Jets in expanding quark-gluon plasmas

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We study the transverse-momentum imbalance (acoplanarity) of a gluon jet propagating in an expanding quark-gluon plasma. Under reasonable assumptions for the gluon cross section for interaction with a quark-gluon plasma, and with a hadron gas, and a proper space-time picture of the time evolution of matter produced in nuclear collisions, a simple formula is established relating the acoplanarity distribution to the jet emission angle, the total multiplicity of produced particles, and the nuclear radius. We find that for reasonable values of the jet-plasma cross section, the acoplanarity distribution stands out significantly beyond experimental cuts.

I. INTRODUCTION

Heavy-ion collisions at ultrarelativistic energies offer us the prospect of studying new states of matter. In particular one expects that if the energy density achieved in such collisions is high enough, a deconfinement transition leading to the formation of a quark-gluon plasma could be observed. An important issue in this field is to pin down observables which could provide unambiguous information on the state of matter produced during a collision. Owing to the complexity of heavy-ion systems, it is likely that no single probe will yield definite conclusions. It is therefore important to investigate as many observables as possible and correlate their predictions.

It has been suggested that the study of the propagation of jets through the plasma could give information on the quark and gluon mean free paths in the plasma, and hence on the properties of the plasma itself. Jets are produced at the very beginning of the reaction by hard collisions of the nucleon's constituents. The typical jet energy is greater than 10 GeV, much more than typical energies of the plasma constituents at the temperatures of interest. Thus they may be expected to escape the plasma as well-identified objects.

In a recent paper, Appel has studied the momentum imbalance of jet pairs which results from the interaction of the jets with the plasma constituents.[1] Appel's analysis strongly suggests that such momentum imbalance or acoplanarity may be quite sensitive to the physical conditions which prevail in the plasma which the jets traverse. The purpose of this paper is to extend Appel's investigation, using a more realistic model for the plasma. In particular we shall take into account the longitudinal expansion of the plasma.[2] We shall also investigate the effects of experimental cuts on the produced particle's momenta along the jet axis. We also allow for a combination of hard and soft scattering of the jet from the surrounding matter. The soft component is phenomenologically parametrized in terms of a total jet matter cross section and an exponential slope typical of soft processes. The hard cross section is taken from perturbative QCD. We determine the relative magnitude of the effects of these two contributions to jet scattering and discuss the prospects for measuring these quantities experimentally.

This paper is organized as follows. In Sec. II we recall the main formulas which allow us to calculate the jet acoplanarity arising from the soft-glue emission prior to the collision. This is the standard bremsstrahlung which takes place already in free space. In Sec. III we describe the model we use for the plasma and the influence of the plasma on the jet propagation. We also implement the effect of experimental cuts which must be made on the momentum of particles along the jet axis in order that an identification of a jet may be made. Section IV contains a presentation of our numerical results. Section V summarizes the conclusions, and discusses what is required of an experiment which might attempt to measure this acoplanarity.

II. JET ACOPLANARITY

For simplicity of our analysis, we assume that the jets are produced in the plane $z=0$, that is, in the center-of-mass frame of the colliding constituents. The more general situation is treated along the lines we present here. For the situation of interest, the central region of head-on nucleus-nucleus collisions, we expect that the rapidity distribution of produced particles is approximately boost invariant, and therefore the jet distribution when expressed in terms of jet mass, transverse momentum, and rapidity is invariant under longitudinal boosts.[2] Since the jet mass Q is very large, and the formation time of the leading particles in the jet is of order $1/Q$, we shall assume that the leading particles in the jet form at $t=0$. In terms of the space-time rapidity y and proper time τ, we are assuming that $\tau=y=0$ are the initial space-time coordinates of the jet.

The geometry of the problem is illustrated in Figs. 1(a) and 1(b). We assume head-on collisions, and therefore cylindrical symmetry of the collision. The two jets ori-

