Hard Probes in High-energy Heavy-ion Collisions

Xin-Nian Wang
Nuclear Science Division, Mailstop 70A-3307
Lawrence Berkeley National Laboratory
University of California, Berkeley, CA 94720 USA

ABSTRACT

Hard QCD processes in ultrarelativistic heavy-ion collisions become increasingly relevant and they can be used as probes of the dense matter formed during the violent scatterings. We will discuss how one can use these hard probes to study the properties of the dense matter and the associated phenomenologies. In particular, we study the effect of jet quenching due to medium-induced energy loss on inclusive particle $p_T$ distributions and investigate how one can improve the measurement of parton energy loss in direct photon events.

1 Introduction

Quantum Chromodynamics (QCD) has been established as the underlying theory of strong interactions. This theory is probably best demonstrated in $e^+e^-$ annihilation processes in which a quark-antiquark pair is produced and then fragmented into hadrons. In the earliest time or at the shortest distance, one can use perturbative QCD (pQCD) to describe the interaction, i.e., parton radiation. Later on, the produced partons will then combine with each via nonperturbative interaction and finally hadronize into hadrons. Therefore, one can consider that there exists an interacting parton system during the prehadronization stage in $e^+e^-$ annihilation processes, which is, however, limited to a space-time region characterized by the confinement scale $\Lambda_{QCD}$. The characteristic particle spectrum (in $p_T$ and rapidity) and the ratios of produced particles should then be determined by the physics of pQCD and nonperturbative hadronization. One seeks to produce a similar interacting parton system or a quark-gluon plasma but at a much larger scale of the order of a nucleus size in ultrarelativistic heavy-ion collisions. Therefore, one should study those experimental observables which are unique to the large size and long life-time of an interacting partonic system as signals of a quark-gluon plasma.

There are many proposed signals of a quark-gluon plasma \cite{1}. Among them, hard probes associated with hard processes are especially useful because they are produced in the earliest stage of the interaction and their abilities to probe the dense matter are less complicated by the hadronization physics. The merits of hard probes are even more apparent at high energies because those processes are also dominating the underlying collision dynamics which will determine the initial conditions of the produced partonic system \cite{2, 3}. Study of them will then enable us to probe the early
parton dynamics and the evolution of the quark-gluon plasma. In general, one can divide the hard probes into two categories: thermal emission and particle suppression by the medium. Particle production, like photon/dilepton and charm particles, from thermal emission can be considered as the thermometers of the dense medium. Their background comes from the direct production in the initial collision processes. On the other hand, suppression of particles produced in the initial hard processes, like high-$p_T$ particles from jets and $J/\Psi$, can reveal evidences of the parton energy loss in dense matter and the deconfinement of the partonic system. Thermal production of these particles is expected to be negligible. Therefore, in both cases, one needs to know the initial production rate accurately enough. One advantage of these hard probes is that the initial production rate can be calculated via pQCD, especially if we understand the slightly nuclear modification one would expect to happen. In this talk, I will only concentrate on high-$p_T$ particle suppression due to jet quenching.

Medium-induced radiative energy loss of a high-energy parton traversing a dense QCD medium is interesting because it depends sensitively on the density of the medium and thus can be used as a probe of the dense matter formed in ultrarelativistic heavy-ion collisions. As recent studies demonstrated [4, 5, 6], it is very important to take into account the destructive interference effect in the calculation of radiation spectrum from a fast parton induced by multiple scatterings. The so-called Landau-Pomeranchuk-Migdal effect can lead to very interesting, and sometimes nonintuitive results for the energy loss of a fast parton in a QCD medium. For example, Baier et al. recent showed [6] that the energy loss per distance, $dE/dx$, is proportional to the total length that the parton has traveled. Because of the unique interference effect, the parton somehow knows its history of propagation. Another feature of the induced energy loss is that it depends on the parton density of the medium it is traversing via the final transverse momentum broadening that the parton receives during its propagation. One can therefore determine the parton density of the produced dense matter by measuring the energy loss of a fast parton when it propagates through the medium.

