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HADRONIZATION MODEL FOR QUARK-GLUON PLASMA IN ULTRA-RELATIVISTIC COLLISIONS

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ABSTRACT

We propose a model for the expansion and hadronization of small droplets of quark-gluon plasma assumed to be produced in ultra-relativistic collisions of hadrons and/or nuclei, and we discuss its experimental consequences. The deconfinement transition is described in terms of formation of QCD strings in the expanding plasma droplet, the string tension stopping the expansion and string breaking leading to possible subdivision of the droplet into smaller ones. The bulk of plasma hadronization is predicted to occur by deflagration of the plasma droplets at transition, through their outer surface and with ejection of low pressure hadron gas. In the rapidity distribution dn/dy of hadrons, this predicts isolated maxima of width $\delta y \sim 1$ or high bumpy regions a few units of rapidity wide, on an event-by-event basis, the corresponding particles having larger than normal transverse momentum. It is in these rapidity intervals that one expects the "traditional" signals for plasma formation (direct dileptons, direct photons, strange particles). We review possible evidence for plasma formation in existing cosmic ray and hadron collider data. Although it appears to be inconclusive, the observed features of the rapidity distribution and the estimated energy densities in the central region are quite encouraging for a systematic search using all proposed signals.

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1. INTRODUCTION

There are compelling reasons to believe that hadronic matter at high energy density (perhaps $\gtrsim 2$ GeV fm^{-3}) takes the form of a quark-gluon plasma, i.e., a dense fluid of quarks, antiquarks and gluons which are not confined in well-separated hadrons. In recent years, many theoretical papers have addressed the possibility that this new state of matter could occur in high energy collisions of hadrons and/or nuclei, and various experimental tests have been proposed[1]. All this work suffers from many uncertainties, however, because it has not been possible so far to extract from Quantum Chromodynamics (QCD) reliable predictions on hadronic matter in the deconfinement transition region between the dense phase (quark-gluon plasma) and the dilute phase (hadron gas, i.e., gas of hadrons separated by QCD vacuum). In particular, Monte-Carlo lattice calculations, which are quite successful for QCD without quarks and indicate for this case a first order phase transition[2], are very difficult to extend to the realistic case of QCD with light quarks. Recently, on the basis of universality arguments and calculations with heavy quarks, some of the experts suggested that the presence of light quarks may change the nature of the deconfinement transition and make it smooth[3], but the matter is undecided. Furthermore, the situation is quite unclear concerning the chiral transition (high density restoration of chiral symmetry for massless quarks).

Despite the uncertainties on the nature of the deconfinement transition, there is little doubt that it must exist. It is likely to be characterized by substantial changes of entropy and energy densities without comparable changes of pressure and temperature, these properties being due to the melting of many quark and gluon degrees of freedom which are frozen in the hadron gas phase and get liberated in the plasma. By combining these qualitative properties with the familiar QCD concept of strings (also called flux tubes), we are led to propose what we believe to be a plausible model for the hadronization of droplets of quark-gluon plasma as could be formed in high-energy collisions. This process concerns very small blobs of hadronic matter (at most 5-10 fm across in transverse as well as longitudinal directions) in a state of rapid expansion and cooling (typical time scale of a few

fm/c). For such situations, the detailed analytic nature of the transition is less relevant than for a large medium in slow expansion (as must have been the case for the quark-gluon plasma which occurred in the early Universe if the Hot Big Bang model is correct). What matters is whether strong supercooling of the plasma is to be expected; in our model this is not the case, in line with the indications that the transition may be smooth (a plasma hadronization mechanism assuming strong supercooling was considered in Ref. 4). Since little hadronization is expected at the surface of the plasma droplet before it reaches the transition region[5,6], we predict that the bulk of hadronization takes place near transition where, according to relativistic combustion theory, the plasma can deflagrate by ejecting a flow of hadron gas. The reason is that this deflagration process, first described in Ref. 7 (to be referred to hereafter as I), turns out to give a much more efficient release of the large energy contents of the plasma than other hadronization mechanisms[5,6]. In view of the small droplet size, the deflagration is expected to occur mostly through the outer surface. The resulting picture of plasma hadronization, despite its qualitative nature, leads to rather specific experimental consequences, especially concerning the rapidity distribution and mean transverse momentum of secondaries resulting from a plasma droplet. Although our qualitative findings should be more general, we concentrate our treatment on the case of zero chemical potential, i.e., zero baryon number, which is appropriate for the central rapidity region of ultra-relativistic collisions.

The plan of the paper is as follows. We describe in Section 2 the thermodynamic properties of hot hadronic matter which we shall use later on. Section 3 develops the picture we propose for the behaviour of a quark-gluon plasma droplet, assumed to be formed in a collision, as it expands and cools down toward the transition region. We associate the transition with the formation of QCD strings between quarks and gluons. After a purely qualitative description we derive some estimates from the simplest model of longitudinal expansion with Lorentz-boost invariance. The hadronization of plasma droplets is discussed in Section 4. After recalling other forms of hadronization of a plasma droplet through the outer surface, we concentrate on the deflagration mechanism of Ref. 7; it only operates when the plasma has reached transition but is then very effective in liberating the high energy contents. This is studied numerically for the simplest case of a plane surface using two examples of equation of state, one with a first order transition and the other with a smooth transition. The experimental consequences resulting from our model are discussed in Section 5. The last Section presents concluding remarks, especially on possible evidence for plasma formation in cosmic ray events and at the CERN colliders. It ends with a brief summary.

2. HOT HADRONIC MATTER

In view of the difficulties of the lattice approach for QCD with light quarks, one must rely on the continuum approach to obtain estimates for the thermodynamics of hadronic matter. For the quark-gluon plasma phase, one uses perturbative QCD improved by inclusion of plasmon effects and of a vacuum energy shift. For the hadron gas phase, one relies on phenomenological knowledge of the hadron spectrum and the hadronic interactions, simulating the latter by bootstrap, mean-field or potential methods[8]. For net baryon number zero, the only case we shall consider in this paper, one finds for the two phases pressure - temperature curves of the type depicted in Fig. 1a. These p-T curves cannot be trusted at temperatures in the region of their crossing point nor beyond (dashed parts of the curves) where the approximations used loose their validity (large higher order terms on the plasma side, unknown hadron properties on the gas side). The transition could be of 1st order (full curve of Fig. 1a), or of second order, or it could be smooth as depicted in Fig. 1b. In the latter case we define the transition region as the interval of rapid variation of the slope dp/dT. By the thermodynamic relation

$$dp/dT = s = (\varepsilon + p)/T \qquad (1)$$

the entropy density s and the energy density ε also vary rapidly in this region. The rapid increase of s has a clear physical origin; as mentioned earlier it is due to the melting of the numerous degrees of freedom of the QCD partons (quarks,

antiquarks and gluons) which in the gas phase are frozen inside the hadrons and get liberated in the plasma phase.