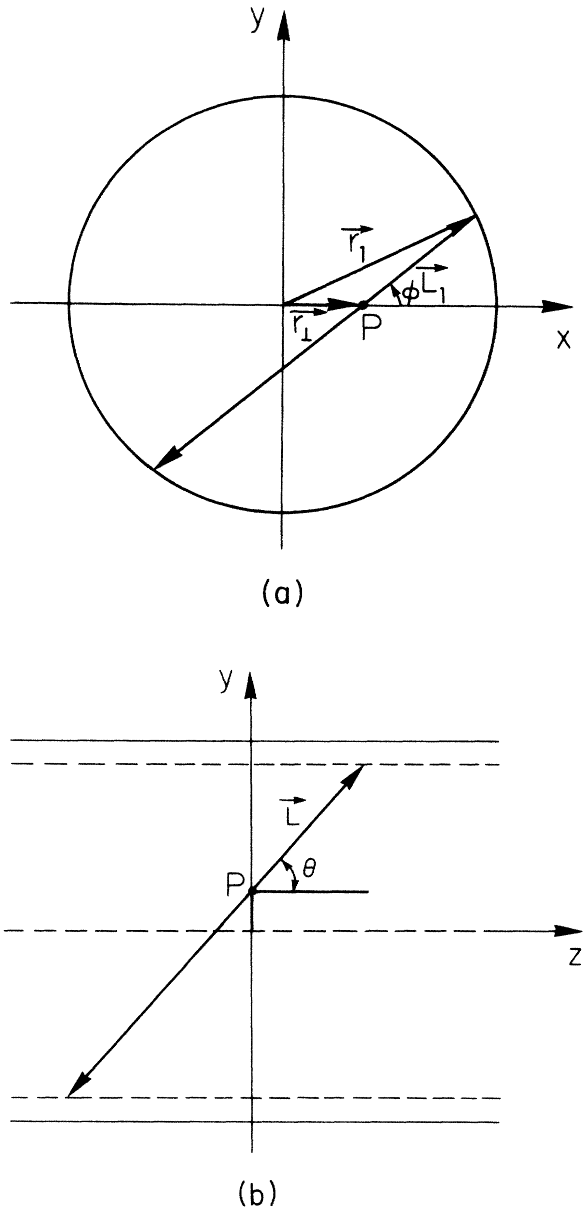

FIG. 1. These figures illustrate the geometry of the problem. The z axis is the collision axis. (a) is drawn in the plane $z=0$. P is the point where the leading particles of the jets are produced. (b) is drawn in the plane of the jets. The solid lines limit the region occupied by the plasma. The dashed lines are projections on the plane of the jets of the parallels to the collision axis drawn through the points where the jets leave the plasma.

ginate from the production point P, back to back. Parallel to the collision axis going through P and the jet axis, a plane is defined to which we refer as the plane of the jets. Φ is the angle between the plane and the plane containing the collision axis and the production point P. θ is the angle of the jets with respect to the collision axis, in the plane of the jets. θ is related to the longitudinal rapidity of the jets through the formula

$$\cosh(y_{\text{jet}})=Q/2P_t=1/\sin(\theta)\ , \tag{1}$$

where Q is the invariant mass of the pair of jets and P_t is the transverse momentum of one jet.

The situation described so far is somewhat idealized. In fact, several effects contribute to give to the jet momenta some nonvanishing component along an axis η perpendicular to the ideal scattering plane defined above. As a result, the two jets do not lie in this ideal scattering plane. The total component $k_{t\eta}$ of the momenta of the two jets along the axis η, referred to as the acoplanarity, may be generated by production of particles in the initial state which are not included in the jet distribution, or by the production of or scattering from low-transverse-momentum particles which are not included in the jet.[3,4] Of course, the ideal scattering plane defined above is not observable, nor is the production point P. However, as discussed by Appel, it can be determined, to within an error in angle of order $k_{t\eta}/Q$, by minimizing over all possible planes the quantity $\sum_i k_{t\eta i}$, where i runs over all the particles in the final state.[1]

Let $B(k_t)$ be the probability density for the leading particle of one jet to acquire transverse momentum k_t due to a single gluon bremsstrahlung prior to the collision. Then the probability $dP/dk_{t\eta}$ density to observe a total momentum imbalance $k_{t\eta}$ is given by

$$dP/kt_{t\eta}=\sum_{n=0}\frac{1}{n!}\left[\prod_{i=1}^{n}\int d^2k_{ti}B(\mathbf{k}_{ti})\right]\times\delta\left[k_{t\eta}-\sum_{i=1}^{n}\mathbf{k}_{ti}\cdot\hat{\mathbf{k}}_{t\eta}\right]\ . \tag{2}$$

Note that the term with $n=0$ is simply $\delta(k_{t\eta})$. This formula acquires a simple form when expressed in terms of the Fourier transform of B with respect to impact parameter b:

$$\overline{B}(b)=\int d^2k_t e^{i\mathbf{k}_t\cdot\hat{\mathbf{k}}_{t\eta}b}B(\mathbf{k}_t)\ . \tag{3}$$

Notice that in Eq. (2), the direction of $k_{t\eta}$ has already been specified by determining the plane of the jet axis. Therefore, the distribution depends only upon the magnitude of $k_{t\eta}$. The distribution of $dP/dk_{t\eta}$ is, therefore, in terms of a one-dimensional impact-parameter representation,

$$dP/dk_{t\eta}=\frac{1}{\pi}\int_0^\infty db\cos(k_{t\eta}b)\exp[\overline{B}(b)]\ . \tag{4}$$

