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Jet quenching at intermediate RHIC energies

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Abstract

The final state energy loss of fast partons penetrating a longitudinally expanding quark–gluon plasma of effective gluon rapidity density $dN^g/dy = 650$–800 is evaluated and incorporated together with the multiple initial state Cronin scattering in the lowest order perturbative QCD hadron production formalism. Predictions for the neutral pion attenuation in central Au + Au collisions at the intermediate RHIC energy of $\sqrt{s_{NN}} = 62$ GeV relative to the binary collision scaled $p + p$ result are given. The quenching is found to be a factor of 2–3 with a moderate transverse momentum dependence and the attenuation of the away side di-hadron correlation function is estimated to be 3–5-fold.

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1. Introduction

Recent combined experimental measurements of the nuclear modification to the moderate and large transverse momentum hadron production in Au + Au [1,2] and d + Au [3] reactions have provided strong evidence in support of the dominance of multiple final state interactions [4] over the initial state Cronin scattering [5] and possible nuclear wavefunction effects [6] in relativistic heavy ion collisions. These findings pave the way for detailed studies of derivative jet quenching observables [7,8] such as the high p_T azimuthal anisotropy [9], the broadening and disappearance of di-jet correlations [10], extensions of the correlation analysis with respect to the reaction plane [11] and possibilities for full jet and lost energy reconstruction [12]. The theory and phenomenology of multiparton dynamics in ultra-dense nuclear matter will, however, remain incomplete without a thorough investigation of the

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center of mass energy and system size dependence [13,14] of medium induced non-Abelian gluon bremsstrahlung and the corresponding hadron attenuation. Possible onset of pion suppression at the SPS energy of $\sqrt{s_{NN}} = 17$ GeV has been recently discussed [15]. The intermediate $\sqrt{s_{NN}} = 62$ GeV RHIC run is the next critical step in mapping out the elastic and inelastic scattering properties of a quark–gluon plasma via jet tomography [8,16].

The magnitude of the energy loss driven nuclear quenching of moderate and high p_T pions is controlled by the soft parton rapidity density [8,16]. To relate the experimentally measured $dN^{\mathrm{ch}}/d\eta$ [17] to the effective dN^g/dy we use $|d\eta/dy| \approx 1.2$ at $\eta = y = 0$, ignoring the $\sqrt{s_{NN}}$ dependent changes in particle composition. The estimated effective

$$\frac{dN^g}{dy} \approx \frac{3}{2}\left|\frac{d\eta}{dy}\right|\frac{dN^{\mathrm{ch}}}{d\eta} \tag{1}$$

follows from the isospin symmetry of strong interactions and the approximate parton–hadron duality [18]. Straightforward application of Eq. (1) for central, $N_{\mathrm{part}} = 340$, Au + Au collisions and $dN^{\mathrm{ch}}/d\eta$ constrained from the data [17] yields $dN^g/dy \approx 550, 850, 1150$ at SPS, the intermediate and maximum RHIC energies, respectively. Such rapidity densities will likely be compatible with a soft participant scaling phenomenology [19]. The predicted $\sqrt{s_{NN}} = 200$ GeV pion quenching [13] was found to be in good agreement with the experimental measurements [1] but the original SPS π^0 data [20] would tend to disfavor parton energy loss in variance with the theoretical expectations [4]. A more recent analysis of the low energy $p + p$ baseline cross section [15] shows that the WA98 [20] and CERES [21] data are not inconsistent with jet quenching calculations with $dN^g/dy = 400$ when the Cronin effect is taken into account [13]. The extracted effective gluon rapidity density, however, still falls short of the expectation from the measured hadron multiplicities [17]. This deviation can be related to the uncertainties in the perturbative calculation in a theory with strong coupling and a non-negligible quark contribution[1] to the bulk soft partons at $\sqrt{s_{NN}} = 17$ GeV at midrapidity. Therefore, the importance of studying the sensitivity of the observable spectral modification to a range of parton densities at a fixed center of mass energy should not be underestimated in theoretical calculations.

The purpose of this Letter is to investigate the interplay of the initial state multiple Cronin scattering and the final state energy loss in a thermalized QCD media with effective $dN^g/dy = 650$–800 in moderate and large transverse momentum hadron production. Section 2 outlines the method for evaluating the medium induced gluon bremsstrahlung off fast partons in dense, dynamically expanding and finite nuclear matter. Section 3 presents the calculated depletion of pion multiplicities in central Au + Au reactions at $\sqrt{s_{NN}} = 62$ GeV relative to the binary collision scaled $p + p$ baseline. Summary and discussion is given in Section 4.

