The 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 ultimate goal of the Relativistic Heavy-Ion Collider (RHIC) at Brookhaven National Laboratory is to produce the quark-gluon plasma (QGP) by smashing two gold nuclei at the speed of light. The discovery of jet quenching effect [1, 2, 3, 4, 5] in central Au+Au collisions together with the observation of parton recombination [6, 7, 8, 9, 10, 11, 12] and the early thermalization of the dense matter [13, 14, 15, 16, 17, 18, 19, 20, 21, 22] has provided clear evidence for the formation of strongly interacting partonic matter [23, 24]. The observed jet quenching effect manifests itself in several aspects of high-energy heavy-ion collisions. The absence of these jet quenching phenomena in pp collisions [25] shows that they are due to final state interactions with the produced strongly interacting matter. Detailed analyses indicate that parton energy loss is the source for the observed jet quenching [22, 26, 27, 28]. The initial gluon density, which the parton energy loss is proportional to, 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 [29, 30].

The radiative parton energy loss incorporated in previous studies within a leading order (LO) perturbative QCD (pQCD) model [31, 32, 33, 34, 35, 36, 37, 38, 39, 40] has two basic non-Abelian features. One of them is the quadratic dependence on the total distance traversed by the propagating parton due to the non-Abelian Landau-Pomeranchuk-Migdal (LPM) interference effect in gluon bremsstrahlung induced by multiple scatterings in a static medium [41, 42, 43, 44, 45, 46]. 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 a quark. Previous works have investigated the consequences of the second non-Abelian feature in the flavor dependence of the high-$p_T$ hadron spectra [47, 48]. 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 more 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 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 more realistic Woods-Saxon nuclear distribution and in addition 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
\[
\frac{d\sigma_{AB}}{dy d^2p_T} = K \sum_{abcd} \int d^4p d^4q d^2r d^2k d^2k' t_A(r) t_B(|b - r|) g_A(k_{aT}, r) g_B(k_{bT}, |b - r|) \times f_{a/A}(x_a, Q^2, r) f_{b/B}(x_b, Q^2, |b - r|) \frac{D_{h/c}(z_c, Q^2, \Delta E_c)}{\pi z_c} \frac{d\sigma(ab \to cd)}{dt},
\]
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 CTEQ parameterization \(49, 50\) and the impact-parameter dependent nuclear modification factor given by the new HIJING parameterization: \(f_{a/A}(x_a, Q^2, r) = R_A^a(x, Q^2)(Z/A)f_{a/p}(x, Q^2) + (1 - Z/A)f_{a/n}(x, Q^2)\) 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: \(g_A(k_T, Q^2, b) = e^{-k_T^2/(\langle k_T^2 \rangle_A)} \sigma_N^2 \sigma_t\). The impact-parameter dependent broadened variance is given by \(\langle k_T^2 \rangle_A(Q^2) = \langle k_T^2 \rangle_N(Q^2) + \sigma^2(\langle Q^2 \rangle)\), where the number of scatterings \(\nu_A(b)\) the projectile suffers inside the the nucleus is \(\nu_A(b) = \sigma_N N_t A(b)\) with the nuclear thickness function \(t_A(b)\) defined as follows, and the scale-dependent \(\sigma^2(\langle Q^2 \rangle)\) is chosen as \(\sigma^2(\langle Q^2 \rangle) = 0.225 \text{ln}^2(Q/\text{GeV})/[1 + \text{ln}(Q/\text{GeV})] \text{GeV}^2/\alpha^2\). The average initial intrinsic transverse momentum in nucleon-nucleon collision is \(\langle k_T^2 \rangle_N(Q^2) = 1.2 + 0.2Q^2\). The scale which 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. \(51\). In this paper we use the Woods-Saxon nucleon distribution \(F_{WS}(r) = N_A/[1 + \exp((r - R_A)/a)]\) to replace the simplified hard-sphere PDFs used in previous papers. Here \(R_A\) is the radius of the nucleus and \(N_A = 1.12\Delta A^{0.8}/A - 0.86\Delta A^{1.0}/A\). \(a = 0.54\) fm is a radius parameter and \(N_A\) is the normalization constant. It can be further written as a function of the coordinate component \(z\) along the beam direction of the nucleus and \(b\) that is perpendicular to it by \(r = \sqrt{z^2 + b^2}\). The nuclear thickness function \(t_A(b)\) is then \(t_A(b) = f_{WS}(z, b)\) with the normalization condition \(f_{WS}(b, b) = A\).

