Recombination of shower partons at high p_T in heavy-ion collisions

Rudolph C. Hwa[1] and C. B. Yang[1,2]

[1]*Institute of Theoretical Science and Department of Physics, University of Oregon, Eugene, Oregon 97403-5203, USA*

[2]*Institute of Particle Physics, Hua-Zhong Normal University, Wuhan 430079, People's Republic of China*

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A formalism for hadron production at high p_T in heavy-ion collisions has been developed such that all partons hadronize by recombination. The fragmentation of a hard parton is accounted for by the recombination of shower partons that it creates. Such shower partons can also recombine with the thermal partons to form particles that dominate over all other possible modes of hadronization in the $3 < p_T < 8$ GeV range. The results for the high p_T spectra of pion, kaon, and proton agree well with experiments. Energy loss of partons in the dense medium is taken into account on the average by an effective parameter by fitting data, and is found to be universal independent of the type of particles produced, as it should. Due to the recombination of thermal and shower partons, the structure of jets produced in nuclear collisions is different from that in *pp* collisions. The consequence on same-side correlations is discussed.

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I. INTRODUCTION

In the production of hadrons at high p_T in heavy-ion collisions there are by now three theoretical collaborations that have shown the importance of quark recombination [1–3]. While there are some differences among the three approaches, they all agree in the basics and in the successful interpretation of the experimental data, among which the most outstanding ones are the p/π ratio being around 1 in the $3 < p_T < 4$ GeV range [4] and the scaling law of elliptic flow in the number of constituents [5,6]. The differences are concerned mostly with the treatment of hadronization in the $4 < p_T < 8$ GeV range. In this paper we show how the particles produced in that range arise from the recombination of the thermal partons and the shower partons created by hard partons at high p_T. The determination of the shower parton distributions (SPD) has recently been achieved by studying the fragmentation functions (FF) in the framework of the recombination model [7]. This paper contains the first application of the SPD outside the realm of parton fragmentation. As a consequence we find a new component that stands between the recombination of soft thermal partons at lower p_T and the fragmentation of hard partons at higher p_T [3], and is also different from the direct recombination of soft and hard partons [2].

In Ref. [1] we have avoided the need to specify the origin of the partons before recombination in order to be independent of the models describing the early and intermediate phases of the evolution in heavy-ion collisions. We use the measured pion spectra to infer the quark and antiquark distributions just before hadronization, and then use those distributions to determine the proton spectrum. In that way the calculated p/π ratio is a direct consequence of the recombination model with essentially no dependence on the other aspects of the dynamics such as the separation into soft and hard components. In this paper we do enter into the origins of the partons. The main difference between our approach and the conventional treatment of particle production at high p_T is that we do not use FF to represent the hadronization of hard partons produced at higher p_T. In our view all hadrons are produced by recombination at any p_T. FFs are phenomenological functions that do not specify the hadronization mechanism. Although string models can be useful in the description of hadronization of a pair of q and $\bar{q}$ receding from each other after being created in vacuum in, for example, e^+e^- annihilation, they cannot be applied to heavy-ion collisions where the abundance of color charges renders invalid any notion of stretched color flux tubes between pairs of partons produced in hard collisions [8]. In the absence of such models to give a conceptual basis for fragmentation, it is necessary to find a meaningful hadronization scheme for hard partons at high p_T.

Our view is that a hard parton creates a shower of partons that recombine subsequently to form hadrons. Although the branching process of gluon radiation and pair creation that eventually lead to shower partons at low virtuality cannot be calculated, the distributions of the shower partons can nevertheless be determined from known FFs in the framework of the recombination model, analogous to how the q and $\bar{q}$ distributions are determined from the experimental distribution $dN/p_T dp_T$ of pions, as done in Ref. [1]. In Ref. [7] a variety of SPDs are given as functions of the momentum fractions of the shower partons. In principle, there should be dependences on Q^2; however, for use in heavy-ion collisions at the relativistic heavy ion collider (RHIC), the SPDs given in Ref. [7] for $Q^2 = 100$ GeV^2 is adequate, since the dependence on Q^2 is not severe.

With the SPDs at hand, we can now consider their role in heavy-ion collisions. Hard scattering of partons gives rise to partons with high p_T, which undergo energy degradation as they traverse the dense medium [9,10]. Instead of tracing the spatial coordinates of the hard partons, we shall use an effective parameter ξ to describe the average fraction of partons that escape the dense medium and are able to hadronize outside. The value of ξ will be determined phenomenologically, and will represent an independent check on the degree of jet quenching [11]. A more important issue is the p_T dependences of the hadrons detected and the origins of the partons that contribute to the formation of those hadrons at different regions of p_T. The recombination of thermal and shower partons will be shown to be important with the consequence that the structure of jets produced in heavy-ion collisions is different from that produced in *pp* collisions. That difference will manifest itself in the correlation between particles in the same jet.

The recombination of two shower partons in the same jet is the same as the usual fragmentation of hard partons and becomes important at higher p_T. The recombination of two shower partons arising from two neighboring hard partons is also possible, but will not become important until the collision energy is very high, e.g., at the large hadron collider (LHC).

In what follows the production of mesons (pions and kaons) will be considered first, and then the baryons (specifically proton). Same-side correlation will be discussed, but only qualitatively to explain some apparent puzzle in the data.

II. MESON PRODUCTION BY RECOMBINATION

A. General considerations

In an overview of a heavy-ion collision we can divide the process into two stages: (1) the initial evolutionary phase, and (2) the final hadronization phase. Our concern in this paper is mainly on the latter phase, although what partons recombine depends on the former. A proper formulation of the formation of a dense medium must be done in full three-dimensional (3D) space and time, but to describe hadronization at large transverse momentum can be much simpler. Since the spatial volume in which recombination can occur is small, any two partons that are not collinear cannot recombine. Thus, collinear parton momenta in one dimension (1D) are all that we need to consider for the production of a particle with a particular momentum $\boldsymbol{p}$. The spatial extension along that direction is also constrained to the size of the hadron, so an integration of the remaining spatial variable results in a momentum-space description in which the parton momenta should reflect the wave function in momentum-space representation of the produced hadron. We can therefore formulate the recombination process in the one-dimensional (1D) momentum space only. Starting with $\boldsymbol{p}$ we investigate the partons that can recombine to form the particle at that $\boldsymbol{p}$. Of course, by not starting with a 3D space from the outset, we cannot calculate the density and expansion properties of the soft thermal partons, which will have to be introduced phenomenologically. A more elaborate formulation can be given in terms of Wigner functions in six-dimensional (6D) coordinate-momentum space [2,3]. Such formulations do provide a more complete picture of the collision process with the added advantage of being able to specify the soft component dynamically. It should be emphasized that whereas the recombination process is considered by us in 1D, it does not mean that the formalism is insensitive to the realistic problem of heavy-ion collisions and the 3D nature of the colliding nuclei. We shall return to this point when we discuss centrality dependence. As far as the hadronization part of the problem is concerned, our formulation in the 1D momentum space is consistent with the result of the formulation in 6D coordinate-momentum space after integration over the other five variables.

In our present problem of heavy-ion collisions we shall consider particle production in the transverse plane at rapidity $y=0$ only. Although the transverse plane is two dimensional, we shall only consider the direction in which a hadron is detected so that all partons relevant to the formation of such a particle move in the same direction. Thus, it is unnecessary to carry the subscript T on all our momentum labels to denote "transverse." The invariant phase space element is therefore dp/p^0, which we shall approximate by dp/p for relativistic particles.

