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ABSTRACT

This paper reviews some aspects of the dynamics of the quark-gluon plasmas which may be produced in ultra-relativistic heavy ion collisions. A space-time description of the central rapidity region is presented. It is shown that the hydrodynamical flow induces correlations between particle transverse momenta and multiplicities. One discusses to which extent these correlations could signal the occurrence of a phase transition in heavy ion collisions.

1. INTRODUCTION

One of the main goals of the study of ultra-relativistic heavy ion collisions, such as those being presently performed at CERN or BNL, is to understand the behaviour of extended hadronic systems under extreme conditions of temperature or baryon density. In particular one hopes, if the energy densities achieved in the collisions are high enough, to induce the so-called deconfinement transition leading to the formation of a quark-gluon plasma (For a general review of the field see [1]).

Our theoretical knowledge of this quark-gluon plasma, a new state of matter predicted by quantum chromodynamics, is still very primitive. The qualitative picture which emerges from lattice gauge calculations is summarized in Fig.1 which displays the entropy density of baryonless matter as a function of the temperature. The entropy density (in units of T^3) exibits a characteristic increase within a narrow temperature interval centered around T_c, and quickly reaches a saturating value above T_c. In all calculations done so far, it appears that the value of the energy density (or the entropy density) at high temperature is compatible with the Stefan-Boltzmann law for a gas of non interacting quarks and gluons [2]. This could be naively expected on the basis of asymptotic freedom and it provides the simplest picture one may have of a quark-gluon plasma, namely that of a gas of weakly interacting particles.

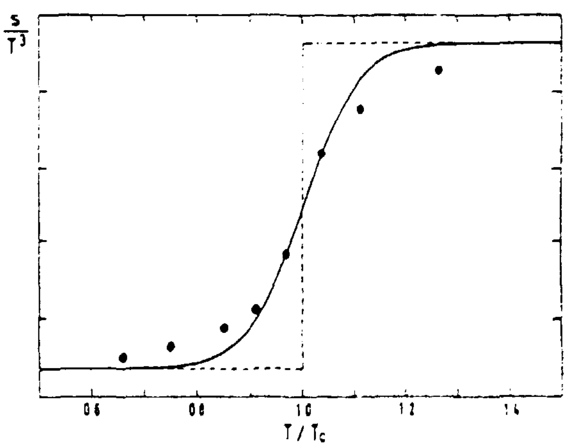

Fig.1: The entropy density of baryonless hadronic matter as a function of the temperature. Dashed line: eq.(1). Dots: predictions of the lattice gauge calculation (from ref.[3]).

The behaviour of the entropy density exhibited in Fig.1 may be understood simply in terms of an equation of state inspired from the bag model of hadron structure. This gives the entropy density as:

$$s(T) = \frac{2\pi^2}{45} T^3 \{\nu_h[1-\theta(T-T_c)]+\nu_{pl}\theta(T-T_c)\} \quad (1)$$

where ν_h (resp. ν_{pl}) is the number of degrees of freedom in the hadronic phase (resp. in the quark-gluon plasma). Eq.(1) implies a first order phase transition at $T=T_c$ between a phase of massless hadrons (e.g.pions) and a phase of massless quarks and gluons. Eq.(1) may actually be generalized so as to allow for a continous transition[3]. In any case, the main feature of the equation of state, independently of whether it is continous or discontinous,is the rapid increase of the number of degrees of freedom accross T_c. Typically, for a plasma made of up and down quarks, one has $\nu_{pl}=37$, while ν_h is only 3 for a massless pion gas.

Many uncertainties remain regarding the nature of the transition as well as the precise value of the parameters such as T_c. But in spite of those uncertainties, it is now commonly accepted that a quark-gluon plasma may be produced whenever the energy density exceeds a critical value of the order of a few GeV per fm^3. To put this number in perspective, let us recall that the energy density of ordinary nuclear matter is only about 0.15 GeV/fm^3.

