 m^2}{d e} - \frac{x^2 m^2}{d^3} - \frac{x^2 m^2 (k_T^2 + q_T^2 + x^2 m^2)}{d^3} \right],
$$

with $d \triangleq \sqrt{(k_T^2 - q_T^2 + x^2 m^2)^2 + 4x^2 m^2 q_T^2}$. For $q_T \gg x m$, one observes a dip of the radiation at $k_T \lesssim x m$ as well as suppression as compared to light quarks (as can be seen from instance from the value of $P^{\text{SQCD}}$ at $k_T = 0$, $\propto \frac{q_T^2}{(q_T^2 + x^2 m^2)^2}$). This is the celebrated “dead - cone” phenomenon [29]. For smaller $q_T$ values ($q_T \lesssim x m$), both cones in (6) interfere and the dead-cone structure disappears, although the radiation of heavy quark stays suppressed. This effect is not identified in [29] – where the radiation from the heavy quark is assumed to be independent of its diffusion angle – but indeed corresponds to
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the most frequent situation encountered in induced radiative energy loss as \( \frac{d\sigma^{qQ \rightarrow qQ}}{dq_T} \) is peaked at small \( q_T \). After \( k_T \) integration of \( P_{SQCD}(m,q_T,k_T,x) \), one obtains a simple formula which can be further approximated by

\[
x \frac{d\sigma^{qQ \rightarrow qQ}}{dx} \approx \frac{d\sigma_{el}^{qQ \rightarrow qQ}}{dx} = \frac{2N_c\alpha_s(1-x)}{\pi} \ln \left( 1 + \frac{q_T^2}{3x^2m^2} \right).
\]

(12)

In fact, the limitation imposed by the phase space boundary becomes drastic only when \( k_{t,\text{max}} \) drops below \( \text{max}(x m, q_T) \), i.e. for \( x \lesssim \frac{q_T^2}{s} \sim \frac{\mu}{E} \), what justifies (12) in the energy range in which we are mostly interested. From (12), one sees that the radiative factor \( P_{SQCD} \) partly cures the IR divergence of the elastic cross section at small \( q_T \) but it still suffers from a logarithmic divergence. As we consider the radiation of heavy quarks propagating through the QGP, we introduce an infrared regulator \( \mu^2 \) of the order of the Debye mass \( m_D^2 \) in (8) and can proceed to the \( \vec{q}_T \) integration, assuming \( \mu^2 \ll m^2 \):

\[
x \frac{d\sigma^{qQ \rightarrow qQ}}{dx} = \frac{4C_F\alpha_s^3}{3} \begin{cases} 
\ln \left( \frac{\mu^2}{3x^2m^2} \right) & \text{for } x \lesssim \frac{\mu}{m} \\
\ln \left( \frac{\mu^2}{x^2m^2} \right) & \text{for } x \gtrsim \frac{\mu}{m}
\end{cases},
\]

(13)

which describes the radiation spectra in the mass range of interest better than 15%, as seen in fig. 2.

![Figure 2](image_url) (Color online) Left: ratio of the the approximate radiation (eq. 13) at \( \mu = 0.3 \text{ GeV} \) and of the exact \((q_T,k_T)\) integration of (7) when \( s \rightarrow +\infty \) for radiation from strange (full), charm (dashed) and bottom (dot-dashed) quarks. Right: radiative energy spectra per unit length with (continuous) and without (dashed) gluon mass; thin, plain and thick lines correspond to \( m = 0.15, 1.5 \) and 5 GeV.

In the QGP environment the radiated gluons polarize the medium, an effect that can be incorporated phenomenologically to the formalism [7] by imposing a thermal gluon mass \( m_g \) of the order of the Debye mass. In this case the previous calculations still hold if \( x^2m^2 \) is replaced by \( x^2m^2 + (1-x)m_g^2 \). On fig. 2 (right) we illustrate the effect of a finite gluon mass on the energy spectra per unit length, defined as \( \frac{\lambda_Q}{\sigma_{el}} \frac{dx}{dx} \frac{d\sigma^{qQ \rightarrow qQ}}{dx} \), where \( \lambda_Q \) is the mean free path of the heavy quark w.r.t. \( \sigma_{el} \). In the region of large \( x \) which dominates the average energy loss, the effect of the gluon mass on the radiative cross section of heavy quark is moderate, but a finite gluon mass has a drastic effect on
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the radiation for small and intermediate $x$-values which are usually thought to dominate the quenching [30] observed in $R_{AA}$, i.e. $x \lesssim 0.2$.

