Non-Abelian feature of parton energy loss in energy dependence of jet quenching in high-energy heavy-ion collisions

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One of the non-Abelian features of parton energy loss is the ratio $\Delta E_g/\Delta E_q = 9/4$ between gluon and quark jets. Since jet production rate is dominated by quark jets at high $x_T = 2p_T/\sqrt{s}$ and by gluon jets at low x_T, high p_T hadron suppression in high-energy heavy-ion collisions should reflect such a non-Abelian feature. Within a leading-order perturbative QCD parton model that incorporates transverse expansion and Woods-Saxon nuclear distribution, the energy dependence of large $p_T \sim 5$–20 GeV/c hadron suppression is found to be sensitive to the non-Abelian feature of parton energy loss and could be tested by data from low-energy runs at Brookhaven National Laboratory's relativistic heavy-ion collider or data from the large hadron collider.

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One of the ultimate goals of the relativistic heavy-ion collider (RHIC) program at Brookhaven National Laboratory is to produce a quark-gluon plasma by smashing two gold nuclei at the speed of light. The discovery of the jet quenching effect [1–4] in central Au+Au collisions together with the observation of parton recombination [5–12] and the early thermalization of dense matter [13–22] has provided clear evidence of the formation of strongly interacting partonic matter [23,24]. The observed jet quenching effect manifests itself in several aspects of high p_T hadron spectra, which include suppression of inclusive spectra in Au+Au relative to pp collisions [1,2], disappearance of back-to-back correlations [3], and azimuthal anisotropy in noncentral Au+Au collisions [4]. The absence of these jet quenching phenomena in d + Au collisions [25–28] shows that they are due to final-state interactions with the strongly interacting matter produced. Detailed analyses indicate that parton energy loss is the source of the observed jet quenching [29–34]. The initial gluon density, to which the parton energy loss is proportional, has been extracted from RHIC data of central Au+Au collisions at $\sqrt{s} = 200$ AGeV and is about 30 times higher than that in a cold nucleus [35,36].

The radiative parton energy loss incorporated in previous studies within a leading-order (LO) perturbative QCD (pQCD) model [32,33,37–40] has two basic non-Abelian features. One is the quadratic dependence on the total distance traversed by the propagating parton due to the non-Abelian Landau-Pomeranchuk-Midal (LPM) interference effect in gluon bremsstrahlung induced by multiple scatterings in a *static* medium [41–48]. The second feature of the parton energy loss is its dependence on the color representation of the propagating parton. Therefore, energy loss for a gluon is 9/4 times larger than for a quark. Previous works have investigated the consequences of the second non-Abelian feature in the flavor dependence of the high p_T hadron suppression [38]. In this paper we study the effect of the non-Abelian parton energy loss on the energy dependence of the inclusive hadron spectra suppression. We exploit the well-known feature of the initial parton distributions in nucleons (or nuclei) that quarks dominate at large fractional momentum x while gluons dominate at small x. Jet or large p_T hadron production as a result of hard scatterings of initial partons will be dominated by quarks at large $x_T = 2p_T/\sqrt{s}$ and by gluons at small x_T. Since gluons lose 9/4 times as much energy as quarks, the energy dependence of the large (and fixed) p_T hadron spectra suppression due to parton energy loss should reflect the transition from quark-dominated jet production at low energy to gluon-dominated jet production at high energy. Such a unique energy dependence of the high-p_T hadron suppression can be tested by combining $\sqrt{s} = 200$ AGeV data with lower energy data or future data from large hadron collider (LHC) experiments.

We will work within a LO pQCD parton model incorporating the non-Abelian QCD parton energy loss in high-energy heavy-ion collisions. We will study the energy dependence of the high p_T hadron suppression and compare the effect of QCD energy loss with that of a non-QCD one where gluons and quarks are chosen to have the same amount of energy loss. In both cases, we will assume that parton energy loss is proportional to the initial gluon density of the system which in turn is assumed to be proportional to the measured total charge hadron multiplicity in the central rapidity region.

In comparison to previous studies within the LO pQCD parton model that employed the hard-sphere model of nuclear distribution and assumed only longitudinal expansion, we will use the more realistic Woods-Saxon nuclear distribution and include the transverse expansion of the dense medium.