2. Medium induced gluon bremsstrahlung in finite dynamical plasmas

The full solution for the medium induced gluon radiation off jets produced in a hard collisions at early times $\tau_{\mathrm{jet}} \simeq 1/E$ inside a nuclear medium of length L can be obtained to all orders in the correlations between the multiple scattering centers via the reaction operator approach [22]. Other existing techniques have been reviewed in [4]. The double differential bremsstrahlung intensity for gluons with momentum $k = [xp^+, \mathbf{k}^2/xp^+, \mathbf{k}]$ resulting from the sequential interactions of a fast parton with momentum $p = [p^+, 0, \mathbf{0}]$ can be written as

$$x\frac{dN_g}{dx\,d^2\mathbf{k}} = \sum_{n=1}^{\infty} x\frac{dN_g^{(n)}}{dx\,d^2\mathbf{k}}$$

[1] Hard jets have a $C_F/C_A = 4/9$ smaller elastic scattering cross section with quarks relative to gluons. If nuclear matter is approximated by effective gluons, the inferred density will be correspondingly smaller while still yielding the same mean σ_{el} and opacity.

$$= \sum_{n=1}^{\infty} \frac{C_R \alpha_s}{\pi^2} \prod_{i=1}^{n} \int_0^{L-\sum_{a=1}^{i-1} \Delta z_a} \frac{d\Delta z_i}{\lambda_g(i)} \int d^2\mathbf{q}_i \left[\sigma_{\rm el}^{-1}(i) \frac{d\sigma_{\rm el}(i)}{d^2\mathbf{q}_i} - \delta^2(\mathbf{q}_i) \right]$$
$$\times \left(-2\mathbf{C}_{(1,\ldots,n)} \cdot \sum_{m=1}^{n} \mathbf{B}_{(m+1,\ldots,n)(m,\ldots,n)} \left[\cos\left(\sum_{k=2}^{m} \omega_{(k,\ldots,n)} \Delta z_k \right) - \cos\left(\sum_{k=1}^{m} \omega_{(k,\ldots,n)} \Delta z_k \right) \right] \right), \tag{2}$$

where $\sum_2^1 \equiv 0$ is understood. In the small angle eikonal limit $x = k^+/p^+ \approx \omega/E$. In Eq. (2) the color current propagators are denoted by

$$\mathbf{C}_{(m,\ldots,n)} = \frac{1}{2} \nabla_{\mathbf{k}} \ln(\mathbf{k} - \mathbf{q}_m - \cdots - \mathbf{q}_n)^2, \qquad \mathbf{B}_{(m+1,\ldots,n)(m,\ldots,n)} = \mathbf{C}_{(m+1,\ldots,n)} - \mathbf{C}_{(m,\ldots,n)}. \tag{3}$$

The momentum transfers $\mathbf{q}_i$ are distributed according to a normalized elastic differential cross section,

$$\sigma_{\rm el}(i)^{-1} \frac{d\sigma_{\rm el}(i)}{d^2\mathbf{q}_i} = \frac{\mu^2(i)}{\pi(\mathbf{q}_i^2 + \mu^2(i))^2}, \tag{4}$$

which models scattering by soft partons with a thermally generated Debye screening mass $\mu(i)$. It has been shown [22,23] via the cancellation of direct and virtual diagrams that in the eikonal limit only the gluon mean free path $\lambda_g(i)$ enters the medium-induced bremsstrahlung spectrum in Eq. (2). For gluon-dominated bulk soft matter $\sigma_{\rm el}(i) \approx \frac{9}{2}\pi\alpha_s^2/\mu^2(i)$ and $\lambda_g(i) = 1/\sigma_{\rm el}(i)\rho(i)$. The characteristic path length dependence of the non-Abelian energy loss in Eq. (2) comes from the interference phases and is differentially controlled by the inverse formation times,

$$\omega_{(m,\ldots,n)} = \frac{(\mathbf{k} - \mathbf{q}_m - \cdots - \mathbf{q}_n)^2}{2xE}, \tag{5}$$

and the separations of the subsequent scattering centers $\Delta z_k = z_k - z_{k-1}$. It is the non-Abelian analogue of the Landau–Pomeranchuk–Migdal destructive interference effect in QED [24].

