Jet and Leading Hadron Production in High-energy Heavy-ion Collisions

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Jet tomography has become a powerful tool for the study of properties of dense matter in high-energy heavy-ion collisions. I will discuss recent progresses in the phenomenological study of jet quenching, including momentum, colliding energy and nuclear size dependence of single hadron suppression, modification of dihadron correlations and the soft hadron distribution associated with a quenched jet.

1. Introduction

The most important consequence of the discovery of the asymptotic freedom of QCD [1] is the small value of strong coupling constant $\alpha_{\rm s}$ at short distances or in hard processes involving large energy and momentum transfer. It makes the perturbative expansion in $\alpha_{\rm s}$ a reliable technique for calculations of many physical observables in these hard processes. Hadronic interaction and production often involve strong interaction at long distance, which is not calculable within the framework of the perturbative expansion. However, it has been proven to the leading power correction $(1/Q^2)$ that the cross section can be factorized into short-distance parts calculable in perturbative QCD (pQCD) and non-perturbative long distance parts [2]. These long distance parts can be expressed as matrix elements that are universal. Therefore, they can be measured in one process and used in another process; therein lies the predictive power of pQCD. It has been extremely successful in the study of hard processes, from electron-positron annihilation, to deeply inelastic scattering (DIS) of leptons and nucleons, Drell-Yan dilepton production and large transverse momentum jet production in hadronic collisions.

Hard processes can also happen in high-energy nucleus-nucleus collisions. They include production of large transverse momentum jets, direct photons and dileptons with large invariant-mass and heavy quarks. Since hard processes happen on a short time scale in the earliest stage of high-energy heavy-ion collisions, they can probe the bulk matter that is formed shortly after the collision. The pQCD parton model serves as a reliable and tested framework for the study of these hard probes. In this talk, I will focus on the physics of jet propagation in the dense medium and recent progresses in the phenomenological study of jet and high p_T hadron production in heavy-ion collisions.

The study of jet production in heavy-ion collisions exploits the attenuation of parton jets or jet quenching during their propagation in dense medium. Such an idea was first proposed by Bjorken [3] to study the space-time structure of high-energy hadron-hadron collisions via elastic energy loss of partons in medium [3]. But it was soon realized that elastic energy loss may be very small relative to the radiative energy loss induced by multiple scatterings [4]. The effect of jet attenuation in medium and its utilization as a probe was taken more seriously only after a Monte Carlo study within the HIJING model [5] that demonstrated significant suppression of high p_T single inclusive hadron spectra due to jet quenching. Such large suppression can be easily measured in experiments without jet reconstruction and detailed study of medium modification of jet structure. It is quite amazing how well the calculated high p_T hadron suppression with a very crude estimate (or rather a guess) and simulation of average parton energy loss, $dE/dx = 1$ GeV/fm, agrees qualitatively with the first RHIC data in central $Au + Au$ collisions [6]. However, later theoretical studies show that radiative parton energy loss involves quantum field treatment of induced radiation and the non-Abelian Landau-Pomeranchuk-Migdal interference effects in QCD [7]. The non-Abelian energy loss depends on both the local parton density and the total distance of the parton propagation.

It was realized from the beginning that one has to understand the cold-nuclear effects on high p_T hadron spectra in p+A collisions in order to extract the genuine suppression caused by the hot medium in heavy-ion collisions [5]. These cold-nuclear effects include nuclear modification of the parton distributions in nuclei and change of high p_T hadron spectra caused by initial and final state multiple scattering in cold nuclei. Nuclear shadowing or depletion of effective parton distributions in nuclei is limited to small fractional momentum $x < 0.1$ and thus to low p_T hadron spectra $p_T < 2 \sim 4$ GeV/c at $\sqrt{s} = 200$ GeV [5]. Multiple initial parton scattering and the consequent intrinsic transverse momentum broadening will partially compensate the effect of nuclear shadowing. The final hadron spectra in $p + A$ collisions at mid-rapidity were predicted to be enhanced slightly, known as the Cronin effect, at intermediate p_T and then become nearly identical to that in $p + p$ collisions at large p_T [8]. These predictions were indeed verified by experimental data in $d + Au$ collisions at RHIC [9]. Therefore, the final state interaction and jet quenching have been established as the true cause of the observed suppression of high p_T hadron spectra in $Au + Au$ collisions. Since parton energy loss depends on both the local parton density and total propagation length, it will depend on the azimuthal angle of the jet propagation relative to the reaction plane in non-central heavy-ion collisions. Such azimuthal angle dependence was predicted [10] to give rise to an azimuthal anisotropy of the suppressed large p_T hadron spectra which was also observed at RHIC [11]. The most striking consequence of jet quenching observed so far is the suppression of the away-side jets in two-hadron correlation measurements [12], providing a clear illustration of the jet quenching picture in heavy-ion collisions.