The only laboratory systems in which one can possibly produce such high energy densities are colliding heavy ions. We shall see explicitly that the energy density achieved in ultra-relativistic nucleus-nucleus collisions may indeed grow significantly with the size of the nuclei involved. One expects high energy nucleus-nucleus collisions to exhibit two different regimes, depending upon the value of the initial kinetic energy. For moderate energies ($E_{lab} \sim 10GeV/A$), one expects substantial stopping of the two nuclei in their center of mass. This leads to systems with large baryon density. At extremely high energy, transparancy sets in. During the collision, the two ions go through each other, leaving between them a highly excited region which may contain a large energy density,but little baryonic charge (see Fig.2). For the rest of this talk, we shall discuss only this central region which is simpler to analyze.

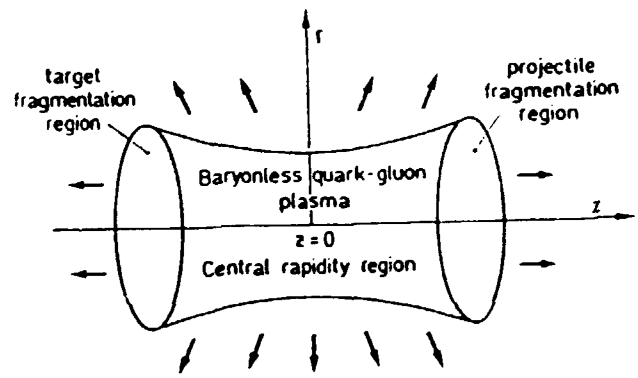

Fig. 2: Representation of an idealized central collision between two identical nuclei. The baryons populate the Lorentz contracted fragmentation regions. The central rapidity region is baryon free.

We shall actually describe a fairly idealized situation, that of a perfectly central collision at an extremely high energy. Such a model study may be somewhat unrealistic as far as comparison with experiment is concerned. Its main virtue is to help developing our intuition and to provide a rather clear conconceptual framework within which many interesting physical questions may be formulated.

2.SPACE-TIME DESCRIPTION OF THE CENTRAL REGION.

Colliding heavy nuclei are rather complicated dynamical systems and in order to make progress symplifying assumptions concerning the geometry and the dynamics are necessary. An obvious simplification occurs if one restricts oneself to central collisions of identical nuclei, since then one may take advantage of the cylindrical symmetry. A further simplification comes from assuming that the particle production process in the central region is invariant under Lorentz boosts along the longitudinal direction[4], as predicted by most particle production models. Cylindrical symmetry and longitudinal boost invariance reduce our problem effectively to a 1+1 dimensional problem.

It is convenient to describe the longitudinal evolution of the system in a space-time diagram, such as the one displayed in Fig.3. In this space-time diagram, the various stages of the collision are bordered by hyperbolae of constant proper time $\tau=\sqrt{t^2-z^2}$, as a result of the longitudinal boost invariance. Let us describe briefly this various stages.

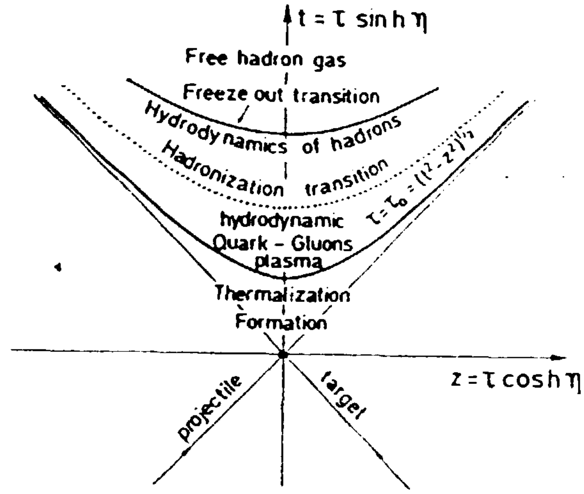

Fig. 3 : Space-time diagram representing various stages of a central collision. These various stages are separated by hyperbolae of constant proper time, as implied by the longitudinal boost invariance.