The values of the thermodynamic variables in the transition region are not known precisely. Reasonable guesses are (we put $\hbar = c = k_B = 1$)

$$T_0 \sim 150-250 \text{ MeV}, \quad p_0 \sim 0.5-1 \text{ GeV fm}^{-3} \qquad (2)$$

the estimate of T_0 being more reliable than the one for p_0. In order of magnitude, the energy density ε may perhaps grow from $\varepsilon_0 \sim 2$ GeV fm^{-3} to $\varepsilon_0' \sim 4$ GeV fm^{-3} across the transition. The growth of ε over the width ΔT of the transition region is assumed to be much more rapid than it would be in either phase. The single phase variation of ε with T is probably not too different from the ultra-relativistic ideal gas law (Stefan-Boltzmann behaviour) which gives

$$\varepsilon = 3p \propto T^4, \quad s \propto T^3 \qquad (3)$$

The rapid variation assumption therefore means

$$(1 + \Delta T/T_0)^4 \ll \varepsilon_0'/\varepsilon_0 \sim 2 \quad \text{or} \quad \Delta T \ll 0.2\, T_0 \sim 40 \text{ MeV}.$$

In the example we shall study in Section 3, $\Delta T/T_0$ will be a few percent and the corresponding $\Delta p/p_0$ of order 15%. Since p is small compared to ε, the entropy density s increases across the transition region by about the same factor as ε.

3. EXPANDING DROPLET OF PLASMA

We now begin our discussion of what we believe to be the likely evolution of a droplet of quark-gluon plasma assumed to be formed in a high energy nuclear collision. We imagine the droplet to be transversely and longitudinally a few fermi across (possibly after some expansion). As long as the energy density ε is larger than ε_0', the plasma expands, mainly in the longitudinal direction, and its losses by radiation of hadrons through its outer surface are small (see next section). In this expansion, ε decreases dominantly due to a decreasing number density n of QCD partons, and also due to a decreasing mean energy $\bar{E}$ per parton (the Stefan-Boltzmann behaviour is $n \propto T^3$, $\bar{E} \propto T$).

The transition begins when ε reaches ε_0', and we picture it as follows. Above transition, the number density n is high enough for the colour fields induced by the colour charges of the partons to be spread all over the space occupied by the plasma (Fig. 2a). As n drops to values $\lesssim$ 2-3 fm^{-3}, the properties generally accepted for QCD imply that small islands of QCD vacuum should form in the plasma, the colour fields between partons gradually collapsing into "flux tubes" or "strings" of the type commonly considered in QCD phenomenology (Fig. 2b). QCD strings have thickness $\sim$ 1 GeV fm^{-1}. A string carries at one end an SU(3) colour index α belonging to the representation 3 and at the other end an index $\bar{\alpha}$ belonging to $\bar{3}$. Some of these indices saturate with the colour indices of the partons (one for an (anti)quark which ends a string, two for a gluon which is a kink in a string). Three strings can also join in a vertex with the indices saturated by the fully antisymmetric tensor $\varepsilon_{\alpha\beta\gamma}$ or $\varepsilon_{\bar{\alpha}\bar{\beta}\bar{\gamma}}$ (see Fig. 3 for an example). As the strings appear in the expanding plasma, they will tend to form networks involving many partons, because complete colour neutralization of a few neighbouring partons is unlikely. After string formation the plasma has the structure recently invoked by Patel in an interesting discussion of the deconfinement transition[9].

Once a string network has formed in the plasma droplet, it is stretched by the expansion, and this stretching in turn slows down the expansion. Some strings will also break by the familiar mechanism of pair creation of light quarks (Fig. 4), but the break up is not fast, as we know from the string model of high spin hadrons (leading Regge trajectory) where string break up is slow enough to give a small width/mass ratio. The likely course of events is therefore that by these processes of string formation, stretching and breaking, a droplet of expanding plasma will either stop expanding, or will break up in a few smaller droplets which stop expanding and separate from each other. The resulting *non-expanding droplets* should have an energy density ε still around ε_0'. While ε decreases in the expansion, it tends to stabilize momentarily at the value ε_0' when the expansion stops. Collective energy of expansion is then converted into internal energy (an irreversible, entropy-producing process). As mentioned already, a plasma droplet formed in a high energy

collision is expected to expand mainly in longitudinal direction. It will therefore get elongated and will tend to break up in droplets of more spherical shape, each with a different longitudinal rapidity y, the likely sizes being again a few fm across.

To make the above picture a little more concrete, let us consider*) the simplest model for an expanding plasma produced in the central region of a very high energy collision, i.e., a uniform cylinder a few fm thick expanding longitudinally according to the well known scaling law[10]) characterized by longitudinal Lorentz-boost invariance

$$t = \tau \cosh y \;, \quad z = \tau \sinh y . \tag{4}$$

Here y is the longitudinal rapidity at time t of an infinitesimal transverse slice of plasma at longitudinal position z, τ being its proper time. The transverse expansion is neglected. One assumes local equilibrium with the thermodynamic state depending on τ. Energy-momentum conservation thus implies that the expansion is adiabatic (constant entropy per comoving volume).

Any hadronic system is expected to be $\gtrsim$ 1 fm across, and a plasma droplet somewhat larger. We therefore want the plasma cylinder to have a length $z \gtrsim 2$ fm in the rest frame of each slice; since (4) implies $|z| < t$, our description can only hold for $t \gtrsim 1$ fm/c in each of these frames, i.e., it only holds for $\tau \gtrsim 1$ fm/c. If two slices are separated by δz in the rest frame of one of them and at its proper time τ, they have a rapidity difference δy given by

$$\sinh \delta y = \delta z / \tau . \tag{5}$$

Since $\tau \gtrsim 1$ fm, we have $\delta y \lesssim 0.9$ for $\delta z = 1$ fm. This means that when strings of length $\sim$ 1 fm form as described above, the difference in rapidity of the collective flow at their endpoints is $\sim$ 1 if the string is longitudinal and $\tau \sim 1$ fm, and it is smaller in all other cases (strings not oriented in the longitudinal direction, larger τ at string formation). This supports our expectation that string breaking should be rather exceptional, the stretching of most strings having the effect to stop the expansion. For a long cylinder in longitudinal expansion, however, break

*) These considerations resulted from a discussion with M. Gyulassy.

up in shorter pieces is likely, each piece becoming a non-expanding droplet. If the break up occurs at average longitudinal separation δz, presumably several fm, the corresponding δy in (5) gives us an estimate for the rapidity difference between adjacent non-expanding droplets. Since δz can now be significantly larger than τ, this rapidity difference can be several units, especially if the initial plasma is not very hot (small τ at string formation).

