PHYSICAL REVIEW C **67**, 061901(R) (2003)

Phenomenology of jet quenching in heavy ion collisions

Berndt Müller
Department of Physics, Duke University, Durham, North Carolina 27708

(Received 21 August 2002; revised manuscript received 24 March 2003; published 20 June 2003)

We derive an analytical expression for the quenching factor in the strong quenching limit where the p_T spectrum of hard partons is dominated by surface emission. We explore the phenomenological consequences of different scaling laws for the energy loss and calculate the additional suppression of the away-side jet.

DOI: 10.1103/PhysRevC.67.061901 PACS number(s): 25.75.−q

It is commonly believed that the yield for ''hard'' observables in high-energy nuclear reactions scales as the number of binary nucleon-nucleon (NN) collisions occurring during the encounter of the two nuclei. This expectation applies to processes that are characterized by a high virtuality q^2, for which final state interactions are negligible, such as lepton pair production or the total yield of heavy flavor quarks. There is no *a priori* reason to expect this rule to hold for high-q^2 processes, in which the final state can be strongly modified by interactions with comovers. A well documented example is the production of heavy vector mesons, such as the J/Ψ, which is found to be ''anomalously'' suppressed in collisions of heavy nuclei at the CERN-SPS [1].

The yield of hadrons produced with high transverse momentum p_T in Au+Au collisions at the Relativistic Heavy Ion Collider (RHIC) has recently been shown to be significantly suppressed in comparison with the cumulative yield of NN collisions [2,3]. This effect, called ''jet quenching,'' was predicted to occur as a result of energy loss by the hard scattered partons due to interactions with the surrounding dense medium [4–6]. The theory of this energy loss has been a topic of intense research over the past few years [7–11]. The present consensus is that the dominant mechanism for the energy loss in QCD is collisionally induced radiation of gluons by the fast parton.

It is difficult to measure the energy loss of a scattered parton directly in heavy ion reactions, because the large multiplicity of emitted hadrons makes it almost impossible to isolate the resulting jet by kinematic cuts. However, the energy loss of the parton is imprinted as an equivalent loss of energy of the leading hadron produced in its fragmentation [12]. This is what has been observed in the RHIC experiments. Generally, it is assumed that the fragmentation occurs after the parton has left the comoving medium, and thus is described by the measured vacuum fragmentation functions. We will, therefore, not be concerned with the conversion from partons to hadrons, but focus directly on the p_T spectrum of scattered partons.

Preliminary data from run 2 at RHIC (Au+Au at a center-of-mass energy of 200 GeV per nucleon pair) confirm the effect observed in run 1 and its interpretation as jet quenching [13]. For pions with $p_T \approx 5$ GeV/c the measured suppression factor $Q(p_T)$ is about 1/5. There is not yet complete agreement between different experiments about the scaling of the hadron yields with the number of nucleons participating in the reaction, $N_{\rm part}$, or with the binary collision number $N_{\rm coll}$. The PHOBOS Collaboration has presented evidence that the yield of charged hadrons with $p_T \approx 4$ GeV/c scales like $N_{\rm part}$ [14]. This finding is contradicted by data from the PHENIX Collaboration [15]. It is not clear whether the discrepancy is due to different normalization methods, different experimental acceptance, or other reasons.

In addition, the experiments have found that the suppression factors for mesons and baryons are quite different up to transverse momenta of 5 GeV/c. It was recently proposed that this phenomenon can be attributed to a competition between different hadron formation processes [16,17], with parton recombination dominating at lower p_T and fragmentation at higher p_T. The participant or binary collision scaling may be complicated in the transition region. The situation is predicted to simplify for $p_T > 6$ GeV/c, where hadrons are overwhelmingly produced by fragmentation of an energetic parton.

We will first show analytically that the spectrum of high-p_T partons, in the limit of large energy loss, is dominated by partons emitted from the surface of the collision zone and thus scales like the surface rather than the volume of the interaction region [18]. We will then explore the scaling of the quenching factor $Q(p_T)$ with participant number and p_T. We finally calculate the additional suppression of the away-side jet and the azimuthal anisotropy of the parton yield in noncentral collisions.