Unlike in the QED case, where one can measure directly the radiative energy loss of a fast electron, one cannot measure directly the energy loss of a fast leading parton in QCD. Since a parton is normally studied via a jet, a cluster of hadrons in the phase space, an identified jet can contain particles both from the fragmentation of the leading parton and from the radiated partons. If we neglect the $p_T$ broadening effect, the total energy of the jet should not change even if the leading parton suffers radiative energy loss. What should be changed by the energy loss are the particle distributions inside the jet or the fragmentation function and the jet profile. Therefore, one can only measure parton energy loss indirectly via the modification of the jet fragmentation function and jet profile.

In principle, one can measure the parton energy loss by directly measuring the fragmentation function and profile of a jet with a determined transverse energy. However, because of the huge background and its fluctuation in high-energy heavy-ion collisions [7], the conventional calorimetric study of a jet cannot determine the jet energy thus the energy loss very well. In Ref. [8], Gyulassy and I proposed that
single-particle spectrum can be used to study the effect of jet energy loss, since the suppression of large $E_T$ jets naturally leads to the suppression of large $p_T$ particles. However, since the single-particle spectrum is a convolution of the jet cross section and jet fragmentation function, the suppression of produced particles with a given $p_T$ results from jet quenching with a wide range of initial transverse energies. One, therefore, cannot measure directly, from the single-particle $p_T$ spectrum, the energy loss of a jet with known initial transverse energy. Recently, Huang, Sarcevic and I proposed to study the jet quenching by measuring the $p_T$ distribution of charged hadrons in the opposite direction of a tagged direct photon [9]. Since a direct photon in the central rapidity region ($y = 0$) is always accompanied by a jet in the opposite transverse direction with roughly equal transverse energy, the $p_T$ distribution of charged hadrons in the opposite direction of the tagged direct photon is directly related to the jet fragmentation function with known initial energy. One can thus directly measure the modification of the jet fragmentation and then determine the energy loss suffered by the leading parton.

In this talk, I will review the effects of energy loss on single-particle distributions both in the normal central $A + A$ collisions and in events with a tagged direct photon with known transverse energy. I will discuss the energy and $A$ dependences of the energy loss and jet quenching. In the case of jet quenching in $\gamma + jet$ events, the $E_T$ smearing of jet due to initial state radiation will be included. The change of jet profile function in the azimuthal angle due to $p_T$ broadening of the jet will also be discussed. Finally, discussions will be given on the feasibility of measuring the energy loss in $\gamma + jet$ events at RHIC.

2 Modified jet fragmentation functions

Jet fragmentation functions have been studied extensively in $e^+e^-$, $ep$ and $p\bar{p}$ collisions [10]. These functions describe the particle distributions in the fractional energy, $z = E_h/E_{jet}$, in the direction of a jet. The measured dependence of the fragmentation functions on the momentum scale is shown to satisfy the QCD evolution equations very well. We will use the parametrizations of the most recent analysis [11] in both $z$ and $Q^2$ for jet fragmentation functions $D_{h/a}^0(z, Q^2)$ to describe jet ($a$) fragmentation into hadrons ($h$) in the vacuum.

In principle, one should study the modification of jet fragmentation functions in a perturbative QCD calculation in which induced radiation of a propagating parton in a medium and Landau-Pomeranchuk-Migdal interference effect can be dynamically taken into account. However, for the purpose of our current study, we can use a phenomenological model to describe the modification of the jet fragmentation function due to an effective energy loss $dE/dx$ of the parton. In this model we assume: (1) A quark-gluon plasma (QGP) is formed with a transverse size of the colliding nuclei, $R_A$. A parton with a reduced energy will only hadronize outside the deconfined phase and the fragmentation can be described as in $e^+e^-$ collisions. (2) The inelastic scattering mean-free-path for the parton $a$ inside the QGP is $\lambda_a$. The radiative energy loss per
scattering is $\epsilon_a$. The energy loss per distance is thus $dE/a/dx = \epsilon_a/\lambda_a$. The probability for a parton to scatter $n$ times within a distance $\Delta L$ is given by a Poisson distribution,

$$P_a(n, \Delta L) = \frac{(\Delta L/\lambda_a)^n}{n!}e^{-\Delta L/\lambda_a}. \hspace{1cm} (1)$$

