Energy Dependence of Jet Quenching and Life-time of the Dense Matter in High-energy Heavy-ion Collisions

Xin-Nian Wang
Nuclear Science Division, MS70R0319, Lawrence Berkeley National Laboratory, Berkeley, CA 94720
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Suppression of high $p_T$ hadron spectra in high-energy heavy-ion collisions at different energies is studied within a pQCD parton model incorporating medium induced parton energy loss. The $p_T$ dependence of the nuclear modification factor $R_{AA}(p_T)$ is found to depend on both the energy dependence of the parton energy loss and the power-law behavior of the initial jet spectra. The high $p_T$ hadron suppression at $\sqrt{s} = 62.4$ GeV and its centrality dependence are studied in detail. The overall values of the modification factor are found to provide strong constraints on the lifetime of the dense matter.

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I. INTRODUCTION

The discovery of jet quenching in central $Au + Au$ collisions at the Relativistic Heavy-ion Collider (RHIC) at Brookhaven National Laboratory has provided clear evidence for the formation of strongly interacting dense matter. The observed jet quenching includes suppression of single hadron spectra at high $p_T$ [1,2], disappearance of back-to-back correlation of high $p_T$ hadrons [3] and the azimuthal anisotropy of high $p_T$ hadron spectra in non-central $Au + Au$ collisions [4]. The absence of these jet quenching phenomena in $d + Au$ collisions [5–8] proves that they are indeed due to final state interaction. Detailed analyses [9] further indicate that they are caused by parton energy loss as predicted [10–13]. Using the parton energy loss extracted from experimental data, one can conclude that the initial gluon density in central $Au + Au$ collisions [5] is about 4 GeV/c and its centrality dependence are studied in detail. Th e additional physics is at play in the energy dependence of the hadron suppression at large $p_T$ (up to $\sqrt{s} = 62.4$ GeV) and its centrality dependence are studied in detail. The overall values of the modification factor are found to provide strong constraints on the lifetime of the dense matter.

II. ENERGY DEPENDENCE OF HIGH $p_T$ HADRON SUPPRESSION

We will use a LO pQCD model [18] to calculate the inclusive high-$p_T$ hadron cross section in $A + A$ collisions,

$$\frac{d\sigma_{AA}}{dy dp_T} = K \sum_{abcd} \int d^2b d^2r dx_a dx_b d^2k_a d^2k_b d^2k_c d^2k_d$$

$$t_A(r) t_A(|b - r|) g_A(k_a, r) g_A(k_b, |b - r|) f_{a/A}(x_a, Q^2, r) f_{b/A}(x_b, Q^2, |b - r|)$$

$$\frac{D_{h/c}(z_c, Q^2, \Delta E_c)}{\pi z_c} \frac{d\sigma}{dt}(ab \to cd),$$

where $\sigma(ab \to cd)$ are elementary parton scattering cross sections and $t_A(b)$ is the nuclear thickness function normalized to $\int d^2b t_A(b) = A$. We will use a hard-sphere model of nuclear distribution in this paper. The $K \approx 1.5 - 2$ factor is used to account for higher order QCD corrections. The hadron is assumed to have the same rapidity as the parton, $y = y_c$ and its fractional momentum is defined as $z_c = p_T/p_T$. The parton distributions per nucleon $f_{a/A}(x_a, Q^2, r)$ inside the nucleus are assumed to be factorizable into the parton distributions in a free nucleon given by the MRS D′ parameterization [21] and the impact-parameter dependent nuclear
modification factor given by the new HIJING parameterization [22]. The initial transverse momentum distribution $p_A(k_T, Q^2, b)$ is assumed to have a Gaussian form with a width that includes both an intrinsic part in a nucleon and nuclear broadening. Detailed description of this model and systematic comparisons with experimental data can be found in Ref. [18].