To implement the radiative processes in MC@sHQ [21], – designed to simulate the heavy quark transport in an expanding medium – we have evaluated the total number of radiated gluon candidates in a given elastic collision of momentum transfer $q_T$ as

$$ P_{\text{SQCD}}(m, q_T) \triangleq \int_{x_{\text{min}}}^1 \frac{d x}{x} P_{\text{SQCD}}(m, q_T, x), $$

with $P_{\text{SQCD}}(m, q_T, x)$ defined by (12), as well as the average number of radiated gluon ensuing an elastic collision as $P_{\text{SQCD}}(s) = \frac{1}{\sigma_{\text{el}}} \int d^2 q_T \frac{d \sigma_{\text{el}}}{d^2 q_T} P_{\text{SQCD}}(m, q_T)$, where the elastic cross section has been taken according with an effective running coupling (model “E” of [21]), although the gluon emission embedded in $P(m, q_T)$ has been taken at a fixed $\alpha_s$ value in this preliminary study.

For each elastic collision happening in the MC evolution, we then generate the gluon candidates according to Poisson statistics with an average of $P_{\text{SQCD}}(s)$, which happens to be $< 1$ for reasonable values of $\alpha_s$. We then successively sample the $\frac{d \sigma_{\text{el}}}{d^2 q_T} P_{\text{SQCD}}(m, q_T)$, $P_{\text{SQCD}}(m, q_T, x)$, $P_{\text{SQCD}}(m, q_T, k_T, x)$ and $P_{\text{SQCD}}(m, q_T, k_T, x)$ weights for $q_T$, $x$, $k_T$ and $\phi(q_T, k_T)$, so generating the full kinematics of the $2 \to 3$ process, and only accept the event if phase-space constrain (7) is satisfied. A similar method is used for the $gQ \to gQg'$ case. The thermal gluon mass was taken as $m_g = 2T$ at each space-time point of the QGP evolution.

3. Results and Discussion

Fig. 3 displays essential results in our approach. On the top-left panel, we show $R_{AA}$ as a function of $p_T$ in comparison with experimental Au-Au data, with and without radiative energy loss. As already discussed in [21], without radiative energy loss $R_{AA}$ is about 0.5 for large $p_T$ values and therefore well above the data. Radiative energy loss alone (shown as the shaded area bounded by calculation for $\alpha_s = 0.2$ and 0.3 for the gluon emission vertex) reproduces almost the observed values of $R_{AA}$, while our results are slightly below the data if we add both types of energy losses. These observations are in contradiction with the established literature [11, 13, 12]. The reasons for this are twofold: a) the running coupling approach presented in [21] naturally leads to $\hat{q}$ values (see fig. 6 of [22]) that exceed the fixed $\alpha_s$ pQCD prediction [31] and b) we limit ourselves, in this first study, to incoherent GB radiation and do not incorporate finite length nor LPM-like effects (see discussion below) which obviously quench the radiation spectra. In our view, a) constitutes a global improvement, while b) is a shortcoming of our present description of radiative energy loss.

On the right hand side of the top panel we display $R_{AA}(p_T)$ separately for leptons from $D$ and $B$ meson decay. Leptons from D mesons are practically not present anymore at large $p_T$ values. $c$-quarks are indeed stronger quenched than $b$-quarks and contribute less to high $p_T$ leptons due to their softer fragmentation function. On the left-bottom panel, we proceed along the strategy of [25] and optimize our multiplicative coefficient

$\xi$ the $x_{\text{min}}$ cut off, chosen as $x_{\text{min}} = 0.05$, permits to discard ultrasoft gluons that lead to IR catastrophe due to the “$x^{-1}$” weighting of the cross section but which contribute little to the stopping.
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Figure 3. (Color online) top left: nuclear modification factor $R_{AA}$ of single leptons resulting from the decay of heavy mesons (B and D) produced in central Au+Au ($\sqrt{s_{NN}} = 200$ GeV), for 3 different energy loss scenarios; top right: same for coll. + rad. case, separated for leptons resulting from $B$ and $D$ mesons decay; bottom left: same as top left in case of pure collisional and coll. + rad. mixture with respective multiplicative factors, $K = 2$ and $K = 0.55$, of the interaction rates; bottom right: same as bottom left for $B$ and $D$ mesons.