In our ignorance of what will control the average separation δz between successive break up points, we shall mention two possibilities. If δz is controlled by the thickness of the cylinder, δz would be about the same for various initial energy densities of the plasma at formation time (only the small transverse expansion would make it somewhat larger for a hotter initial plasma); then an initially hotter plasma would have larger τ at break up and by (5) the resulting droplets would have smaller rapidity differences δy. Another possibility is that δz gets larger for an initially hotter plasma since, as explained above, the larger τ at string formation implies a smaller rapidity difference between the end points of a string and hence a smaller probability of string break up; an initially hotter plasma would then tend to give larger droplets with comparable rapidity separation.

The cooling of the plasma during the scaling expansion (4) is very simple to derive if one neglects transverse expansion and assumes constant $w = p/\varepsilon$ (the velocity of sound is then $w^{1/2}$). As is well known, eq. (1) then gives $s \propto T^{1/w}$ and adiabaticity $s \propto \tau^{-1}$, so that $T \propto \tau^{-w}$, $\varepsilon \propto \tau^{-1-w}$, and the internal energy per comoving volume varies as $\varepsilon\,\tau \propto \tau^{-w}$. The value of w is expected to be small, $w \lesssim 1/3$; the decrease of $\varepsilon\,\tau$ is therefore very slow. By overall energy conservation, this decrease of internal energy drives the transverse expansion which is thereby confirmed to be a small effect. The number density n of partons in the plasma is expected to vary roughly as the entropy density s, in the sense that their ratio should vary much less with temperature than either of them. At high T, asymptotic freedom gives a Stefan-Boltzmann behaviour, the values being

$$\varepsilon = 12.2\,T^4 , \quad s = 16.2\,T^3, \quad n = 4.1\,T^3 \tag{6}$$

for gluons and two light quark flavours. Hence $s/n \approx 4$ at high temperature where the partons carry all energy and entropy. At lower T, in particular for T near the transition temperature T_0 where the strings form, part of the energy and entropy will be carried by the colour fields between the partons (as also discussed by Patel[9]) who, however, neglects kinks, i.e. gluons, and colour degrees of freedom in the entropy) and we therefore expect s/n to be larger. As explained before, our estimates for the plasma near transition are

$$T_0 \sim 200\ \mathrm{MeV},\ p_0 \sim 0.5\text{-}1\ \mathrm{GeV\,fm^{-3}},\ \varepsilon_0' \sim 4\ \mathrm{GeV\,fm^{-3}},\ n_0 \sim 2\text{-}3\ \mathrm{fm^{-3}} \tag{7}$$

They give

$$s_0' = (\varepsilon_0' + p_0)/T_0 \sim 25\ \mathrm{fm^{-3}}, \quad s_0'/n_0 \sim 10. \tag{8}$$

This suggests an increase of s/n by about a factor 2 as a plasma cools down from high temperatures to transiton. The estimates (7) are very uncertain, however, and there would be little increase of s/n if $\varepsilon_0' + p_0$ would be substantially smaller. At any rate, we are led to expect that the decrease of n in the longitudinal expansion (4) is somewhat faster than $s \propto \tau^{-1}$.

4. HADRONIZATION OF PLASMA DROPLETS

As long as a plasma droplet is above transition, the only way it can hadronize is by evaporation of hadrons (mostly mesons) at its outer surface. This process has been studied in various ways. Banerjee et al[5]) considered the breaking of single QCD strings pulled out of the plasma by (anti)quarks punching through its surface. The resulting energy fluxes are found to be (see Fig. 3 of Ref. 5)

$$F_\varepsilon \sim 10^{22}\ \mathrm{GeV\,fm^{-2}\,sec^{-1}} \sim 30\ \mathrm{MeV\,fm^{-3}}\,c \quad \text{at } T \sim 250\ \mathrm{MeV},$$

$$F_\varepsilon \sim 10\ \mathrm{MeV\,fm^{-3}}\,c \quad \text{at } T \sim 200\ \mathrm{MeV}.$$

These values are very low if one remembers that the energy density of the plasma is $\geq \varepsilon_0' \sim 4$ GeV fm^{-3}. Higher fluxes are obtained when considering emission of pions into free space by (anti)-quarks reaching the plasma surface, either pion emission by an (anti)-quark reflected at the surface, or $q\bar{q}$ annihilation into a pion at the surface. These effects have been estimated by Müller[6] in two approximations leading to

$$F_\varepsilon \sim 0.2\, T^4 \quad \text{and} \quad \sim 0.07\, T^6/f_\pi^2 \tag{9}$$

with f_π = 93 MeV the pion decay constant, see eqs. (4.6) and (4.13) of Ref. 6. This gives respectively

$F_\varepsilon \sim 100$ and 250 MeV $fm^{-3}c$ at $T \sim 250$ MeV,

$F_\varepsilon \sim 40$ and 65 MeV $fm^{-3}c$ at $T \sim 200$ MeV.

Such fluxes are still quite small and they cause little energy loss of the plasma during its expansion. We disregard earlier pion emission calculations by Danos and Rafelski[11] because they contain adjustable parameters which cannot be reliably estimated.

When the plasma droplet is close to transition, it can deflagrate through its outer surface, producing an outflow of hadron gas, as shown in Section 3 of I[7]. It should be noted that this process of surface deflagration is forbidden by the condition of entropy increase as long as the plasma is at higher temperature; it can only occur if T is very close to the transition temperature T_0. We shall find that the energy fluxes produced by surface deflagration are much larger than (9) at $T \sim T_0 \sim 200$ MeV. Not knowing of any process liberating more rapidly the high energy and entropy contents of the plasma, we therefore propose that surface deflagration is the main hadronization mechanism of the plasma droplets. It should take place at the end of their expansion and possible break up. It may be triggered by the slowing down and stopping of the expansion. Since the deflagration is slow (see the low value of v_s below), it will mostly occur for non-expanding droplets of plasma. Since the droplets are small, a few fm across, the hadronization by deflagration is expected to occur mainly at the outer surface of the droplets.

Things would be different in a larger volume of expanding plasma where the formation and stretching of the string network and the breaking of some strings are expected to create inside the plasma bubbles which expand and fill up with hadron gas (possibly again by deflagration). Such internal bubble formation should be

unimportant for the small droplets that could be formed in high energy collisions.

We now proceed with a more quantitative discussion of deflagration, based on section 3 of I and on the simple geometry we considered there (deflagration through a planar surface). We treat two cases, case A with a first order transition and case B with a smooth transition. For case A, we take the deflagrating plasma to be at the transition point with temperature T_0, pressure p_0, energy density ε_0' and entropy density $s_0' = (\varepsilon_0' + p_0)/T_0$. For the hadron gas we take the equation of state

$$p_g = \varepsilon_g/3 = p_0\,(T/T_0)^4 \;, \quad T \leq T_0 \tag{10}$$

The value of $s_0' = (dp/dT)_{T_0}$ of the plasma at transition is taken to be twice the entropy density $s_g = (\varepsilon_g + p_g)/T = 4\, p_g/T$ of the gas at T_0, which amounts to putting

$$s_0' = 8\, p_0/T_0 \;, \quad \varepsilon_0' = 7\, p_0 . \tag{11}$$

This corresponds to the value 0.5 of the parameter α of I (the plasma variables had there the index L, the gas variables the index R). The corresponding equation of state is plotted in Fig. 5. As shown in I (Figs. 3-4) the entropy condition gives that deflagration can occur into gas with pressure p_g in the wide range

$$0.2\, p_0 \leq p_g \leq p_0 \tag{12}$$

with a very flat probability distribution over most of the range (see the near constancy of the curve in Fig. 4 of I over the above interval where it is > 1 and the near constancy of the velocity v_s defined below).