We begin by considering the loss of energy by an energetic parton traversing a homogeneous, static medium of thickness L. We assume that the geometry is given by a cylinder with radius R, as in the boost-invariant Bjorken model [19] for a nuclear collision with impact parameter $b=0$, and that the parton moves in the transverse plane in the local rest frame of the medium. We further assume that the effective energy loss, defined as the shift of the momentum spectrum of fast partons, depends on p_T and L in the following general way:

$$\Delta p_T = \eta p_T^{\mu} L, \tag{1}$$

where μ is a scaling exponent. The linear dependence on L holds when the loss occurs in subsequent, independent interactions with the medium. It has also been shown to be valid, when multiple interactions are suppressed by the Landau-Pomeranchuk-Migdal effect, due to the steep fall-off of the parton spectrum with p_T [10]. The traversed path length inside the medium for a parton created at transverse position $\mathbf{r}$ with an angle ϕ relative to the radial direction is

$$L(\phi)=(R^2-r^2\sin^2\phi)^{1/2}-r\cos\phi\approx\frac{z^2}{2R\cos\phi}, \qquad (2)$$

where $r=|\mathbf{r}|$, $z^2=R^2-r^2$, and the approximation is valid near the surface ($z\ll R$). Perturbative QCD predicts that the parton spectrum at moderately large values of p_T has the form [20]

$$\frac{dN}{d^2p_T}=N_0\left(1+\frac{p_T}{p_0}\right)^{-\nu} \qquad (3)$$

with a power $\nu\approx 8$ and $p_0\approx 1.75$ GeV/c. The quenched spectrum is given by

$$\frac{d\tilde{N}}{d^2p_T}=Q(p_T)\frac{dN}{d^2p_T}=\frac{1}{2\pi^2R^2}\int_0^{2\pi}d\phi\int_0^R d^2r\frac{dN(p_T+\Delta p_T)}{d^2p_T}. \qquad (4)$$

Replacing the integration over r with an integration over z, using approximation (2), and formally extending the range of the integration to infinity, one finds

$$Q(p_T)\approx\frac{2(p_0+p_T)}{\pi R\,\eta(\nu-1)p_T^{\mu}}. \qquad (5)$$

Two things are remarkable about this result. First, the factor R in the denominator reduces the scaling of the parton yield with the size of the reaction zone by one power of R, from a volume to a surface dependence. Second, the dependence of $Q(p_T)$ on p_T is determined by the power μ governing the p_T dependence of the energy loss. For $\mu=1$ the quenching factor Q falls slowly with increasing p_T; for $\mu=1/2$ it grows with p_T, implying less quenching at higher p_T.

We have confirmed the range of validity of the approximate analytical expression (5) by comparing it with the exact integral (4). As seen in Fig. 1, the analytical approximation $Q_{\rm anal}$ deviates from the exact result by less than 5% when the quenching factor $Q\leqslant 0.2$. An even better agreement is found when the analytical result is divided by the correction factor $(1+Q_{\rm anal}^2)$. The excellent agreement suggests that we may extend the calculation to noncentral collisions. Generalizing the surface-to-volume ratio of the cylindrical geometry ($S/V=2/R$) to the geometry of a collision with impact parameter b, one finds

$$Q(p_T,b)=Q(p_T,0)\alpha_b/(\alpha_b-\sin\alpha_b) \qquad (6)$$

with $\alpha_b=2\arccos(b/2R)$. This approximation is valid with about the same accuracy as Eq. (5), except for very peripheral collisions.

In order to be able to address the experimental data, we need to relax some of the geometrical oversimplifications used above.

(1) The transverse profile of the primary jet yield is not homogeneous, but proportional to the binary NN collision profile,

$$T(\mathbf{r},\mathbf{b})=\rho_1(\mathbf{r})\rho_2(\mathbf{r}-\mathbf{b}), \qquad (7)$$