In order that $dP/dk_{t\eta}$ be normalized to one, upon integration over all $k_{t\eta}$, we require that $\overline{B}(b)$ vanish when $b\rightarrow 0$. This is accomplished by subtracting $\overline{B}(0)$ from $\overline{B}(b)$. We emphasize that this subtraction is only done to make $dP/dk_{t\eta}$ a normalized probability distribution. We should not confuse $\mathbf{k}_t$, the component of the jet momentum perpendicular to the jet axis, with $k_{t\eta}$, the projection of $\mathbf{k}_t$ along the axis η perpendicular to the plane of the jets ($k_{t\eta}=k_t\cdot\hat{\mathbf{k}}_{t\eta}$). For $\overline{B}(b)$, we use the same function as Appel:[5]

$$\overline{B}(b) = -\frac{C_F}{\pi\beta}(\{\ln\ln(Q^2/\Lambda^2) - \ln\ln[(b_0/b)^2/\Lambda^2]\} \times [\ln(Q^2/\Lambda^2) - \tfrac{3}{2}] , - \ln[Q^2/(b_0/b)^2]) , \quad (5)$$

where $C_F = \frac{4}{3}$ and $\beta = (33 - 2N_F)/12\pi$ for N_F flavors. Λ is the QCD scale parameter, (0.2 GeV), $b_0 = 1.123$. Q is the invariant mass of the pair of jets. The validity of the formula above is restricted to the region $b \ll \Lambda^{-1}$. Furthermore, we shall take $\overline{B} = 0$ when $b < b_0/Q$. This ensures in particular that $\overline{B}(0) = 0$.

As we shall see in Sec. III, the scattering of the leading particles off the plasma constituents leads to a very simple modification of Eq. (4).

III. JETS IN THE PLASMA

In this section we derive the formulas which allow the calculation of the acoplanarity arising from the multiple scattering of a gluon jet on the constituents of the plasma. Let us first specify our model for the plasma. We consider only central collisions and assume that the plasma is contained in a cylinder of radius $R = 1.2$ fm $A^{1/3}$, where A is the mass number. We shall also assume that the particle-production process is invariant under Lorentz boosts in the longitudinal direction. Then all the quantities specifying the state of the system depend only on the proper time $\tau = (t^2 - z^2)^{1/2}$, and not on z and t separately. The evolution of the plasma after it is formed is described by the hydrodynamic equations. In the simplified situation which we are considering, this evolution corresponds to a uniform cooling of the plasma, as it expands in the longitudinal direction. This cooling implies a decrease of the entropy density according to the law

$$s(\tau) = s(\tau_0)\tau_0/\tau . \quad (6)$$

Thus, as time goes on, the entropy density, or equivalently the temperature, decreases and eventually reaches the value at which the phase transition from the quark-gluon plasma into ordinary hadronic matter is expected to take place. In order to be able to follow the system through the phase transition, further assumptions have to be made. In this paper, we shall investigate a plausible scenario in which the plasma adiabatically converts into hadrons. Let us call $s_{\rm pl}$ and $s_{\rm h}$ the entropy densities of the plasma and of the hadronic matter, respectively, at the transition. We have

$$s_{\rm pl}/s_{\rm h} = N_{\rm pl}/N_{\rm h} = r , \quad (7)$$

where $N_{\rm pl}$ and $N_{\rm h}$ count the number of degrees of freedom in the plasma and in the hadron gas, respectively, and we have assumed a bag-model equation of state for both phases. Typically, r is of the order of 10. From the fact that the total entropy is conserved, and the fact that the entropy density decreases with respect to time according to Eq. (6), one easily derives the formula

$$s(\tau) = x s_{\rm pl} + (1 - x) s_{\rm h} , \quad (8)$$

where x, the fraction of hadronic matter in the mixed phase is given by

$$x = [r\tau_{\rm pl}/(\tau - 1)]/(r - 1) , \quad (9)$$

$\tau_{\rm pl}$ being the time at which the hadronization of the plasma starts. $\tau_{\rm pl}$ may be obtained as a function of the initial time τ_0, where the hydrodynamic evolution starts from the entropy equation, Eq. (6). An important feature of the present scenario is that it implies that the system spends a lot of time in the mixed phase. If $\tau_{\rm h}$ denotes the time at which the hadronization ends, we have $\tau_{\rm h}/\tau_{\rm pl} = r$ as can be obtained from Eq. (7).

The jets are produced at very short times and are supposed to propagate at the speed of light. Before they leave the interaction region, whose thickness is typically a few fm, they will scatter on the plasma particles. The net effect of this multiple scattering is to increase the momentum imbalance. As shown by Appel, the effect of the plasma can be accounted for by a simple modification of the formula Eq. (4), leading to the simple replacement of the function $\overline{B}(b)$ by the function $\overline{B}(b) + \overline{F}(b)$. The new function $\overline{F}(b)$ is the Fourier transform [as defined by Eq. (3)] of the probability density that the gluon jet scatters elastically off the plasma constituents with transverse momentum k_t. This function $\overline{F}$ has been subtracted so that it vanishes at $b = 0$, and that $dP/dk_{t\eta}$ is properly normalized. We take $F(k_t)$ to be the integral of the inverse mean free path throughout its space-time transversal of the plasma:

$$F(k_T) = \sum_i \int dx\, n_i\, d^2\sigma/d^2k_t , \quad (10)$$

where n_i is the number density of plasma constituents of type i and $d^2\sigma/d^2k_t$ is the differential cross section for jet scattering with transverse-momentum transfer k_t. This latter cross section is assumed to be the sum of two contributions. The first is due to soft hadronic processes, which we parametrize as