In the recombination model [12] the invariant inclusive distribution for a produced meson with momentum p is

$$p\frac{dN_M}{dp} = \int \frac{dp_1}{p_1}\frac{dp_2}{p_2}F_{q\overline{q}'}(p_1,p_2)R_M(p_1,p_2,p), \tag{1}$$

where $F_{q\overline{q}'}(p_1,p_2)$ is the joint distribution of a quark q at p_1 and an antiquark $\overline{q}'$ at p_2, and $R_M(p_1,p_2,p)$ is the recombination function (RF) for $q\overline{q}'\to M$. For central collisions it is immaterial which direction $\boldsymbol{p}$ points. Although we only consider $\boldsymbol{p}$ in the transverse plane here, Eq. (1) has been used in the longitudinal direction in Refs. [12,13], where the RF is specified. It is

$$R_M(p_1,p_2,p) = \frac{1}{B(a+1,b+1)}\left(\frac{p_1}{p}\right)^{a+1}\left(\frac{p_2}{p}\right)^{b+1} \times \delta\left(\frac{p_1}{p}+\frac{p_2}{p}-1\right), \tag{2}$$

where $B(m,n)$ is the beta function. For pion it is shown in Ref. [13] from the analysis of Drell-Yan production data in pion-initiated process [14] in the framework of the valon model [12] that $a=b=0$. Thus, we have explicitly

$$R_\pi(p_1,p_2,p) = \frac{p_1p_2}{p^2}\delta\left(\frac{p_1}{p}+\frac{p_2}{p}-1\right), \tag{3}$$

The statistical factor for the recombination process is 1 [1]. For kaon it follows from the constituent quark masses of the valons and kaon-initiated inclusive production [13] that $a=1$ and $b=2$.

The δ function in Eq. (2) guarantees the conservation of momentum in the recombination process. R_M is an invariant distribution that is related to the non-invariant probability density $G_M(y_1,y_2)$ of finding the two valons in M with momentum fractions y_1 and y_2 by

$$R_M(p_1,p_2,p) = y_1y_2G_M(y_1y_2), \quad y_i = p_i/p. \tag{4}$$

Thus, for pion $G_\pi(y_1y_2)$ is a constant, apart from the δ function. It means that in the momentum space the wave function of pion in terms of the valons is very broad, corresponding to the pion being a tightly bound state of its constituent quarks. That is not the case for kaon. The broadness of $G_\pi(x_1x_2)$ is an important reason why the thermal-shower recombination in the formation of pions makes a dominant contribution in the intermediate p_T range, as we shall see below.

It should be noted that R_π given in Eq. (3) differs from the RF given in Ref. [1], where our formulation of the high-p_T problem is in the two-dimensional (2D) transverse plane.

Even the dimension of R_π in Ref. [1] is equivalently p_T^{-2}, whereas it is dimensionless in Eq. (3). Thus, caution should be exercised in relating our treatment here to that in Ref. [1].

It is important to recognize that the RF describes the probability of recombination of quarks (and antiquarks). In the recombination model the gluons are not regarded as partons that can directly hadronize [12,13]. They hadronize through the q and $\bar{q}$ channels by conversion to $q\bar{q}$ pairs. Thus, $F_{q\bar{q}'}$ in Eq. (1) must include all the q and $\bar{q}'$ generated by gluon conversion for the purpose of hadronization through the use of R_M in Eq. (1). The sea is therefore saturated by complete conversion from the gluons, a procedure that leads to the correct inclusive cross section for the hadronic production of pions, both in normalization [12] and in the momentum distribution [12,13]. The saturation of the sea of q and $\bar{q}'$ makes possible that the momenta of the gluons (which carry roughly half of the momentum of each nucleon) can be properly accounted for in the produced hadrons in any hadronic or nuclear collision. This procedure of saturating the sea is also followed implicitly in Ref. [7] for the determination of the shower partons in the fragmentation of an initiating parton. We also note that the conversion of gluons to quark pairs increases the number of degrees of freedom of all the partons in the medium and thereby overcomes the decrease of entropy that one might otherwise conclude by focusing on individual recombination processes.

Last, we mention that soft gluon emission and absorption are always possible in a recombination process; in fact, that is how color mutation can take place for a $q\bar{q}'$ pair to become colorless, while the partons dress themselves to become valons before forming a meson. Such soft processes do not change the momenta of valons and therefore do not influence the momentum consideration in Eq. (2). We do not consider quark-antiquark-gluon recombination with large momentum fraction for the gluon because the valon representation is complete. That is precisely the reason why the valon model was constructed in the beginning to describe hadron structure on the one hand as well as the recombination function (for the time-reversed process) on the other [12].

B. Pion distribution

Restricting our attention to pion production for now, we obtain from Eqs. (1) and (3):

$$\frac{dN_\pi}{pdp} = \frac{1}{p^3}\int_0^p dp_1 F_{q\bar{q}'}(p_1, p-p_1), \tag{5}$$

which clearly exhibits the simple dependence of the pion spectrum on the momentum distributions of the q and $\bar{q}'$ that recombine. It should be recognized that, since we work in 1D, dN_π/pdp in Eq. (5) is actually the number density of pions in $p_T\,dp_T\,dy\,d\phi$ evaluated at $y=0$. With the assumption that it is independent of ϕ in central collisions, what we denote as dN_π/pdp is equivalent to the experimental $dN/(2\pi p_T\,dp_T)$, where the number N refers to the integrated result over all ϕ so that when divided by 2π the average density in 2D is the same as what we calculate in 1D.

Note that in Eq. (5) p_1 is integrated over a wide range due to the broadness of the RF. There are two components of parton sources that contribute to $F_{q\bar{q}'}$. One is thermal $(\mathcal{T})$ and the other shower $(\mathcal{S})$. Before giving the specifics of what they are, let us first express $F_{q\bar{q}'}$ in terms of them in a schematic way

$$F_{q\bar{q}'} = \mathcal{TT} + \mathcal{TS} + (\mathcal{SS})_1 + (\mathcal{SS})_2. \tag{6}$$

All four terms make contributions at all p_T, although each is important in only restricted regions of p_T. $(\mathcal{SS})_1$ denotes two shower partons arising from one hard parton (hence, within one jet), and can be related through the use of R to the usual fragmentation that is described by the D function. $(\mathcal{SS})_2$ denotes two shower partons that are from two separate but nearby hard partons, and are therefore associated with two overlapping jets. $(\mathcal{SS})_2$ is not expected to be important unless the density of hard partons is extremely high, such as that possibly at LHC. For brevity, we shall write $(\mathcal{SS})_1$ simply as $\mathcal{SS}$, when no confusion is likely to arise. $\mathcal{TT}$ signifies two thermal partons whose recombination yields the thermal hadrons, usually referred to as the soft component. $\mathcal{TS}$ denotes thermal-shower pairing and is the new component that has never been considered before. It turns out to be important in the $3 < p_T < 8$ GeV range. We emphasize that the $\mathcal{TS}$ term would be absent if we do not treat the fragmentation of a hard parton as the recombination of shower partons as done in Ref. [7]. It is now evident that by considering $\mathcal{S}$ as the showering effect of hard partons all hadrons are produced by recombination, as is made explicit by Eqs. (5) and (6) in the case of pions. Let us now specify what $\mathcal{T}$ and $\mathcal{S}$ represent.