We also assume that the mean-free-path of a gluon is half that of a quark, and the energy loss $dE/dx$ is twice that of a quark. (3) The emitted gluons, each carrying energy $\epsilon_a$ on the average, will also hadronize according to the fragmentation function with the minimum scale $Q_0^2 = 2.0 \text{ GeV}^2$. We will also neglect the energy fluctuation given by the radiation spectrum for the emitted gluons. Since the emitted gluons only produce hadrons with very small fractional energy, the final modified fragmentation functions in the moderately large $z$ region are not very sensitive to the actual radiation spectrum and the scale dependence of the fragmentation functions for the emitted gluons.

We will consider the central rapidity region of high-energy heavy-ion collisions. We assume that a parton with initial transverse energy $E_T$ will travel in the transverse direction in a cylindrical system. With the above assumptions, the modified fragmentation functions for a parton traveling a distance $\Delta L$ can be approximated as,

$$D_{h/a}(z, Q^2, \Delta L) = \frac{1}{C_N^a} \sum_{n=0}^{N} P_a(n, \Delta L) \frac{z_a^n}{z} D_{h/a}^0(z_a^n, Q^2) + \langle n_a \rangle \frac{z_a'}{z} D_{h/g}^0(z_a', Q_0^2), \hspace{1cm} (2)$$

where $z_a^n = z/(1 - n\epsilon_a/E_T)$, $z_a' = zE_T/\epsilon_a$ and $C_N^a = \sum_{n=0}^{N} P_a(n)$. We limit the number of inelastic scatterings to $N = E_T/\epsilon_a$ by energy conservation. For large values of $N$, the average number of scatterings within a distance $\Delta L$ is approximately $\langle n_a \rangle \approx \Delta L/\lambda_a$. The first term corresponds to the fragmentation of the leading partons with reduced energy $E_T - n\epsilon_a$ and the second term comes from the emitted gluons each having energy $\epsilon_a$ on the average.

3 Energy loss and single-particle $p_T$ spectrum

To calculate the $p_T$ distribution of particles from jet fragmentation in a normal central heavy-ion collision, one simply convolutes the fragmentation functions with the jet cross section [12],

$$\frac{dN_{AA}^{\text{hard}}}{dyd^2p_T} = K \int d^2r \sum_{abcdh} \int_{x_{amin}}^{1} dx_a \int_{x_{bmin}}^{1} dx_b f_{a/A}(x_a, Q^2, r) f_{b/A}(x_b, Q^2, r).$$
where $z_c = x_T(e^y/x_a + e^{-y}/x_b)/2$, $x_{min} = x_a x T e^{-y}/(2 x_a - x_T e^{-y})$, $x_{min} = x_T e^y/(2 - x_T e^{-y})$, and $x_T = 2p_T/\sqrt{s}$. The K ≈ 2 factor accounts for higher order corrections [13]. The parton distribution density in a nucleus, $f_{a/A}(x, Q^2 , r) = t_A(r) S_{a/A}(x, r) f_{a/N}(x, Q^2)$, is assumed to be factorizable into the nuclear thickness function $t_A(r)$ (with normalization $\int d^2 r t_A(r) = A$), parton distribution in a nucleon $f_{a/N}(x, Q^2)$ and the parton shadowing factor $S_{a/A}(x, r)$ which we take the parametrization used in HIJING model [14]. Neglecting the transverse expansion, the transverse distance a parton produced at $(r, \phi)$ will travel is $\Delta L(r, \phi) = \sqrt{R_A^2 - r^2(1 - \cos^2\phi) - r \cos\phi}$.