$$d^2\sigma_s^i/d^2k_t = \frac{\sigma_{gi}}{2\pi m^2}\exp(-k_t/M) . \quad (11)$$

This differential cross section has been parametrized by an exponential, as is typical of low-p_t phenomenon. We do not assume that this cross section has the form extracted from QCD under the assumption of large momentum transfer, since for the processes we consider, the momentum transfer is on the average quite small. In this equation, i denotes either σ_{gg}, or σ_{gq}, or σ_{gh}. We are assuming that the leading particle in the jet is a gluon, and do not consider here the scattering of quark jets, although they could be treated by the same methods. The parameter M is taken to be of the order of 0.4 GeV. For the cross section σ_i we use

$$\sigma_{gg} = \tfrac{9}{4}\sigma_{gq} = (\tfrac{9}{4})^2\sigma_{qq} \quad (12)$$

with σ_{qq} taken from the additive quark model, i.e., $\sigma_{qq} = 4.5$ mb. This relation between cross sections is only derivable in perturbative QCD, and is useful to use for making estimates. Again, this assumption might be relaxed to obtain a more general computation, but in this preliminary analysis, we will use Eq. (12).

The hard-scattering contribution is taken from perturbative QCD (Ref. 6). We only include this contribution for $k_t > 1$ GeV. In this limit, the relations between differential cross sections are as in Eq. (12). The differential cross section for glue-glue scattering is

$$d^2\sigma_{gg}^{\mathrm{h}}/d^2k_t = \tfrac{9}{2}\alpha_s{}^2\frac{1}{k_t{}^2} \ . \tag{13}$$

Here, α_s is the QCD running coupling constant:

$$\alpha_s = 1/\beta\ln(k_t{}^2/\Lambda^2) \ . \tag{14}$$

In order to evaluate the function F, we need to integrate along the jet path. This integral is conveniently done by changing the variable, using the proper time τ as the new integration variable. Noting that $dx = v\,dt = dt$ (since the jet propagates with the speed of light) and that $\tau = t\sin(\theta)$, one gets

$$F = \frac{1}{\sin(\theta)}\sum_i \int_{\tau_0}^{\tau_1} d\tau\,(d^2\sigma_i/d^2k_t)n_i(\tau) \ , \tag{15}$$

where τ_L denotes the proper time at which the jet leaves the plasma. It is easily shown that τ_L is related to the path length of the jet in the plasma by

$$\tau_L = L\sin(\theta) \tag{16}$$

with

$$L = [(R^2 - r_t{}^2\sin^2\Phi)^{1/2} - r_t\cos(\Phi)]/\sin(\theta) \ . \tag{17}$$

Here r_t is the distance of the origin of the jet from the collision axis.

It is convenient at this stage to look at the space-time diagram (Fig. 2) which summarizes the evolution of the system. In this diagram, the jet trajectory is represented by a straight line which with the z axis makes an angle α related to the angle θ by

$$\tan\alpha = 1/\cos\theta \ . \tag{18}$$

The various stages of the plasma evolution are delimited by the hyperbolas corresponding to constant proper times. The hyperbola τ_0 corresponds to the plasma formation. The next hyperbola is labeled τ_{pl} which is the proper time at which the plasma enters the mixed phase. The mixed phase lasts until proper time τ_{h}. As we already mentioned, τ_{h} may be quite large. We shall assume that τ_{h} is always larger than τ_L, so that the jet escapes from the plasma before the mixed phase has totally disappeared.

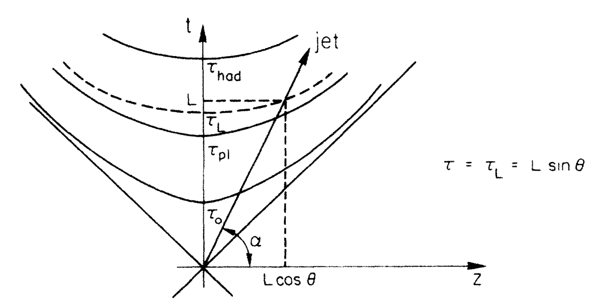

FIG. 2. Space-time diagram. The hyperbolas of constant proper-time delimit the various stages of the plasma history. The straight line labeled "jet" is the jet trajectory. τ_L is the proper time at which the jet leaves the plasma.