2. Jet quenching and modified jet fragmentation

Though jet quenching can be intuitively related to parton energy loss, the experimentally measurable consequences can only be found in the modification of the final hadron spectra from jet fragmentation in medium relative to that in vacuum. In addition, there are many partonic processes, such as quark-anti-quark annihilation, in which the leading parton can lose their identities (flavor or gluon versus quark). In these processes, the concept of parton energy loss become very ambiguous. Furthermore, hard processes are normally followed by final state radiations with a short time even in the vacuum. The

DGLAP evolution and induced radiation due to final state multiple scattering should be considered together in the same framework of jet fragmentation.

Within a framework of twist-expansion in the collinear factorized parton model, one can study such medium modification of the jet fragmentation function via modified DGLAP evolution equations [13]. For quark propagation in a cold nuclear medium in deeply inelastic lepton scattering off a nuclear target, the modified fragmentation function,

$$\widetilde{D}_{q\to h}(z_h,Q^2) \equiv D_{q\to h}(z_h,Q^2) + \int_0^{Q^2} \frac{d\ell_T^2}{\ell_T^2}\frac{\alpha_s}{2\pi}\int_{z_h}^1 \frac{dz}{z}\left[\Delta\gamma_{q\to qg}(z,x,x_L,\ell_T^2)\, D_{q\to h}(z_h/z) + \Delta\gamma_{q\to qg}(1-z,x,x_L,\ell_T^2) D_{g\to h}(z_h/z)\right], \quad (1)$$

has the same form as the DGLAP correction in vacuum, which determines the evolution of $D_{q\to h}(z_h,Q^2)$ and $D_{g\to h}(z_h,Q^2)$ as the leading-twist quark and gluon fragmentation functions. The difference from the vacuum DGLAP evolution lies in the modified splitting functions

$$\Delta\gamma_{q\to qg}(z,x,x_L,\ell_T^2) = \frac{1+z^2}{(1-z)_+} T^A_{qg}(x,x_L)\frac{2\pi\alpha_s C_A}{(\ell_T^2+\langle k_T^2\rangle)N_c f_q^A(x,\mu_I^2)} + (\text{virtual corr.})\,, \quad (2)$$

which depends on the properties of the medium through the twist-four parton matrix element of the nucleus,

$$T^A_{qg}(x,x_L) = \int\frac{dy^-}{2\pi}\, dy_1^- dy_2^- e^{i(x+x_L)p^+y^-}(1-e^{-ix_Lp^+y_2^-})(1-e^{-ix_Lp^+(y^- - y_1^-)}) \times\frac{1}{2}\langle A|\bar{\psi}_q(0)\,\gamma^+\, F_\sigma^{\ +}(y_2^-)\, F^{+\sigma}(y_1^-)\,\psi_q(y^-)|A\rangle\theta(-y_2^-)\theta(y^- - y_1^-) \ . \quad (3)$$

Here, the fractional momentum $x=x_B$ is carried by the initial quark and $x_L=\ell_T^2/2p^+q^-z(1-z)$ is the additional momentum fraction required for induced gluon radiation. The dipole-like structure in the above twist-four parton matrix element is a result of the LPM interference in gluon bremsstrahlung. Assuming a factorized form of the twist-four matrix element, the above modified fragmentation function can describe the suppression of leading hadron spectra in HERMES data of DIS [14] very well, including the quadratic nuclear size dependence.