We will use the MRS D−I parametrization of the parton distributions [15] in a nucleon. The resultant $p_T$-spectra of charged hadrons ($\pi^\pm, K^\pm$) for $pp$ and $p\bar{p}$ collisions are shown in Fig. 1 together with the experimental data [16, 17, 18] for $\sqrt{s} = 63, 200, 900$ and 1800 GeV. The calculations (dot-dashed line) from Eq. (3) with the jet fragmentation functions given by Ref. [11] agree with the experimental data remarkably well, especially at large $p_T$. However, the calculations are consistently below the experimental data at low $p_T$, where we believe particle production from soft processes is very important. To account for particle production at small $p_T$, we introduce a soft component to the particle spectra in an exponential form whose parameters are fixed by total $dN/dy$ [15]. The total $p_T$ spectra including both soft and hard component are shown in Fig. 1 as solid lines.

To calculate the $p_T$ spectrum in $AA$ collisions, one has to take into account both the parton shadowing effect and the modification of the jet fragmentation func-
Figure 2: The ratio of charged particle $p_T$ spectrum in central $Au+Au$ collisions at $\sqrt{s} = 200$ GeV over that of $pp$ collisions, normalized by the total binary nucleon-nucleon collisions in central $Au+Au$ collisions. The mean-free-path of a quark inside the medium is assumed to be 1 fm.

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The ratio $R_{AA}(p_T)$ is given by:

$$R_{AA}(p_T) = \frac{dN_{AA}/dy/d^2p_T}{\sigma_{pp}T_{AA}(0)dN_{pp}/dy/d^2p_T},$$

(4)

between the spectrum in central $AA$ and $pp$ collisions. The ratio is normalized to the effective total number of binary $pp$ collisions in a central $AA$ collision. If none of the nuclear effects (shadowing and jet quenching) are taken into account, this ratio should be unity at large transverse momentum. Shown in Fig. 2 are the results for central $Au+Au$ collisions at the RHIC energy with $dE/dx = 1, 2$ GeV/fm, respectively. As we have argued before, jet energy loss will result in the suppression of high $p_T$ particles as compared to $pp$ collisions. Therefore, the ratio at large $p_T$ in Fig. 2 is smaller than one due to the energy loss suffered by the jet partons. It, however, increases with $p_T$ because of the constant energy loss (or even some weak energy dependent energy loss). At hypothetically large $p_T$ when the total energy loss is negligible compared to the initial jet energy, the ratio should approach to one.

At small $p_T$, particles from soft interaction (or from hadronization of QGP) dominate. The ratio $R_{AA}(p_T)$ is highly sensitive to the $A$-scaling behavior of the soft particle production. Since we assumed a linear scaling for the soft particle production, the ratio should approach to $A/\sigma_{pp}T_{AA}(0) = 0.149$ at small $p_T$ for central $Au+Au$ at the RHIC energy, as shown in Fig. 2.
In this framework, one can also study the effect of energy and $A$ dependence of the energy loss and the effect of energy loss on particle production of different flavors [14].

4 Jet quenching in $\gamma+$jet events

Since hadron production at a fixed large $p_T$ comes from fragmentation of jets with different transverse energies, the suppression factor in Eq. (4) only provides the information about the effect of energy loss on jet fragmentation at an averaged value of $z = p_T/E_T^{\text{jet}}$. In order to study the modification of the fragmentation function due to energy loss, one might in principle measure the inclusive $p_T$ spectrum in the direction of a triggered jet. However, with the large background and its fluctuation due to hadrons from many other minijets and soft processes, the determination of the jet energy is almost impossible. To overcome this difficulty, we proposed the study of high $p_T$ particle spectrum in the opposite direction of a tagged direct photon [9]. Direct photons are always accompanied by a jet in the opposite transverse direction. Even though taking into account of the initial state radiation, the average energy of the jet is approximately that of the tagged photon. One can therefore relate the $p_T$ distribution of hadrons from the jet fragmentation to the fragmentation function of a jet with known initial energy and study the modification of the fragmentation function due to parton energy loss.