1. Thermal component

$\mathcal{T}$ is the thermal component including hydrodynamical flow. It is not our intention to derive soft parton distribution from hydrodynamics. We shall simply assume what is necessary to give rise to the observed distribution of pions for $p_T < 2$ GeV. Since the observed dN_π/pdp at low p is exponential, the necessary invariant parton distribution for the thermal component is

$$\mathcal{T}(p_1) = p_1\frac{dN_q^{\rm th}}{dp_1} = Cp_1\exp(-p_1/T), \tag{7}$$

where T is the inverse slope enhanced by flow. C is the normalization factor to be adjusted to fit the pion data at low p, and has the dimension $[\text{momentum}]^{-1}$. For noncentral collisions, which we do not consider in this paper, C would depend on centrality. We assume that the thermal partons are uncorrelated, except in special circumstances, so that we may write $F_{q\bar{q}'}^{\rm th}$ in the factorizable form

$$F_{q\bar{q}'}^{\rm th}(p_1,p_2) = \mathcal{T}(p_1)\mathcal{T}(p_2) = C^2 p_1 p_2\,\exp[-(p_1+p_2)/T]. \tag{8}$$

Substituting this into Eq. (5) yields for the thermal component of the pion distribution

$$\frac{dN_\pi^{\rm th}}{pdp} = \frac{C^2}{6}\exp(-p/T), \tag{9}$$

which is the exponential form aimed for. To fit the data the parameters are

$$C = 23.2\ \text{GeV}^{-1}, \quad T = 0.317\ \text{GeV}, \tag{10}$$

as we shall see later.

2. Shower partons

Next we consider the shower distribution. In a heavy-ion collision let the parton i be scattered into the transverse plane at rapidity $y=0$ with probability $f_i(k)$, i.e.:

$$\left.\frac{dN_i^{\rm hard}}{kdkdy}\right|_{y=0} = f_i(k), \tag{11}$$

where k is the transverse momenta of the parton. We shall use the parametrization of $f_i(k)$ given in Ref. [15], obtained for the study of dilepton production in central collision of gold nuclei at $\sqrt{s_{NN}}=200$ GeV. Due to energy loss of the partons in the dense medium, not all partons emerge from the reaction zone to hadronize outside. Only a fraction of them do, and we shall use ξ to represent the effective fraction after averaging over all central events such that $\xi f_i(k)$ denotes the number of unquenched partons with momentum k that are to hadronize. Note that we do not refer to them as jets, since the notion of jets presupposes that a hard parton fragments into a jet of hadrons. That supposition is, of course, just what we want to avoid. We regard ξ as an effective fraction because we do not consider its dependence on k and do not delve into the space-time properties of the hard scattering and subsequent evolution. The value of ξ will be determined phenomenologically, and be regarded as an empirical quantification of the degree of energy loss.

The hard parton i at momentum k creates a parton shower. Since there are various types of shower partons for each type of initiating parton i, let us use S_i^j to denote the matrix of SPDs for $i \to j$. Although i can be u, d, s, $\bar{u}$, $\bar{d}$, $\bar{s}$, and g, j is allowed to be quark and antiquarks, but not gluon. The role of gluons in the recombination process has already been discussed earlier at the end of Sec. II A. The sea partons in the shower are saturated by gluon conversion. In the notation of Ref. [7] we use K_{NS} to denote the valence quark in the shower, $L(L_s)$ the light (strange) sea quarks in a quark-initiated shower, and $G(G_s)$ the light (strange) sea quarks in a gluon-initiated shower. Thus, the shower matrix S_i^j has the form

$$S_i^j = \begin{pmatrix} K & L & L_s \\ L & K & L_s \\ L & L & K_s \\ G & G & G_s \end{pmatrix}, \quad i = u,d,s,g, \quad j = u,d,s, \tag{12}$$

where $K=K_{NS}+L$ and $K_s=K_{NS}+L_s$. The antiquarks $\bar{u}$, $\bar{d}$, and $\bar{s}$ have the same structure, and are related to u, d, and s as sea, and vice-versa. The parametrizations for these SPDs have been completely determined in Ref. [7] as functions of the momentum fraction z of parton j in parton i. Note that in Eq. (12) there is no fourth column corresponding to gluons in the shower. The determination of the five essential SPDs is carried out on the condition of no shower gluons so that all hadrons specified by the fragmentation functions are formed by $q\bar{q}'$ recombination without gluons. The underlying physics for this has already been discussed in the last two paragraphs of Sec. II A.

The distribution of shower parton j with transverse momentum p_1 in central heavy-ion collisions is then

$$\mathcal{S}(p_1) = \xi \sum_i \int_{k_0}^{\infty} dk k f_i(k) S_i^j(p_1/k), \tag{13}$$

where p_1 and k are collinear. Since the input on hard parton distribution $f_i(k)$ cannot be valid at low k, we shall consider the earlier integral only for $k > k_0$, for which we set the minimum at $k_0=3$ GeV. In practice we cutoff the upper limit of integration at 20 GeV.

3. Thermal-shower recombination

With the shower partons specified we can now combine them with the thermal partons to describe the $\mathcal{TS}$ term in Eq. (6). We have

$$\mathcal{T}(p_1)\mathcal{S}(p_2) = \xi\, C\, p_1 e^{-p_1/T} \sum_i \int dk\, k f_i(k) S_i^j(p_2/k), \tag{14}$$

where the distributions of thermal light quarks are assumed to be flavor independent and the appropriate one is implied to pair off with j to form the meson under consideration. The contribution to the pion spectrum from thermal-shower recombination is then, using Eq. (5):

$$\frac{dN_\pi^{\mathcal{TS}}}{pdp} = \frac{1}{p^3}\int_0^p dp_1 \mathcal{T}(p_1)\ \mathcal{S}(p-p_1), \tag{15}$$

where a sum over j is implied to match the flavors of the valence quarks of the detected pion. Only the overall normalization of this term depends on ξ. The p dependence is our prediction, which is used to fit the data and thereby determine ξ.

4. Shower-shower recombination

For two shower partons in the same jet we have

$$(\mathcal{SS})_1(p_1,p_2) = \xi \sum_i \int dk k f_i(k) \left\{ S_i^j\left(\frac{p_1}{k}\right), S_i^{j'}\left(\frac{p_2}{k-p_1}\right) \right\}, \tag{16}$$

where the curly brackets signify the symmetrization of the leading parton momentum fraction

$$\left\{ S_i^j(z_1), S_i^{j'}\left(\frac{z_2}{1-z_1}\right) \right\} = \frac{1}{2}\left[S_i^j(z_1) S_i^{j'}\left(\frac{z_2}{1-z_1}\right) + S_i^j\left(\frac{z_1}{1-z_2}\right) S_i^{j'}(z_2) \right]. \tag{17}$$

We have shown in Ref. [7] that the SPDs can be determined from the recombination formula for the fragmentation function

$$xD_i^M(x) = \int \frac{dx_1}{x_1}\frac{dx_2}{x_2}\left\{S_i^j(x_1), S_i^{j'}\left(\frac{x_2}{1-x_1}\right)\right\}R_M(x_1,x_2,x). \tag{18}$$

A sum over j and j' is implied in consort with the j and j' labels hidden in the RF that are relevant for M. The substitution of Eq. (16) for the $(\mathcal{SS})_1$ term in Eq. (6) into Eq. (1) clearly yields

$$p\frac{dN_M^{\text{frag}}}{dp} = \xi\sum_i \int dk k f_i(k)\frac{p}{k}D_i^M\left(\frac{p}{k}\right), \tag{19}$$

which is the usual formula for the production of a meson at high p_T in the fragmentation model, except for the presence of ξ here for reasons that have already been discussed earlier.