The parton energy loss is encoded in an effective modified fragmentation function \(53, 54\)
\[
D_{h/c}(z_c, Q^2, \Delta E_c) = 1 - e^{-\langle \Delta L/\lambda \rangle} \left[ z_c' D^0_{h/c}(z_c', Q^2) + \frac{\Delta L}{\lambda} z_c' D^0_{h/g}(z_g', Q^2) \right] + e^{-\langle \Delta L/\lambda \rangle} D^0_{h/c}(z_c, Q^2).
\]
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 the above formula. Here \(z_c' = p_T/(p_{Tc} - \Delta E_c)\), \(z_g' = (\Delta L/\lambda) p_T/\Delta E_c\) 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 fragmentation functions in free space \(D^0_{h/c}(z_c, Q^2)\) are given by the BBK parameterization \(57\).

In contrast to previous calculations where only longitudinal expansion was considered, 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 use hard-sphere nuclear distribution again. Let us denote the gluon number \(N_g\) and assume that \(N_g\) is a slowly varying function of rapidity \(y\) and proper time \(\tau\) at central rapidity region \(y = 0\), then we have \(d^2N_g/d\tau dy = 0\). Noting that \(dN_g/dy = p dV/dy\) and \(dV/dy = d\tau R_T^2\), we obtain
\[
\frac{d\rho}{d\tau} \frac{dV}{dy} + \rho \left[ \pi R_T^2 + 2\pi \tau R_T \frac{dR_T}{d\tau} \right] = 0,
\]
which is
\[
\frac{d\rho}{d\tau} + \rho \left[ \frac{1}{\tau} + 2 \frac{dR_T}{R_T} \frac{d\tau}{d\tau} \right] = 0.
\]
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 the above equation is then
\[
\tau \rho [R_A + (\tau - \tau_0)c_s^2]^2 = \tau_0 \rho_0 R_A^2.
\]
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

$$\Delta E(b, r, \phi) \approx \frac{dE}{dL} \int_{\tau_0}^{\tau_{\text{max}}} d\tau \frac{\langle R_{\text{min}} + (\tau - \tau_0) c_s^2 \rangle}{\tau_0 R_{\text{min}}^2 \rho_0} \rho_g(\tau, b, r + n\phi), \quad (6)$$

where $R_{\text{min}} = \text{Min}(R_A, R_B)$, $n$ is the direction where a parton is propagating. The upper limit $\tau_{\text{max}} = \text{Min}(\Delta L, \tau_f)$ is the longest time for the parton to propagate in the dense medium, where $\tau_f$ is the lifetime of the dense matter before breakup. $\Delta L(b, r, \phi)$ is the distance the parton, produced at $r$, travels along $n$ at the azimuthal angle $\phi$ 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 $\tau_{\text{max}} = \Delta L$. $\langle dE/dL \rangle_{1d}$ is the average parton energy loss over a distance $R_A$ in a 1-dimensional expanding medium with an initial uniform gluon density $\rho_0$.

The gluon density $\rho_g$ in the longitudinally and transversely expanding medium is then given by

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

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

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

The energy loss function can be parameterized as

$$\frac{dE}{dL} = \frac{E/\mu_0 - 1.6}{7.5 + E/\mu_0} \quad (9)$$

according to a study of parton energy loss [48] that include both included bremsstrahlung and thermal aborption of gluons. For $\sqrt{s}=200$ AGeV, we find following set of parameters $\epsilon_0 = 1.2$ GeV, $\mu_0 = 1.6$ GeV and $l_0 = 0.2$ fm ($l_0$ appears in the formula of $\langle \Delta L/\Lambda \rangle$) can fit the data. In Ref. [33, 34], these parameters are set to slightly different values $\epsilon_0 = 1.07$ GeV, $\mu_0 = 1.5$ GeV and $l_0 = 0.3$ 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 previous values $\epsilon_0 = 1.07$ GeV, $\mu_0 = 1.5$ GeV [33, 34]. 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 $l_0$ in this paper is inversely proportional to the average number of scatterings undergone by the propagating energetic parton. The smaller value $l_0 = 0.2$ fm than previously used one $l_0 = 0.3$ fm means that the average number of scatterings is tuned larger to make more energy loss by compensating the effect caused by transverse expansion which makes the medium more rapidly diluted. Note that the parameter $\epsilon_0$ is proportional and $l_0$ 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 [55]

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

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

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

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 one. Hadron suppression due to parton energy loss leads to $R_{AB} < 1$.