For the case of local thermal equilibrium we relate all dimensional scales in the problem to the temperature of the medium $T(i)$, for example, $\mu^2(i) = 4\pi\alpha_s T(i)^2$ and the elastic scattering cross section indicated above. The gluon radiative spectrum is evaluated numerically in an ideal $(1+1)$D Bjorken expanding plasma [25], $\rho(i) = \rho_0(\tau_0/\tau_i)^\alpha$, $\epsilon_i = \epsilon_0(\tau_0/\tau_i)^{\alpha+v_s^2}$ with $\alpha = 1$ and v_s being the speed of sound. The scaling with proper time of the temperature, the Debye screening mass $\mu(i)$ and the gluon mean free path $\lambda_g(i)$ naturally follow. To leading power, the initial equilibration time τ_0 cancels in the evaluation of the bremsstrahlung integrals [4] since

$$\int_0^L d\tau_i\, \rho(i)\tau_i = \frac{1}{A_\perp} \frac{dN^g}{dy} L. \tag{6}$$

In Eq. (6) $A_\perp$ is the transverse size of the medium and dN^g/dy is the relevant physical quantity that controls the attenuation of the final state partonic flux.

Transverse, $\beta_T \neq 0$, expansion was not explicitly included in the energy loss calculation. Hence, we first clarify its impact on the medium induced non-Abelian bremsstrahlung. Numerically, the contributions to the radiative spectrum in Eq. (2) have to be evaluated in the background of the dynamically evolving soft parton multiplicity. In practice, for realistic $(3+1)$D hydrodynamic simulation only the dominant $n = 1$ term [22] has been considered in the mean energy loss approximation [26]. Analytic treatments of transverse expansion must therefore provide important guidance to its effect on the radiative spectra and $\langle \Delta E \rangle$.

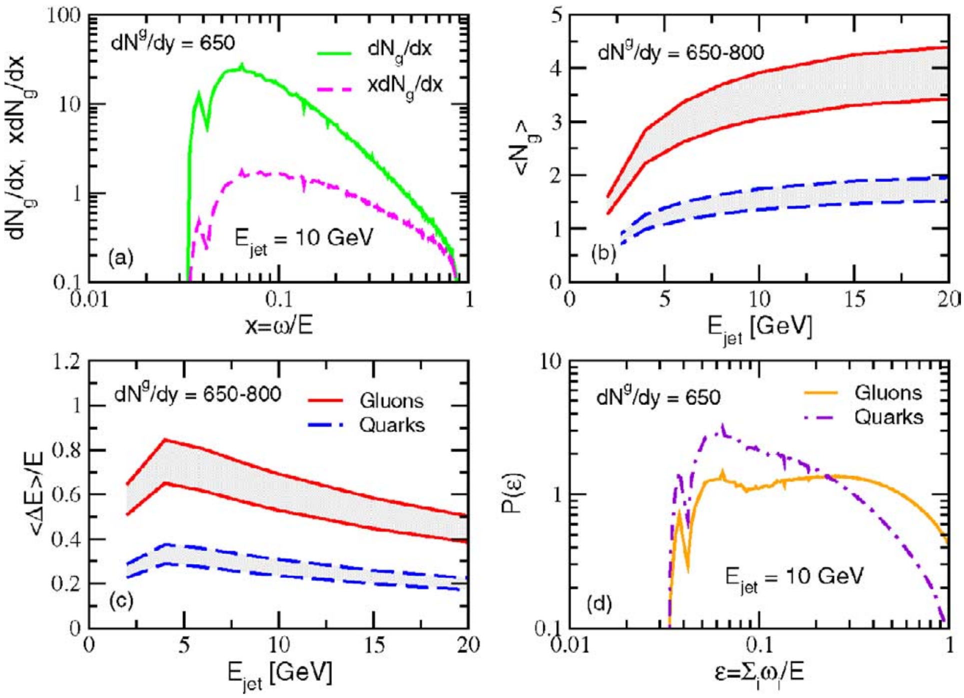

Fig. 1. (a) Medium-induced single inclusive gluon spectrum dN_g/dx and fractional intensity $x\,dN_g/dx$ versus $x \approx \omega/E$ for a $E = 10$ GeV gluon jet. (b) Mean medium-induced gluon number $\langle N_g \rangle$ per jet for initial effective soft parton rapidity density $dN^g/dy = 650$–800. (c) Corresponding mean fractional energy loss $\langle \Delta E \rangle/E$. (d) Probability distribution $P(\epsilon)$ of the fractional energy loss of $E = 10$ GeV quarks and gluons via multiple independent gluon bremsstrahlung, $\epsilon = \sum_i \omega_i/E$.