Case B is obtained by smoothing out the first order transition of case A in the following way. For pressure $p \geq p_0$ we lower by 1% the plasma temperature T and divide by 0.99 its entropy density $s = dp/dT$, so that $Ts = \varepsilon + p$ and hence its energy density ε remain the same function of p. We take the deflagrating plasma to be at the same pressure p_0 as in case A, so that its energy density is again $\varepsilon_0' = 7\, p_0$, whereas its temperature is 0.99 T_0 and its entropy density $s_0' = (8/0.99)p_0/T_0$. For p below p_0, the T(p) curve of case B is then connected smoothly with the curve of A, and we make them identical for $p < 0.85\, p_0$. On the scale of Fig. 5 the curves of A and B are indistinguishable. The difference is shown in Fig. 6 which gives an enlargement of the transition region. One sees in this example that the transition region is very narrow in temperature ($\Delta T/T_0 \sim 3\%$) and wider in pressure ($\Delta p/p_0 \sim 15\%$). The condition of entropy increase is easily worked out for case B. It gives that deflagration is possible for gas pressures in the range

$$0.23\, p_0 \leq p_g \leq 0.83\, p_0 \tag{13}$$

with again a flat probability distribution over most of the range. In both cases A and B, for any reasonable equation of state of the plasma, the entropy condition forbids deflagration if the plasma pressure is significantly above p_0. For the smooth transition of case B, it is important to note in (13) that the deflagration produces gas with $p_g \leq 0.83\ p_0$, so that it "jumps" over the region where p is just below p_0. This shows that our analysis does not depend on the detailed properties of hadronic matter just below transition.

Using the equations of I, section 3, we calculate the dependence on the hadron gas pressure p_g of the following quantities, all of which we define in the rest frame of the deflagrating plasma :

i) the inward velocity v_s of the plasma surface (equal to the velocity v_L in I);

ii) the flow velocity v_g of the gas (called v in I) and the corresponding rapidity y_g;

iii) the energy flux F_ε in the gas in units of its maximum value $\varepsilon_0' v_s$, i.e., the quantity

$$\frac{F_\varepsilon}{\varepsilon_0' v_s} = \frac{1}{\varepsilon_0' v_s} \frac{(\varepsilon_g + p_g) v_g}{1 - v_g^2} . \tag{14}$$

In both cases A and B, for the p_g ranges (12) and (13) respectively, these quantities are the following functions of $\beta = p_g/p_0$:

$$v_s = (1 + 2\beta - 3\beta^2)^{1/2} \Big/ (49 - 14\beta - 3\beta^2)^{1/2} \tag{15}$$

$$v_g = \tanh y_g = (7 - 10\beta + 3\beta^2)^{1/2} \Big/ (7 + 22\beta + 3\beta^2)^{1/2} \tag{16}$$

$$F_\varepsilon = (49 - 14\beta - 3\beta^2) / 56. \tag{17}$$

They are plotted in Figs. 7 and 8.

These curves have interesting qualitative properties. The velocity v_s is always small, $v_s \lesssim 1/6$, i.e., it takes a time $\gtrsim 6$ fm/c for the plasma surface to recede by 1 fm. This smallness of v_s is understandable; for sensible values of y_g ($y_g \lesssim 1$), the gas can only take up slowly the high energy and entropy contents of the plasma. We also note that v_s is rather constant and close to its maximum over the range (13). As to the gas velocity v_g, it drops from a maximum of 2/3 to 0 as p_g ranges over (12). That there is a strong decrease is again understandable; at low p_g and ε_g most of the internal energy of the plasma (density ε_0') goes into collective flow of the gas, whereas more goes into internal energy of the gas when p_g is larger.

The fact that $0.6 < F_\varepsilon/\varepsilon_0' v_s \lesssim 0.8$ in Fig. 8 means that, as a layer of stationary plasma transforms into gas, some 60 to 80% of its energy contents transforms into energy flux of the gas, the balance being contributed by the small energy density of the gas. For $\varepsilon_0' = 4$ GeV fm^{-3}; $v_s = 1/6$ and $F_\varepsilon/\varepsilon_0' v_s = 0.75$ one finds an energy flux $F_\varepsilon = 0.5$ GeV fm^{-3}c, an order of magnitude larger than (9) at $T = T_0 = 200$ MeV. Such a flux is in fact quite high in absolute value; it amounts to $\varepsilon_0'/8$, whereas one would get $\varepsilon_0'/4$ if all partons of a Stefan-Boltzmann plasma which fly towards its surface would escape through it without deflection nor energy loss. The high flux is due to the fact that the energy density ε drops by a large factor in the process, about a factor 4 in the middle of the range (12). This can be compared with the rate of decrease of ε in the scaling expansion considered at the end of the previous section; for $w = 1/3$ we there had $\varepsilon \propto \tau^{-4/3}$ for $\tau \gtrsim 1$ fm/c, which gives a drop of ε by a factor 4 over the proper time intervals $\tau = 1 - 2.8$ fm/c, $2 - 5.7$ fm/c and $3 - 8.5$ fm/c (see Ref. 12 for more elaborate calculations of longitudinal expansion in the hydrodynamical model). Returning to our deflagration model, we note from Figs. 7 and 8 that the energy flux F_ε varies little over the wide range of pressure p_g where v_s is almost constant; the reason is that cold gas (low p_g, ε_g) escapes with higher velocity than hot gas (high p_g, ε_g).

While our results were calculated with very simple examples of equation of state, it is their qualitative nature which is of interest, in particular the possibility of deflagration, the high energy flux F_ε of the gas and the low inward velocity v_s of the plasma surface. These properties result from the general shape of the equation of state rather than from its detailed mathematical expression; they remain true in cases where the sound velocity of the gas varies with pressure, as will happen for realistic equations of state. In such cases, however, the range of gas pressures p_g reached with substantial probability by the deflagration may be narrower than in our examples, while still falling well below the pressure p_0 of the deflagrating plasma.

For given properties of plasma and gas, our very elementary treatment of the deflagration process uses only energy-momentum conservation and the requirement of entropy increase, i.e. of transitions to states of larger probability[7]. Hence, the qualitative characteristics of the process are not likely to be changed by effects due to thickness and curvature of the plasma surface. Also incomplete thermalization of the plasma before and during the deflagration should not cause qualitative changes; all what is required is that the energy-momentum tensor and the number of states of given internal energy (exponential of entropy) be rather close to equilibrium values. Nevertheless, the considerations of the previous section suggest that the end phase of the expansion of the plasma droplet and its possible break up in smaller droplets are rather slow processes, favouring an approach to local equilibrium before the deflagration starts. In addition, since the deflagration is slow (small v_s), there is some time for the plasma to thermalize further during the deflagration. Still, one should remember that in our model the stopping of the expansion and the deflagration are irreversible processes producing entropy, so that the overall evolution cannot be entirely adiabatic.