There is finally the possibility of recombination of two shower partons from two different but partially overlapping showers. The corresponding $(\mathcal{SS})_2$ term in Eq. (6) should then be

$$(\mathcal{SS})_2(p_1,p_2) = \delta_{y\phi}\xi^2\sum_{i,i'} \times \int dkdk' kk' f_i(k) f_{i'}(k') S_i^j\left(\frac{p_1}{k}\right) S_{i'}^{j'}\left(\frac{p_2}{k'}\right), \tag{20}$$

where a multiplicative factor $\delta_{y\phi}$ is included to reflect the probability of overlap in y and ϕ of the two showers in order for collinear recombination of the partons j and j' to take place. The value of $\delta_{y\phi}$ can be estimated by studying the size of the jet cone, and is expected to be small. Thus, this mode of recombination is not likely to be important at RHIC. However, at very high energy, such as at LHC, where $f_i(k)$ is orders of magnitude higher, the $(\mathcal{SS})_2$ term may well become significant.

5. Result on the pion spectrum

Collecting all the pieces of Eq. (6) together and substituting them in Eq. (5), we obtain the four contributions to the pion spectrum. The parameters C and T in Eq. (10) are determined by fitting the low-p_T data. Ignoring the $(\mathcal{SS})_2$ contribution on the basis that $\delta_{y\phi}$ is very small, there is only one parameter, ξ, to adjust to fit the data for $p_T > 2$ GeV. The result is shown in Fig. 1. With the value

$$\xi = 0.07 \tag{21}$$

the fit of the π^0 data from PHENIX [16] on central Au-Au collisions at $\sqrt{s_{NN}}=200$ GeV is excellent up to $p_T \simeq 10$ GeV. Note that in the region $3 < p_T < 8$ GeV, the dominant contribution is from thermal-shower recombination (line with crosses). The conventional jet fragmentation is from shower-shower recombination in one jet (line with circles); it becomes more important than the $\mathcal{TS}$ contribution only for $p_T > 9$ GeV. To show the relative size of the $(\mathcal{SS})_2$ recombination from two jets, we assume $\delta_{y\phi}=0.01$ (a very rough estimate) just to put its contribution on the figure. It is indicated by the line with squares, which is much lower than all others.

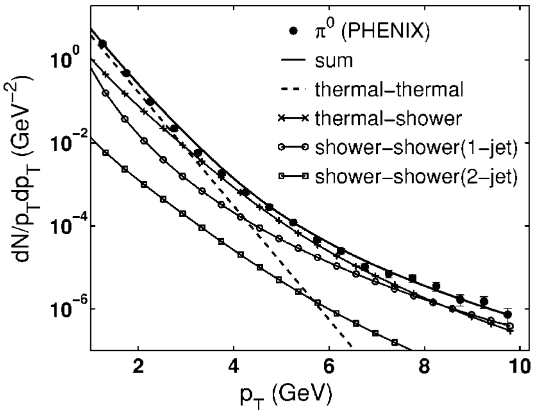

FIG. 1. Transverse momentum distribution of π^0 in Au-Au collisions. Data are from Ref. [16]. The solid line is the sum of four contributions to the recombination of partons: $\mathcal{TT}$ (dashed line); $\mathcal{TS}$ (line with crosses); $(\mathcal{SS})_1$, two shower partons in one jet (line with open circles); and $(\mathcal{SS})_2$, two shower partons from two overlapping jets (line with squares).

The result that thermal-shower dominates over shower-shower (1-jet) recombination for $p_T < 8$ GeV is our main finding in this work. It shows the importance of considering the interaction between the thermal partons and the partons created by hard scattering. That interaction becomes particularly significant at the hadronization scale where recombination occurs.

One way to see why $\mathcal{TS}$ is greater than $\mathcal{SS}$ in their contributions to Eq. (1) via Eq. (6) is to examine their sum in the following form:

$$\mathcal{TS} + \mathcal{SS} = \xi\sum_i \int dk\ k f_i(k)\ S_i^{j'}\left(\frac{p_2}{k}\right) \times\left[Cp_1 e^{-p_1/T} + S_i^j\left(\frac{p_1}{k-p_2}\right)\right], \tag{22}$$

putting aside the other symmetrizing term in $\mathcal{SS}$ without any impediment to our argument. The first (thermal) term inside the square bracket is much larger than the second (shower) term when p_1 is small (but not infinitesimal), which is a region of p_1 that is relevant for the soft parton to recombine with a shower parton at $p_2 > 3$ GeV only if $R_\pi(p_1,p_2,p)$ is broad enough to encompass both. This linking between soft and semihard partons not only enhances the spectrum over simple fragmentation, but also modifies the structure of what is usually regarded as minijet. Since there are no thermal soft partons in pp collisions, the pion distribution at intermediate and high p_T in pp collisions must differ from that in AA collisions. Furthermore, the same-side correlations in the two cases are necessarily different, as we shall discuss in Sec. IV.

6. Dependence on the recombination function

To see how the thermal-shower recombination depends on the broadness of RF, consider the general form of the valon distribution

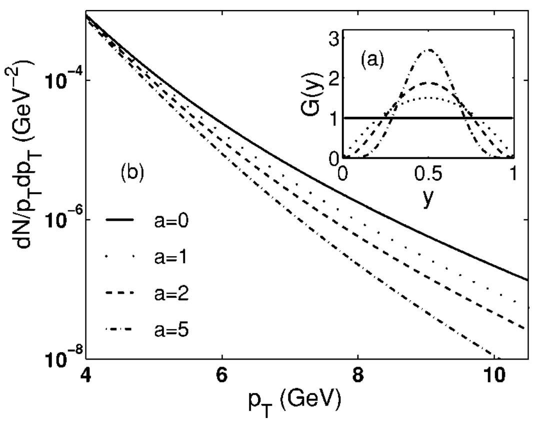

FIG. 2. (a) Various valon distributions in momentum fraction y according to Eq. (24). (b) The corresponding p_T distributions from $\mathcal{TS}$ recombination.

$$G(y_1,y_2) = \frac{1}{B(a+1,a+1)}(y_1 y_2)^a \; \delta(y_1+y_2-1). \quad (23)$$

Equations (3) and (4) follow from Eq. (23) for $a=0$. If a were larger, as one would expect for the ρ meson (since small a corresponds to a tightly bound state), then the valon distribution would be more sharply peaked at $y_1=y_2=1/2$. In that case the recombination of a thermal parton at low p_1 and a shower parton at intermediate p_2 is suppressed. To make this point transparent, we show in Fig. 2(a) several possible widths of a single-valon distribution $G(y)$ obtained from $G(y_1,y_2)$ by one integration, i.e.:

$$G(y) = \frac{1}{B(a+1,a+1)}[y(1-y)]^a, \quad (24)$$

which is normalized to 1 by one more integration. The corresponding RF is given by Eqs. (4) and (23). When we use that RF in Eq. (1) and calculate the distribution dN/pdp with only the $\mathcal{TS}$ contribution to $F_{q\bar{q}}$ taken into account, the result is shown in Fig. 2(b) for $a=0,1,2,5$. It is evident that the recombination of thermal-shower partons is significantly suppressed at high p_T when a is large, and becomes negligible when $R(p_1,p_2,p)$ tends toward being proportional to $\delta(p_1-p/2)\delta(p_2-p/2)$.