The effect of parton energy loss is implemented through an effective modified fragmentation function [23],

$$D_{h/c}(z_c, Q^2, \Delta E_c) = (1 - e^{-\langle D_{h/c}/c \rangle}) \left[ \frac{z_c}{z_c} D^0_{h/c}(z_c, Q^2) + \left(\frac{\Delta L}{\lambda} \cdot \frac{z_c}{z_c} D^0_{h/c}(z_c, Q^2) \right) + e^{-\langle D_h/c \rangle} D^0_{h/c}(z_c, Q^2) \right]. \quad (2)$$

This effective form is a good approximation to the actual calculated medium modification in the multiple parton scattering formalism [24], given that the actual energy loss should be about 1.6 times the input value in the above formula. Here $z_c' = p_T/(p_T - \Delta E_c)$, $z_g' = \langle\Delta L/\lambda \rangle p_T/\Delta E_c$ are the rescaled momentum fractions and $\Delta E_c$ is the total energy loss during an average number of inelastic scatterings $\langle\Delta L/\lambda\rangle$. The FF’s in free space $D^0_{h/c}(z_c, Q^2)$ are given by the BBK parameterization [25].

In this study, we assume a 1-dimensional expanding medium with a gluon density $\rho_g(\tau, r)$ whose initial distribution is proportional to the transverse profile of participant nucleons. The total energy loss for a parton propagating this medium is

$$\langle \Delta E(\tau, r) \rangle \approx \langle \frac{dE}{dL} \rangle_{1d} \int_{\tau_0}^{\tau_{\text{max}}} \frac{d\tau - \tau_0}{\tau_0 \rho_0} \rho_g(\tau, b, \vec{r} + \vec{n}\tau), \quad (3)$$

according to recent theoretical studies [13,26,27], where $\rho_0$ is the averaged initial gluon density at $\tau_0$ in a central collision. $\langle \frac{dE}{dL} \rangle_{1d}$ is the average parton energy loss over a distance $R_A$ in a 1-d expanding medium with an initial uniform gluon density $\rho_0$. Similarly, the average number of scatterings along the path of parton propagation is

$$\langle \Delta L/\lambda \rangle = \int_{\tau_0}^{\tau_{\text{max}}} \frac{d\tau}{\tau_0 \rho_0} \rho_g(\tau, b, \vec{r} + \vec{n}\tau). \quad (4)$$

Inclusion of transverse radial expansion generates a faster dilution of the gluon density relative to 1-d expansion, but also results in longer propagation time in medium. These effects offset each other and the final total energy loss in the cases of 1-d and 3-d expansion is found very similar [28]. With a hard-sphere nuclear distribution, the gluon density profile $\rho_g(\tau, r)$ can be expressed in a simple form and the analytic expressions of $\Delta E(\tau, b, r)$ and $\langle \Delta L/\lambda \rangle$ are given in the Appendix.

The energy dependence of the energy loss is parameterized as

$$\langle \frac{dE}{dL} \rangle_{1d} = e_0(E/\mu - 1.6)^{1.2}(7.5 + E/\mu), \quad (5)$$

according to the numerical results in Ref. [29] in which thermal gluon absorption is also taken into account in the calculation of parton energy loss. Fit to the most central Au + Au collisions at $\sqrt{s} = 200$ GeV in $e_0 = 1.07$ GeV/fm, $\mu = 1.5$ GeV and $\lambda_0 = 1/\sigma\rho_0 = 0.3$ fm. The corresponding energy loss in a static medium with parton density $\rho_0$ over a distance $R_A$ is [27] $dE_0/dL = (R_A/2\tau_0)(dE/dL)_{1d} \approx 14$ GeV/fm. This is about 30 times higher than the parton energy loss in a cold nucleus [14].

At different energies, we assume that $e_0$ and $\rho_0$ are proportional to the measured hadron multiplicity as given in Ref. [30]. Since there is no experimental measurement at the LHC energy yet, we will use the model calculation [22] to extrapolate and assume $(N_{ch}/d\eta)$ at $\sqrt{s} = 2.4$.