The calculation of the integral giving F can then be split into two pieces: one which corresponds to the propagation of the jet in the plasma, the other one corresponding to the propagation of the jet in the mixed phase. We shall write the Fourier transform of F in the following form:

$$\bar{F}(b) = \frac{1}{\sin(\theta)}(\bar{F}_s + \bar{F}_h)(f_{\mathrm{pl}} + f_{\mathrm{mix}}) \ . \tag{19}$$

In this equation, $\bar{F}_s$ and $\bar{F}_h$ are the impact-parameter transforms of

$$\bar{F}_s = \frac{1}{\sigma_{gg}}\int d^2k_t(e^{i\mathbf{k}_t\cdot\hat{\mathbf{k}}_\eta b} - 1)d^2\sigma_{gg}^s/d^2k_t \tag{20}$$

and

$$\bar{F}_h = \frac{1}{\sigma_{gg}}\int d^2k_t(e^{i\mathbf{k}_t\cdot\hat{\mathbf{k}}_\eta b} - 1)d^2\sigma_{gg}^h/d^2k_t \ , \tag{21}$$

where s and h refer to hard and soft scattering and where f_{pl} and f_{mix} are two constants which we will soon calculate.

First, however, consider $\bar{F}$. For $\bar{F}_s$, a short computation gives

$$\bar{F}_s = (1 + b^2M^2)^{-3/2} - 1 \ . \tag{22}$$

A tedious, but straightforward computation of $\bar{F}_h$ gives, to leading order in α_s, the expression

$$\bar{F}_h = (2\pi)\frac{1}{\sigma_{gg}}\frac{9}{2}\frac{1}{\Lambda_0{}^2}\alpha_s{}^2(b\Lambda) \times \left[(\Lambda_0 b/2)^2[\Psi(2) - \ln(\Lambda_0 b/2)] + \sum_{p=0} (-)^p(\Lambda_0 b/2)^{2p+4} \times \frac{1}{2p+2}\frac{1}{\Gamma(p+3)} \right] . \tag{23}$$

In this expression, Λ_0 is a cutoff below which we cannot trust perturbative QCD, which we choose to be $\Lambda_0 = 1$ GeV. Λ is the QCD Λ parameter. The running coupling constant is taken to be a constant equal to its value for $b = 1\ \mathrm{GeV}^{-1} = 0.2$ fm for $b > 0.2$ fm.

The quantities f_{pl} and f_{mix} are given as integrals over the trajectory of the jet in the plasma phase and the mixed phase, respectively, of

$$f = \int dx\,\sigma_{gg}(n_g + \tfrac{4}{9}n_q) \ . \tag{24}$$

Consider first f_{pl}, which is given by

$$f_{\mathrm{pl}} = \int_{\tau_0}^{\tau_{\mathrm{pl}}} d\tau\,\sigma_{gg}(n_g + \tfrac{4}{9}n_q) \ . \tag{25}$$

The particle densities are directly related to the entropy densities for an ideal gas, which in turn can be calculated from Eq. (6) as a function of the proper time τ. A simple calculation then shows that

$$f_{\rm pl}=\sigma_{\rm pl}(dN/dy)\ln(\tau_L/\tau_0)/\pi R^2 \;, \quad (26)$$

where τ_L is the smaller of the two times τ_L and $\tau_{\rm pl}$. To derive this expression, we have used the relation

$$\frac{dN}{dy}\frac{1}{\pi R^2}=[n_g(\tau_0)+\tfrac{7}{6}n_q(\tau_0)]\tau_0 \;. \quad (27)$$

This expression follows essentially from entropy conservation. The peculiar factor of $\frac{7}{6}$ comes from the difference in statistics between bosons and fermions. The factor of 1 for gluons arises because pions and gluons are both bosons. Combining Eqs. (25)—(27) gives

$$\sigma_{\rm pl}=\sigma_{gg}(16+4N_F)/(16+21N_F/2)\sim 14 \text{ mb} \;, \quad (28)$$

where N_F is the number of active quark flavors, which for our purposes, we choose to be $N_F=3$.

The calculation of the mixed-phase contribution implies evaluating the densities of particles in the mixed phase. This may be done easily using the entropy equation. One then finds

$$f_{\rm mix}=\frac{dN}{dy}\frac{1}{\pi R^2}\left[\frac{(r\sigma_{\rm pl}-\sigma_{gh})}{(r-1)}\ln(\tau_L/\tau_{\rm pl}) + \frac{(\sigma_{gh}-\sigma_{\rm pl})}{(r-1)}\frac{(\tau_L-\tau_{\rm pl})}{\tau_{\rm pl}}\right] . \quad (29)$$

This formula assumes of course that $\tau_L > \tau_{\rm pl}$.

A remarkable simplification takes place if one assumes that the scattering cross section of a gluon from hadrons is the same as the cross section for scattering on the plasma, i.e., if one sets

$$\sigma_{gh}=\sigma_{\rm pl} \;. \quad (30)$$

This equality is probably a good approximation. In the additive quark parton model of cross sections, we have approximately

$$\sigma_{gh}\sim\tfrac{9}{4}\sigma_{q\pi}=\tfrac{1}{2}\sigma_{pp}=19 \text{ mb} \;. \quad (31)$$

This number is remarkably close to the average value in the plasma, 14 mb. A more detailed computation would of course include better estimates, but in the crude analysis which we present, such refinement is inconsequential. Then $f_{\rm pl}$ and $f_{\rm mix}$ add up to the simple expression

$$f_{\rm pl}+f_{\rm mix}=\sigma_{\rm pl}\frac{dN}{dy}\frac{1}{\pi R^2}\ln(\tau_L/\tau_0) \;. \quad (32)$$