C. Energy loss

In the absence of a space-time study of the problem that includes the locations where hard collisions occur, it is not possible to consider the medium effect on each and every hard parton that traverses the medium. We have used the parameter ξ to represent the overall effect after averaging over all events in central collisions. The value $\xi=0.07$ given in Eq. (21) is a quantity deduced from fitting the pion spectrum. It clearly indicates that not all hard partons created in a heavy-ion collisions can get out of the dense medium to hadronize. As suggested by the STAR data [17], only those near the surface can escape the quenching effect. Our result on the value of ξ seems low by comparison to the nuclear modification factor $R_{AA}(p_T)$, which is roughly 0.2 for $p_T>6$ GeV in central Au-Au collisions. To understand their difference, let us examine what they measure, respectively.

$R_{AA}(p_T)$ is defined by the ratio

$$R_{AA}(p_T) = \frac{dN/p_T\,dp_T\;(AA)}{N_{\rm coll} dN/p_T\,dp_T\;(pp)}, \quad (25)$$

where $N_{\rm coll}$ is the average number of binary collisions. In the denominator the production of particles at high p_T in pp collisions can be well described by the fragmentation of hard partons. However, we have seen that the usual fragmentation corresponds to the shower-shower recombination in 1-jet, which is not the important part of hadronization in AA collisions for $3<p_T<8$ GeV. The suppression factor ξ may, in the spirit of Eq. (25), be written in the schematic form

$$\xi = \left\langle \frac{dN/p_T\,dp_T\;(AA)}{\int \mathcal{T}\hat{\mathcal{S}}R\;(AA)} \right\rangle, \quad (26)$$

where $\hat{\mathcal{S}}$ is the shower component expressed in Eq. (13), but without the ξ factor. The angular brackets denote an average over all p_T. The denominator is the thermal-shower recombination in AA collisions, if all hard partons get out of the medium to hadronize. It is then clear why ξ is smaller than R_{AA}. It should not only account for the suppressed number of hard partons in the numerator that get out from the dense medium to hadronize (which R_{AA} does also), but is also made smaller by the denominator where the $\mathcal{T}\hat{\mathcal{S}}$ recombination is larger than the scaled contribution from pp collisions.

Another way to exhibit the effect of energy loss on our result is to calculate $R_{AA}(p_T)$ directly from our pion distribution, assuming that the pions dominate over all other particles. That assumption is invalid for $p_T<4$ GeV, since the production of proton is known to be roughly equal to that of pion in the $3<p_T<4$ GeV region. Nevertheless, it is illuminating to see what the pion component gives (call it R_{AA}^{π}) relative to the experimental R_{AA} for $p_T>4$ GeV. Since $\mathcal{SS}$ recombination is just the fragmentation component in AA collisions as stated in Eq. (19), we can obtain $N_{\rm coll} dN_\pi/p_T\,dp_T(pp)$ simply by omitting the factor ξ from $dN_\pi^{\mathcal{SS}}/pdp$. That is, we have

$$R_{AA}^{\pi}(p) = \frac{dN_\pi/pdp}{\xi^{-1} dN_\pi^{\mathcal{SS}}/pdp}, \quad (27)$$

where the numerator corresponds to the solid line in Fig. 1 and $dN_\pi^{\mathcal{SS}}/pdp$ corresponds to the line with open circles in the same figure. The result for R_{AA}^{π} is shown in Fig. 3 by the solid line, which is in reasonable agreement with the data on R_{AA} [18] that is not limited to the pions. It is now clear that the ratio of those two lines is equal to the ratio R_{AA}^{π}/ξ, which in turn is approximately equal to R_{AA}/ξ. Since the thermal-thermal recombination is negligible for $p_T>4$ GeV, the ratio $(dN_\pi/pdp)/(dN_\pi^{\mathcal{SS}}/pdp)$ is essentially independent of ξ. Thus, R_{AA}^{π}/ξ is roughly independent of energy loss and therefore of centrality.

It is worth remarking on how the nuclear-size effect enters into our one-dimensional treatment of the hadronization process. We emphasize that hadronization occurs outside the

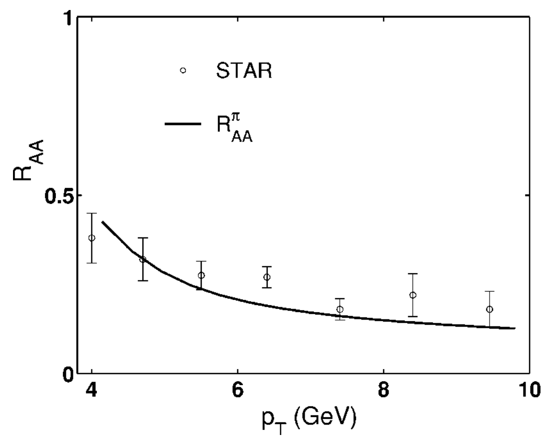

FIG. 3. Nuclear modification factor R_{AA} compared to the calculated result using only pions, R^{π}_{AA}.

volume of dense matter and along the direction of the detected particle. The 3D nuclear-size effect influences the properties of the partons *before* they enter into the interaction region for hadronization, and is therefore not explicitly present in the recombination formula. The recombining partons are either the soft thermal ones or the semihard shower partons. Centrality obviously affects the magnitude of the thermal source, i.e., C, and the degree of energy loss, i.e., ξ. We have made a preliminary study of the centrality dependence of dN_π/pdp and found agreement with the data by the appropriate adjustment of C and ξ in essentially the same way as we have done in this paper for the most central collisions. We mention this only to point out that the 1D description of recombination does not preclude the possibility of accounting for the nuclear-size effect, which is manifestly 3D, but occurs prior to the hadronization process.

The suppression factor ξ, as expressed in Eq. (26), cannot be measured directly. There is, however, a measurable quantity that is closely related to ξ. Since only the hard partons that are created near the surface of the collision region can get out, most jets detected in AA collisions do not have a partner in the opposite direction. That does not mean the nonexistence of back-to-back jets. A pair of hard partons that are created near the edge, but directed at around 90° relative to the radial position of creation measured from the center, and yet remain in the transverse plane (i.e., roughly tangent to the cylinder), do not go through the bulk of the medium. Those two hard partons can then lead to back-to-back jets that are detectable. It is then of interest to measure the total number of events that contain back-to-back jets. Let $R_{jj/j}(\overline{p_T})$ be the ratio of events with back-to-back jets to the total number of events containing any jets, where $\overline{p_T}$ is the minimum value of p_T that any particle must have to be counted as part of a jet. Thus, for example, for $\overline{p_T}=5$ GeV, $R_{jj/j}(\overline{p_T})$ refers to the fraction of events containing two particles with $p_T>5$ GeV in opposite directions out of all events with any particle having $p_T>5$ GeV. $R_{jj/j}(\overline{p_T})$ is then a measure of the fraction of space near the surface that can give rise to particles at high p_T through thermal-shower recombination. It can have dependence on centrality. The precise relationship between ξ and $R_{jj/j}$ is a separate problem worthy of detailed investigation, especially if the experimental determination of $R_{jj/j}$ is forthcoming.

In a Monte Carlo calculation, such as that employed in Ref. [2], where space-time trajectories of the hard partons are tracked, it is possible to compute the parton momenta after energy losses are taken into account. Those emergent partons then generate shower partons whose recombination with thermal partons presumably results in a yield that can check the value of our mean suppression factor ξ. In our treatment in the momentum space only, we determine ξ by fitting the pion data, but once fixed the relative magnitude between $\mathcal{TS}$ and $\mathcal{SS}$ is also fixed. Moreover, there is no more freedom in the determination of the inclusive distribution of other particles, such as kaon and proton.