In conclusion of this section, we sketch in Fig. 9 the world lines of collective flow predicted by our model in the plane of the time t and the longitudinal position z already considered at the end of Section 3. Below the broken line, a droplet of plasma expands according to the scaling law (4). Near the broken line

which corresponds to constant proper time, the droplet breaks and gives rise to two non-expanding droplets of plasma near transition, which begin to deflagrate. The dotted line gives the movement of the deflagration fronts. Below that line one has plasma and above it the hadron gas formed by the deflagration.

5. EXPERIMENTAL CONSEQUENCES

Despite its qualitative nature, the mechanism we have proposed for the expansion and hadronization of quark-gluon plasma droplets which could be formed in high energy collisions has rather specific experimental consequences which we discuss in this section. In our model, if plasma is created in a collision, its evolution goes through two distinct phases:

Phase A - The plasma expands, strings form and some of them may break, the tension of the remaining strings stop the expansion, and the plasma takes the form of one or more small non-expanding droplets in a state just above transition. There is little hadronization during this phase.

Phase B - Each non-expanding droplet resulting from phase A hadronizes by deflagration processes of the type described in I and in the previous section, i.e., by ejection of low pressure hadron gas through the outer surface. This is an irreversible process presumably triggered by fluctuations at the end of the expansion and has therefore a stochastic character.

The most direct prediction of our model is the occurrence of isolated peaks or high bumpy intervals in the longitudinal rapidity distribution dn/dy of the outgoing hadrons (all hadrons, or all charged hadrons). This can be seen as follows. The hadronization of a non-expanding plasma droplet of longitudinal rapidity y_0 will give hadrons in a limited rapidity interval around y_0, of order $|y - y_0| \lesssim 0.5 - 1$, because the hadron gas will escape in all directions with velocities $v_g \lesssim 0.6$ relative to the plasma droplet. As explained in section 4, v_g may vary considerably from direction to direction and from droplet to droplet, higher v_g going with colder gas (lower p_g, see Fig. 7). The contribution of a droplet to dn/dy should be large because the multiplicity reflects the entropy contents which is high in the plasma and can only increase further in the deflagration and the later transformation of the expanding hadron gas in a vapour of non-interacting hadrons (freeze-out). We note in this connection that the entropy density in the plasma at transition can be as high as 20-30 fm^{-3} with our previous estimates, see eq. (8).

The contribution of a non-expanding droplet to dn/dy is hard to determine quantitatively because large fluctuations can occur from case to case and because the observed pions will often be decay products of heavier mesons. Qualitatively, we expect that a single plasma droplet will contribute to dn/dy a peak of width $\delta y_0 \lesssim 1$ above the background contributed by the rest of the particle production process. The particles in the peak will tend to be distributed over all azimuths around the longitudinal direction, with again the possibility of substantial fluctuations. The peak in dn/dy should be isolated if no other plasma droplet occurs within a few units of rapidity. This raises the question of the rapidity separation δy between adjacent non-expanding plasma droplets resulting from the break up of a larger expanding droplet. We have seen in section 3 that δy could be ~ 1 or significantly larger. In the latter case, or if δy_0 is well below 1, the peaks in dn/dy will be well separated, otherwise there will be considerable overlap and the overall effect of plasma formation will be a high and bumpy region in dn/dy, extending over several units of rapidity. The two cases are illustrated in Fig. 10. It should be stressed that all these considerations are meant on an event-by-event basis; inclusive distributions of one or a few particles would reveal very little of the above properties. One should also note that the dn/dy properties just discussed depend mainly on the features of phase A of the plasma evolution, and only weakly on those of phase B.

Also the transverse momentum distribution should reflect plasma formation. Although a quantitative determination of the p_T distribution is difficult for the reasons mentioned above in connection with dn/dy, the qualitative expectations are rather simple*). The thermal motion in the hadron gas, characterized by its temperature T_g, is boosted by the flow velocity v_g of the gas with respect to the

*) The author is indebted to R. Hagedorn for a discussion on this point.

plasma droplet (the original transverse velocity of the droplet should be too small to have significant effects). While the gas pressure p_g may vary widely (see Section 4), T_g which goes roughly as ${p_g}^{1/4}$ should remain in a rather narrow range around $T_0/2^{1/4}$. For $T_0 \sim 200$ MeV this is ~ 170 MeV, somewhat larger than the Hagedorn temperature which fits normal p_T distributions in terms of thermal motion alone. Furthermore, lower T_g and p_g go with larger v_g (Fig. 7) and hence with a larger boost effect due to the flow velocity of the gas. At least if $T_0 \gtrsim 200$ MeV, we therefore expect a widening of the p_T distribution to be a second signal for plasma formation. It should occur in the same events and rapidity intervals as the dn/dy signal discussed above.

To go beyond these crude considerations, the best approach may be to concentrate on heavy mesons (ρ, K, K*) rather than to consider the overall p_T distribution which mainly concerns pions (only part of the detected pions belong to the hadron gas, the others being decay products of resonances). The distribution of such heavy mesons in y and p_T could be measured in the rapidity intervals where dn/dy indicates likely plasma formation (isolated maxima or high bumpy regions). More directly than for pions, it should reflect the properties of the flowing hadron gas ejected by deflagration of the plasma, i.e., the characteristics of phase B of the plasma evolution.

The above implications of our surface deflagration model for dn/dy and p_T distributions should also be considered in conjunction with the "traditional" signals of plasma formation discussed by many authors and based on direct dileptons, direct photons and strange particles[1]. These signals reflect the properties of the quark-gluon plasma before it hadronizes, and they are therefore complementary to the features discussed above which mainly concern the way the plasma hadronizes. But there is an obvious correlation, in the sense that we predict the traditional signals to appear in the same events, and per event in the same rapidity intervals, as the dn/dy signal proposed above.

In our opinion, experimental work in this domain should take into account the possibility that for given beam, target and incident energy, plasma formation would occur (if at all!) in a fraction of events, and that in such events only a fraction of the secondaries would originate from a plasma droplet. If our model is right, the dn/dy signal we propose can be used for filtering out candidate events and localizing in rapidity where their possible plasma phase was likely to be situated. Finding in these events and rapidity intervals not only higher p_T but also an abnormal behaviour of the traditional signals (presumably an abnormally high relative abundance) would then suggest that a plasma droplet may have been formed. Caution would be needed, however, in comparing quantitatively the data with the existing theoretical calculations for the traditional signals. The reason is that the latter are usually performed in perturbative approximation, whereas the plasma, if formed in a collision, will spend most of its lifetime near the transition temperature, where in our view it should already be outside the perturbative regime (especially due to string formation, see Section 3). This difficulty is less serious if the initial plasma temperature is $\gg T_0$, especially for strange particle production[1,13,14]) which should be mainly controlled by the high temperature part of the plasma evolution; it is then that most strange (anti)quarks s, $\bar{s}$ will be created and their disappearance during plasma expansion by $s\bar{s}$ annihilation will be very small due to their low density (disappearance by weak decay during the plasma phase is negligible).