An approximation that illustrates the effects of $\beta_T \neq 0$ has been developed in [27]. In this scenario, for a medium of mean initial radius L

$$\rho(\tau) = \frac{1}{\pi}\frac{dN^g}{dy}\frac{1}{\tau}\frac{1}{(L+\beta_T\tau)^2}. \tag{7}$$

Additional dilution relative to the $(1+1)$D Bjorken case is generated by the transverse motion of the soft partons and the increasing transverse size $A_\perp$ of the system. For a discussion of non-central collisions, $R_x \neq R_y$ and $\beta_{Tx} \neq \beta_{Ty}$, see [27]. It follows from Eq. (7) that the propagating jets interact with the bulk matter at midrapidity over an increased time period $\Delta\tau \simeq L/(1-\beta_T) > L$ that largely compensates for the rarefaction of the medium. It has been demonstrated [27] that for small and moderate expansion velocities β_T the $(1+3)$D angular (ϕ) averaged energy loss approximates well the $(1+1)$D result. Logarithmic corrections arise only in the $\beta_T \to 1$ limit. It is physically intuitive that the two effects associated with transverse expansion, the increase of the medium size and the decrease of the medium density, tend to cancel each other and leave the Bjorken expansion part as a good approximation for the calculation of quenching in the single and double inclusive hadron spectra.

An important aspect of the application of the theory of medium induced non-Abelian bremsstrahlung is the inclusion of finite kinematic bounds. These are particularly relevant at RHIC and SPS energies over the full accessible p_T range. Additional details on the evaluation of the radiative energy loss are given in [22].

Fig. 1(a) shows the differential gluon spectrum dN_g/dx and the fractional radiation intensity $x\,dN_g/dx$ computed to 3rd order in opacity ($n = 1, 2, 3$) from Eq. (2) for a 10 GeV gluon jet. The fluctuations in the calculation reflect the numerical accuracy of the cancellation between the propagator poles, Eq. (3), and the interference phases, Eq. (5), and come predominantly from higher orders in opacity. These, however, do not affect the evaluation of the nuclear modification factor. The jet quenching strength is largely set by the $n = 1$ term [22].

In a thermalized medium the plasmon frequency $\omega_{pl}(i) \sim \mu(i)$ regulates the infrared modes and the medium induced bremsstrahlung results in a few semihard gluons. Numerical results for quark jets are not given since they differ by a simple $C_F/C_A = 4/9$ color factor. Fig. 1(b) and (c) show the mean induced gluon number $\langle N_g \rangle$ and the mean fractional energy loss $\langle \Delta E \rangle / E$ calculated directly from Eq. (2) for a range of soft parton rapidity densities $dN^g/dy = 650$–800. Comparing the shape of $\langle \Delta E \rangle / E$ in Fig. 1(c) to the naive analytic expectation in the infinite kinematic limit

$$\frac{\langle \Delta E \rangle}{E} \approx \frac{9 C_R \pi \alpha_s^3}{4} \frac{1}{A_\perp} \frac{dN^g}{dy} L \frac{1}{E} \ln \frac{2E}{\mu^2 L} + \cdots \tag{8}$$

and observing the deviation from the hyperbolic $1/E$ dependence it is easy to recognize the need for a careful treatment of phase space. It is also instructive to note that the average energy per gluon $\langle \Delta E \rangle / \langle N_g \rangle \simeq 0.8$–$1.9$ GeV, which implies that the radiative quanta might be experimentally observable with finite p_T cuts [8,12].

Applications that extend beyond the mean energy loss and invoke a probabilistic treatment with multiple gluon fluctuations require additional assumptions [28]. So far even the case of two gluon emission with scattering has not been calculated. The Poisson approximation assumes independent radiation and equivalence of the single inclusive and the single exclusive gluon spectrum. Finite kinematics alone is guaranteed to violate this ansatz and so is angular ordering. Nevertheless, the usefulness of such probabilistic treatment is in allowing the system to maximize the observable cross section and correspondingly minimize, to the extent to which this is possible, the effect of the non-Abelian medium-induced bremsstrahlung.