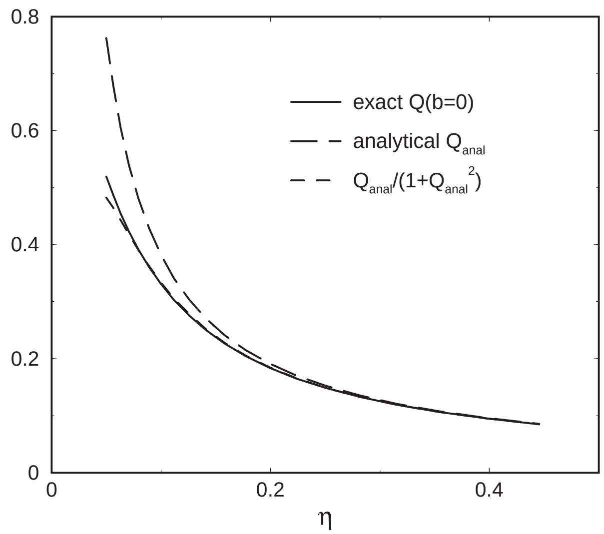

FIG. 1. Comparison between the analytical approximation (5) and the exact result (4) for the quenching factor Q. The curves are for Au+Au at $b=0$ and $p_T=6$ GeV/c with $\mu=1/2$. The calculation assumes a homogeneous transverse profile for jet production and quenching medium. Also shown is the expression $Q_{\rm anal}/(1+Q_{\rm anal}^2)$.

where $\rho_i(\mathbf{r})$ is the longitudinally integrated density of nucleus i.

(2) The density of the comoving medium is also not homogeneous in the transverse plane. We assume that it is proportional to the local density of participant nucleons, which, in the Glauber model, is given by

$$\rho_{\rm part}(\mathbf{r})=\rho_1(\mathbf{r})(1-e^{-\sigma\rho_2(\mathbf{r}-\mathbf{b})})+\rho_2(\mathbf{r}-\mathbf{b})(1-e^{-\sigma\rho_1(\mathbf{r})}), \qquad (8)$$

where σ denotes the inelastic NN cross section.

(3) The comover medium expands and its density decreases with time. Here we assume

$$\rho(\mathbf{r},\tau)=C\rho_{\rm part}(\mathbf{r})/(\tau+\tau_0), \qquad (9)$$

which is modeled to represent a longitudinal, boost invariant expansion. As an approximation to the phenomenology at full RHIC energy we use the values $C\approx 3$ and $\tau_0=1$ fm/c.

Finally, we need to return to Eq. (1) for the energy loss. Perturbative QCD predicts that the radiative energy loss depends quadratically on the medium thickness L [8]. As pointed out by Baier *et al.* [10], this holds for the average energy loss $\overline{\Delta E}$ of a given parton, but the average energy loss of observed partons with fixed transverse momentum p_T has a different scaling. This is so, because the energy loss distribution $D(\epsilon)$ is strongly skewed toward small values of ϵ by the steeply falling p_T spectrum of fast partons. In fact, the average shift of the spectrum due to the energy loss, here called the effective energy loss, is given by [10]

$$\Delta p_T\approx\alpha_s L\sqrt{\pi\hat{q}p_T/\nu}, \qquad (10)$$

where $\hat{q}$ encodes the "scattering power" of the medium, which is proportional to the density. For an expanding medium, the expression $\hat{q}L^2$ must be replaced with

$$\hat{q}_0 L_{\rm eff}^2 = 2\hat{q}_0 \int_0^L \tau d\tau \frac{\rho[\mathbf{r}(\tau),\tau]}{\rho(\mathbf{r},0)}, \tag{11}$$

where $\mathbf{r}(\tau) = \mathbf{r} + \mathbf{v}\tau$ denotes the position of the fast parton in the medium at time τ, and $\hat{q}_0$ is a function of the transverse position $\mathbf{r}$ at which the jet is produced. We thus can make contact with Eq. (1) by writing

$$\Delta p_T = \eta' L_{\rm eff}\sqrt{\rho(\mathbf{r},0) p_T/\nu} \tag{12}$$

with the constant $\eta' = \alpha_s \sqrt{\pi \hat{q}_0/\rho(0)}$, which does not depend on $\mathbf{r}$.

We will denote the scaling law (12) for the energy loss as BDMS. In our following numerical study we have explored two other scaling laws. The first one is the Bethe-Heitler (BH) scaling law [21,22],

$$\Delta p_T = \eta p_T \int_0^L d\tau \rho[\mathbf{r}(\tau),\tau] \equiv \eta p_T (L\rho)_{\rm eff}, \tag{13}$$

corresponding to $\mu = 1$. The second scaling law is

$$\Delta p_T = \eta p_T \sqrt{(L\rho)_{\rm eff}}, \tag{14}$$

which we will call the RW scaling law. It could be interpreted as describing a random walk in p_T as the fast parton traverses the medium, with some interactions resulting in an energy gain and others in a loss of energy.