In a LO pQCD model [37], the inclusive invariant differential cross section for high p_T hadrons in $A + B$ collisions is given by

$$\frac{d\sigma_{AB}^h}{dyd^2p_T} = K\sum_{abcd}\int d^2\mathbf{b}d^2\mathbf{r}dx_a dx_b d^2\mathbf{k}_{aT}d^2\mathbf{k}_{bT}t_A(r)t_B \times (|\mathbf{b}-\mathbf{r}|)g_A(k_{aT},r)g_B(k_{bT},|\mathbf{b}-\mathbf{r}|)$$

$$\times f_{a/A}(x_a, Q^2, r) f_{b/B}(x_b, Q^2, |\mathbf{b}-\mathbf{r}|)$$
$$\times \frac{D_{h/c}(z_c, Q^2, \Delta E_c)}{\pi z_c} \frac{d\sigma(ab \to cd)}{d\hat{t}}, \qquad (1)$$

where $\sigma(ab \to cd)$ are elementary parton scattering cross sections. The factor $K \approx 1.0$–2.0 is used to account for higher order QCD corrections and is set to be the same for both $p + p$ and $A + B$ collisions at the same energy. The hadron is assumed to have the same rapidity as the parton, i.e., $y = y_c$, and its fractional momentum is defined by $z_c = p_T/p_{Tc}$. The parton distributions per nucleon $f_{a/A}(x_a, Q^2, r)$ inside the nucleus can be factorized into the parton distributions in a free nucleon given by the Coordinated Theoretical-Experimental Project on QCD (CTEQ) parametrization [49,50] and the impact-parameter-dependent nuclear modification factor given by the new (HIJING) parametrization: $f_{a/A}(x_a, Q^2, r) = R_a^A(x, Q^2)[(Z/A) f_{a/p}(x, Q^2) + (1 - Z/A) f_{a/n}(x, Q^2)]$ with $R_a^A(x, Q^2)$ given by Eqs. (8) and (9) of Ref. [51]. We assume that the initial transverse momentum distribution $g_A(k_T, Q^2, b)$ has a Gaussian form [37,52] with a width that includes both an intrinsic k_T in a nucleon and the nuclear broadening due to initial multiple scattering in a nucleus: $g_A(k_T, Q^2, b) = e^{-k_T^2/\langle k_T^2\rangle_A}/(\pi\langle k_T^2\rangle_A)$. The impact-parameter-dependent broadened variance is given by $\langle k_T^2\rangle_A(Q^2) = \langle k_T^2\rangle_N(Q^2) + \delta^2(Q^2)[\nu_A(b) - 1]$, where the number of scatterings $\nu_A(b)$ the projectile suffers inside the nucleus is $\nu_A(b) = \sigma_{NN} t_A(b)$ with the nuclear thickness function $t_A(b)$ defined as follows, and the scale-dependent $\delta^2(Q^2)$ chosen as $\delta^2(Q^2) = 0.225 \ln^2(Q/\text{GeV})/[1 + \ln(Q/\text{GeV})]\ \text{GeV}^2/c^2$. The average initial intrinsic transverse momentum in nucleon-nucleon collisions is $\langle k_T^2\rangle_N(Q^2) = 1.2 + 0.2 Q^2 \alpha_s(Q^2)$. The scale that characterizes the partonic process is chosen to be $Q = p_T$, where p_T is the transverse momentum of the final-state partons in a partonic scattering. Detailed description of this model and systematic comparisons with experimental data can be found in Ref. [37]. In this paper we use the Woods-Saxon nuclear distribution $F_{\text{WS}}(r) = N_A/[1 + \exp((r - R_A)/a)]$ to replace the simplified hard-sphere one used in previous papers. Here R_A is the radius of the nucleus given by $R_A = 1.12 A^{1.0/3.0} - 0.86 A^{-1.0/3.0}$, $a = 0.54$ fm is a radius parameter, and N_A is the normalization constant. The Woods-Saxon distribution can be further written as a function of the coordinate component z along the beam direction of the nucleus and $\mathbf{b}$ that is perpendicular to it by $r = \sqrt{z^2 + b^2}$. The nuclear thickness function $t_A(b)$ is then $t_A(b) = \int_{-\infty}^{\infty} dz F_{\text{WS}}(z, b)$ with the normalization condition $\int d^2b\, t_A(b) = A$.