The probability $P(\epsilon)$ for fractional energy loss $\epsilon = \sum_i \omega_i / E$ due to multiple gluon emission is given in Fig. 1(d). The $\epsilon = 0$ bin contribution $\delta(\epsilon) e^{-\langle N_g \rangle}$ is not shown. Details on the calculation of $P(\epsilon)$ are given in [28], where a fast iterative procedure for its evaluation form the gluon radiative spectrum dN_g/dx was developed. If a non-zero fraction of the probability $P(\epsilon)$ falls in the $\epsilon > 1$ region, it is uniformly redistributed in the physical $\epsilon \in [0, 1]$ interval. This approach ensures that $\langle \Delta E \rangle / E \leqslant 1$ in the large energy loss limit and provides corrections relative to its direct evaluation from Eq. (2) that are different for quarks and gluons. The properties of $P(\epsilon)$ can be summarized by the normalization of the first two ϵ^k, $k = 0, 1$ moments:

$$\int_0^1 P(\epsilon)\, d\epsilon = 1, \qquad \int_0^1 \epsilon P(\epsilon)\, d\epsilon = \frac{\langle \Delta E \rangle}{E}. \tag{9}$$

We note the distinct difference in the manifestation of the large energy loss relative to the single inclusive spectra. In the former case the overall scale changes and in the probabilistic interpretation the distribution is shifted toward large ϵ values, see Fig. 1(d).

3. Phenomenological application at intermediate energies

Dynamical nuclear effects in $p + A$ and $A + A$ reactions are most readily detectable through the nuclear modification ratio

$$R_{AA}^{h_1}(\mathbf{p}_{T1}) = \frac{dN^{AA}(\mathbf{b})}{dy_1\, d^2\mathbf{p}_{T1}} \left[\frac{T_{AA}(\mathbf{b})\, d\sigma^{pp}}{dy_1\, d^2\mathbf{p}_{T1}} \right]^{-1} \quad \text{in } A + A, \tag{10}$$

where $T_{AA}(\mathbf{b}) = \int d^2\mathbf{r}\, T_A(\mathbf{r}) T_B(\mathbf{r} - \mathbf{b})$ is calculable in terms of the nuclear thickness functions $T_A(\mathbf{r}) = \int dz\, \rho_A(\mathbf{r}, z)$. In $R_{AA}^{h_1}(\mathbf{p}_{T1})$ the uncertainty associated with the next-to leading order factor, K_{NLO}, drops out.