We begin the discussion of our numerical results for the quenching factor Q with its dependence on the transverse momentum of the fast parton, shown in Fig. 2. The QCD-motivated BDMS law (solid line) and the other two scaling laws exhibit clearly different behaviors. This reflects the different p_T scaling of the energy loss in these models (linear for BH and RW; square root for BDMS). The data from the RHIC experiments [2,3,15] suggest that the quenching first becomes stronger with increasing momentum, reaches a minimum, and finally begins to diminish. This would indicate that the BDMS law only applies at high p_T, and that other laws govern the energy loss at lower p_T [21] or hadron production is not dominated by parton fragmentation in this kinematic region. We note that the dependence of Q on p_T is well described by the analytical formula (5).

The impact parameter dependence of the quenching factor is shown in Fig. 3, plotted as the yield per half the number of participant nucleons against the participant number $N_{\rm part}$. The unquenched jet yield, scaling with the number of binary NN collisions, would increase relative to $N_{\rm part}$. As the figure shows, the quenching counteracts this increase, and the yield per participant actually falls for the BDMS and the BH laws as the collision centrality increases. An approximately flat behavior, as observed in the PHOBOS experiment [14], is only found for the RW scaling law.

For noncentral collisions, the quenching factor Q is a function of the azimuthal emission angle, because the geometry is not axially symmetric. This is known to lead to an angular asymmetry of a quadrupole shape in the spectra of high-p_T particles [23,24]. The elliptic flow parameter v_2 is defined as the Fourier component proportional to cos(2ϕ) of the angular distribution of particles with respect to the scattering plane [25]. We find (see Fig. 4) that the values of $v_2 \leqslant 0.1$ obtained for all three scaling laws are significantly smaller than the measured values ($v_2 \approx 0.2$) for semicentral and peripheral collisions [26]. However, the calculated v_2 for partons would be large enough to explain the measured elliptic flow of hadrons, if the hadrons were produced by re-

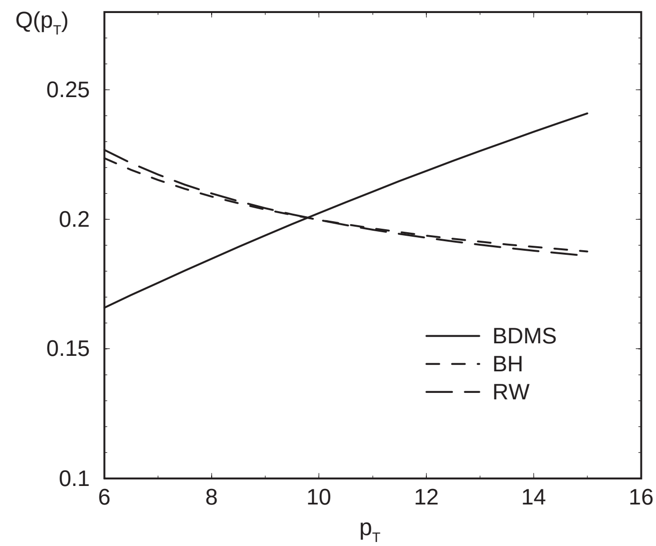

FIG. 2. Dependence of the quenching factor Q on p_T for central collisions. The parameter η is chosen such that $Q(p_T) \approx 0.2$ for $p_T = 10$ GeV/c in each case. The scaling laws (BH, RW) exhibit stronger quenching with increasing p_T, in agreement with preliminary RHIC data, in contrast to the BDMS law. Equation (5) provides a good description of the dependence on p_T seen here.