To study the dependence of the high $p_T$ hadron suppression on the colliding energy, let us assume first that the lifetime of the dense matter is larger than the system size. Shown in Fig. 1 are the nuclear modification factors

$$R_{AB}(p_T) = \frac{\langle d\sigma_A^h/dy \rangle}{\langle N_{\text{binary}} \rangle d\sigma_{pp}^h}, \quad (6)$$

for charged hadron (solid lines) and neutral pions (dashed lines) in central Au + Au (Pb + Pb at the SPS energy) collisions at different energies, from SPS $\sqrt{s} = 17.2$ GeV to LHC $\sqrt{s} = 5.5$ TeV. Here,

$$\langle N_{\text{binary}} \rangle = \int d^2 b d^2 r A(r)a_A(\vec{b} - \vec{r}) \quad (7)$$

is the number of geometrical binary collisions at a given range of impact parameters. At the SPS energy, the observed nuclear modification factor in central Pb + Pb collisions is consistently about 1 due to strong Cronin effect via initial multiple parton scattering, leaving not much room for large parton energy loss [19]. We shall return to this point later.

In central Au + Au collisions at RHIC, however, strong suppression of high $p_T$ hadrons is observed. This can be attributed to large parton energy loss that overcomes the modest Cronin enhancement as observed in d + Au collisions [5-8] and gives rise to the large hadron suppression. The energy-dependence of the parton energy loss in Eq. (5) describes well the flat $p_T$ dependence of the nuclear modification factor $R_{AA}(p_T)$ for neutral pions at large $p_T$ in Au + Au collisions at $\sqrt{s} = 200$ GeV, as shown in the figure. For charged hadrons there are complications from other medium effects as we will discuss later. The flatness of the modification factor at this energy is more clearly illustrated by the ratio of central to peripheral collisions [31,32]. Such a flat $p_T$ dependence is actually a coincidence, as the combined effect of the
energy dependence of the parton energy loss in Eq. (5) and the power-law behavior of the initial jet spectra. Using the same energy dependence of the parton energy loss but with a reduced amplitude due to smaller initial gluon density at $\sqrt{s} = 62.4$ GeV, the nuclear modification factor is found to decrease with $p_T$ and even becomes smaller than the modification factor at $\sqrt{s} = 200$ GeV at high $p_T > 10$ GeV/$c$. This is simply a consequence of the energy dependence of jet spectrum shape. The initial jet spectra at $\sqrt{s} = 62.3$ GeV are much steeper than those at 200 GeV. The same amount of energy loss leads to a larger suppression of the final hadron spectra at 62.3 GeV than at 200 GeV. As one increases the colliding energy, the power-law spectra for the initial jet production become flatter, and the same parton energy loss will lead to less suppression of the final hadrons. As shown in the same figure, the nuclear modification factor at the LHC energy $\sqrt{s} = 5.5$ TeV is smaller than at 200 GeV in the intermediate $p_T$ region, due to larger initial gluon density. However, the modification factor $R_{AA}(p_T)$ increases with $p_T$ due to the flatter power-law spectra of jet production at LHC.

The parameterized energy dependence of the parton energy loss is in part due to the detailed balance effect in induced gluon radiation and absorption [27]. This effect is most important in the intermediate $p_T$ region. In this region, one expects the jet fragmentation process to be modified by other non-perturbative processes such as parton recombination or coalescence [33–35]. The observed flavor dependence of the hadron suppression and of the azimuthal anisotropy clearly points to the effect of parton recombination that enhances both baryon and kaon spectra in the presence of dense medium. To include this effect in the current parton model, we have added a soft component to kaon and baryon FF’s that is proportional to the pion FF with a weight $\sim \langle N_{\text{bin}}(b, \phi) \rangle / \langle 1 + \exp(2p_T\phi - 15) \rangle$. The functional form and parameters are adjusted so that $(K + p)/\pi \approx 2$ at $p_T \sim 3$ GeV/$c$ in the most central $Au + Au$ collisions at $\sqrt{s} = 200$ GeV and approaches its $p + p$ value at $p_T > 5$ GeV/$c$. This gives rise to the splitting of the suppression factor for charged hadrons and $\pi^0$ in the calculation. Because of the steeper power-law spectra of jet production at 62.4 GeV, the effect of the non-perturbative parton recombination persists to higher $p_T$ than in $Au + Au$ collisions at 200 GeV. In this region of $p_T$, the non-perturbative recombination effects dominate the nuclear modification of the charged hadron spectra. As a consequence, the $p_T$ dependence of the modification factors $R_{AA}(p_T)$ at $\sqrt{s} = 62.4$ and 200 GeV are similar. They only diverge at high $p_T > 8$ GeV/$c$ where the recombination effects are negligible at both energies. Since the soft components due to parton recombination are not closely related to medium thermalization, they will still contribute to the final hadron spectra in peripheral $A + A$ and $p + A$ collisions. This will lead to flavor dependence of the Cronin effect in $p + A$ collisions [36].