The parton energy loss is encoded in an effective modified fragmentation function [53,54]

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

This effective form is a good approximation to the actual calculated medium modification in the multiple parton scattering formalism [55,56], given that the actual energy loss should be about 1.6 times the input value in formula (2). Here $z_c' = p_T/(p_{Tc} - \Delta E_c)$, $z_g' = \langle\Delta L/\lambda\rangle p_T/\Delta E_c$ are the rescaled momentum fractions and ΔE_c is the total energy loss during an average number of inelastic scatterings $\langle\Delta L/\lambda\rangle$. The fragmentation functions in free space $D^0_{h/c}(z_c, Q^2)$ are given by the BBK parametrization [57].

In contrast to previous calculations that considered only longitudinal expansion, we incorporate in this paper both longitudinal and transverse expansion of the medium in the calculation of parton energy loss. To simplify the calculation, we again use hard-sphere nuclear distribution. Let us assume the gluon number N_g is a slowly varying function of rapidity y and proper time τ at central rapidity region $y = 0$, then we have $d^2N_g/d\tau dy = 0$. Noting that $dN_g/dy = \rho dV/dy$ and $dV/dy = \tau\pi R_T^2$, we obtain

$$\frac{d\rho}{d\tau}\frac{dV}{dy} + \rho\left[\pi R_T^2 + 2\pi\tau R_T \frac{dR_T}{d\tau}\right] = 0\ , \qquad (3)$$

which is

$$\frac{d\rho}{d\tau} + \rho\left[\frac{1}{\tau} + \frac{2}{R_T}\frac{dR_T}{d\tau}\right] = 0\ . \qquad (4)$$

The radius has the form $R_T(\tau) = R_A + (\tau - \tau_0)c_s^2$ where c_s is the speed of sound in the medium given by $c_s^2 = \partial P/\partial e$ (1/3 for ideal gas). The solution of Eq. (4) is then

$$\tau\rho\left[R_A + (\tau - \tau_0)c_s^2\right]^2 = \tau_0\rho_0 R_A^2\ . \qquad (5)$$

The expansion is characterized by the gluon density $\rho_g(\tau, r)$ whose initial distribution is proportional to the transverse profile of participant nucleons. We can write the total energy loss for a parton traversing the medium as

$$\Delta E(b, \mathbf{r}, \phi) \approx \left\langle\frac{dE}{dL}\right\rangle_{1d} \times \int_{\tau_0}^{\tau_{\max}} d\tau \frac{\tau\left[R_{\min} + (\tau - \tau_0)c_s^2\right]^2 - \tau_0 R_{\min}^2}{\tau_0 R_{\min}^2 \rho_0} \times \rho_g(\tau, b, \mathbf{r} + \mathbf{n}\tau), \qquad (6)$$

where $R_{\min} = \text{Min}(R_A, R_B)$ and $\mathbf{n}$ is the direction where a parton is propagating. The upper limit $\tau_{\max} = \text{Min}(\Delta L, \tau_f)$ is the longest time for the parton to propagate in the dense medium, where τ_f is the lifetime of the dense matter before breakup. $\Delta L(b, \mathbf{r}, \phi)$ is the distance the parton, produced at $\mathbf{r}$, travels along $\mathbf{n}$ at the azimuthal angle ϕ relative to the reaction plane in a collision with impact parameter b. Since the formation time of a hadron fragmented from a parton is proportional to the energy of the parton, very high energetic partons generally hadronize after the dense medium breaks up, or they hadronize outside the medium. In this case we have $\tau_{\max} = \Delta L$. Function $\langle dE/dL\rangle_{1d}$ is the average parton energy loss over a distance R_A in a one-dimensional expanding medium with an initial uniform gluon density ρ_0. The gluon density ρ_g in the longitudinally and transversely expanding medium is then given by