Let us select events which has a direct photon with transverse energy $E_\gamma^T$ in the central rapidity region, $|y| \leq \Delta y/2$, $\Delta y = 1$. For sufficiently large $E_\gamma^T$ of the photon, the rapidity distribution of the associated jet is also centered around zero rapidity with a comparable width. If the azimuthal angle of the photon is $\phi_\gamma$ and $\tilde{\phi}_\gamma = \phi_\gamma + \pi$, most of the hadrons from the jet fragmentation will fall into the kinematic region, $(|y| \leq \Delta y/2, |\phi - \tilde{\phi}_\gamma| \leq \Delta \phi/2)$, where one can take $\Delta \phi = 2$ according to the jet profile as measured in high-energy $p\bar{p}$ collisions [20]. Given the jet fragmentation functions $D_{h/a}(z)$, with $z$ the fractions of momenta of the jet carried by hadrons, one can calculate the differential $p_T$ distribution of hadrons from jet fragmentation in the kinematical region $(\Delta y, \Delta \phi)$,

$$\frac{dN_{\gamma-jet}^{pp}}{dy d^2p_T} = \sum_{a,h} r_a(E_T^\gamma) \int dE_T^{\text{jet}} g(E_T^{\text{jet}}, E_\gamma^T) \frac{D_{h/a}(p_T/E_T^{\text{jet}}) C(\Delta y, \Delta \phi)}{p_T E_T^{\text{jet}}} \frac{C(\Delta y, \Delta \phi)}{\Delta y \Delta \phi},$$

where $C(\Delta y, \Delta \phi) = \int_{|y| \leq \Delta y/2} dy \int_{|\phi - \tilde{\phi}_\gamma| \leq \Delta \phi/2} d\phi f(y, \phi - \tilde{\phi}_\gamma)$ is an overall acceptance factor and $f(y, \phi)$ is the normalized hadron profile around the jet axis. The summation is over both jet ($a$) and hadron species ($h$), and $r_a(E_T^\gamma)$ is the fractional production cross section of $a$-type jet associated with the direct photon. $C(\Delta y, \Delta \phi)$ is the acceptance factor for finding the jet fragments in the given kinematic range. We find $C(\Delta y, \Delta \phi) \approx 0.5$ at $\sqrt{s} = 200$ GeV, independent of the photon energy $E_\gamma^T$, using HIJING [14] Monte Carlo simulations for the given kinematic cuts. The normalized
Figure 3: The $E_T$ distributions of the jet, caused by the initial state radiations, accompanying a triggered direct photon with $E_T^\gamma = 10, 15$ GeV, respectively at $\sqrt{s} = 200$ GeV.

As shown in Fig. 3, is the $E_T$ distribution of the jet with a given $E_T^\gamma$ of the tagged direct photon. As we can see that the transverse energy of the jet has a wide smearing around $E_T^\gamma$ due to the initial state radiation associated with the hard processes. Because of the rapidly decrease in $E_T$ of the cross section of direct photon production, the distribution is biased toward smaller $E_T^\gamma$ than $E_T$. The average $E_T^\gamma$ is thus smaller than $E_T$. Since one only triggers a direct photon with a given $E_T^\gamma$, one should average over the $E_T$ smearing of the jet. Such a smearing is important especially for hadrons with $p_T$ comparable or larger than $E_T^\gamma$.

If we define the inclusive fragmentation function associated with a direct pho-


\begin{align}
D^\gamma(z, E_T^\gamma) &= \sum_{ah} r_a(E_T^\gamma) \int dE_T^{jet} g(E_T^{jet}, E_T^\gamma) \frac{E_T^\gamma}{E_T^{jet}} D_{h/a}(z \frac{E_T^\gamma}{E_T^{jet}}), \quad (7)
\end{align}

we can rewrite the \( p_T \) spectrum [Eq. (5)] in the opposite direction of a tagged photon as

\begin{align}
\frac{dN^{\gamma\rightarrow jet}_{pp}}{dyd^2p_T} &= \frac{D^\gamma(p_T/E_T^\gamma, E_T^\gamma)}{p_T E_T^\gamma} C(\Delta y, \Delta \phi) \Delta y \Delta \phi . \quad (8)
\end{align}