One can quantify the modification of the fragmentation by the quark energy loss which can be defined as the momentum fraction carried by the radiated gluon,

$$\langle\Delta z_g\rangle = \int_0^{Q^2}\frac{d\ell_T^2}{\ell_T^2}\int_0^1 dz\frac{\alpha_s}{2\pi} z\,\Delta\gamma_{q\to gq}(z,x_B,x_L,\ell_T^2). \quad (4)$$

To extend the study of modified fragmentation functions to jets in heavy-ion collisions, one can assume $\langle k_T^2\rangle\approx\mu^2$ (the Debye screening mass) and a gluon density profile $\rho(y)=(\tau_0/\tau)\theta(R_A-y)\rho_0$ for a 1-dimensional expanding system. Since the initial jet production rate is independent of the final gluon density, which can be related to the parton-gluon scattering cross section $[\alpha_s x_T G(x_T)\simeq(N_c/2\pi^2)\mu^2\sigma_g]$, one has then

$$\frac{\alpha_s T^A_{qg}(x_B,x_L)}{f_q^A(x_B)} \approx \frac{N_c}{\pi}\mu^2\int dy\sigma_g\rho(y)[1-\cos(y/\tau_f)], \quad (5)$$

where $\tau_f = 2Ez(1-z)/\ell_T^2$ is the gluon formation time. In the limit of high initial jet energy E, the total energy loss becomes [14]

$$\langle \Delta E \rangle \approx \pi C_a C_A \alpha_s^3 \int_{\tau_0}^{R_A} d\tau (\tau - \tau_0) \rho(\tau, \vec{r}_0 + \vec{n}(\tau - \tau_0)) \ln \frac{2E}{\tau \mu^2}, \tag{6}$$

where $\sigma_g \approx C_a 2\pi\alpha_s^2/\mu^2$ (C_a=1 for qg and 9/4 for gg scattering) is assumed. Taking into account the kinematic limits in the integration, the total energy loss has a strong energy dependence for finite values of E [15].

3. Jet quenching in heavy-ion collisions

Working in the same framework of twist expansion in the collinear factorized parton model, one can similarly obtain the single inclusive hadron spectra at high p_T [8],

$$\begin{aligned} \frac{d\sigma^h_{AA}}{dyd^2p_T} &= K \sum_{abcd} \int d^2b d^2r dx_a dx_b d^2k_{aT} d^2k_{bT} t_A(r) t_A(|\mathbf{b}-\mathbf{r}|) g_A(k_{aT}, r) g_A(k_{bT}, |\mathbf{b}-\mathbf{r}|) \\ &\times f_{a/A}(x_a, Q^2, r) f_{b/A}(x_b, Q^2, |\mathbf{b}-\mathbf{r}|) \frac{d\sigma}{d\hat{t}}(ab \to cd) \frac{D'_{h/c}(z_c, Q^2, \Delta E_c)}{\pi z_c}, \end{aligned} \tag{7}$$

with medium modified fragmentation functions $\widetilde{D}_{h/c}$. Here, $z_c = p_T/p_{Tc}$, $y = y_c$, $\sigma(ab \to cd)$ are elementary parton scattering cross sections and $t_A(b)$ is the nuclear thickness function normalized to $\int d^2b t_A(b) = A$. The $K \approx 1.5$–2 factor is used to account for higher order pQCD corrections. For simplification, many studies assume the modified fragmentation functions as given by the vacuum ones with the fractional momentum rescaled by $1/(1-\Delta z)$. Such effective description is found to reproduce the full calculation very well, but only when $\Delta z = \Delta E_c/E$ is set to be $\Delta z \approx 0.6\langle z_g \rangle$. One can similarly calculate back-to-back dihadron spectra in the same parton model, assuming two jets fragment independently as described by medium modified fragmentation functions [16]. Such a parton model with modified fragmentation functions is the framework for many phenomenological studies of jet quenching in heavy-ion collisions [8,17]. It can describe well the suppression of the single hadron spectra, back-to-back hadron correlation and azimuthal angle anisotropy [16], which are three different consequences of jet quenching. Since jet quenching depends on the initial gluon density of the hot medium which is not known within the parton model, one has to fit the data on the suppression of single hadron spectra in the most central $Au + Au$ collisions at a given colliding energy, *e.g.*, $\sqrt{s} = 200$ GeV. Then the centrality dependence, the energy dependence, the back-side suppression and the azimuthal anisotropy are all predictions within the parton model. Combining measurements of the above three different effects of jet quenching and compare with the results from jet quenching in deeply inelastic $e + A$ collisions, one can conclude that the initial gluon density (at initial time $\tau_0 = 0.2$ fm/c) reached in central $Au + Au$ collisions at $\sqrt{s} = 200$ GeV is about 30 times higher than in a cold Au nuclei [16,14].