Assuming that one day the signals will have been found and the tests are positive, can one then go further and extract from high energy collisions more detailed information on the quark gluon plasma and the hadronic phase transition? The answer is probably yes for some features, e.g. the size of the jump in energy density which strongly affects the deflagration, and no for some others, e.g. the true nature of the transition (first or second order, continuous) which would require larger and longer-lived plasmas than nuclear collisions can provide. Unfortunately the answer is also no for the property that should differentiate between quark-gluon plasma and hadron gas in the most spectacular way, to wit the colour-electric conductivity which should be large in the plasma and zero in the gas. It is true that pair production of very heavy quarks could provide a slowly

varying colour-electric field over time intervals of several fm/c. But even if a plasma would be produced in the same collision in the same rapidity interval, the field would be too weak (a single triplet charge per heavy quark) for the colour conductivity to have significant effects.

6. CONCLUDING REMARKS AND SUMMARY

An obvious question we have not addressed so far is whether evidence already exists for plasma formation in high energy collisions. We discuss this first for cosmic ray events, then for pp and $p\bar{p}$ processes at the CERN colliders.

Strong fluctuations in $dn/d\eta$ (η = pseudorapidity) have been observed in high multiplicity cosmic ray collisions involving nuclei[15]). It is tempting to relate them to formation of quark-gluon plasma, as already suggested by Gyulassy[16]). Some theoretical models of particle production are able to reproduce the high observed multiplicities[17]). In the energy region of these events ($\gtrsim$ 1 TeV/nucleon of cosmic ray nucleus) one estimates energy densities of several GeV fm^{-3} in the central rapidity region, high enough for plasma formation. The models do not reproduce the fluctuations in $dn/d\eta$, but this could be remedied by assuming that the objects produced are clusters instead of single hadrons*). The JACEE events[15]) have larger than normal values for the average p_T, but the correlation in η between higher $\langle p_T \rangle$ and $dn/d\eta$ peaks is not clear. While intriguing and encouraging, this evidence is insufficient to draw conclusions on plasma formation.

On the question of possible plasma formation in hadron-hadron collisions, opposite views are found in the literature. While Kajantie and McLerran[18]) advocate that plasma formation is only to be expected in sufficiently heavy nuclei, Hagedorn[19]) argues that positive evidence already exists at CERN in pp collisions at the ISR (centre of mass energy E_{cm} up to 63 GeV) and $p\bar{p}$ collisions at the SPS collider (E_{cm} = 540 GeV). Hagedorn's arguments imply very high temperatures, up to 400-500 MeV at the ISR and $\sim$ 700 MeV at the $p\bar{p}$ collider, so that the corresponding energy densities would be well above 100 GeV fm^{-3}. At the ISR, this means by energy conservation that the plasma volume would be $\ll 1$ fm^3, which is much too

*) This was suggested by G. Ekspong on the basis of similar findings of the UA5 Collaboration on the $p\bar{p}$ collider (see below).

small in our view for the application of thermodynamical concepts. At the $p\bar{p}$ collider, however, Hagedorn's temperature estimate cannot be criticized on these grounds.

The UA5 Collaboration has analyzed the pseudorapidity distribution $dn/d\eta$ of single events at the $p\bar{p}$ collider. They find[20] that 0.5 - 1% of their streamer chamber events have spikes in $dn/d\eta$ about 3-4 times the average background level, with a width $\sim$ 0.5 in η (the biggest observed spike contains 18 tracks in half a unit of pseudorapidity; assuming the p_T's to be normal one estimates an energy density of $\sim$ 10 GeV fm^{-3}). If the p_T's are normal, such spikes can be reproduced by fluctuations in a rather conventional cluster model. While again intriguing and encouraging, this evidence is inconclusive. Abnormally large p_T's, abnormally high relative abundances of strange particles, direct dileptons, direct photons should be looked for; if they occur in the spikes, they would provide very serious indications for plasma formation.

In our view, an important fact in assessing the situation at the ISR and SPS collider is that the latter shows a strong increase of average p_T with multiplicity in the central region of rapidity[21], whereas at the ISR this increase is weak at E_{cm} = 63 GeV and is not seen at E_{cm} = 31 GeV[22]. Clearly, from the ISR energy range onward, an increasing amount of incident energy is transferred to transverse motion without creation of separated transverse jets, the simplest interpretation being an increasing amount of thermalization[23,24]. The dual parton model, which successfully accounts for multiplicities, requires implausible ad hoc assumptions to incorporate this effect[25]. We noted in Ref. 23 that the shape of the average p_T versus multiplicity correlation could reveal plasma formation if it occurred in most collisions, an argument that does not require a first order phase transition (Fig. 2 or Ref. 23 is unchanged if the transition in its Fig. 1 is smoothed out over a small temperature interval). We also noted that the available data were inconclusive in this respect, a situation which still prevails today.

Although we know of no experimental data already indicating plasma formation, we believe that a systematic search is warranted. At ultra-relativistic energies, it should involve the study of fluctuations in dn/dy, increase in p_T and the

"traditional" plasma formation signals for hadron-hadron, hadron-nucleus and nucleus-nucleus collisions. Interpretation of the data will have to rely on qualitative features, in particular on qualitative differences between the expected manifestations of plasma formation and the predictions of conventional particle production mechanisms. This is due to the fact that no quantitative theory exists for the proposed signals, neither in the plasma near transition, nor in conventional production mechanisms. The present paper has been written in this spirit; it shows that qualitative signals and correlations are indeed available for guiding the experimental work.