As we mentioned earlier that 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 slow 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 $R_{AA}$ results for Au+Au collisions are shown in Figs. 1 and 2 with Fig. 1 for $\pi^0$ and Fig. 2 for charged hadrons. The results for three collision energies $\sqrt{s}=62.4$, 200 and 5500 AGeV are given. The parameters $c_0$ and $\lambda_0$ at these energies are set to appropriate values based on the ratios of charged particle (or gluon) multiplicity density $dN_{ch}/dy$ at 62.4, 200, and 5500 AGeV and their values $c_0=1.2$ and $\lambda_0=0.2$ at 200 AGeV. In the figures one can see different transverse momentum behaviors of the nuclear modification factor at these energies. Similar behaviors have been seen in recent studies [32, 33, 64]. 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. [32] 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 non-perturbative processes such as parton recombination or coalescence [9, 14, 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 have added a soft component to kaon and baryon fragmentation function that is proportional to the pion fragmentation function with a weight $\sim (N_{bin}(b, r))/[1 + \exp(2p_T - 15)]$ where $p_Tc$ is the transverse momentum for parton $c$ (actually we have also found that the similar effect can be also achieved by using a function of the hadron transverse momentum $p_T$: $(N_{bin}(b, r))/[1 + \exp(p_T - 5)]$. 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 $\pi^0$ 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 the energy loss is chosen to be the same for both gluons and quarks. Similarly, the average number of inelastic scatterings obeys $(\Delta L_g)/(\Delta L_q) = 9/4$ in the QCD case. For the non-QCD case we are considering, the above ratio is set to one. In Figs. 1 and 2 one can see that the difference 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 $R_{AA}$ 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 not sensitive 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 both to quark and gluon energy loss. At LHC energy, the results are only sensitive to gluon energy loss in the $p_T$ range we calculated. Such energy dependence pattern is a direct consequence of the non-Abelian feature of the energy loss.

In order to demonstrate the colliding energy dependence of the nuclear modification factor and illustrate the difference between QCD and non-QCD energy loss, we compute the $R_{AA}$ for neutral pions at fixed $p_T = 6$ GeV in Au+Au collisions as a function of $\sqrt{s}$ from 20 AGeV 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 $c_0$ and $\lambda_0$ which are relevant to the energy loss are determined according to the gluon number or the charged particle multiplicity per rapidity $55, 61$. One can see that due to the dominant gluon bremsstrahlung or gluon energy loss at high energy the $R_{AA}$ for the QCD case is more suppressed than the non-QCD case where the gluon energy loss is assumed to take an equal role as the quark one. In the calculation, we have assumed that the lifetime of the dense matter is equal 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 at around $\sqrt{s} = 20$ AGeV. If one takes short lifetime, the suppression factor $R_{AA}$ is much larger than 1 due to strong Cronin effect [33]. The dashed box around $\sqrt{s} = 20$ AGeV in the Fig. 3 assumes a lifetime $\tau_f = 0 - 2$ fm/c and thus provides an estimate of the uncertainty due to lifetime of the dense matter. Since finite lifetime reduce the effect of full parton energy loss, the difference between QCD and non-QCD energy loss effect in $R_{AA}$ should be smaller. The difference in Fig. 3 is therefore the upper limit.

Another interesting feature with the energy dependence of $R_{AA}$ is the change of slope around $\sqrt{s} = 1300$ GeV. The rapid decrease of $R_{AA}$ at $\sqrt{s} = 20 - 1300$ GeV is mainly due to increase of parton energy loss due to increased initial gluon density and also 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 will survive. This will be like surface emission with a finite depth. The suppression factor $R_{AA}$ will then be determined by the width of the outlayer which is just the averaged mean-free-path $\langle \lambda \rangle$. As a consequence, $R_{AA}$ will then have 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 a LO perturbative QCD model with medium induced parton energy loss. The previous calculations based on the
FIG. 1: Nuclear modification factor $R_{AA}$ for neutral pions at $\sqrt{s} = 62.4$, 200 and 5500 AGeV. We choose the values of corresponding parameters at 62.4 and 5500 AGeV based on their values at 200 AGeV and the ratio $(dN_{ch}/dy)_{62}/(dN_{ch}/dy)_{200}$ and $(dN_{ch}/dy)_{5500}/(dN_{ch}/dy)_{200}$.

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 difference at higher colliding energies and thus can be tested by the energy dependence of the hadron suppression factor $R_{AA}$ in the range between $\sqrt{s} = 20 - 1000$ GeV. We also found a weaker energy dependence above $\sqrt{s} = 1000$ GeV due to surface emission with finite depth.

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

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FIG. 3: Nuclear modification factor $R_{AA}$ 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 one where the energy loss is identical for quarks and gluons.

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