Since parton degradation in medium depends both on the local gluon density and the propagation length (which is non-linear due to non-Abelian LPM interference effect), the centrality dependence of the single inclusive hadron suppression reflects a combination of these two dependencies. One can extend the study of the density and length dependence by varying the nuclear size at a fixed energy. Shown in Fig. 1 are the parton model

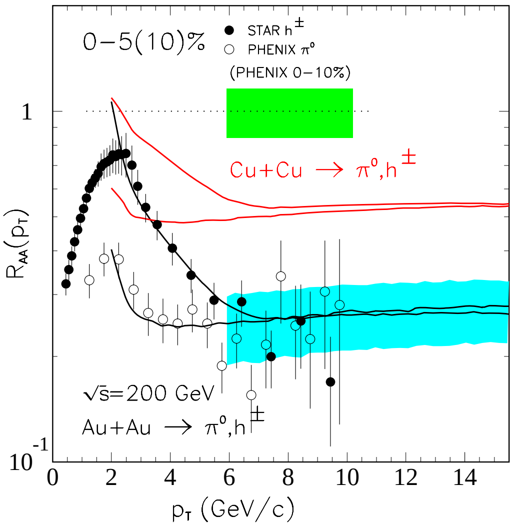

Figure 1. Hadron suppression factor $R_{AA}(p_T)$ for the most central (0-10%) $Au + Au$ and $Cu + Cu$ collisions at $\sqrt{s} =$ 200 GeV.

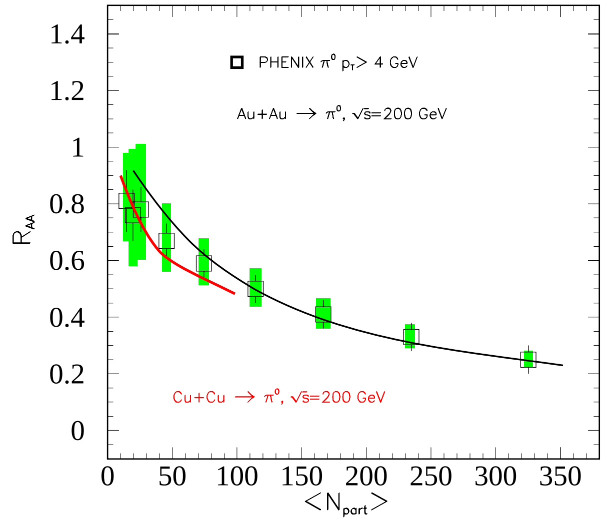

Figure 2. The hadron suppression factor $R_{AA}(p_T)$ at fixed p_T and $\sqrt{s} = 200$ GeV as a function of $N_{\rm part}$ in $Au + Au$ and $Cu + Cu$ collisions.

calculations of the suppression factor $R_{AA}(p_T)$ for the most 0-10% central $Cu + Cu$ collisions at $\sqrt{s} = 200$ GeV together with the calculation and experimental data of central $Au + Au$ collisions. As expected, the suppression is very similar to semi-peripheral (30-40%) $Au + Au$ collisions with the same p_T dependence. This is exactly what is observed by the experiments on $Cu + Cu$ collisions [18,19]. In principle, the suppression factor should only be a function of the total energy loss. At fixed energy and p_T, it is proportional to the path integral in Eq. (6). As shown in Fig. 2 the suppression factor at fixed $p_T \approx 4$ GeV and energy $\sqrt{s} = 200$ GeV is only approximately a function of $N_{\rm part}$. The new data on $Cu + Cu$ collisions indeed indicate small but finite deviation from $N_{\rm part}$ scaling [18,19]. However, the statistical and most importantly the systematic errors are still too big to quantify the small deviation.