D. Kaon production

For the production of kaon the RF has the explicit form that follows from Eq. (2) with $a=1$ and $b=2$ [13]:

$$R_K(p_1,p_2,p) = 12\frac{p_1^2p_2^3}{p^4}\delta(p_1+p_2-p). \tag{28}$$

Using this in Eq. (1) yields for K^+:

$$\frac{dN_{K^+}}{pdp} = \frac{12}{p^6}\int_0^p dp_1\, p_1(p-p_1)^2\, F_{u\bar{s}}(p_1,p-p_1). \tag{29}$$

Note that because the kaon is not as tightly bound as pion, the RF is not as broad as that for pion, with the consequence that the factor $p_1(p-p_1)^2$ in the integrand forces the u and $\bar{s}$ quarks to have closer momenta than those in Eq. (1).

There are four terms for $F_{u\bar{s}}$ as in Eq. (6). The calculational procedure is basically the same as for pion. A slight complication arises from the strange quark being different from the light quarks. In the thermal component we shall simply attach a multiplicative factor λ_s for the s sector

$$\mathcal{T}_s = \lambda_s\,\mathcal{T}, \tag{30}$$

where we set λ_s to be the Wròblewski factor at $\lambda_s=0.5$ [19]. The thermal-shower recombination now has two terms

$$\mathcal{T}(p_1)\mathcal{S}_s(p_2) + \mathcal{S}(p_1)\mathcal{T}_s(p_2) = \xi\, C\sum_i \int dk\, kf_i(k)\left[p_1e^{-p_1/T}\, S_i^{\bar{s}}\left(\frac{p_2}{k}\right) + \lambda_s p_2 e^{-p_2/T}S_i^{u}\left(\frac{p_1}{k}\right)\right]. \tag{31}$$

The $(\mathcal{SS})_1$ and $(\mathcal{SS})_2$ contributions are as expressed in Eqs. (16) and (20), respectively, with $\{j,j'\}$ identified as $\{u,\bar{s}\}$.

The four contributions to the kaon spectrum are shown in Fig. 4. No adjustable parameter has been used beyond what has already been determined in the previous section. The agreement between the sum (solid line) and the data is evidently very good. The data are for K_s^0 at 0–5% centrality and extend to $p_T\sim 6$ GeV [20], farther than for K^+ [4,21]. As in the case of pions the thermal-shower recombination is more important than both thermal-thermal and shower-shower (1-jet) for $3<p_T<8$ GeV. As in Fig. 1 the shower-shower (2-

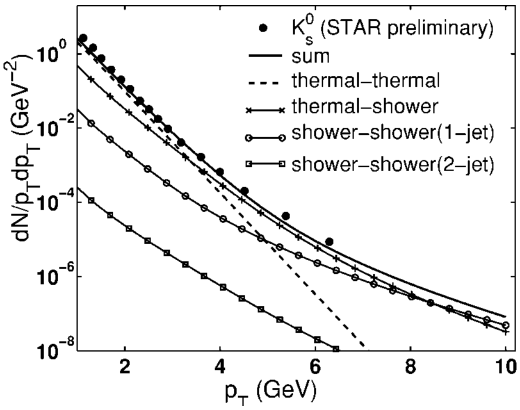

FIG. 4. Transverse momentum distribution of K^+ in Au-Au collisions. Data for K_s^0 are preliminary from Ref. [20]. The symbols for the four contributions are the same as in Fig. 1.

jet) curve is shown for $\delta_{y\phi}=0.01$ and is negligible. There are some small differences in the various components of the pion and kaon spectra, but on the whole the two are basically similar.

III. BARYON PRODUCTION BY RECOMBINATION

A. General considerations

The extension of the treatment in the preceding section to baryon production is conceptually straightforward. The generalization of Eq. (1) is clearly

$$p\frac{dN_B}{dp}=\int\frac{dp_1}{p_1}\frac{dp_2}{p_2}\frac{dp_3}{p_3}F(p_1,p_2,p_3)\ R_B(p_1,p_2,p_3,p), \tag{32}$$

where $F(p_1,p_2,p_3)$ is the joint distribution of three relevant quarks to form the baryon B. The RF is related to the noninvariant valon distribution by

$$R_B(p_1,p_2,p_3,p)=g_{\rm st}\ y_1y_2y_3\ G_B(y_1,y_2,y_3),y_i=p_i/p, \tag{33}$$

where $g_{\rm st}$ is a statistical factor, and $G_B(y_1,y_2,y_3)$ has the general form

$$G_B(y_1,y_2,y_3)=g_B\ y_1^{\alpha}y_2^{\beta}y_3^{\gamma}\ \delta(y_1+y_2+y_3-1), \tag{34}$$

$$g_B=[B(\alpha+1,\beta+\gamma+2)\ B(\beta+1,\gamma+1)]^{-1}. \tag{35}$$

For proton, y_1 and y_2 refer to the momentum fractions of the U valons, and y_3 to that of the D valon. The exponents α,β, and γ have been determined in Ref. [22] to be

$$\alpha=\beta=1.75, \quad \gamma=1.05 \tag{36}$$

from parton distribution functions of the proton. Note that although the U and D valons have the same constituent quark masses, they have different average momentum fractions calculable from Eq. (34) [22]:

$$\langle y\rangle_U=0.3644, \quad \langle y\rangle_D=0.2712, \tag{37}$$

which satisfies the sum rule

$$2\ \langle y\rangle_U+\langle y\rangle_D=1. \tag{38}$$

For hyperons, the lack of information about their parton distribution functions deprives us of any such detail knowledge of their valon distributions, and consequently of their RFs. Nevertheless, on the basis that the constituent quarks in all baryons are not tightly bound, we expect the exponents α,β and γ for hyperons to be also in the range between 1 and 2.

The 3-quark distribution now has more terms in the various possible contributions from the thermal and shower partons. Schematically, it takes the form

$$F_{qq'q''}=\mathcal{TTT}+\mathcal{TTS}+\mathcal{T}(\mathcal{SS})_1+(\mathcal{SSS})_1+\mathcal{T}(\mathcal{SS})_2 + (\mathcal{S}(\mathcal{SS})_1)_2+(\mathcal{SSS})_3. \tag{39}$$

They are arranged in increasing order of the number of hard partons involved: the first term has none, the next three one, the following two two, and the last term three. We shall consider only the first four terms, since they involve thermal partons and the shower of only one hard parton. The fifth term involves partons from two overlapping jets, and is ignored at RHIC energy despite the enhancement by $\mathcal{T}$.

B. Proton production

Let us focus our attention on proton production at high p_T. As before, we omit the subscript $\mathcal{T}$ in referring to momenta in the transverse plane, and obtain from Eqs. (32)–(35):

$$\frac{dN_p}{pdp}=\frac{g_p\ g_{st}}{p^{2\alpha+\gamma+4}}\int_0^p dp_1\int_0^{p-p_1}dp_2\ (p_1p_2)^{\alpha}(p-p_1-p_2)^{\gamma} \times F(p_1,p_2,p-p_1-p_2). \tag{40}$$

The $\mathcal{TTT}$ contribution can be computed analytically. Using Eq. (7) for each $\mathcal{T}$, we obtain for the thermal spectrum

$$\frac{dN_p^{\rm th}}{pdp}=\frac{C^3}{6}pe^{-p/T}\frac{B(\alpha+2,\gamma+2)B(\alpha+2,\alpha+\gamma+4)}{B(\alpha+1,\gamma+1)B(\alpha+1,\alpha+\gamma+2)}, \tag{41}$$

where C and T are given in Eq. (10), and α,γ in Eq. (36). The statistical factor $g_{\rm st}$ is 1/6 when the spin-flavor consideration is taken into account [1]. Equation (41) is not reliable at very small p, since the proton mass effect invalidates our essentially scale-invariant formulation for relativistic particles. Furthermore, at small p_T the problem becomes 3D, so our treatment should be modified accordingly.