III. EFFECT OF FINITE LIFETIME

In the calculation of the parton energy loss in Eqs. (3) and (4), the upper limit of the path integral should be

$$\tau_{\text{max}} = \min(\Delta L, \tau_f),$$

where $\tau_f$ is the lifetime of the dense matter before breakup, $\Delta L(b, \vec{r}, \phi)$ is the distance the parton, produced at $\vec{r}$, has to travel along $\vec{n}$ at an azimuthal angle $\phi$ relative to the reaction plane in a collision with impact-parameter $b$. According to the analyses of experimental data on high $p_T$ hadron suppression and suppression of away-side correlation in $Au + Au$ collisions at $\sqrt{s} = 200$ GeV, the extracted energy loss points to an initial gluon density of about 30/fm$^3$ at an initial time $\tau_0 = 0.2$ fm. Given the measured transverse energy per charged hadron of 0.8 GeV [37], this gives a lower bound on the initial energy density of about 25 GeV/fm$^3$. In 1-d expansion with the equation of state of an ideal fluid, the energy density decreases with time, $\epsilon(\tau) = \epsilon_0(\tau/\tau_0)^{4/3}$. Assuming the 1-d hydrodynamics expansion starts at $\tau_i = 1$ fm/$c$ (free-streaming before that), the lifetime of the plasma or the duration for the parton energy loss should be about $\tau_f \sim 5$ fm, before the phase transition with a critical energy density $\epsilon_c \sim 1$ GeV/fm$^3$. The early stage of the mixed phase or crossover could also contribute to the jet quenching and thus extend the effective time duration for parton energy loss. When this time is larger than the
average path length, the total parton energy loss is then limited only by the system size. In the previous analysis [14,29], such an assumption for central Au + Au collisions is justfied given the high initial energy density.

In central Pb + Pb collisions at the SPS energy $\sqrt{s} = 17.3$, the rapidity density of charged particles is about half of that in central Au + Au collisions at $\sqrt{s} = 200$ GeV [30]. The average transverse energy per charged particle is about the same at SPS and RHIC energy [37]. One can then assume the initial energy density at SPS to be half of that in central Au + Au collisions at $\sqrt{s} = 200$ GeV. The lifetime of the plasma in central Pb + Pb collisions at SPS, $\tau_s \sim 2 - 3$, if it were formed, is then considered longer than that at the highest energy of RHIC. The longer thermalization time at SPS could make the effective lifetime even shorter. To demonstrate the dependence on the lifetime, we show in Fig. 1 the nuclear modification factor of $\pi^0$ at the SPS energy with short lifetime $\tau_f = 0, 1, 2$ fm/$c$. The modification factor is quite sensitive to the lifetime $\tau_f$. Even if one takes into account the uncertainty in the reference $p + p$ data [29], which gives a systematic error of about a factor of 2, the data still point to a short lived dense system with $\tau_f < 2$ fm/$c$. This might also explain the observed elliptic flow $v_2$ at SPS that is much smaller than the hydrodynamic limit [38]. On the other hand, the longer lifetime of the plasma allows development of a full hydrodynamic flow, giving rise to a large $v_2$ that saturates the hydrodynamic limit [39].

To study the sensitivity of the high $p_T$ hadron suppression to the lifetime of the plasma at $\sqrt{s} = 62.4$ GeV, we show in Fig. 2 the nuclear modification factors for both charged hadrons and neutral pions with different values of $\tau_f$ in Au + Au collisions at $\sqrt{s} = 62.4$ GeV. The hadron suppression in the large $p_T$ region in the most central collisions is very sensitive to the lifetime of the plasma in this calculation. In peripheral collisions, the small size of the dense medium limits the parton energy loss. As a result, the hadron suppression is only sensitive to values of $\tau_f$ that are smaller than the average medium size. In reality, the values of $\tau_f$ should decrease from central to peripheral collisions.