For negative times, the two ions are moving towards each other at a speed close to that of light, and suffer a Lorentz contraction which reduces their apparent longitudinal size to $2R / \gamma$ where $\gamma=1/\sqrt{1-v^2}$ and R is the radius of the ions.(Strickly, this Lorentz contraction does not apply to the momentum components of the nuclear wavefunctions [4], defined in the center of mass frame). The nuclei collide at z=t=0 , and a lot of quanta get produced. The detailed mechanisms by which these quanta materialize as well identified particles are still poorly understood. A recent calculation [5] based on a parton model gives the following predictions for the average transverse momentum of the partons which contribute dominantly to the energy density:

$$p_T^2 = \frac{9\pi}{16} \alpha C_A \frac{AxG}{R^2} \quad \alpha\ A^{1/3} \tag{2}$$

where G is the gluon density of a nucleon as measured for example in deep inelastic scattering. Typically, $xG \gtrsim 3$ and $\alpha C_A \approx 1$.The number and the energy densities of these partons are respectively:

$$n = \frac{3}{2} \sqrt{\frac{\alpha C_A}{\pi}} \left(\frac{AxG}{R^2}\right)^{3/2} \quad \alpha\ A^{1/2} \tag{3}$$

$$\epsilon = \frac{9}{8}\,\alpha\, C_A \left(\frac{AxG}{R^2}\right)^2 \; \sim \alpha\, A^{2/3} \qquad (4)$$

In order to get orders of magnitude, let us take $R=1.2A^{1/3}$, $A^{1/3}=6$, $xG=3$, $\alpha C_A=1$. Then one gets $p_T \simeq 0.94$GeV, $n \simeq 37 fm^{-3}$, $\epsilon \simeq 35 GeV/fm^3$. These numbers represent the typical transverse momentum and the densities of the partons which are expected to contribute dominantly to the energy density at a time $\tau_0 = 1/p_T \simeq .2$fm/c. At this time, these distributions resemble those of free particles with velocities $\tanh y = z/t$.

A further major step in arriving at a simple picture is to assume that the newly formed quarks and gluons quickly thermalize, i.e. their distributions turn into local equilibrium distributions:

$$n(r,p,t) = \frac{1}{\exp(\epsilon_p / T(r,t)) \mp 1} \qquad \epsilon_p = \sqrt{p^2+m^2} \qquad (5)$$

where $T(r,t)$ is the local temperature of the system . From there on the system is then described by the equations of hydrodynamics. There are two basic ingredients in hydrodynamics. One is a statement of local conservation of energy and momentum, which can be expressed in the form:

$$\partial_\mu T^{\mu\nu} = 0 \qquad (6)$$

This equation is much more general, of course, than hydrodynamics. Any local field theory would lead to eq.(2.1). Specific aspects of hydrodynamics enter in the prescription of the form of the energy-momentum tensor $T^{\mu\nu}$. If one ignores dissipative processes as is done in most calculations, one has:

$$T^{\mu\nu} = (\epsilon+P)u^\mu u^\nu - P g^{\mu\nu} \qquad (7)$$

where $g^{\mu\nu}$ is the metric tensor, u^μ the fluid four-velocity, ϵ the energy density and P the pressure. The pressure and the energy density are related by an equation of state:

$$P = P(\epsilon) \qquad (8)$$

The eqs.(6-8) certainly constitute the simplest set of dynamical equations capable of describing the evolution of the baryonless plasma. Also, they directly incorporate our available microscopic information about plasma properties through the equation of state. Finally, since they deal mostly with basic conservation laws, they should provide reasonable orientation for average properties.