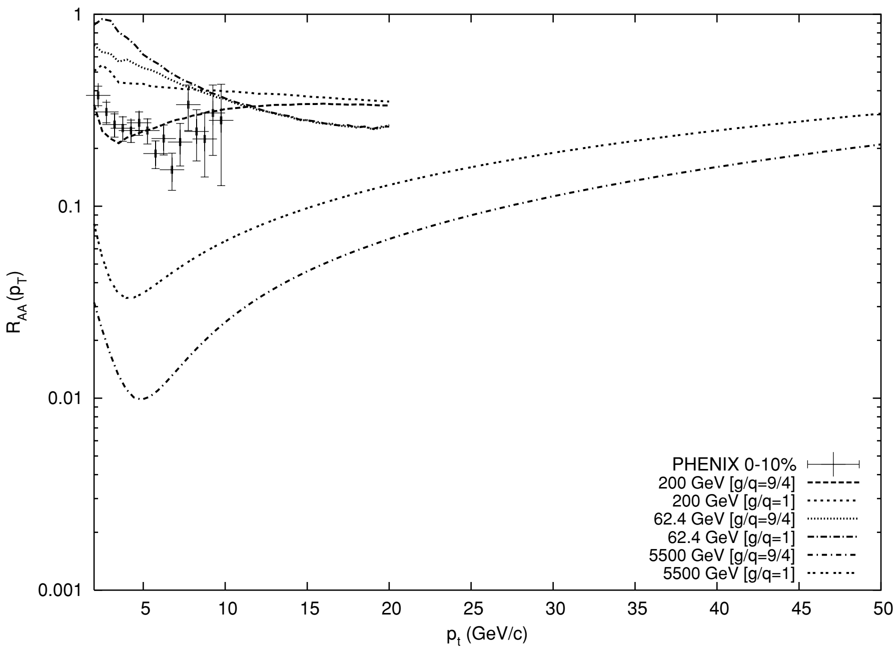

FIG. 1. Nuclear modification factor R_{AuAu} for neutral pions at $\sqrt{s} = 62.4$, 200, and 5500 AGeV. We chose the values of corresponding parameters at 62.4 and 5500 AGeV based on their values at 200 AGeV and the ratio $(dN_{\mathrm{ch}}/dy)_{62}/(dN_{\mathrm{ch}}/dy)_{200}$ and $(dN_{\mathrm{ch}}/dy)_{5500}/(dN_{\mathrm{ch}}/dy)_{200}$.

$$\rho_g(\tau, b, \mathbf{r} + \mathbf{n}\tau) = \frac{\tau_0\rho_0}{\tau}\frac{R_{\mathrm{min}}^2}{\left[R_{\mathrm{min}} + (\tau - \tau_0)c_s^2\right]^2}\frac{\pi}{2c_{AB}R_{\mathrm{min}}} \times \left[\frac{R_A^3}{A}t_A(r) + \frac{R_B^3}{B}t_B(|\mathbf{b} - \mathbf{r}|)\right]. \quad (7)$$

The average number of scatterings along the path of parton propagation is

$$\langle \Delta L/\lambda \rangle = \int_{\tau_0}^{\tau_{\mathrm{max}}} d\tau \sigma \rho_g(\tau, b, \mathbf{r} + \mathbf{n}\tau). \quad (8)$$

The energy loss function can be parametrized as

$$\left\langle \frac{dE}{dL} \right\rangle_{1d} = \epsilon_0 \frac{(E/\mu_0 - 1.6)^{1.2}}{7.5 + E/\mu_0}, \quad (9)$$

according to a study of parton energy loss [48] that included both bremstrahlung and thermal aborption of gluons. For $\sqrt{s} =$ 200 AGeV, we find the following set of parameters can fit the data: $\epsilon_0 = 1.2$ GeV, $\mu_0 = 1.6$ GeV, and $\lambda_0 = 0.2$ fm (λ_0 appears in the formula of $\langle \Delta L/\lambda \rangle$). In Ref. [35,39], these parameters are set to slightly different values: $\epsilon_0 = 1.07$ GeV, $\mu_0 =$ 1.5 GeV, and $\lambda_0 = 0.3$ fm. The value of $\langle dE/dL \rangle_{1d}$ with $\epsilon_0 =$ 1.2 GeV, $\mu_0 = 1.6$ GeV used in this paper is almost the same in the energy range $E =$ 5–20 GeV as with the previous values $\epsilon_0 = 1.07$ GeV, $\mu_0 = 1.5$ GeV [35,39]. For example, at $E = 5$ and 20 GeV we have $\langle dE/dL \rangle_{1d}(\epsilon_0 = 1.2, \mu_0 = 1.6) = 0.19$ and 1.05, while $\langle dE/dL \rangle_{1d}(\epsilon_0 = 1.07, \mu_0 = 1.5) = 0.19$ and 0.99. Another modified parameter λ_0 in this paper is inversely proportional to the average number of scatterings undergone by the propagating energetic parton. The value $\lambda_0 = 0.2$ fm is smaller than the previously used $\lambda_0 = 0.3$ fm, which means that the average number of scatterings is tuned larger to make more energy loss by compensating for the effect caused by transverse expansion, which makes the medium more rapidly diluted. Note that the parameter ϵ_0 is proportional and λ_0 is inversely proportional to the gluon or multiplicity density per rapidity. The energy loss in a corresponding static medium is found to be 14 GeV/fm, which is about 30 times as high as in a cold nuclei [35].

