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Abstract

The hydrodynamic evolution of a quark-gluon plasma, produced in the central rapidity region, is studied incorporating the external bag pressure acting on the plasma surface. It is shown that the plasma fluid, which undergoes the scaling longitudinal expansion, generates a non-trivial transverse flow pattern, near the plasma surface, consisting of a rarefaction wave and a compression wave, instead of a simple rarefaction wave in the case of free expansion without a surface boundary condition. We also discuss a possible global phase transition of supercooled plasma, in the interior of the cylinder, to a superheated dense hadron gas through a time-like surface of discontinuity.

1. In the past few years interest in ultrarelativistic nucleus-nucleus collisions has been considerably stimulated by the possibility of creating a new form of matter (quark-gluon plasma) which once pervaded the early universe.[1,2] At relatively low bombarding energies, $E_{CM} = (2 - 5)$ GeV per nucleon, where the nuclear stopping power reaches a maximum, a quark-gluon plasma which is characterized by large baryon number density would be produced.[3] On the other hand, at very high energies, $E_{CM} \gtrsim (30 - 50)$ GeV per nucleon, where the nucleus becomes almost transparent with respect to incident nucleons, a hot plasma with small net baryon number is expected to be formed in the central rapidity region just after two nuclei pass through each other.[4,5] The space-time evolution of the hadronic matter produced in the latter regime was described by Bjorken,[4] using the relativistic hydrodynamics with a Lorentz-boost invariant initial condition on account of the ansatz of the central plateau formation in the rapidity distribution. Bjorken's one-dimensional scaling solution to the subsequent longitudinal hydrodynamic expansion of the central regime was extended recently by Baym et al.[6] to include the transverse flow, assuming cylindrical symmetry for the collision volume. Although these authors considered that the hadronic matter is formed in the deconfined quark-gluon plasma phase, the hadronization stage of the plasma has not yet been described in these studies. The aim of this paper is to put forward the above hydrodynamic picture of ultrarelativistic nuclear collisions, supplemented by two possible mechanisms of the hadronization of the plasma, and to study the dynamical effect of the hadronization on the hydrodynamic evolution of the plasma fluid.

Consider a space-time picture of the transverse plasma evolution, as schematically shown in Fig. 1, where the plasma domain is surrounded by two types of surface; one is a space-like surface with a space-like unit normal vector n_μ ($n^2 \equiv \underset{\sim}{n}^2 - n_0^2 > 0$) and the other a time-like surface with an opposite relation $n^2 < 0$. The space-like surface corresponds to the ordinary surface in three-dimensional coordinate space, which travels with a velocity smaller than the light velocity, while the time-like surface appears when an instantaneous bulk phase transition of the plasma occurs.

A semi-classical microscopic description of the hadronization process at a space-like surface of the quark-gluon plasma was obtained by the present authors by making use of the chromoelectric flux tube model.[7] It was shown that only high momentum quarks ($k_0 \gtrsim$ GeV) in the thermal distribution can penetrate the surface, accompanied by fission of the flux tube (hadronization), and most of the low momentum thermal quarks are reabsorbed by the plasma due to color confinement. We shall use this result to derive a boundary condition at the space-like surface of the plasma fluid. Our problem is then reduced to the problem of gas flow in the cylindrical domain which is covered by an almost impenetrable membrane with the external bag pressure acting on it. It will be shown that the resultant transverse flow, which couples to the scaling longitudinal expansion, consists of an inward moving rarefaction wave and a compression wave generated near the surface. This is in contrast to the result of a free expansion to the vacuum without a surface boundary condition,[6] where a simple rarefaction occurs. We will see that, with reasonable values for the initial plasma temperature, the compression wave dominates the rarefaction wave and, consequently, the plasma shrinks rather than expands. This happens due to the longitudinal expansion which causes a uniform rapid decrease of the interior plasma pressure below the external bag pressure. If the interior region of the plasma is sufficiently supercooled, without forming a bubble of hadron gas,[8] an instantaneous global phase transition to a superheated dense hadron gas may take place through a time-like surface, followed by an explosive transverse expansion of the system. We shall estimate the increase of the temperature and pressure of the system by this dynamical phase transition.

2. The basic equations in the following analysis are obained from the energy-momentum conservation law:

$$\partial_\mu T^{\mu\nu}(x) = 0 \quad . \tag{1}$$

Introducing a step function $\theta(x)$ which becomes $\theta(x) = 1$ in the plasma domain and vanishes outside of this region, and writing the stress-energy tensor as

$$T^{\mu\nu}(x) = T^{\mu\nu}_{IN}(x)\theta(x) + T^{\mu\nu}_{OUT}(x)(1 - \theta(x)) \quad , \tag{2}$$

Eq. (1) splits into three equations:

$$\theta(x)\partial_\mu T^{\mu\nu}_{IN}(x) = 0 \tag{3}$$

$$(1 - \theta(x)) \; {}_\mu T^{\mu\nu}_{OUT}(x) = 0 \tag{4}$$

$$T^{\mu\nu}_{IN}(x) n_\mu \delta(x) = T^{\mu\nu}_{OUT} n_\mu \delta(x) \tag{5}$$

where, in deriving the last equation, we used $\partial_\mu \theta(x) = n_\mu \delta(x)$ with a surface delta function $\delta(x)$ and a surface normal unit vector n_μ.