Using this equation, one can extract the inclusive jet fragmentation function, \( D^\gamma(z, E_T^\gamma) \), from the measured spectrum. Shown in Fig. 4 are the calculated \( p_T \) distributions from the fragmentation of photon-tagged jets with \( E_T^\gamma = 10, 15 \text{ GeV} \) and the underlying background from the rest of a central \( Au + Au \) collisions at the RHIC energy. The points are HIJING simulations of 10K events and solid lines are numerical results from Eqs. (3) and (5) with the fragmentation functions given by the parametrization of \( e^+e^- \) data [11]. The effect of parton energy loss is not included yet. As we can see, the spectra from jet fragmentation are significantly higher than the background at large transverse momenta. One can therefore easily extract the fragmentation function from the experimental data without much statistical errors from the subtraction of the background. One also notice that there are significant number of particles with \( p_T \) larger than the triggered photon, \( E_T^\gamma \), because of the \( E_T \) smearing of the jet caused by initial state radiations.

Consider parton energy loss in central \( AA \) collisions, we model the jet fragmentation functions as given by Eq. (4). Including the \( E_T \) smearing and averaging over the \( \gamma \)-jet production position in the transverse direction, the inclusive fragmentation function of a photon-tagged jet is,

\begin{align}
D_{AA}^\gamma(z) &= \int \frac{d^2r_T^2}{T_{AA}(0)} \sum_{ah} r_a(E_T^\gamma) \int dE_T^{jet} g(E_T^{jet}, E_T^\gamma) \frac{E_T^\gamma}{E_T^{jet}} D_{h/a}(z \frac{E_T^\gamma}{E_T^{jet}}, \Delta L), \quad (9)
\end{align}

where \( T_{AA}(0) = \int d^2r_T^2(r) \) is the overlap function of \( AA \) collisions at zero impact-parameter. We assume that jet production rate is proportional to the number of binary nucleon-nucleon collisions.

Shown in Fig. 5 are the ratios of the inclusive fragmentation function in a central \( Au + Au \) collisions with energy loss \( dE/q/dx = 1 \text{ GeV/fm} \), over the ones in \( pp \) collisions without energy loss. The enhancement of soft particle production due to induced emissions is important only at small fractional energy \( z \). The fragmentation function is suppressed for large and intermediate \( z \) due to parton energy loss. For a fixed \( dE/q/dx \), the suppression becomes less as \( E_T^\gamma \) increases. The optimal case is when the average total energy loss is significant as compared to the initial jet energy, and yet the \( p_T \) spectrum from jet fragmentation is still much larger than the underlying background. Notice that we now define \( z \) as the hadron’s fractional energy of the triggered photon. Because of the \( E_T \) smearing of the jet caused by initial state radiations, hadrons can have \( p_T \) larger than \( E_T^\gamma \). Therefore, the effective inclusive jet fragmentation function does not vanish at \( z = p_T/E_T^\gamma > 1 \).
Figure 4: The differential $p_T$ spectrum of charged particles from the fragmentation of a photon-tagged jet with $E_T^\gamma = 10, 15$ GeV and the underlying background in central $Au + Au$ collisions at $\sqrt{s} = 200$ GeV. The direct photon is restricted to $|y| \leq \Delta y/2 = 0.5$. Charged particles are limited to the same rapidity range and in the opposite direction of the photon, $|\phi - \phi_\gamma - \pi| \leq \Delta \phi/2 = 1.0$. Solid lines are from the jet fragmentation function and points are HIJING simulations of 10K events. Parton energy loss is not included yet.
As compared to our earlier results [9] where we did not take into account the $E_T$ smearing of the jet, the modification of the averaged fragmentation function due to energy loss is quite sensitive to the value of the mean-free-path for $dE_q/dx = 1$ GeV. To study the sensitivity of the suppression to the energy loss, we plot in Fig. 6 the same ratio at a fixed value of $z = 0.4$ as functions of $dE_q/dx$. The ratio in general decreases with $dE_q/dx$ as more large $p_T$ particles are suppressed when leading partons loss more energy. For small values of $dE_q/dx$, the suppression factor is more or less independent of the mean-free-path. However, for large values of $dE_q/dx \geq 1$ GeV/fm, the ratio is sensitive to the mean-free-path. One thus needs additional information or a global fit to determine both the energy loss $dE/dx$ and the mean-free-path from the experimental data.