The jet quenching effect can be shown by the nuclear modification factor defined as [58]

$$R_{AB} = \frac{d\sigma_{AB}^h/dyd^2p_T}{\langle N_{\mathrm{binary}}\rangle d\sigma_{pp}^h/dyd^2p_T}, \quad (10)$$

where N_{binary} is the average number of geometrical binary collisions at a given range of impact parameters

$$\langle N_{\mathrm{binary}}\rangle = \int d^2\mathbf{b}d^2\mathbf{r}t_A(r)t_B(|\mathbf{b} - \mathbf{r}|). \quad (11)$$

If there is no energy loss, the cross section for nucleus-nucleus collisions is a simple sum of that for elementary binary nucleon-nucleon collisions, so the nuclear modification factor R_{AB} is Fig. 1. Hadron suppression due to parton energy loss leads to $R_{AB} < 1$.

As we mentioned earlier, we use the Woods-Saxon nuclear distribution in the parton model calculation. The numerical difficulty with the Woods-Saxon distribution is that one cannot simply put the analytical formula into the program because that would substantially reduce the speed of the calculation and make the numerical calculation practically impossible. One trick to overcome this problem is to calculate the distribution before hand and then store the results in tables whose entries can be called in the run time of the program. The calculated R_{AB} results for Au+Au collisions are shown in Figs. 1 and 2 with Fig. 1 for π^0 and Fig. 2 for charged hadrons. The results for three collision energies $\sqrt{s} = 62.4$, 200, and 5500 AGeV are given. The parameters ϵ_0 and λ_0 at these energies are set to appropriate values based on the ratios of charged particle (or gluon) multiplicity density [51,61] $(dN_{\mathrm{ch}}/dy)_{5500}/(dN_{\mathrm{ch}}/dy)_{200}$, $(dN_{\mathrm{ch}}/dy)_{62.4}/(dN_{\mathrm{ch}}/dy)_{200}$ and their values $\epsilon_0 = 1.2$ and $\lambda_0 = 0.2$ at 200 AGeV. In the figures, we can see different transverse momentum behaviors of the nuclear modification