It is then easy to add up the contributions of the two jets leaving the plasma at Φ and $\Phi+\pi$. Note that the Φ dependence comes through the time τ_L. One then gets

$$f_{\rm pl}+f_{\rm mix}=\sigma_{\rm pl}\frac{dN}{dy}\frac{1}{\pi R^2}\ln[(R^2-r_t^2)/\tau_0^2] \;. \quad (33)$$

This expression should now be inserted into Eq. (19) for $\bar{F}(b)$. The resulting expression still depends on the unmeasured variable r_t, the distance from the collision axis at which the jets appear. One should in principle average $\exp[\bar{F}(b,r_t)]$ over r_t. However, in order to get a simple expression, we have merely replaced the quantity $R^2-r_t^2$ by its average over r_t, which is $R^2/2$. Since this quantity sits in a logarithm, one may expect this to give a fair approximation. The resulting formula is then

$$\bar{F}(b)=\frac{1}{\sin(\theta)}\sigma_{\rm pl}\frac{dN}{dy}\frac{1}{\pi R^2}(\bar{F}_s+\bar{F}_h)\ln\left[\frac{R^2}{2\tau_0{}^2}\right] . \quad (34)$$

In an experimental environment, in order to identify particles as associated with jets, those particles with small momentum along the jet axis one must explicitly subtract out. These particles primarily are associated with the hadron matter distribution, and have no origin in the jet itself. We therefore subtract out particles with a rapidity along the jet axis less than a specified amount determined by experimental cuts. We expect that a cutoff of about two units of rapidity along each jet axis should be sufficient.

We now will determine the minimum momentum and therefore the minimum rapidity along the jet axis for which particles may be unambiguously identified as belonging to the jet, and not as belonging to the tail of the low-p_t distribution associated with the typical soft-hadron production. To estimate the magnitude of this cut, we first notice that all of the particles with momentum along the jet axis satisfying $p>p_0$ are within a rapidity interval of M/p_0, with $M\sim 0.4$ GeV. The distribution of particles along the jet axis must exceed the background from low-p_t particle production to be detectable. We have, therefore, that

$$\frac{M}{p}\frac{dN^{AA}}{dy}e^{-p/M}<\frac{M}{p}\frac{dN^{\rm jet}}{dy} \;. \quad (35)$$

The factor of M/p on the right-hand side comes from converting dN/dp along the jet axis to dN/dy along the jet axis and multiplication by a factor of M on both sides of the equation to compensate for the $1/M$ normalization factor in dN/dp on the left-hand side of the equation. Assuming that the nuclear multiplicity is about $2A$ times that of the jets, we crudely obtain therefore that the cutoff in rapidity along the jet axis, y_0, is

$$y_0\sim\ln[\ln(2A)] \;. \quad (36)$$

The assumption that the jet multiplicity is $1/2A$ that of the nuclear should be adequate for the crude order-of-magnitude estimate we make here. Even for uranium, this cutoff is only about 2.

Because of momentum conservation, the transverse momentum lost by the remainder of the jet is therefore given by

$$C(k_t)=\int_0^{y_0}dy\frac{dN}{dy\,dk_t}\sim 4y_0\frac{e^{-k_t/M}}{M} \;. \quad (37)$$

The distribution $dP/dk_{t\eta}$ can then be obtained from Eq. (4) with $\bar{B}$ replaced by $\bar{B}+\bar{F}+\bar{C}$.

IV. RESULTS AND DISCUSSION

We have investigated the contributions to the jet acoplanarity due to the scattering of the jets in the plasma and the effects of the experimental cuts which must be imposed in order to be able to identify the jets without ambiguity. In our calculations, we have taken $\sigma_{\rm pl}=14$ mb. We have considered a variety of A and Q values.

The effect of the plasma is contained in the term $\bar{F}(b)$. $\bar{F}(b)$ depends explicitly on three variables: A, the mass number of the colliding nuclei, the jet emission angle θ, and the initial entropy density proportional to dN/dy. The A dependence is fairly weak and comes mostly from the ratio $(dN/dy)/\pi R^2$ which grows like $A^{1/3}$ (there is an extra $\ln A$ dependence buried in the term $\ln R^2$). This effect is however amplified when the expotential of F is taken to generate the acoplanarity. To generate curves we have assumed that $dN/dy = 2A\, dN/dy\,|_{pp} \sim 8A$.

The effect of varying A is shown for 20-GeV jets in Figs. 3(a)—3(d). In these figures, the value of the acoplanarity is plotted for the case of rescattering from the plasma with experimental cuts, and experimental cuts only. In both curves, the preemission bremsstrahlung contribution is taken into account. All of the curves are for jet production at 90° relative to the collision axis, so that the effect of the plasma is minimized. Notice that for all values of A in the range $20 < A < 200$, there is a significant effect due to the plasma, which broadens the distribution and lowers it at small $k_{t\eta}$. For A of 200, the distribution is broadest, about a factor of 2 broader than is the case with no plasma. For A of 20, the distribution seems only a factor of 20% broader. The effects of A are therefore significant.