We conclude with a brief summary. The most important properties of hadronic matter assumed in this work are the large differences of entropy density s and energy density ε between the gas and plasma phases, the localization of the transition in a narrow interval of temperature T, and the smallness of the pressure p compared to $Ts = \varepsilon + p$. The hadronization of a plasma droplet assumed to be created in a high energy collision must therefore liberate much entropy and energy corresponding to freezing out of many of the QCD degrees of freedom excited in the plasma. We have proposed what we believe to be a plausible model for this hadronization process. It describes the end of the expansion of a plasma droplet when transition is reached (formation of QCD strings and possible break up of the droplet, the string tension stopping the expansion), followed by the bulk of hadronization in the form of ejection of hadron gas through the outer plasma surface (a process with the kinematic and thermodynamic properties of a deflagration). Studying a very simple model, we found that the deflagration is slow in the sense that the plasma surface recedes at a small fraction of the velocity of light ($\lesssim c/6$), but the resulting energy fluxes in the hadron gas are found to be high, about half of what a free escape of QCD partons would give. The flow velocity v_g of the emitted hadron gas ranges up to $\sim$ 0.6 c and its pressure p_g ranges down to a small fraction of the pressure at transition, larger v_g going with smaller p_g (see Fig. 7). The overall result is predicted to be the appearance on an event-by-event basis of one or more peaks of width $\lesssim 1$ in the rapidity distribution dn/dy of final state hadrons, either isolated or overlapping into a high bumpy region a few rapidity units wide, and p_T values larger than normal, all this with large fluctuations from event to event and presumably in a fraction only of all events. These predictions concerning dn/dy and p_T should be considered in conjunction with the traditionally proposed signals (direct dileptons and photons, strange particles), the latter being expected to occur in the rapidity intervals indicated by the former.

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Figure Captions

Fig. 1 : The qualitative shape of the equation of state for hot hadronic matter at zero chemical potential. Fig. 1a refers to a first order phase transition with metastable states (dashed parts of the curves), Fig. 1b corresponds to a smooth transition.

Fig. 2 : A quark-gluon plasma well above transition (Fig. 2a) and near transition (Fig. 2b). The QCD partons are represented by triangles for quarks and antiquarks, by circular dots for gluons. The hatched area represents the region of space occupied by the colour fields.

Fig. 3 : A piece of QCD string network involving quarks (triangles) and gluons (circular dots). The SU(3) index structure is indicated (see text).

Fig. 4 : A string under tension (above) can break by quark pair creation (below).

Fig. 5 : The equation of state of cases A and B (see text).

Fig. 6 : Enlargment of Fig. 5 in the transition region.

Fig. 7 : The inward velocity v_s of the plasma surface, the velocity v_g and rapidity y_g of the ejected gas, all defined in the plasma rest frame, are plotted versus gas pressure.

Fig. 8 : The energy flux F_ε in the gas defined in the plasma rest frame and divided by its maximum value is plotted versus the gas pressure.

Fig. 9 : The evolution of an expanding droplet of plasma in longitudinal space and time according to our model. The plasma reaches transition and stops expanding at the broken line, it deflagrates into hadron gas at the dotted line (see also the text).

Fig. 10 : Qualitative shape of the longitudinal rapidity distribution of hadrons (or of charged hadrons) predicted for the case of plasma formation. Non-expanding plasma droplets may give well separated peaks (a) or a high bumpy region resulting from overlapping peaks (b).