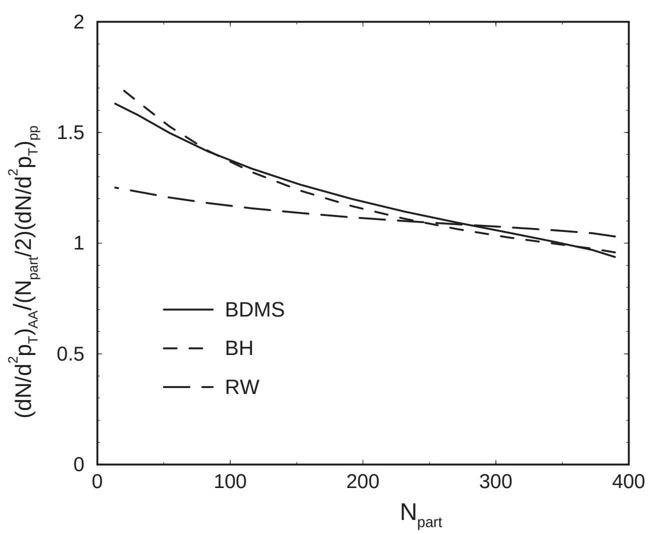

FIG. 3. Quenched hard parton yield divided by half the number of participant nucleons as a function of $N_{\rm part}$ for $p_T = 10$ GeV/c. The values of the stopping power strength parameters are $\eta = 0.06$ (RW), $\eta = 0.017$ (BH), and $\eta' = 0.78$ (BDMS).

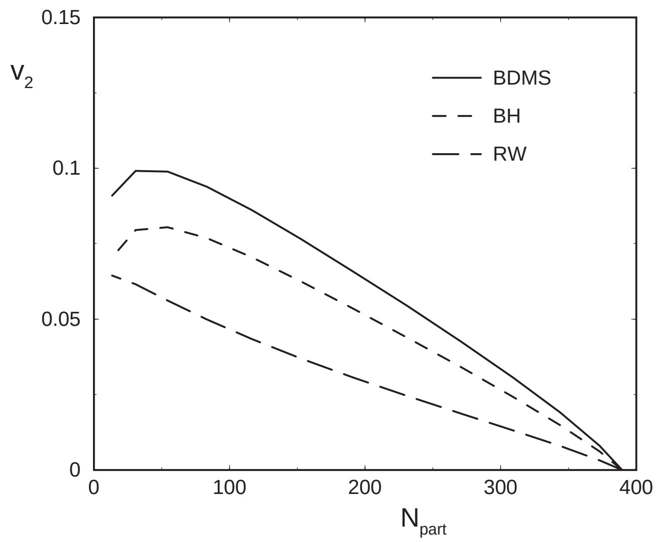

FIG. 4. Elliptic flow parameter v_2 as a function of the number of participants for the three energy loss models. The values of the parameters as the same as in Fig. 3.

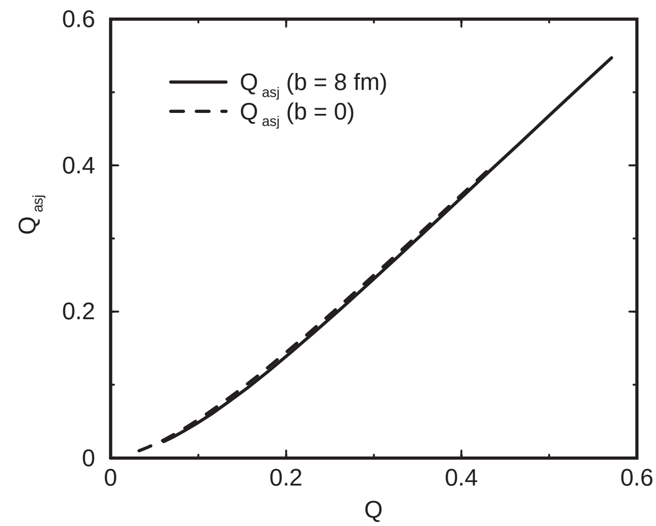

FIG. 5. Incremental away-side jet suppression factor Q_{asj} as function of the "same-side" jet suppression factor Q for the BDMS energy loss law.

combination in the p_T range of the RHIC data [27].

As observed hadrons from hard partons preferentially originate from the surface region facing the detector, the parton emitted in the opposite direction has to traverse more material and thus endures an even larger energy loss. This leads to an additional suppression of the away-side jet and its leading hadron spectrum. The dependence of the incremental away-side hadron suppression factor Q_{asj} on the primary suppression factor Q is shown in Fig. 5 for the BDMS scaling law for two impact parameters ($b=0$ and $b=8$ fm). There appears to be a universal relationship, which is linear for $Q \geqslant 0.2$.