In Figs. 4(a)—4(c) the effects of Q of the jet are determined. Here we have the same plots as in Figs. 3(a)—3(d) except that three different values of $Q=10$, 20, and 40 GeV are considered for $A=100$. For $Q=10$ GeV, the acoplanarity distribution is almost flat, suggesting that the plasma has destroyed any jettiness in the distribution. The computed distribution is probably not quantitatively correct since our computations are strictly speaking only valid if $k_{t\eta}/Q \ll 1$, a condition not satisfied for this case. For $Q=20$ GeV, the distribution is jetlike, but the distribution with the plasma is about twice as broad as that of the distribution without the plasma. For $Q=40$, the distribution without the plasma becomes closer to that of the case with the plasma, and is perhaps about 30—40% broader.

It appears that there is a significant contribution due to rescattering in the plasma. A good measurement of these distributions may ultimately give a measure of the gluon-plasma cross section. Our computations show a small shift in the distributions as a function of cross section, but other contributions, such as the rescattering of inelastically produced particles from the jet in the plasma must be added in to make a precise comparison. A difference of a factor of 2 in the cross section effectively modifies the distribution as a shift in A by a factor of 8. The difference of a cross section by a factor of 2 therefore produces a change equivalent to changing $A=20$ to $A=200$. From Fig. 3 we see that this is in fact a quite significant shift.

Finally, we have studied the dependence of the jet acoplanarity upon jet opening angle. We find little dependence upon the opening angle for reasonable value of the

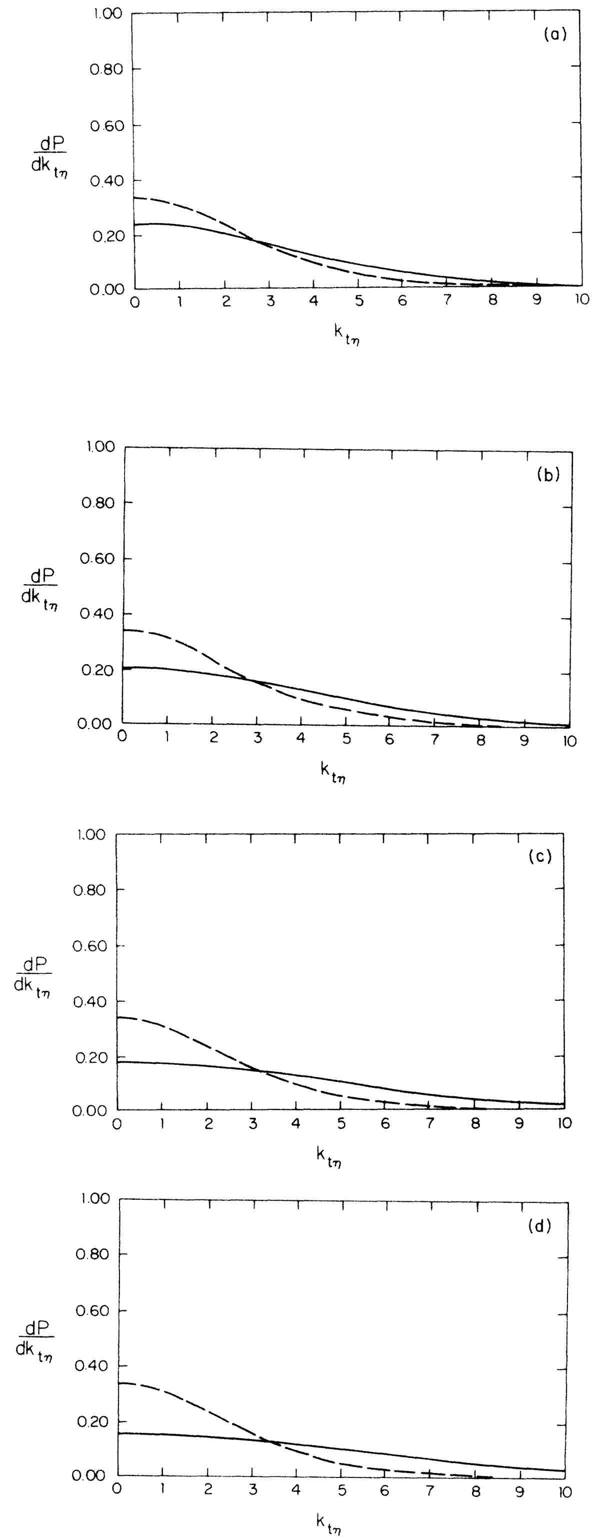

FIG. 3. These figures illustrate how the distribution $dP/dk_{t\eta}$ depends on the baryon number of the colliding nuclei. The dashed curve in each plot represents the distribution computed in the absence of a plasma with experimental cuts included. The solid curve contains the contribution from the plasma plus cuts. In both curves the effects of gluon bremsstrahlung are taken into account. The jet mass is 20 GeV and the assumed cross section is 14 mb. The values of A are (a) 20, (b) 50, (c) 100, (d) 200.