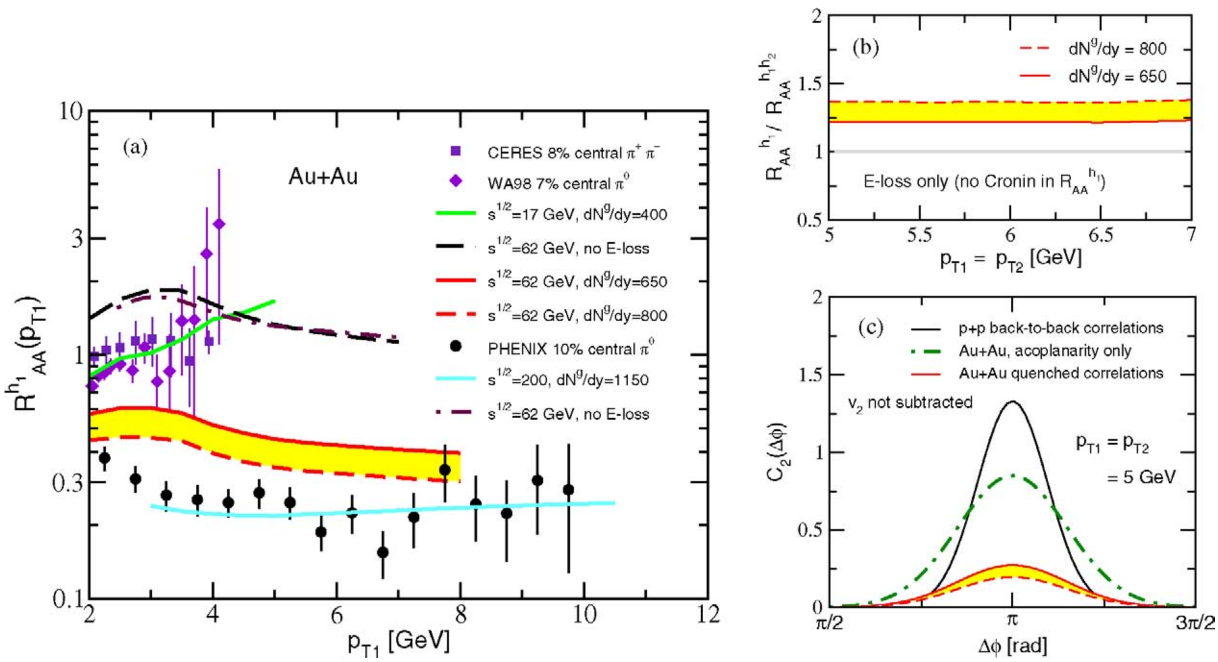

Fig. 2. (a) Calculated nuclear modification factor $R_{AA}(\mathbf{p}_{T1})$ versus the center of mass energy $\sqrt{s_{NN}} = 17, 62, 200$ GeV for central Au + Au collisions. The yellow band represents the perturbative QCD expectation for the depletion of the neutral pion multiplicity at $\sqrt{s_{NN}} = 62$ GeV. Enhancement, arising from transverse momentum diffusion in cold nuclear matter without final state energy loss is given for comparison. PHENIX [2], WA98 [20] and CERES [21] data as presented in [15] are shown. (b) The ratio of the attenuation of the single and double inclusive π^0 cross sections $R^{h_1}_{AA}(p_{T1})/R^{h_1h_2}_{AA}(p_{T_1} = p_{T_2})$ in central Au + Au collisions at $\sqrt{s_{NN}} = 62$ GeV and (c) the manifestation of the double inclusive hadron suppression in the quenching of the away-side di-hadron correlation function $C_2(\Delta\phi)$ [37].

The standard lowest order perturbative expression for the hadron multiplicity in $A + A$ reactions including the effects of multiple initial state scattering and the final state energy loss reads [13,29]:

$$\frac{1}{T_{AA}(\mathbf{b})}\frac{dN^{\pi}_{AA}(\mathbf{b})}{dy_1\,d^2\mathbf{p}_{T1}} = K_{\mathrm{NLO}}\sum_{abcd}\int_{x_{a\,\min}}^{1} dx_a \int_{x_{b\,\min}}^{1} dx_b \int d^2\mathbf{k}_a \int d^2\mathbf{k}_b\, g(\mathbf{k}_a)g(\mathbf{k}_b)\, f_{a/A}\big(x_a, Q^2\big) f_{b/A}\big(x_b, Q^2\big)$$
$$\times\, S_{a/A}\big(x_a, Q^2\big) S_{b/A}\big(x_b, Q^2\big)\frac{\alpha_s^2}{\hat{s}^2}\big|\bar{M}_{ab\to cd}(\hat{s},\hat{t},\hat{u})\big|^2 \int_0^1 d\epsilon\, P(\epsilon)\frac{z_c^*}{z_c}\frac{D_{\pi/c}(z_c^*, Q_c^2)}{z_c}. \tag{11}$$

In Eq. (11) $f_{a/A}(x, Q^2)$ are the isospin corrected ($A = Z + N$) lowest order parton distribution functions [30], $S_{a/A}(x, Q^2)$ is the leading twist shadowing parameterization [31], and $D(z)_{\pi/c}(z, Q^2)$ is the fragmentation function into pions [32]. Vacuum and medium-induced initial state parton broadening, $\langle \mathbf{k}_T^2\rangle = \langle \mathbf{k}_T^2\rangle_{pp} + \langle \mathbf{k}_T^2\rangle_{\mathrm{nucl.}}$, is incorporated via a normalized Gaussian $\mathbf{k}_T$ smearing function [29]. If the final state parton looses a fraction ϵ of its energy the correspondingly rescaled fragmentation momentum fraction reads $z^* = z/(1-\epsilon)$. For consistency the calculation is performed in the same way as in [13] where additional details can be found.