Several conclusions can be drawn from our results. First, the momentum dependence of the BDMS energy loss formula does not seem to agree with some of the RHIC data. A linear dependence of the average energy loss on p_T is in better agreement with the data in the p_T region explored so far [21]. Also, the dependence of the effective momentum loss Δp_T on the medium thickness L_{eff} predicted by BDMS does not yield a scaling of the charged hadron yield with participant number as seen in the PHOBOS data up to $p_T = 4.25$ GeV/c; only the RW scaling law yields such a dependence. This may indicate that energy gain and loss mechanisms are competing in this momentum range [28]. Another explanation may be that hadrons at intermediate momenta are not produced by fragmentation of fast partons, but by other processes, such as parton recombination. The observed magnitude of the elliptic flow lends support to this interpretation. Finally, we have found a universal relation between the same-side and away-side suppression factors for the BDMS law, which can be tested experimentally.

This work was supported in part by a grant from the U.S. Department of Energy (Contract No. DE-FG02-96ER90945). I thank S. A. Bass, X.-N. Wang, P. Jacobs, and K. Rajagopal for valuable comments on the manuscript.

[1] G. Gerschel and J. Hüfner, Annu. Rev. Nucl. Part. Sci. **49**, 255 (1999).
[2] K. Adcox *et al.*, PHENIX Collaboration, Phys. Rev. Lett. **88**, 022301 (2002).
[3] C. Adler *et al.*, STAR Collaboration, Phys. Rev. Lett. **89**, 202301 (2002); **90**, 082302 (2002).
[4] J. D. Bjorken, Fermilab-PUB-82-59-T (unpublished).
[5] M.H. Thoma and M. Gyulassy, Nucl. Phys. **B351**, 491 (1991).
[6] M. Gyulassy and X.-N. Wang, Nucl. Phys. **B420**, 583 (1994).
[7] B.G. Zakharov, JETP Lett. **63**, 952 (1996); **65**, 615 (1997).
[8] R. Baier, Yu.L. Dokshitzer, A.H. Mueller, S. Peigne, and D. Schiff, Nucl. Phys. **B483**, 291 (1997); **B484**, 265 (1997).
[9] U.A. Wiedemann, Nucl. Phys. **B588**, 303 (2000).
[10] R. Baier, Yu.L. Dokshitzer, A.H. Mueller, and D. Schiff, J. High Energy Phys. **109**, 033 (2001).
[11] R. Baier, D. Schiff, and B.G. Zakharov, Annu. Rev. Nucl. Part. Sci. **50**, 37 (2000).
[12] X.-N. Wang and M. Gyulassy, Phys. Rev. Lett. **68**, 1480 (1992).
[13] Presentations by the BRAHMS, PHENIX, PHOBOS, and STAR Collaborations at the Quark Matter 2002 Conference: C. E. Jørgensen, S. Mioduszewski, N. N. Ajitanand, D. d'Enterria, T. Sakaguchi, C. Roland, D. Hardtke, K. Filimonov, and J. Klay, Nucl. Phys. (to be published), http://alice-france.in2p3.fr/qm2002
[14] C. Roland, PHOBOS Collaboration, hep-ex/0212006.
[15] D. d'Enterria, PHENIX Collaboration, hep-ex/0209051.
[16] R.J. Fries, B. Müller, C. Nonaka, and S.A. Bass, nucl-th/0301087, Phys. Rev. Lett. (to be published).
[17] V. Greco, C.M. Ko, and P. Levai, nucl-th/0301093, Phys. Rev. Lett. (to be published).

[18] E.V. Shuryak, Phys. Rev. D **66**, 027902 (2002).
[19] J.D. Bjorken, Phys. Rev. D **27**, 140 (1983).
[20] R.J. Fries, B. Müller, and D.K. Srivastava, Phys. Rev. Lett. **90**, 132301 (2003).
[21] S. Jeon, J. Jalilian-Marian, and I. Sarcevic, Phys. Lett. B **562**, 45 (2003).
[22] M.G. Mustafa, D. Pal, D.K. Srivastava, and M. Thoma, Phys. Lett. B **428**, 234 (1998).
[23] X.-N. Wang, Phys. Rev. C **63**, 054902 (2001).
[24] M. Gyulassy, I. Vitev, and X.-N. Wang, Phys. Rev. Lett. **86**, 2537 (2001).
[25] A.M. Poskanzer and S.A. Voloshin, Phys. Rev. D **58**, 1671 (1998).
[26] C. Adler *et al.*, STAR Collaboration, Phys. Rev. Lett. **89**, 132301 (2002); **90**, 032301 (2003).
[27] D. Molnar and S.A. Voloshin, nucl-th/0302014.
[28] E. Wang and X.-N. Wang, Phys. Rev. Lett. **87**, 142301 (2001).