The most important assumption done in writing the equations (6-8) concerns the thermalization. Whether this thermalization does occur, and over a sufficiently short period of time for the whole picture to make sense is still an unproved conjecture. One can give however some plausibility arguments based on mean free path estimates. The mean free path is given by $\lambda = 1/\sigma n$, where σ is the quark or gluon cross section in the plasma and n the number density of the plasma. If ones assumes thermalization at a temperature T in a baryonless plasma made of u and d quarks in addition to gluons, one can calculate the number of quanta to be $n \sim 4.14T^3$. The average cross section for a gluon propagating through the plasma may be estimated to be of the order of 14 mb.Thus in a plasma at temperature $T \sim 200$MeV one has $n \sim 4\text{fm}^{-3}$ and the mean free path of a gluon is $\lambda \sim 0.2$fm . This mean free path is small compared to a typical size of the system, e.g. $2R \sim 15$ fm for big nuclei.

The hydrodynamical evolution of the system may be understood as the superposition of two collective motions [7]. The first one is the longitudinal expansion of the fluid, which because of the boost invariance reduces to a simple scaling mode [6] . Thus for example the longitudinal expansion alone would make the entropy density drop like $1/\tau$:

$$s(\tau)=s(\tau_0)\ \tau_0/\tau \qquad (9)$$

This result reflects simply the fact that the entropy is conserved during the evolution and that the proper volume of the system increases like the proper time τ. The longitudinal expansion of the system implies a rapid cooling of the plasma. Superimposed to this longitudinal expansion is the transverse motion which begins with the inwards propagation of a rarefaction wave (see Fig.4). If no phase transition takes place in the system the rarefaction wave will reach the collision axis in a time $R/c_s=R\sqrt{3} \sim 12$fm for a large nucleus (we have used the fact that the speed of sound in an ultrarelativistic medium is $c/\sqrt{3}$). This is a fairly long time scale compared to the time scale of the longitudinal expansion. Note for example that in a time of order 12fm the central temperature drops by some 40% due to the longitudinal cooling alone (this is obtained from (9) which implies that temperature and time are related by: $T^3\tau \sim T_0^3\tau_0$, and τ_0 is taken to be 1fm) . Thus the overall hydrodynamics is by far dominated by the longitudinal expansion. However the interesting physics is contained in the transverse motion. This is because the transverse motion may be quite modified by the occurrence of a phase transition and this affects the final distribution of particles[8-13].

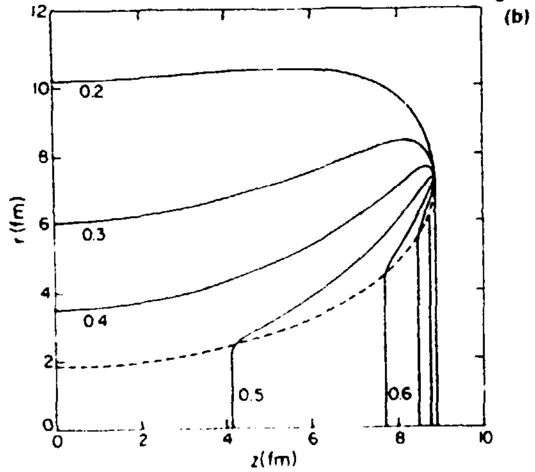

Fig. 4 : Cylindrical expansion of an ultra-relativistic ideal gas.Snapshots of isotherms in the z,r plane . The dashed curve shows the inwardly moving rarefaction front (from ref. [7]).

As the plasma cools down, it eventually reaches the critical temperature at which it may hadronize. In the framework of the hydrodynamic description, this hadronization is treated as a phase transition (the reverse of the deconfinement transition). When the hadronization is over, the system, mostly composed of pions, is still strongly interacting and its evolution described by hydrodynamics until the hadron mean-free-path becomes comparable to the size of the system. This is the so-called freeze-out transition where the particles decouple. In the framework of the hydrodynamical model, one usually assumes that freeze-out takes place when the temperature has dropped down to some value $T_{f.o.}$ called the freeze-out temperature [14]. This is usually taken to be of the order of the pion mass for the reason that the pion mean-free-path, which is inversely proportional to the pion density, increases rapidly as the temperature decreases below the pion mass. From the freeze-out transition, the particles evolve as free particles.