The recent experimental results from PHOBOS [40] on nuclear modification factors for charged hadrons in Au + Au collisions at 62.4 GeV only extend to $p_T \sim 4$ GeV/$c$. In this region, the suppression of charged hadron is indeed much smaller than at 200 GeV. However, charged hadrons in this region are also dominated by non-perturbative recombination effects, though our results are still sensitive to the lifetime. Experimental measurements of $\pi^0$ and high $p_T$ charged hadrons, both are less influenced by the parton recombination effect, should provide more stringent constraints on the lifetime of the dense matter.

IV. SUMMARY AND DISCUSSION

Within a parton model incorporating medium induced parton energy loss, we have studied in this brief report the suppression of inclusive hadron spectra at high $p_T$ in heavy-ion collisions at different energies. We found that the $p_T$ dependence of the nuclear modification factor $R_{AA}(p_T)$ is determined by the energy dependence of the parton energy loss and the power-law behavior of the initial jet spectra. With the onset of parton energy loss and the change of the power-law jet spectra, the $p_T$ dependence of the modification factor changes from monotonic decrease at $\sqrt{s} = 62.4$ GeV to monotonic increase with $p_T$ at the LHC energy $\sqrt{s} = 5.5$ TeV. The flat $p_T$ dependence observed at $\sqrt{s} = 200$ GeV is just a coincidence.

We also studied the sensitivity of the hadron suppression factor to the lifetime of the plasma or the duration of parton energy loss. We found that the hadron suppression factor at intermediate and large $p_T$ is sensitive to the lifetime if it is comparable or smaller than the system size. The experimental measurements could provide important constraints. Together with the measurement of elliptic flow, which is also sensitive to the lifetime and thermalization time of the plasma, one can gain additional information on the dynamic evolution of the produced quark-gluon plasma.

One can also calculate the back-to-back dihadron correlation at different energies. We find that suppression of the back-to-back dihadron correlation at 62.4 GeV in
central Au + Au collisions is almost identical to that at 200 GeV. This is partly due to the trigger bias that selects dihadron production close to the surface and results in completely suppression the back-side jets that traverse the whole length of the dense matter. The suppression due to $k_T$ broadening of initial multiple parton scattering is also independent of the colliding energy.

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V. APPENDIX

In this appendix, we give the basic analytic formula for calculating the path integral in the parton energy loss in Eqs. (3) and (4), assuming a hard-sphere nuclear distribution.

Given two overlapping nuclei as illustrated in Fig. 3, we want to calculate a path integral over the path $\Delta L$. Let $r_1$ and $r_2$ be the radial coordinates of the jet production point as measured from the center of the two nuclei. For given $\vec{b}$ and $\vec{r}_1$

$$r_2 = \sqrt{b^2 + r_1^2 - 2br_1 \cos \phi_b}$$  (9)

For a hard-sphere distribution, $r_1 \leq R_A$ and $r_2 \leq R_B$.

The jet travels at an angle $\phi_p$ with respect to the impact-parameter $\vec{b}$ (which determines the reaction plane with the beam direction). $\phi_b$ is the angle between $\vec{b}$ and $\vec{r}_1$. The distance the jet travels is then the $\Delta L = \min(\tau_1, \tau_2)$;

$$\tau_1 = \sqrt{R_A^2 - r_1^2 \sin^2 \phi_r - r_1 \cos \phi_r};$$  (10)

$$\tau_2 = \sqrt{R_B^2 - r_1^2 + (r_1 \cos \phi_r - b \cos \phi_p)^2} - (r_1 \cos \phi_r - b \cos \phi_p),$$  (11)

where $\phi_r = \phi_p - \phi_b$. These are solutions of

$$R_A = |\vec{r}_1 + \vec{n}\tau_1|,$$  (12)

$$R_B = |\vec{r}_2 - \vec{b} + \vec{n}\tau_2|.$$  (13)

If $b < |R_A - R_B|$

$$\Delta L = \sqrt{R_{\min}^2 - r_1^2 \sin^2 \phi_r - r_1 \cos \phi_r},$$  (14)

where $R_{\min} = \min(R_A, R_B)$. We also define $R_{\max} = \max(R_A, R_B)$.