In the $p_T \leqslant 5$ GeV range, experimentally accessible at $\sqrt{s_{NN}} = 62$ GeV, sizable non-perturbative effects in baryon production, manifest in enhanced p/π and Λ/K ratios, will likely be observed. Discussion of the moderate p_T baryon phenomenology is beyond the scope of this Letter and details are given in [33]. In the limit of vanishing baryon masses, $m_B \to 0$, perturbative calculations of the nuclear modification factor $R^{h_1}_{AA}(\mathbf{p}_{T1})$ yield results comparable to the one for neutral and charged pions. We finally note that the dominant contribution to the nuclear shadowing may come from enhanced dynamical power corrections [34,35], resulting from the multiple initial and final state interactions of the partons on a nucleus. However, in the calculated p_T range ($x \geqslant 0.08$, $Q^2 \geqslant 8\ \text{GeV}^2$) their effect was found to be small.

Results form the perturbative calculation of the nuclear modification to the neutral and charged pion production in central Au + Au collisions are given in the left-hand side of Fig. 2(a). At all energies there is a strong cancellation between the Cronin enhancement, which arises from the transverse momentum diffusion of fast partons in cold nuclear matter [7,22], and the subsequent inelastic jet attenuation in the final state. This interplay is most pronounced at the low SPS $\sqrt{s_{NN}} = 17$ GeV where, in the absence of quenching, the corresponding enhancement could reach a factor of 3–4 [13]. With final state energy loss taken into account, $R^{h_1}_{AA}(\mathbf{p}_{T1})$ is shown versus the reanalyzed [15] WA98 [20] and the CERES [21] Pb + Pb and Pb + Au data. We note that even the previously extracted nuclear modification at the SPS allows for inelastic final state interactions in a medium of $dN^g/dy = 200$ [13].

At present, there are strong indications [8,16] that the highest nuclear matter density reached in the early stages ($\tau_0 = 0.6$ fm) of the $\sqrt{s_{NN}} = 200$ GeV central Au + Au collisions at RHIC may be on the order of 100 times cold nuclear matter density, $\epsilon_{\text{cold}} = 0.14\ \text{GeV}/\text{fm}^3$. Perturbative calculations with a corresponding $dN^g/dy = 1150$ are compatible with the measured π^0 attenuation [1], as illustrated in Fig. 2(a). Quenching ratios of similar magnitude, $R^{h_1}_{AA}(\mathbf{p}_{T1})$, $R^{h_1}_{CP}(\mathbf{p}_{T1}) \sim 0.2$–$0.25$, have also been predicted and measured for inclusive charged hadrons at $p_T \geqslant 5$ GeV [1]. It is interesting to note that in spite of the very large initial energy density the nuclear attenuation is "only" a factor of 4–5. The reason for this somewhat unintuitive result is the strong longitudinal expansion in the absence of which energetic jets could have been completely absorbed.

The yellow band represents a calculation of the π^0 attenuation at $\sqrt{s_{NN}} = 62$ GeV from Eqs. (10), (11). The Cronin enhancement alone in central Au + Au reactions was found to be still sizable with $R^{h_1}_{AA}(\mathbf{p}_{T1})_{\max} = 1.8$–$2$ at $p_T = 3$–3.5 GeV with a subsequent decrease at higher transverse momenta. For comparison, at $\sqrt{s_{NN}} = 200$ GeV the Cronin maximum decreases with the center of mass energy and $R^{h_1}_{AA}(\mathbf{p}_{T1})_{\max} = 1.7$. While such enhancement may naively appear large, it is consistent with the system (Au + Au versus d + Au) and $\sqrt{s_{NN}}$ dependence coming from p_T diffusion [7]. We note that even at the maximum RHIC energy the $y = 0$ central d + Au collisions exhibit $\sim 35\%$ maximum pion enhancement [36]. The final state partonic energy loss was computed for a range of initial effective gluon rapidity densities $dN^g/dy = 650$–800 as discussed in Section 2. Fig. 2(a) clearly shows that the perturbative calculation presented here predicts a dominance of the inelastic jet interactions and net π^0 quenching at the intermediate RHIC energy.