The p_T dependence of the hadron suppression factor is sensitive to many aspects of the jet production and parton energy loss. Because of trigger bias, the initial jet energy for fixed hadron p_T varies with colliding energy $\sqrt{s}$ because the shape of the jet spectra change dramatically from the SPS to RHIC and LHC energies. At the same time, the parton energy loss has also a strong energy dependence on the initial jet energy in the kinematic regime of current experimental data. For given amount of parton energy loss, the hadron suppression factor is also sensitive to the slope of the initial jet spectra which changes with $x_T = 2p_T/\sqrt{s}$. In the kinematic regime $p_T \sim 10$ GeV, jet spectra at low energies are very steep at the edge of kinematic limit. Any given parton energy loss will lead to the increase of suppression at larger p_T. In this case $R_{AA}(p_T)$, shown in Fig. 3, decreases with p_T as observed in experiments at $\sqrt{s} = 63$ GeV [18,19]. At $\sqrt{s} = 200$ GeV, such effect due to the shape of jet spectra is compensated by the energy dependence of the parton energy loss. The resulting $R_{AA}(p_T)$ is almost independent of p_T in the range of $p_T = 5 - 20$ GeV. At LHC energy, $p_T \sim 50$ GeV is far away from the kinematic limit. Jet spectra have a nice power-law behavior. In this case, the suppression factor increases slowly with p_T. At the SPS energy, initial parton fractional momentum is quite large ($x_T \sim 0.6$ for $p_T = 5$ GeV). The nucleon Fermi motion will cause the effective parton distribution to rise (EMC effect) and thus the modification factor of hadron spectra also increase with p_T, even with jet quenching.

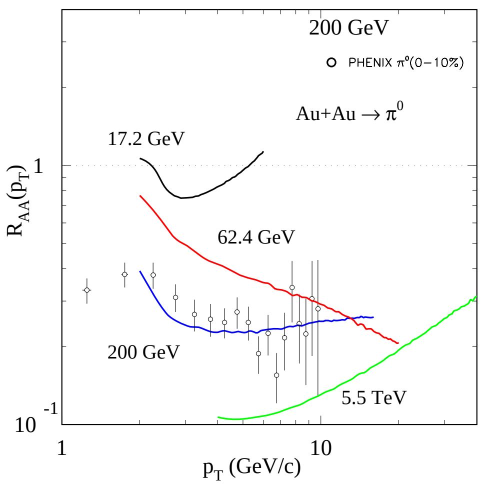

Figure 3. Hadron suppression factor $R_{AA}(p_T)$ for the most central (0-10%) $Au + Au$ collisions at different colliding energies.

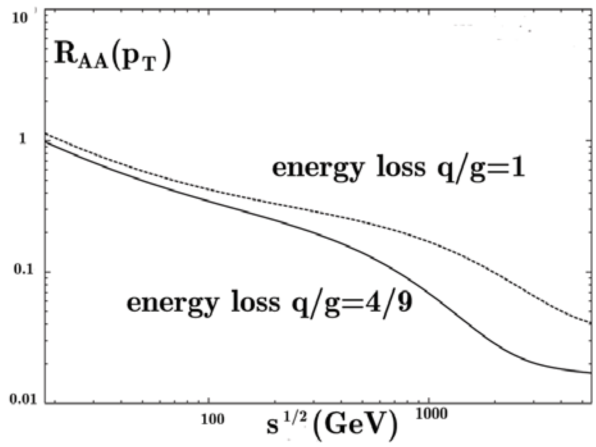

Figure 4. Nuclear modification factor R_{AuAu} for neutral pions as function of collision energy at fixed $p_T = 6$ GeV in central collisions (0-10%) with both the QCD and a non-QCD energy loss.