Assuming that the soft gluon density is proportional to the number of participant nucleons, it is then given by

$$\rho_g(\tau, \vec{b}, \vec{r}) = \frac{\tau_0 \rho_0}{\tau} \left[ \frac{R_A^3}{A} \theta(\tau - |\vec{b} - \vec{r}|) + \frac{R_B^3}{B} \theta(\tau - R_B - |\vec{b} - \vec{r}|) \right].$$  (15)

where $c_{AB} = 1 - (1/2)(1 - R_{\min}^2/R_{\max}^2)^{3/2}$ and $\rho_0$ is defined as the averaged gluon density in central collisions ($b = 0$) at an initial time $\tau_0$:

$$\rho_0 = \frac{1}{\pi R_{\min}^2} \int d^2r \rho(\tau_0, \vec{r}, b = 0).$$  (16)

Using the nuclear thickness function defined as

$$t_A(r) = \frac{3}{2\pi R_A^2} \sqrt{1 - r^2/R_A^2},$$  (17)

the gluon density is then

$$\rho_g(\tau, \vec{b}, \vec{r}) = \frac{3\tau_0 \rho_0}{4\tau c_{AB} R_{\min}} \left[ \theta(\tau - R_A - r) \sqrt{R_A^2 - (\vec{b} - \vec{r})^2} + \theta(\tau - R_B - |\vec{b} - \vec{r}|) \right] \left[ \frac{R_A^2 - (\vec{r} + \vec{n}\tau)^2}{R_A^2 - \vec{b}^2} \right].$$  (18)

According to Eq. (3), the total energy loss along the path $\Delta L$ is

$$\Delta E(b, r, \phi) = \frac{dE}{dL} |_{1d} \int_{\tau_0}^{\tau_0 + \tau_0} d\tau \frac{\tau - \tau_0}{\tau_0 \rho_0} \rho_g(\tau, b, \vec{r} + \vec{n}\tau)$$

$$= \frac{dE}{dL} |_{1d} \int_{\tau_0}^{\tau_0 + \Delta L} d\tau \frac{3}{4\tau c_{AB} R_{\min}} \left[ \frac{R_A^2 - (\vec{r} + \vec{n}\tau)^2}{R_A^2 - \vec{b}^2} \right].$$  (19)

The average number of scatterings is

$$\langle \Delta L/\lambda \rangle = \int_{\tau_0}^{\tau_0 + \Delta L} d\tau \sigma \rho_g(\tau, b, \vec{r} + \vec{n}\tau)$$

$$= \frac{3}{4\lambda_0 R_{\min}} \int_{\tau_0}^{\tau_0 + \Delta L} d\tau \left[ \frac{R_A^2 - (\vec{r} + \vec{n}\tau)^2}{R_A^2 - \vec{b}^2} \right].$$  (20)
The above integrals can be completed analytically. The following are some basic integrals:

\begin{align}
\int d\tau \sqrt{R^2 - (\vec{r} + \vec{\nu}\tau)^2} &= \frac{\tau + \vec{r} \cdot \vec{\nu}}{2} \sqrt{R^2 - (\vec{r} + \vec{\nu}\tau)^2} \\
+ \frac{R^2 - r^2 + (\vec{r} \cdot \vec{\nu})^2}{2} \arcsin \frac{\tau + \vec{r} \cdot \vec{\nu}}{\sqrt{R^2 - r^2 + (\vec{r} \cdot \vec{\nu})^2}},
\end{align}

(21)

\begin{align}
\int \frac{d\tau}{\tau} \sqrt{R^2 - (\vec{r} + \vec{\nu}\tau)^2} &= \sqrt{R^2 - (\vec{r} + \vec{\nu}\tau)^2} \\
- (\vec{r} \cdot \vec{\nu}) \arcsin \frac{\vec{r} \cdot \vec{\nu} + \tau}{\sqrt{R^2 - r^2 + (\vec{r} \cdot \vec{\nu})^2}} \\
- \sqrt{R^2 - r^2} \left\{ \log \left[ R^2 - r^2 - (\vec{r} \cdot \vec{\nu})\tau \right] \\
+ \sqrt{R^2 - r^2} \right\} - \log \tau
\right\}
\end{align}

(22)

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