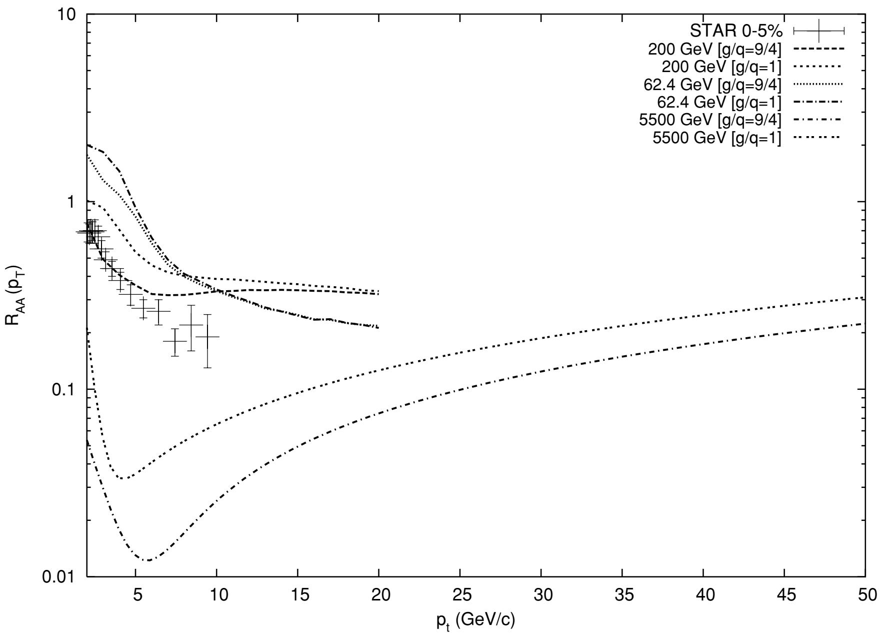

FIG. 2. Nuclear modification factor R_{AuAu} for hadrons at $\sqrt{s} = 62.4$, 200, and 5500 AGeV. We chose the values of corresponding parameters at 62.4 and 5500 AGeV based on their values at 200 AGeV and the ratio $(dN_{\mathrm{ch}}/dy)_{62}/(dN_{\mathrm{ch}}/dy)_{200}$ and $(dN_{\mathrm{ch}}/dy)_{5500}/(dN_{\mathrm{ch}}/dy)_{200}$.

factor at these energies. Similar behaviors have been seen in recent studies [39,59,60]. The nuclear modification factor decreases with p_T at 62.4 AGeV, while it slightly increases with p_T at 200 AGeV. So the nuclear modification factors for neutral pions and charged hadrons at 62.4 AGeV intersect at about $p_T = 11$ and 10 GeV, respectively, with those of 200 AGeV in the QCD case, where the energy loss parameters for gluons and quarks satisfy $g/q = 9/4$ (we will explain this point later). The same feature also occurs in Ref. [39] where the hard-sphere distribution and only the longitudinal expansion are used. In the intermediate p_T region, one expects the jet fragmentation process to be modified by other nonperturbative processes such as parton recombination or coalescence [9–11]. 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 added a soft component to the kaon and baryon fragmentation functions that is proportional to the pion fragmentation function with a weight $\sim\langle N_{\mathrm{bin}}(b,r)\rangle/[1+\exp(2p_{Tc}-15)]$ where p_{Tc} is the transverse momentum for parton c. {Actually we also found that a similar effect can be achieved by using a function of the hadron transverse momentum p_T: $\langle N_{\mathrm{bin}}(b,r)\rangle/[1+\exp(p_T-5)]$.} 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 $\sqrt{s} = 200$ AGeV 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 π^0 in the calculation.

To study the sensitivity of hadron spectra suppression to the non-Abelian parton energy loss, we compare the results with two different cases at each energy: one for the QCD case where the energy loss for a gluon is 9/4 times as large as for a quark, i.e., $\Delta E_g/\Delta E_q = 9/4$; the other is for a non-QCD case where energy loss is chosen to be the same for both gluons and quarks. Similarly, the average number of inelastic scatterings obeys $\langle\frac{\Delta L}{\lambda}\rangle_g \big/ \langle\frac{\Delta L}{\lambda}\rangle_q = 9/4$ in the QCD case. For the non-QCD case, that ratio is set to 1. In Figs. 1 and 2 we can see that the differences between the QCD and non-QCD cases are more significant for higher collision energies. This fact manifests itself at 200 and 5500 AGeV, where the nuclear modification factors R_{AB} are much lower for the QCD energy loss pattern than for the non-QCD one. As shown in the figures, the suppression at 62.4 AGeV is sensitive not to gluon energy loss but only to quark energy loss because of the dominance of quark jets at large p_T. At 200 AGeV, however, the suppression is sensitive to both quark and gluon energy losses. At LHC energy, the results are sensitive only to gluon energy loss in the p_T range we calculated. Such an energy-dependence pattern is a direct consequence of the non-Abelian feature of the energy loss.

To demonstrate the colliding energy dependence of the nuclear modification factor and illustrate the difference between QCD and non-QCD energy loss, we computed the R_{AA} for neutral pions at fixed $p_T = 6$ GeV in Au+Au collisions as a function of $\sqrt{s}$ from 20 to 5500 AGeV. Shown in Fig. 3 are the calculated results with both the QCD energy loss and a non-QCD case where the energy loss is set to be identical for quarks and gluons. Two parameters ϵ_0 and λ_0, which are relevant to the energy loss, are determined according to the gluon number or the charged particle multiplicity per rapidity [51,61]. One can see that because of the dominant gluon bremstrahlung or gluon energy loss at high energy the R_{AA} for the QCD case is more suppressed than for the non-QCD case where the gluon energy loss is assumed to have an influence equal to that of the quark energy loss. In the calculation, we assumed that the lifetime of the dense matter is equal to or longer than the parton propagation time, which is essentially determined by the system size. This might not be the case for heavy-ion collisions at lower energies, in particular around $\sqrt{s} = 20$ AGeV. If one assumes a short lifetime, the suppression factor R_{AA} is much larger than 1 due to a strong Cronin effect [39]. The dashed box around