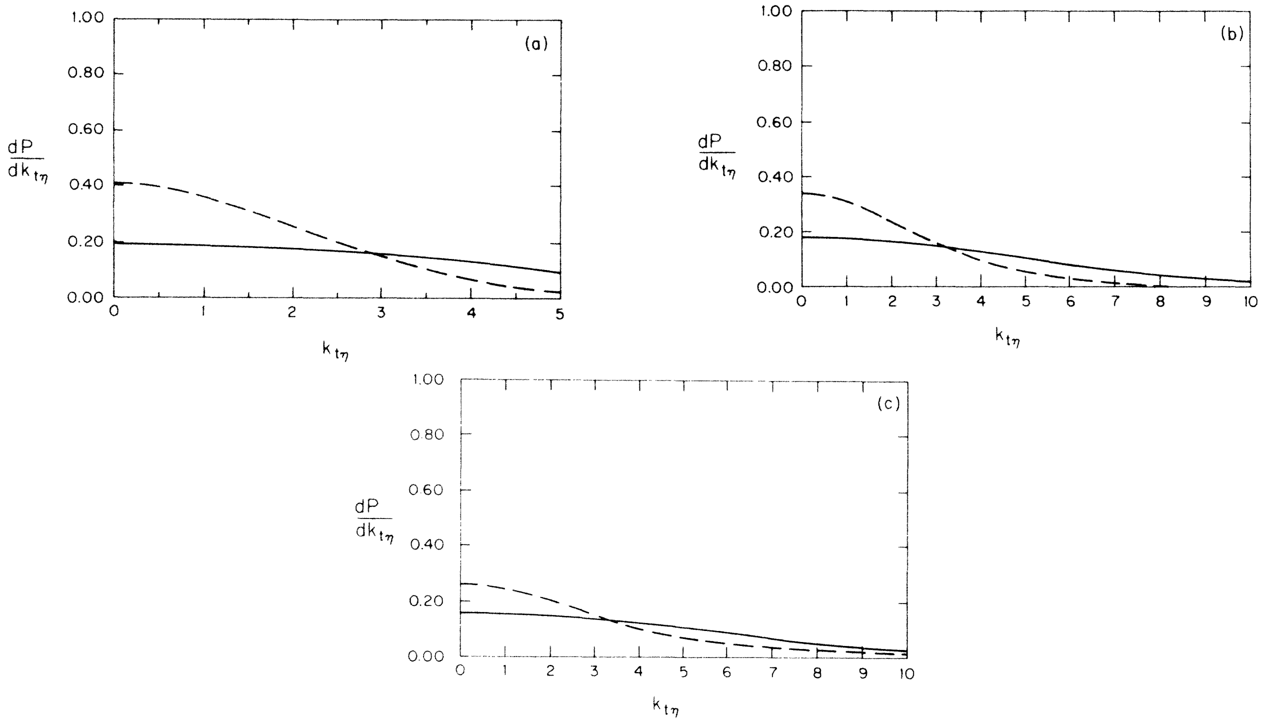

FIG. 4. For $A=100$, the acoplanarity distributions for (a) $Q=20$ GeV, (b) $Q=20$ GeV, (c) $Q=40$ GeV.

angle. This seems to be because the probability distribution is normalized to one, and the width is not too strongly dependent upon the value of θ for $\theta > 30°$. In this range, the value of $\sin(\theta)$ changed by only a factor of 2, and therefore the distribution is not too rapidly varying. A change in θ by this large amount corresponds to an effective increase in the gluon-plasma cross section by a factor of 2. For our choice of $\sigma = 14$ mb, there is not too large a variation if the cross section increases by a factor of 2. Also, most of the variation takes place only at small angles, where experiments are most difficult. If the cross section were chosen to be a little smaller, then there would be a somewhat larger variation. Although the angular variation may be useful for studying the effects of a plasma upon jet acoplanarity, without more knowledge of the range of angles which may be observed with cuts imposed, it is difficult to make any precise statement.

V. CONCLUSIONS

In summary, our analysis supports Appel's conclusion that jets may provide a useful diagnostic tool for studying the quark-gluon plasma. We have shown that in high-energy nuclear collisions, the effects of jet rescattering do in fact appear in the acoplanarity distribution. The cross section for scattering from the plasma may be inferred. We should be careful to note, however, that the existence of acoplanarity does not by itself alone give evidence for a quark-gluon plasma, and may in fact be generated by scattering from a hadronic gas. The jet acoplanarity is therefore not a signal for the plasma, merely a diagonstic tool. The utility of this tool for determining the gluon-plasma cross section, and hence inferring the degree of thermalization of the plasma, if it exists, remains to be established.

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