For fixed p_T and centrality in given $A + A$ collisions, one can also use the energy dependence to study a non-Abelian feature of the parton energy loss, *i.e.*, the dependence on the color representation of the propagating parton. The energy loss for a gluon is 9/4 times larger than a quark. In this case, one can 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 more energy than quarks, the energy dependence of the large (but 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 sensitivity can be illustrated by comparing hadron suppression factors with two different parton energy losses: one for the QCD case where the energy loss for a gluon is 9/4 times as large as that 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 $\langle\frac{\Delta L}{\lambda}\rangle_g/\langle\frac{\Delta L}{\lambda}\rangle_q = 9/4$ in the QCD case. For the non-QCD case we are considering, the above ratio is set to one. Shown in Fig. 4 are the calculated R_{AA} for neutral pions at fixed $p_T = 6$ GeV in central $Au + Au$ collisions as a function of $\sqrt{s}$ from 20 AGeV to 5500 AGeV with both the QCD and non-QCD energy loss [20]. In these calculations, the initial gluon number density and partons' mean-free path is fixed to fit the overall hadron suppression in the most central $Au + Au$ collisions at $\sqrt{s} = 200$ GeV. For any other energy, the initial gluon number is assumed to be proportional to the final measured total charged hadron multiplicity per unit rapidity. 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 be equal to the quark. 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. As pointed out in Ref. [21], the existing data from SPS ($\sqrt{s} = 17$ GeV) and RHIC ($\sqrt{s} = 63$, 130 and 200 GeV) already favor the case of non-Abelian energy loss, which differs from the Abelian case by almost 50% at $\sqrt{s} = 200$ GeV. The difference will grow to about a factor 2 at the LHC energy.

4. Two-hadron correlations

To reduce the sensitivity of jet quenching study to the underlying jet spectra in the suppression of single hadron spectra, one can measure two-hadron correlations at large p_T [12]. It is accomplished by measuring the spectra of hadrons associated with a triggered high p_T hadron. Such two-hadron correlation is effectively the ratio of dihadron and single hadron spectra,

$$D_{AA}(z_T, \phi, p_T^{\rm trig}) = p_T^{\rm trig} \frac{d\sigma_{AA}/dp_T^{\rm asso} dp_T^{\rm trig}}{d\sigma_{AA}/dp_T^{\rm trig}}, \quad (8)$$

where $z_T = p_T^{\rm asso}/p_T^{\rm trig}$ and ϕ is the azimuthal angle between the triggered and associated hadron.

For $\phi < \pi//4$, the same-side two-hadron correlations are determined by the ratio of dihadron and single hadron fragmentation functions. The dihadron fragmentation functions in terms of the overlapping matrix between parton field operators and the final hadron states have been defined and their DGLAP evolution equations have been derived recently [22], which are similar to that of single hadron fragmentation functions. However, there are extra contributions that are proportional to the convolution of two single hadron fragmentation functions. These correspond to independent fragmentation of both daughter partons after the parton split in the radiative processes. Medium modification to the dihadron fragmentation functions due to induced radiation was found to have the identical form as the DGLAP evolution equations [23]. These medium modifications depend on the same gluon correlation functions as in the modification to the single hadron fragmentation functions. Therefore, in the numerical calculation of the medium modification of dihadron fragmentation functions, there are no additional parameters involved. The predicted results for jet quenching in DIS are in good agreement with HERMES data [23]. The nuclear modification is found to manifest mostly in the single hadron fragmentation functions. Since dihadron fragmentation functions already contain the information of single hadron fragmentation function, the modification to the remaining correlated distribution, $D_q^{h_1h_2}(z_1, z_2)/D_q^{h_1}(z_1)$ is very small. This explains why the same-side two-hadron correlation in central heavy-ion collisions remains approximately the same as in $p + p$ collisions [12]. However, trigger bias in heavy-ion collisions could lead to some apparent change of dihadron correlations [23].

For $\phi \sim \pi$, away-side two-hadron spectra are proportional to the product of two single fragmentation functions. Therefore, the away-side two-hadron correlations in Eq. (8) essentially reflect the medium modification of the single fragmentation function in the opposite direction of the triggered jet. For fixed $p_T^{\rm trig}$ and $p_T^{\rm asso}$, the initial transverse momentum of the away-side jet can fluctuate. Such fluctuation leads to a rather flat medium modification of the correlation function, $I_AA(z_T) = D_{AA}(z_T)/D_{pp}(z_T)$ for large values of z_T and it only starts to increase at small $z_T < 0.2$ [16]. Recent measurements from STAR have qualitatively confirmed this feature [24]. For large $p_T^{\rm trig}$, the correlation $I_{AA}(z_T)$ scales approximately with z_T. It will be useful to verify this experimentally.