Additional constraints for the energy loss calculations would arise from the measurement of a suppressed double inclusive hadron production

$$R^{h_1h_2}_{AA} = \frac{dN^{AA}(\mathbf{b})}{dy_1\, dy_2\, d^2\mathbf{p}_{T1}\, d^2\mathbf{p}_{T2}} \left[\frac{T_{AA}(\mathbf{b})\, d\sigma^{pp}}{dy_1\, dy_2\, d^2\mathbf{p}_{T1}\, d^2\mathbf{p}_{T2}} \right]^{-1}. \tag{12}$$

The cross sections in Eq. (12) can be evaluated as in [37] by fragmenting the second hard scattered parton. We note that by unitarity the details of its energy loss and fragmentation do not affect the calculation of the single inclusive spectra. Di-hadron correlations, however, pick only one of the particle states. Since the second parent parton also looses a fraction of its energy in the medium, qualitatively $1 \leqslant R^{h_1}_{AA}/R^{h_1h_2}_{AA} \leqslant 2$. Numerical estimates in Fig. 2(b) from jet quenching alone indicate that the double inclusive hadron suppression is 25–40% larger than the single inclusive quenching in the $5 \leqslant p_{T1} = p_{T1} \leqslant 7$ GeV range at $\sqrt{s_{NN}} = 62$ GeV. The ratio was computed in the mean ΔE approximation, but with an explicit average over the di-jet production point for a realistic Woods–Saxon geometry. Double inclusive cross section quenching is manifest in the attenuation of the away-side correlation function $C_2(\Delta\phi) = (1/N_{\text{trig}})dN^{h_1h_2}/d\Delta\phi$ [37]. In Fig. 2(c) it leads to a 3–5 fold attenuation of the area A_{far} relative to the $p + p$ case. In contrast, elastic transverse momentum diffusion will only result in a broader $C_2(\Delta\phi)$. We note that the experimentally measured suppression value will be sensitive to the subtraction of the elliptic flow v_2 component [39], which may remove part or all of the residual correlations. The results shown in Fig. 2(c), therefore, correspond to the smallest anticipated observable disappearance of the away-side jet at high p_T.

4. Discussion and summary

For $p_T \geqslant 2$ GeV the reduction of the pion multiplicity in central Au + Au reactions relative to the binary collision scaled $p+p$ result at the intermediate RHIC energy of $\sqrt{s_{NN}} = 62$ GeV was found to be a factor of 2–3. This result is consistent with $R_{AA}(p_T = 4 \text{ GeV}) \sim 0.5$ from [38]. The moderate transverse momentum dependence of the quenching ratio arises from the interplay of the Cronin effect, in particular its decrease toward larger transverse momenta, and the jet energy loss. The medium induced energy loss also results in an apparent disappearance of the away-side jet at high p_T, which is similar to the $\sqrt{s_{NN}} = 200$ GeV result [10].

To relate the effective $dN^g/dy = 650$–800 to the temperature and energy density we employ an initial equilibration time $\tau_0 \simeq 0.8$–1 fm suggested by the hydrodynamic description of the SPS data [39]. For $(1+1)$D expansion $\epsilon_0 = 6$–8 GeV/fm^3, $T_0 = 300$–325 MeV and the lifetime of the plasma phase $\tau - \tau_0 = 3.5$–4 fm. Such initial conditions are already significantly above the current expectation for the critical temperature and energy density for a deconfinement phase transition, $\epsilon_c \simeq 1$ GeV and $T_c \simeq 175$ MeV [40].

From a theoretical point of view, the most interesting result would be a strong deviation of the pion attenuation from the pQCD calculation in Fig. 2(a). This may be either a factor of ~ 5 suppression, comparable to the $\sqrt{s_{NN}} = 200$ GeV quenching, or binary and above binary scaling, as seen at the SPS energies. Both cases would indicate a marked non-linear dependence of the non-Abelian energy loss on the soft parton rapidity density dN^g/dy if we assume that parton–hadron duality still holds. Equally striking will be an absence of strong attenuation in the away-side $C_2(\Delta\phi)$. It is important to check whether the established connection between the single and double inclusive hadron suppression [10] persist at the intermediate RHIC energies, see Fig. 2(c).

In summary, the upcoming $\sqrt{s_{NN}} = 62$ GeV measurements and their comparison to theoretical calculations will provide a critical test of jet tomography. They will help to further clarify the properties of the hot and dense nuclear matter crated in relativistic nuclear collisions and the position that it occupies on the QCD phase transition diagram.

Note added

Shortly after the completion of this work preliminary PHENIX data on π_0 production and attenuation at the intermediate RHIC energy of $\sqrt{s_{NN}} = 62$ GeV became available. Comparison between the data and the predictions given in this Letter can be found in [41].

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