As we discussed in the introduction, recent studies [6] of energy loss in a dense medium indicate that the energy loss per distance $dE/dx$ might be proportional to the total distance that the parton has traveled since it was produced. One way to test this experimentally is to study the suppression factor at any given $z$ value for different nucleus-nucleus collisions or for different centrality (impact parameter). Shown in Fig. 7 are the suppression factor for the jet fragmentation function at $z = 0.4$ as functions of $A^{1/3}$. In one case (dashed lines), we assume a constant energy loss $dE/dx=0.5$ GeV/fm. The suppression factor decreases almost linearly with $A^{1/3}$. In another case (solid lines), we assume $dE_q/dx = 0.2(L/fm)$ GeV/fm. The average distance a parton travels in a cylindrical system with transverse size $R_A$ is $\langle L \rangle = 0.905R_A$. We assume $R_A = 1.2A^{1/3}$ fm. We choose the coefficient in $dE_q/dx$ such that it roughly equals to 0.5 GeV/fm for $A = 20$. As we can see, the suppression factor for a distance-dependent $dE/dx$ decreases faster than the one with constant
\[ \lambda_q = 0.5 \text{ fm} \]
\[ \lambda_q = 1.0 \text{ fm} \]
\[ \lambda_q = 1.5 \text{ fm} \]

\[ E_T^\gamma = 10 \text{ GeV} \]
\[ E_T^\gamma = 15 \text{ GeV} \]

\[ A = 197 \quad z = 0.4 \]

\[ dE_q/dx (\text{GeV/fm}) \]

\[ D_A^\gamma(z)/D_P^\gamma(z) \]

Figure 6: Ratio of the inclusive fragmentation function of a photon-tagged jet with and without energy loss in central \( Au + Au \) collisions at \( z = 0.4 \) as a function of \( dE_q/dx \).

\[ dE/dx. \] Unfortunately, we have not found a unique way to extract the average total energy loss so that we can show it is proportional to \( A^{2/3} \) for the distance-dependent \( dE/dx \).

5 \( p_T \) broadening and jet profile

In the above calculation, we have assumed that the jet profile in the opposite direction of the tagged photon remains the same in \( AA \) collisions, since we used the same acceptance factor \( C(\Delta y, \Delta \phi) \). However, due to multiple scatterings suffered by the leading parton, the final jet must acquire additional acoplanarity with respect to its original transverse direction. Such a change to the jet profile could affect the acceptance factor, which will be an overall factor to the measured jet fragmentation function if we assume the jet profile to be the same for particles with different fractional energies.

To demonstrate this, we plot in Fig. 8 as the solid line the azimuthal angle distribution of \( E_T \) (within \( |y| < 0.5 \)) with respect to the opposite direction of the tagged photon with \( E_T^\gamma = 10 \text{ GeV} \). We have subtracted the background so that \( dE_T/d\phi = 0 \) at \( \phi = \pi \). The profile distribution includes both the intrinsic distribution from jet fragmentation and the effect of initial state radiation. The acceptance factor is simply the fractional area within \( |\phi| < \Delta \phi/2 \) region. The \( p_T \) broadening of jets due to multiple scatterings will broaden the profile function. Shown as the dashed line is the profile function for an average \( \Delta p_T^2 = 4 (\text{GeV}/c)^2 \) with a Gaussian distribution. It is clear that with a modest value of the \( p_T \) broadening, the acceptance factor only
Figure 7: Ratio of the inclusive fragmentation function of a photon-tagged jet with and without energy loss in central Au+Au collisions at $z = 0.4$ as a function of $A^{1/3}$. A constant energy loss, $dE_q/dx = 0.5 \text{ GeV/fm}$ (dashed lines) and $dE_q/dx = 0.2(L/\text{fm})$ GeV/fm (solid lines), which is linear in the total distance the jet has traveled, are assumed.
Figure 8: The effective profile function of the photon-tagged jet, defined as the $E_T$ distribution in azimuthal angle respect to the opposite direction of a photon with $E_T^\gamma = 10$ GeV. The dashed line assumes a Gaussian form of $p_T$ broadening for the jet with an average $\Delta p_T^2 = 4$ (GeV/c)$^2$.

changes by around 10%\textsuperscript{[21]}.