3. TRANSVERSE EXPANSION AND OBSERVABLES

We shall now proceed to a discusion of some specific results of hydrodynamical calculations. The geometry of the system which we study is as discussed in sections 1 and 2. Furthermore we assume that the phase transition proceeds smoothly, without any excursion through metastable states. In the case where the deconfinement transition is treated as a first order transition, one assumes that a uniform mixed phase of hadrons and quark-gluon plasma is formed.

The dominant feature of the transverse hydrodynamic flow in this scenario, with a first order phase transition, is the existence of a rarefaction shock converting the mixed phase into hadrons [8-11]. The effects of this shock is most easily analyzed in the one dimensional flow displayed in Fig.5 . This figure shows the energy density profile for a system prepared in a mixed phase. As time goes on, the rarefaction shock slowly moves inwards. The hadrons emerge from the shock at the speed of sound relative to the shock [8] (Chapman-Jouguet point). It is important to observe that the velocity of the shock relative to the mixed phase decreases as the energy density, or equivalently the entropy density, in the mixed phase increases. This can be verified by an explicit calculation but is also easily understood when one knows that very little entropy is produced in the transformation of the mixed phase into hadrons (at most a few percent[8]). The displacement of the shock into the mixed phase liberates a large amount of entropy which has to be carried away by the hadrons. Because the velocity of the hadrons is bounded when they leave the shock, there is a limit to the total amount of entropy they can carry away per unit time. This implies a slowing down of the shock with the increase of the entropy density contained in the mixed phase. In the case where the transition is continous, the flow may become continous, i.e. the shock may disappear. However the slowing down of the transverse expansion coming from the conversion of the large entropy content of the plasma still persists.

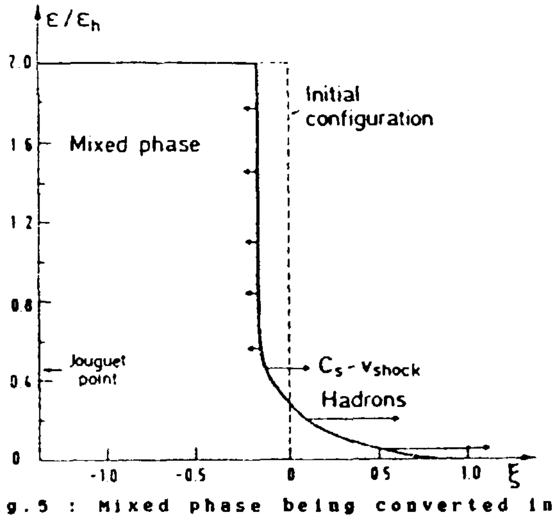

Fig.5 : Mixed phase being converted into hadrons through a rarefaction shock. This figure shows the energy density profile of a one-dimensional system as as a function of $\xi = x/t$. The arrows on on the profile indicate schematically the velocity of the hadrons (to right), or of the shock (on the left).

Let us now consider the expansion of a system initially at a temperature larger than T_c . Fig. 6 displays the temperature profile for the expansion of a one-dimensional slab of size $2R_0$. The various pieces of the flow are clearly visible: rarefaction wave propagating into the plasma; the plateau at $T \sim T_c$ corresponding to the formation of the mixed phase; the rarefaction shock converting the mixed phase into *hadrons . Note also that the* smoothening of the equation of state does not alter the bulk features of the flow pattern. Fig.7 shows the temperature profiles for the full three dimensional expansion of plasma initially contained in a cylinder of radius R_0 with the same temperature as in Fig.5, i.e. $2T_c$. The effect of the longitudinal expansion is clearly visible.