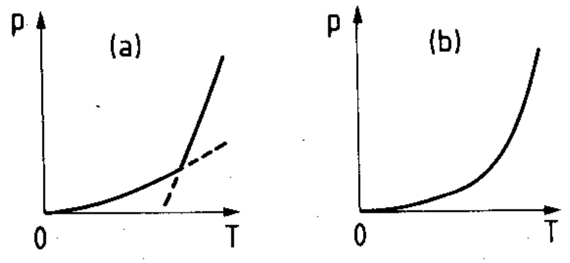

Fig. 1

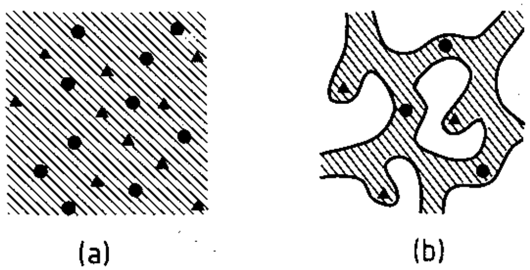

Fig. 2

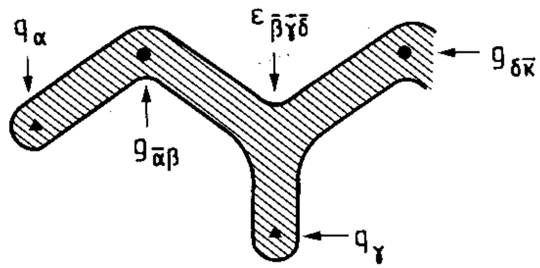

Fig. 3

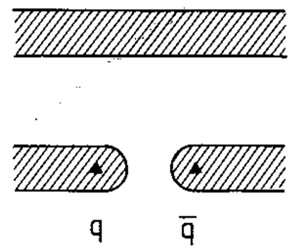

Fig. 4

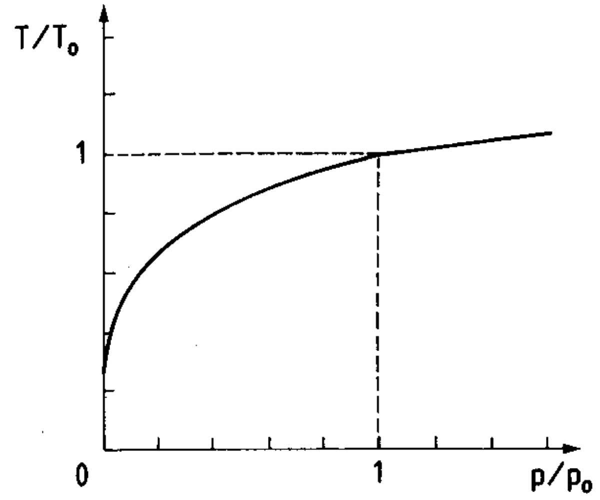

Fig. 5

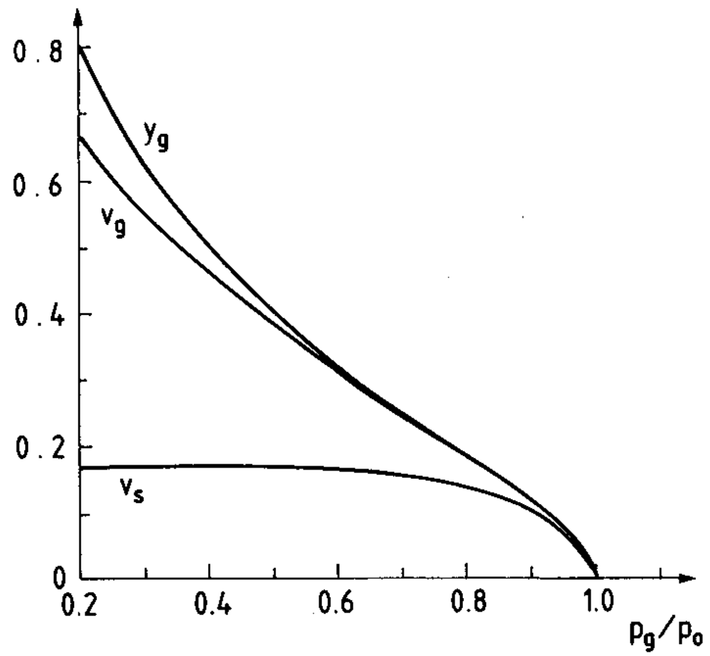

Fig. 7

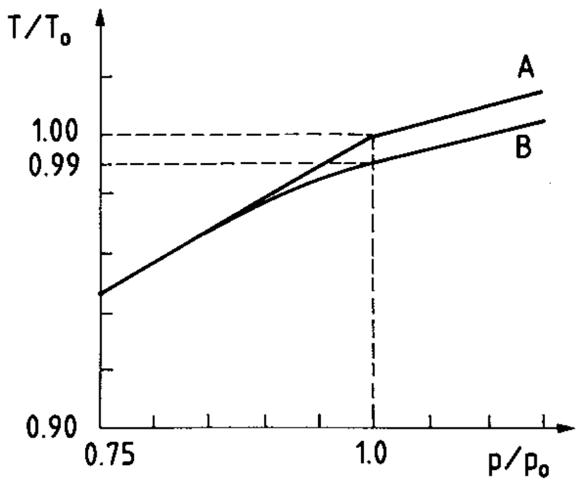

Fig. 6

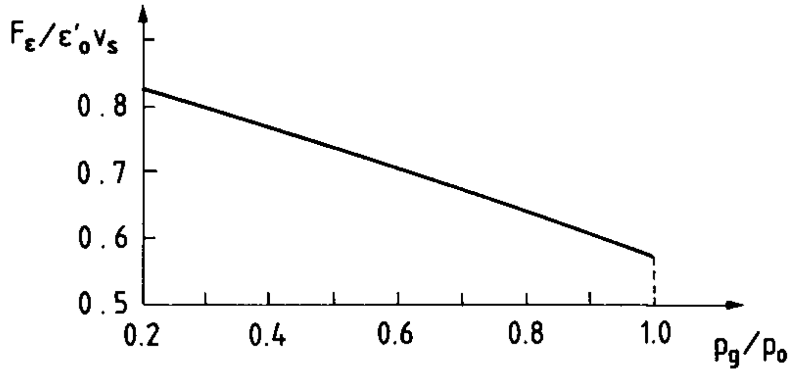

Fig. 8

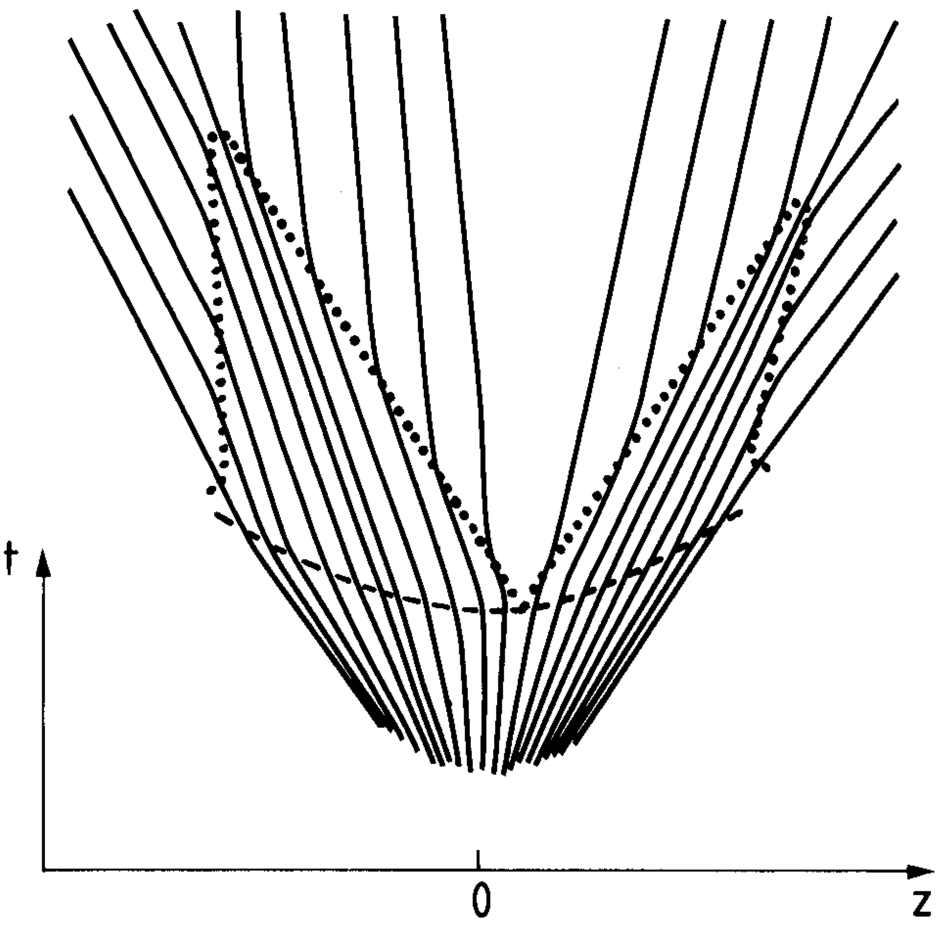

Fig. 9

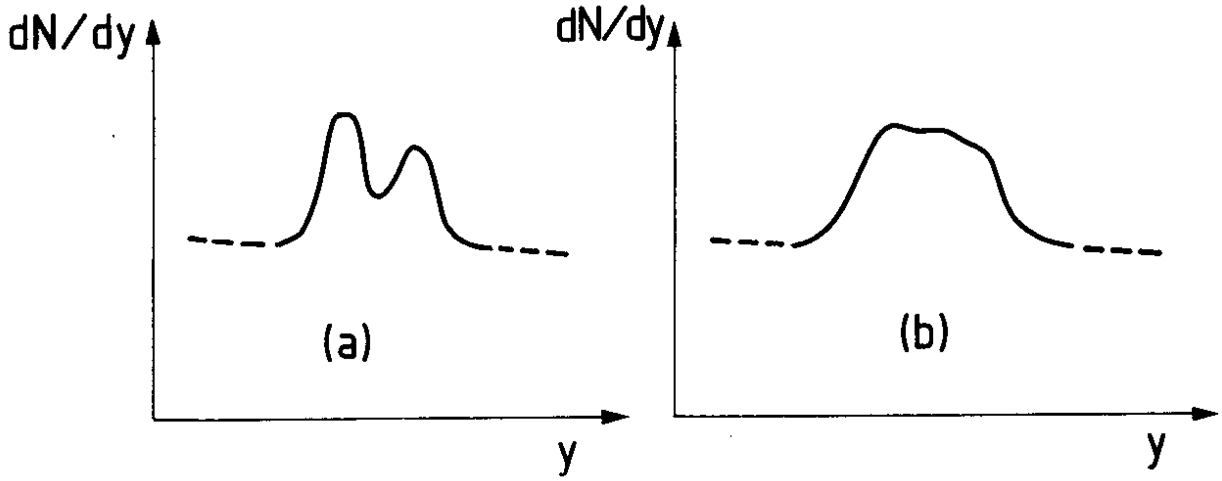

Fig. 10