5. Jet-induced collective excitation

The suppression of high p_T single inclusive hadron spectra and two-hadron correlations can be attributed to modification of jet fragmentation in medium via induced gluon bremsstrahlung. However, most of the studies have neglected the interaction of the radiated gluons with the thermal medium. Such interaction could potentially affect the spectra of soft hadrons associated with a jet. Recent experimental studies of angular correlations of soft hadrons with respect to a quenched jet indeed have revealed a peculiar pattern [25]. In central $Au + Au$ collisions, soft hadrons associated with a quenched jet (in the opposite direction of the triggered hadron) are peaked at a finite angle $\Delta\phi \sim 1$ away from the jet, whereas, they peak in the direction of the jet in peripheral $Au + Au$ or $p + p$ collisions. This observation has lead to suggestions of different scenarios for the interaction between soft partons and the thermal medium as the leading parton propagates through the dense medium.

In the most simplified scenario, one can assume that the soft partons radiated from the leading jet (or equivalently the recoil from elastic scattering) strongly interact with the medium and are immediately thermalized. The energy deposited by the jet through such strong interaction will then propagate through the medium as a sound wave. Because the sound velocity c_s is smaller than the (light) velocity of the leading massless parton, a shock wave would eventually develop [26]. The wave front of the shock will have a Mach cone angle $\cos\theta_M = c_s$. Such formation of shock wave or Mach cone has been demonstrated in hydrodynamical simulations [27].

In the limit of slow thermalization of the radiated partons, a ring structure in the final particle distribution can also be formed via Cerenkov gluon radiation if the soft gluon can develop a space-like dispersion relation through interaction with the thermal medium [28]. Such space-like dispersion is possible if the strongly interactive quark-gluon plasma has many colored resonant structures or partonic bound states. The transitional scattering between a soft gluon and two bound states with different masses has been shown to lead to a space-like dispersion for the soft gluon [29]. However, the total energy loss caused by Cerenkov gluon radiation is very small as compared to radiative energy loss induced by multiple parton scattering. If one employs a space-like dispersion relation for soft gluons in the calculation of induced gluon bremsstrahlung, the LPM interference will produce a final gluon spectra with a peak at a finite angle [30]

$$\cos^2\theta_c = z + \frac{1-z}{\epsilon(\ell)}, \tag{9}$$

which is determined by gluon's dielectric constant $\epsilon(\ell)$, where z is the gluon's fractional energy. In the soft radiation limit $z \sim 0$, this corresponds exactly to the angle of classical Cerenkov radiation. Such Cerenkov-like bremsstrahlung can induce a large energy loss. For a large dielectric constant $\epsilon \gg 1 + 2/z^2LE$, the corresponding total radiative parton energy loss is about twice that from normal gluon bremsstrahlung [30]. The unique feature of Cerenkov-like bremsstrahlung is that the Cerenkov cone size decreases with the momentum of the soft gluon. On the other hand, Mach cone of sonic shock wave is independent of the particle momentum.

6. Summary

In summary, the discovery and detailed study of jet quenching in high-energy heavy-ion collisions at RHIC have provided strong evidence that a strongly interactive quark-gluon plasma has been formed in the central $Au + Au$ collisions. Such sQGP is opaque to energetic parton jets. The initial gluon density at $\tau_0 = 0.2$ fm/c is estimated to be about 30 times higher than in a cold nuclear matter and the energy density about 100 times higher. The peculiar angular distribution of soft hadrons relative to the quenched jet could be an indication of sonic shock wave or Cerenkov-like gluon bremsstrahlung. If confirmed by further experimental test, this could provide another powerful information about properties of the sQGP.

Combined with many other aspects of jet quenching, jet tomography has become a useful and powerful tool to study the properties of dense matter in heavy-ion collisions. With accumulation of data and development of new experimental analysis techniques, the study of jet quenching via direct γ tagged jet will become available. Such study would be ideal because there will not be complicated trigger bias effect as compared to high-p_T hadron triggering. The suppression is expected to be similar to the single hadron spectra [31], except that one knows better the initial jet energy and therefore can study the energy dependence of the jet quenching. Two-hadron correlation in γ-tagged jet is also better due to absence of correlated background because there is no trigger-biased correlation with the reaction plane.

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