$K$ according to the most central class of Au-Au collisions. What is astonishing is the fact that the functional form of $R_{AA}(p_T)$ is very similar for “radiative + collisional” and for “collisional only” energy loss and that they both provide a good description of the data, although the quenching at high $p_T$ is often considered as dominated by radiative processes. Although not displayed here*, an equally good agreement of both rescaled models is found for all Au+Au centrality classes as well as for the elliptical flow $v_2$. In fact, a detailed analysis demonstrates that the $R_{AA}$ observable is mostly sensitive to a succession of moderate energy transfers for which both processes have similar probability distribution of fractional energy loss up to a global factor; this explains why rescaling both types of energy loss scenarios leads to similar agreement for $R_{AA}(p_T)$ and $v_2$. In short terms, the physics of $R_{AA}$ at RHIC is dominated by the Fokker-Planck regime, even at the largest $p_T$. This explains why models lacking radiative energy loss like [17] are able to reproduce the data as well. On the other hand this implies that radiative and collisional energy loss are not distinguishable on the basis of present n.p.s.e. experimental data we have analyzed. This is our main conclusion and the main scientific message that we wanted to deliver at SQM 2009.

* At small and intermediate $p_T$, this claim should be taken with a grain of salt for the “coll. + rad.” scenario, as the transition amplitude was derived in the large energy limit.

* see [link for explicit plots].
To provide some perspective for future experiments, we have shown, on the bottom-right panel, the \( R_{AA} \) for \( B \) and \( D \) mesons for both rescaled scenario (\( K = 2 \) for “coll.” and \( K = 0.55 \) for “coll + rad. (GB)”). For pure collisional and \( p_T \) around 10 GeV, the ratio of heavy mesons \( R_{AA} \), 
\[ R_{cb}(p_T) = \frac{R_{DAA}(p_T)}{R_{BAA}(p_T)} \]
found to be of the order of 0.6 [23]. When radiative energy loss is included, one has \( R_{cb}(10) \approx 0.4 \), which agrees with the value of 0.45 found in [13] and is still twice as large as predicted in the AdS/CFT approach [19]. Thus, those various models could be deciphered with the help of the high-\( p_T \) \( R_{cb} \) ratio, although it is questionable whether upgrades of RHIC experiments will be sufficient to provide the mandatory resolution power, especially for the \( D \)-mesons.

The Landau Pomeranchuck Migdal (LPM) and finite length effects, which are of essential importance for the radiation off light partons, turn out to be less important for the heavy quarks because the finite mass reduces the (effective) formation length, found to be \( l_f \approx \frac{x(1-x)E}{m_\gamma^2 + x^2 m^2} \) for vacuum radiation. For small as well as for large \( x \), \( l_f \) is smaller than the mean free path \( \lambda_Q \); only for the emission off a heavy quarks around \( x \approx m_g/m \), the formation length exceeds substantially \( \lambda_Q \), but we are nevertheless closer to the incoherent regime than for radiation off l.q., as demonstrated in [7]. This was the primary motivation for neglecting LPM and finite length effects in this first implementation of radiative processes in MC@HQ. Looking back, finding the same typical values for \( R_{cb} \) as in [13] – where those effects are taken into account – might be the indication that the main source of discrepancy relies in the \( \hat{q} \) value, that is in the overall interaction rates. Although our approach [21] complemented with radiative energy-loss constitutes a way out of level-2 s.e.p., it is still incomplete in the sense that we have not yet investigated its consequences on the quenching of light hadrons, i.e. the level-2 s.e.p.

In conclusion, we have shown that the combination of radiative and collisional energy loss, both calculated in a running \( \alpha_s \) pQCD-inspired model, allows to describe simultaneously the nuclear modification factor of n.p.s.e. for all centrality classes as well as the elliptical flow observed in Au-Au collisions at \( \sqrt{s} = 200 \) AGeV. We have argued that radiative and collisional energy loss are not distinguishable on the basis of these data only. The influence of some theoretical uncertainties [12, 13] and of different expansion scenarios \( \# \) – i.e. the generalization of [32] for heavy-quark observables – remains to be seen, although we do not expect it will affect our main conclusion. A detailed incorporation of the coherence effects (LPM and finite-length) is also mandatory, although a preliminary implementation of LPM effect in our framework, to be presented in an upcoming publication, was found of mere importance for \( c \)-quarks at RHIC and negligible for \( b \)-quarks.

\( \# \) different initial conditions, different freeze out temperatures and different equations of states in the hydrodynamical calculations.
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