The $\mathcal{TTS}$ contribution is

$$\frac{dN_p^{\mathcal{TTS}}}{pdp}=\frac{g_p\ C^2\ \xi}{6p^{2\alpha+\gamma+4}}\int dp_1dp_2\ (p_1p_2)^{\alpha+1}(p-p_1-p_2)^{\gamma+1} \times\sum_i\int dk\ k\ f_i(k)\ U_i(k,p_1,p_2,p), \tag{42}$$

where

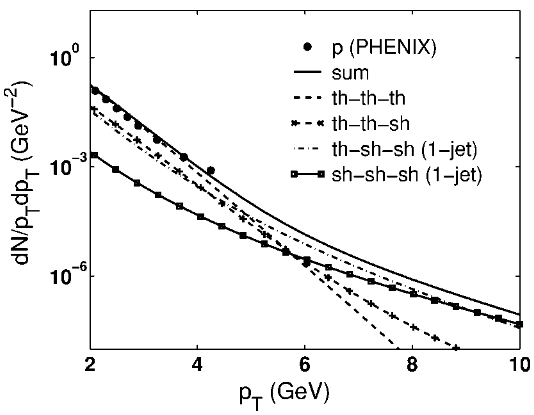

FIG. 5. Transverse momentum distribution of proton in Au-Au collisions. Data are from Ref. [16]. The solid line is the sum of four contributions to the recombination of partons: $\mathcal{TTT}$ (dashed line); $\mathcal{TTS}$ (line with crosses); $\mathcal{TTS}$, one thermal parton with two shower partons in one jet (dashed-dot line); and $\mathcal{SSS}$, three shower partons from one jet (line with squares).

$$U_i(k,p_1,p_2,p) = \frac{1}{p_1}e^{-(p-p_1)/T}\, S_i^u\left(\frac{p_1}{k}\right) + \frac{1}{p_2}e^{-(p-p_2)/T}\, S_i^u\left(\frac{p_2}{k}\right) + \frac{1}{p-p_1-p_2}e^{-(p_1+p_2)/T}\, S_i^d\left(\frac{p-p_1-p_2}{k}\right). \quad (43)$$

The $\mathcal{TSS}$ contribution is

$$\frac{dN_p^{\mathcal{TSS}}}{pdp} = \frac{g_p\, C\, \xi}{6p^{2\alpha+\gamma+4}}\int dp_1 dp_2\, (p_1p_2)^{\alpha}(p-p_1-p_2)^{\gamma} \times \sum_i \int dk\, k\, f_i(k)\, V_i(k,p_1,p_2,p), \quad (44)$$

where

$$V_i(k,p_1,p_2,p) = p_1\, e^{-p_1/T}\left\{S_i^u\left(\frac{p_2}{k}\right), S_i^d\left(\frac{p-p_1-p_2}{k-p_2}\right)\right\} + p_2\, e^{-p_2/T}\left\{S_i^u\left(\frac{p_1}{k}\right), S_i^d\left(\frac{p-p_1-p_2}{k-p_1}\right)\right\} + (p-p_1-p_2)\, e^{-(p-p_1-p_2)/T} \times\left\{S_i^u\left(\frac{p_1}{k}\right), S_i^u\left(\frac{p_2}{k-p_1}\right)\right\}. \quad (45)$$

The curly brackets in Eq. (45) denote symmetrization as in Eq. (17). Finally, we also have $\mathcal{SSS}$ which is simply related to the fragmentation of a hard parton into proton. The latter has already been studied in Ref. [7]. The corresponding FF, $D_i^p(z)$, can be used here as in Eq. (19) to give

$$\frac{dN_p^{\mathcal{SSS}}}{pdp} = \frac{\xi}{p}\sum_i \int dk\, f_i(k)\, D_i^p\left(\frac{p}{k}\right). \quad (46)$$

The result of our calculation for the four types of contributions are shown separately in Fig. 5. Their sum is shown as solid line, and agrees well with the data from PHENIX in Au-Au collisions at $\sqrt{s_{NN}}=200$ GeV for 0–5% centrality [4]. We emphasize that there are no free parameters to adjust to achieve the good agreement. The result below $p_T=2$ GeV is not compared with data, since the proton mass effect becomes important there and invalidates our scale-invariant formulation. All three contributions that involve thermal partons have about the same strength at $p_T \approx 4$ GeV. What dominates at higher p_T is the $\mathcal{TSS}$ component until $p_T > 9$ GeV where the $\mathcal{SSS}$ component takes over. The data available now are only for p_T up to 4.2 GeV and therefore cannot check our prediction in detail. At this point we can only conclude that the departure from the exponential behavior in the data for $p_T > 3.5$ GeV is well accounted for by the two types of thermal-shower recombination, but not by direct hard-parton fragmentation. That may be regarded as empirical support for the role of thermal-shower recombination.

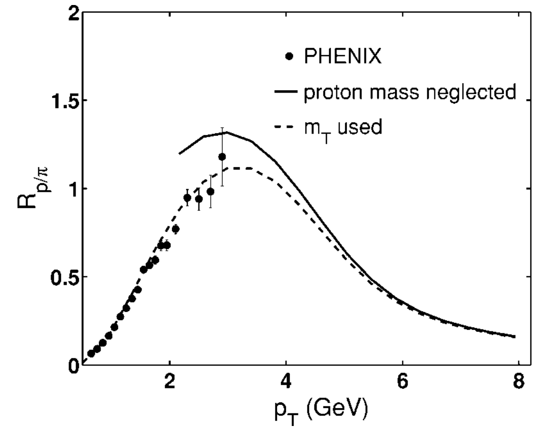

FIG. 6. Comparison of calculated p/π ratio with data from Ref. [4]. The solid line is the ratio of the solid lines in Figs. 5 and 1, without the mass effect taken into account. The dashed line is the result after p is replaced by m_T in the way described in the text.

Having obtained both the proton spectrum here and the pion spectrum in Sec. II, we can now calculate the p/π ratio and compare it to the data. Since the calculated result shown in Fig. 5 is only for $p_T \geq 2$ GeV/c due to our neglect of the proton mass m_p in our formulation, the p/π ratio that is represented by the solid line in Fig. 6 is not shown for $p_T < 2$ GeV/c. It has the broad feature that the ratio is greater than 1 at $p_T \approx 3$ GeV/c. For low p_T where the proton mass effect is important we adopt the Ansatz by replacing the factor $p^{-(2\alpha+\gamma)}$ in Eq. (40), which arises from G_B in Eq. (34), by $m_T^{-(2\alpha+\gamma)}$, where $m_T=(m_p^2+p^2)^{1/2}$. The result is shown as dashed line in Fig. 6 and exhibits excellent agreement with the data [4,23]. It should be noted that the amendment is at best kinematical, since no dynamical effect at low p_T has been considered. What is noteworthy is that the maximum exceeds 1 at the peak, which is the anomaly that the fragmentation model cannot explain. The slow decrease of the ratio as p_T increases is a prediction of our model that can be checked by future data at $p_T > 3$ GeV/c.

IV. SAME-SIDE CORRELATION

We have stated that because of the thermal-shower recombination the structure of jets in AA collisions must be different from that in pp collisions. One way to probe that difference is to study same-side correlation of particles. In Ref. [7] we have calculated the 2-pion correlated distribution in a u-quark initiated jet. That type of calculation can be incorporated into the study of same-side correlation in pp collisions by integrating over hard-parton momentum and summing over all types of hard partons. For AA collisions one would consider $2q$ and $2\,\bar{q}$ from thermal and shower partons and then recombine them to form two pions. Quantitative computation of the process will not be attempted here. Instead, we make some qualitative remarks in light of some data from RHIC that has some relevance.