One of the goals of the hydrodynamical calculations is to provide a framework for the calculation of various global observables. These observables may be reconstructed from the particle distributions at freeze-out which have the form of local equilibrium distributions:

$$f(x,p) = \frac{1}{\exp(p^\mu u_\mu / T_{f.o.}) - 1} \qquad (10)$$

where $u_\mu(x)$ is the four-velocity of the fluid on the frezze-out isotherm at temperature $T_{f.o.}$.

It can be shown that, in the hydrodynamical model which we consider, the properties of the flow at freeze- out, which are needed to calculate the particle distributions depends only on the dimensionless parameter :

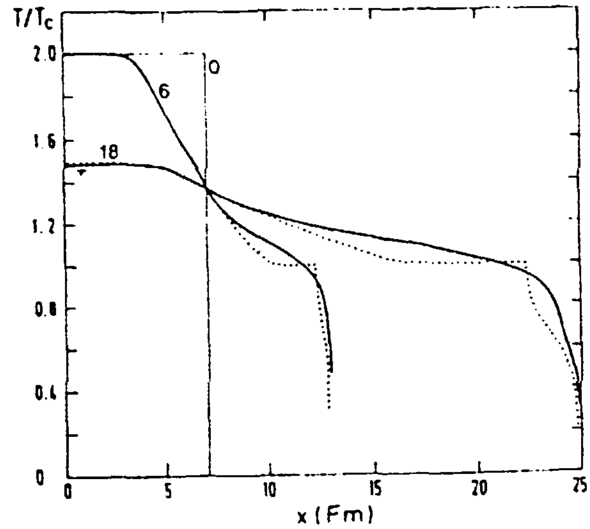

Fig. 6 : Temperature profile for the one-dimensional expansion of a quark-gluon plasma in a slab of initial length $2R_0$ (R_0=7fm). The initial temperature is T_0 = $2T_c$. The curves are labelled by the time in fm/c. The dotted lines correspond to the equation of state (1), while the full lines are for a smoothened version of (1).

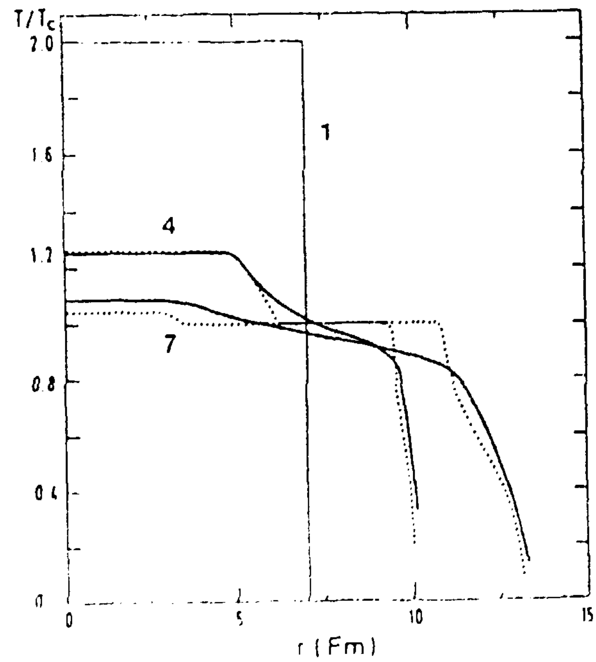

Fig. 7 : Temperature profile corresponding to the three dimensional expansion of a quark-gluon plasma with cylindrical geometry (R_0 = 7fm). The initial proper time is 1fm/c and the initial temperature is $2T_c$. Same convections as in Fig.6 (from ref. [13]).

$$\omega = \frac{s_0 \tau_0}{s_c R_0} \tag{11}$$

where s_0 is the initial entropy density, τ_0 the time at which the hydrodynamical evolution starts, and s_c some reference entropy density. On the other hand, since the hydrodynamical evolution conserves the entropy to a good approximation, one has:

$$\frac{1}{A}\frac{dN}{dy} \propto \frac{s_0 \tau_0}{R_0} \tag{12}$$

This equation relates the initial conditions of the hydrodynamical evolution of the plasma to the multiplicity of the produced particles per nucleon and per unit of rapidity.