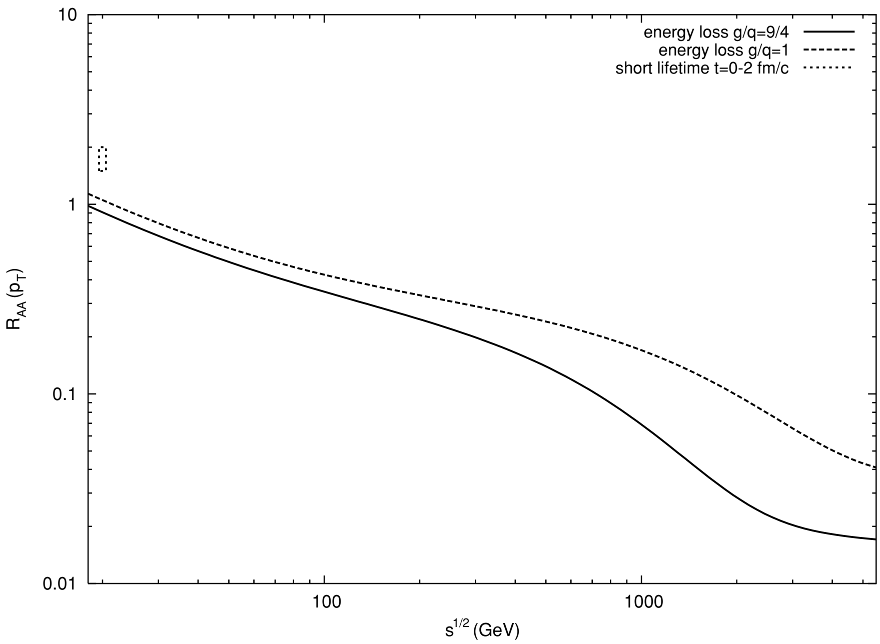

FIG. 3. Nuclear modification factor $R_{\rm AuAu}$ for neutral pions as function of collision energy at fixed $p_T = 6$ GeV in most central collisions (with centrality 10%). Here we compare the QCD energy loss and a non-QCD case where the energy loss is identical for quarks and gluons.

$\sqrt{s} = 20$ AGeV in Fig. 3 assumes a lifetime $\tau_f = 0$–2 fm/c and thus provides an estimate of the uncertainty due to the lifetime of the dense matter. Since a finite lifetime reduces the effect of the full parton energy loss, the difference between the QCD and non-QCD energy loss effects in $R_{\rm AA}$ should be smaller. The difference seen in Fig. 3 is therefore the upper limit.

Another interesting feature of the energy dependence of $R_{\rm AA}$ is the change of slope around $\sqrt{s} = 1300$ GeV. The rapid decrease of $R_{\rm AA}$ at $\sqrt{s} = 20$–1300 GeV is mainly due to the increase in parton energy loss caused by the increased initial gluon density and the change of p_T slope of jet production cross section with $\sqrt{s}$. As the energy loss increases, more jets produced inside the overlapped region are completely suppressed. Only those that are produced within an outlayer in the overlapped region survive. This phenomenon is like surface emission with a finite depth. The suppression factor $R_{\rm AA}$ is then determined by the width of the outlayer, which is just the averaged mean-free-path $\langle\lambda\rangle$. As a consequence, $R_{\rm AA}$ has much weaker $\sqrt{s}$ dependence.

In summary, nuclear modification factors in Au+Au collisions at $\sqrt{s} = 62.4$, 200, and 5500 AGeV are calculated in an LO perturbative QCD model with medium induced parton energy loss. The previous calculations based on the hard-sphere distribution of nucleus and the longitudinal expansion of the dense medium are improved in terms of a more realistic Woods-Saxon distribution and both longitudinal and transverse expansion. The comparison of nuclear modification factors for energy loss patterns in QCD and non-QCD cases shows sizable differences at higher colliding energies. Thus, the energy loss pattern can be tested by the energy dependence of the hadron suppression factor $R_{\rm AA}$ in the range $\sqrt{s} = 20$–1000 GeV. We also found a weaker energy dependence above $\sqrt{s} = 1000$ GeV due to surface emission with finite depth.

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