In Ref. [17] a quantity I_{AA} is defined to be the ratio of all charged particles (above a modeled background) within a range of azimuthal angle $\Delta\phi$ around the trigger particle in Au-Au collisions to the same quantity (without background) in pp collisions. If there is no medium effect and if jets in AA and pp collisions are the same, then I_{AA} should be 1. It is found in Ref. [17] that if the trigger momentum is in the range $4<p_T^{\rm trig}<6$ GeV, the value of I_{AA} in central collisions is consistent with 1, though being more like 1.1. However, for $3<p_T^{\rm trig}<4$ GeV, I_{AA} is significantly higher, roughly 1.5 for $N_{\rm part}$ from 150 to 350. The former case for higher $p_T^{\rm trig}$ has been regarded as evidence that the nearside azimuthal peaks in AA and pp collisions are similar [17,24], but no explanation is given for the discrepancy when $p_T^{\rm trig}$ is lower. The azimuthal distribution of the nearside particles is not shown in Refs. [17,24] for $3<p_T^{\rm trig}<4$ GeV. It is also not clear for that case whether the particles associated with the trigger are integrated over the range $2<p_T<p_T^{\rm trig}$ only, as it is stated explicitly for the case of $4<p_T^{\rm trig}<6$ GeV [24].

In our view there is no reason why I_{AA} should be equal to 1 for the same-side particles. Indeed, we regard the value $I_{AA}\approx 1.5$ for $3<p_T^{\rm trig}<4$ GeV to be an evidence in support of the enhancement due to thermal-shower recombination. The question that we should address is why I_{AA} is lower at higher $p_T^{\rm trig}$.

Although our formalism does not facilitate the calculation of the azimuthal distribution, we can regard the integral over $\Delta\phi$ under the same-side peak that enters the determination of I_{AA} as being roughly equal to an appropriate integral of the two-particle distribution dN/dP_1dP_2 over P_2, with P_1 being assigned the role of the trigger momentum. Despite our restriction to collinear momenta, the integral $\int dP_2$ accounts for all associated particles, as does $\int d\Delta\phi$ in the data analysis by STAR [17]. In our approach we have a hard parton at k creating a shower in which two partons at p_1 and p_2 combine separately with two thermal partons at p_1' and p_2' to form two pions at $P_1=p_1+p_1'$ and $P_2=p_2+p_2'$. The distribution is

$$P_1P_2\frac{dN_{\pi\pi}}{dP_1dP_2}=\int\frac{dp_1}{p_1}\frac{dp_2}{p_2}\frac{dp_1'}{p_1'}\frac{dp_2'}{p_2'}(\mathcal{SS})_1(p_1,p_2)\times\mathcal{T}(p_1')\mathcal{T}(p_2')R_\pi(p_1,p_1',P_1)R_\pi(p_2,p_2',P_2). \quad (47)$$

For pp collisions the thermal partons are replaced by other shower partons initiated by the same hard parton that produces the ones at p_1 and p_2, i.e., in Eq. (47) we make the replacement

$$(\mathcal{SS})_1(p_1,p_2)\mathcal{T}(p_1')\mathcal{T}(p_2')\rightarrow(\mathcal{SSSS})_1(p_1,p_2,p_1',p_2')/\xi, \quad (48)$$

where $(\mathcal{SSSS})$ is a generalization of $(\mathcal{SS})_1$ in Eq. (16) to four shower partons.

We now see that in pp collisions when the trigger momentum, P_1, is high, it forces p_1 or p_1' in $(\mathcal{SSSS})_1(p_1,p_2,p_1',p_2')$ to be high, with the consequence that p_2 and p_2' must be low. Shower partons $S_i^j(z)$ at small momentum fraction z are known to have high density. They are shown in Fig. 2 of Ref. [7]. Indeed, the 2-pion distribution in a hard parton based on $(\mathcal{SSSS})_1$ has been calculated in Ref. [7], where it is shown that the distribution for the momentum fraction $X_2=P_2/k$ of the (nontrigger) second particle becomes very high at low X_2. The high density at low p_2' renders $(\mathcal{SSSS})_1$ to be of the same order as $\mathcal{T}(p_2')$ in AA collision, resulting in an increase of the $\pi\pi$ distribution in pp collisions. For that reason I_{AA} is lower (near 1) at higher $p_T^{\rm trig}$. In other words, since I_{AA} is dominated by the large number of particles in the low p_T range in the jets, it masks the differences in the structures of the jets (at intermediate and high p_T) produced in nuclear and hadronic collisions.

The drawback in using I_{AA} as a measure of same-side correlation is that it involves integrations in both $\Delta\phi$ and P_2, a procedure that is likely to suppress the distinctive features of thermal-shower recombination. It is recommended that the momenta of all particles in a jet are projected along the trigger momentum, P_1, and then the distribution in those projected momenta, P_2, is determined for several values of P_1.

V. CONCLUSION

By showing the importance of considering shower partons created by hard partons, we have called into question the conventional paradigm in particle production at high p_T in heavy-ion collisions. The usual approach is to regard such particles as the products of parton fragmentation. We have shown that all particles are the result of parton recombination, including but not limited to the ones usually regarded as fragments. The important input that makes feasible our approach based entirely on recombination is the shower parton distributions derived in Ref. [7]. Those distributions are determined by analyzing the fragmentation functions for parton-initiated jets in the recombination model. Once they are known, it is conceptually unavoidable to consider the recombination of thermal and shower partons in heavy-ion collisions. The result of such subprocesses turns out to be very important in the $3<p_T<8$ GeV range, as we have shown. The usual subprocesses of hard scattering followed by fragmentation are found to be unimportant until higher p_T. That results in a paradigm shift that has far reaching consequences.

The phenomenon of energy loss of partons traversing dense medium can be related to experimental observables only by means of some valid model of hadronization that connects the partons to hadrons. If fragmentation is not important in the region of p_T under investigation, then serious modification of the quantitative implications of jet quenching must be considered, since the content of a jet has been altered from that produced in *pp* collisions. What such a modification should be has not been studied in this paper. We have only used an effective parameter ξ to account for the fact that not all hard partons can get out of the dense medium to hadronize. We admit that such a method of treating energy loss is very rudimentary in this first attempt to study the effects of shower partons. Yet we have found consistency in being able to reproduce the spectra of pion, kaon, and proton. A more detailed investigation that tracks the space-time history of the produced partons would require an evolution code that is beyond the scope of this paper.

Particles that are formed by thermal-shower recombination are part of a jet produced in heavy-ion collisions but are not present in jets in *pp* collisions. Thus, the structures of jets produced in *AA* and *pp* collisions are different. To find evidence for such differences is of paramount importance in both experimental and theoretical work to follow. The ratio I_{AA} discussed in Sec. IV on same-side correlation reveals a limited glimpse of that difference. We have suggested that the measurement of two-particle inclusive distribution at high p_T would provide a more accurate description of the jet structure and can be used to check the prediction that can be made in our framework.

At this point the single-particle spectra that we have considered provide sufficient encouragement from the agreement with existing data to suggest that the recombination approach has captured the essence of hadronization at any p_T and that shower partons play an important role in the process. More data at higher p_T and on other species will give more stringent tests that the recombination model must pass in order to establish the solidity required for a reliable mechanism of hadronization.

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