In addition, since the change of the jet profile function is related to the average $p_T$ broadening, one can combine the measurement with the measured energy loss to verify the relationship between $dE/dx$ and $\Delta p_T^2$ as suggested by recent theoretical studies \textsuperscript{[9]}.

6 Discussions

To have a feeling of the experimental feasibility of the proposed $\gamma$–jet measurement, we list in Table 1 the number of $\gamma$–jet events per year per unit rapidity and unit (GeV) $E_T$. We assume a central $Au + Au$ cross section of 125 mb with impact-parameters $b < 2$ fm. We have taken the designed luminosity of $\mathcal{L} = 2 \times 10^{26}$ cm$^{-2}$s$^{-1}$ with 100 operation days per year. As we can see, the rate for $E_T^\gamma = 15$, 20 GeV is too small to give any statistically significant measurement of the fragmentation function and its modification in $AA$ collisions. If one can increase the luminosity by a factor of 10, the numbers of events for both $E_T^\gamma = 10$ and 15 GeV are significant enough for
Table 1: Rate of direct photon production in central Au + Au collisions at $\sqrt{s} = 200$ GeV, with luminosity $L = 2 \times 10^{26}$ cm$^{-2}$s$^{-1}$ and 100 operation day per year.

| $E_T$ (GeV) | 7   | 10  | 15  | 20  |
|-------------|-----|-----|-----|-----|
| $dN^{\gamma-jet}/dydE_T$/year | 20500 | 3500 | 400 | 70  |

a reasonable determination of the fragmentation function of the photon-tagged jets.

Given increased luminosity and enough number of events, one still has to overcome the large background of $\pi^0$'s to identify the direct photons. Plotted in Fig. 9 are the production rates of direct photons (solid line) and $\pi^0$'s (dashed and dot-dashed lines). We can see that without parton energy loss, $\pi^0$ production rate is about 20 times larger than the direct photons at $p_T = 10$ GeV/$c$. Fortunately, jet quenching due to parton energy loss can significantly reduce $\pi^0$ rate at large $p_T$ as shown by the dot-dashed line. However, one still has to face $\pi^0$'s about 4 times higher than the direct photons at $p_T = 10$ GeV/$c$. At larger $p_T$, the situation improves, but one loses the production rate. Since the isolation cut method normally employed in $pp$ collisions to reduce the background to direct photons does not work anymore, the only way one can identify them has to be through improvement of detector hardwares.

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Figure 9: The inclusive $p_T$ distribution for direct photons as compared to that of $\pi^0$'s with and without parton energy loss in central $Au + Au$ collisions at $\sqrt{s} = 200$ GeV. $dE_q/dx = 1$ GeV/fm and mean-free-path $\lambda_q = 1$ fm are assumed.
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Fig. 1
A=197 \ dE_q / dx = 1 \text{ GeV/fm}

\text{Fig. 2}

\frac{D_{AA}^\gamma (z)}{D_{pp}^\gamma (z)}

E_T^\gamma = 60 \text{ GeV}

E_T^\gamma = 100 \text{ GeV}

E_T^\gamma = 20 \text{ GeV}

E_T^\gamma = 15 \text{ GeV}

\lambda_q = 1 \text{ fm}

\lambda_q = 2 \text{ fm}
Fig. 3

\[ \frac{D_{AA}^\gamma(z)}{D_{pp}^\gamma(z)} \] vs. \( \frac{dE_q}{dx} \) (GeV/fm)

- \( \lambda_q = 0.5 \text{ fm} \)
- \( \lambda_q = 1.0 \text{ fm} \)
- \( \lambda_q = 1.5 \text{ fm} \)

A=197  z=0.3

\( E_T^\gamma = 15 \text{ GeV} \)