Fig. 8 is a plot of the average transverse momentum carried away by the pions as a fuction of the multiplicity, or equivalently the parameter ω. If collective flow could be ignored, such a plot would provide a direct information on the equation of state [15,16]. Indeed, for a hot pion source at rest, $\langle p_T \rangle$ is proportional to the temperature, while dN/dy is proportional to the entropy density. Thus the plot of $\langle p_T \rangle$ versus dN/dy should ehibit the same characteristic structure as the entropy density versus temperature, see Fig.1. In fact the hydrodynamic flow, especially the longitudinal expansion, washes out any of the structures which could be expected. The dominant effect of the occurence of a phase transition is the softening of the average $\langle p_T \rangle$ as more degrees of freedom come into play (compare the curves labelled pion gas and plasma in Fig.7).

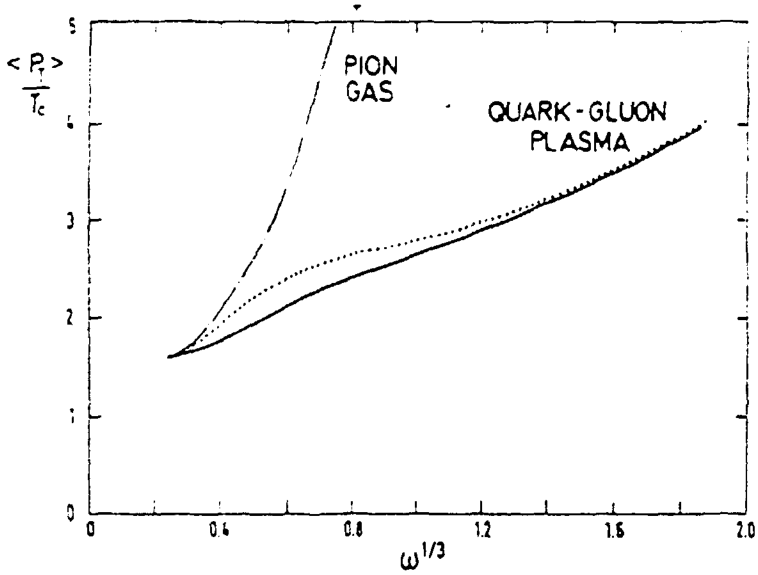

Fig. 8 : Average transverse momentum as a function of the dimensionless parameter ω, eq.(11). The dotted and full lines correspond to variants of the equation of state for the plasma. The dashed dotted line corresponds to a massless pion gas without any phase transition.

7. CONCLUDING REMARKS

We have presented a model study which provides a useful phenomenological guide into the dynamics of ultra-relativistic nucleus -nucleus collisions. We have seen that hydrodynamical calculations are reaching the level of sophistication where they allow for the calculation of observables which can eventually be compared with experiment. We have shown that the correlations between various observables may reflect important features of the underlying equation of state. In particular, we have analyzed in detail the correlations between multiplicities and transverse momenta and discussed how they could possibly signal the occurrence of a phase transition in the evolution of the system. In order to get from such idealized model studies to more realistic situations, many improvements are needed. For example the freeze-out transition is crudely treated in our approach. Also, it would be desirable to get rid of the longitudinal boost invariance, to include finite baryon number and to treat the fragmentation regions. This is especially important in view of understanding experiments at CERN or BNL which will deal mostly with fragmentation regions or with a baryon contaminated central region. Finally, it is worth keeping in mind that one still poorly understands the very beginning of the collisions during which the plasma is expected to be formed.

Acknowledgements

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