In the collision dominated plasma region, the stress-energy tensor may be written in a form of perfect fluid gas in terms of a local pressure $p(x)$, a local energy density $\varepsilon(x)$ and a local four fluid velocity $u^\mu(x)$

$$T^{\mu\nu}_{IN}(x) = (p + \varepsilon)u^\mu u^\nu + pg^{\mu\nu} \tag{6}$$

where we use the convention, $g^{oo} = -1$ and $g^{11} = g^{22} = g^{33} = 1$, for the metric tensor. Upon insertion of an equation of state, Eqs. (3) and (6) describe the hydrodynamics of the plasma fluid. In the following we shall use an ideal gas equation of state (with zero chemical potential) for the quark-gluon plasma. The energy density, pressure and entropy density are given by,

$$\varepsilon = aT^4 \quad , \quad p = \frac{1}{3}\,\varepsilon = \frac{1}{3}\,aT^4 \quad , \quad \sigma = \frac{4}{3}\,aT^3 \tag{7}$$

where

$$a = \frac{\pi^2}{30}\,(16 + \frac{21}{2}\,N_F) = \frac{37}{30}\,\pi^2 \quad , \tag{8}$$

for a plasma with two flavour components $(N_F = 2)$.

3. The stress-energy tensor in the region outside of the plasma takes a more complex form. It contains a vacuum term corresponding to quantum fluctuations in the non-perturbative QCD vacuum, which we represent by the phenomenological bag term $Bg^{\mu\nu}$. In addition, there is a hadronic term whose form depends on the state of the hadrons coming out of the plasma. If the hadron gas is sufficiently dense and locally thermalized, the hadronic term is reduced to the hydrodynamic form as in Eq. (6). This approximation has been taken in ref. 9, 10, 8 to obtain a macroscopic description of the space-like surface of the plasma by making use of relativistic combustion theory. However it was found in a microscopic calculation of the hadronization at the space-like plasma surface[7] that the chemical reaction from the deconfined quarks to mesons proceeds very slowly due to color confinement. This implies that mesons would stream out of the plasma rather freely without strongly interacting with each other once they are formed at the plasma surface. In this case the stress-energy tensor outside of the plasma becomes

$$T^{\mu\nu}_{OUT}(x) = Bg^{\mu\nu} + \int dn(k,x)\frac{k^\mu k^\nu}{E} \qquad (9)$$

where $dn(k,x)$ is the density of hadrons with four momentum $k^\mu = (E,k)$ at x. Near the plasma surface the meson density is obtained from the radiating meson flux normal to the surface as

$$\frac{d^2n}{dk_n dE} = \frac{d^2(\text{Meson Flux})_n}{dk_n dE} \frac{1}{v_n + v_s} \qquad (10)$$

where k_n and $v_n = k_n/E$ are the meson momentum and velocity normal to the surface, measured on the rest frame of the plasma fluid element in the vicinity of the surface. Note here that the plasma surface is moving inward with velocity v_s on this frame because the plasma is being eaten away by the meson radiation. Thus the unit four vector normal to the surface is given by

$$n_\mu = \gamma_s(v_s, \underset{\sim}{n}) \tag{11}$$

where $\gamma_s = (1 - v_s^2)^{-1/2}$ and $\underset{\sim}{n}$ is a unit three vector normal to the surface.

The meson flux from the plasma was given in ref. 7 as

$$\frac{d^2(\text{Meson Flux})_n}{dk_n dE} = \frac{2\gamma}{(2\pi)^3} \int d\underset{\sim}{k}_o \; f(k_o/T_s) v_{no} \frac{1}{k_c^2} \exp\left(\frac{k_o(k_n - k_{no})^2}{2k_{no}k_c^2}\right)$$

$$\theta(k_n < v_{no}E) \cdot \theta(v_{no}E < k_{no}) \cdot \theta(k_n < k_{no}) \cdot \theta(k_n > 0) \; . \tag{12}$$

Here the integration is over the thermal flux of massless quarks and antiquarks with four momentum $k_o^\mu = (k_o, \underset{\sim}{k}_o)$, multiplied by the probability that the flux tube formed by these outgoing quarks will fission to produce a meson of energy E and normal momentum k_n. Also $\gamma = \gamma_c \times \gamma_s \times \gamma_f = 3 \times 2 \times 2 = 12$ is the color, spin and flavor degeneracy factor of massless quarks and antiquarks, $f(k_o/T_s)$ is the thermal distribution function with surface plasma temperature T_s, k_{no} and $v_{no} = k_{no}/k_o$

are the initial quark momentum and velocity normal to the surface, and $k_c = \sqrt{12}/R_{FT}$ is a parameter which characterizes the fission probability of the flux tube with radius R_{FT}. If we adopt the MIT bag model to describe the chromoelectric flux tube, the flux tube radius can be related to the bag constant B and the string tension $\sigma \doteq 0.9$ GeV/fm as $R_{FT} = (\sigma/2\pi B)^{1/2}$,[11] thus yielding

$$k_c = (24\pi B/\sigma)^{1/2} \quad . \tag{13}$$

Using Eqs. (6), (9), (10), (11), one can readily derive the boundary condition at the space-like plasma surface from Eq. (5).

$$p_{pl}(T_s) = B + \int dk_n dE \frac{d^2(\text{Meson Flux})_n}{dk_n dE} k_n \quad , \tag{14}$$

$$\epsilon_{pl}(T_s)v_s = -Bv_s + \int dk_n dE \frac{d^2(\text{Meson Flux})_n}{dk_n dE} E \quad . \tag{15}$$

The integrals contained in these equations are to be carried out with Eq. (12). The leading order term in a power series expansion of T_s/k_o can be easily computed analytically by taking a Boltzman distribution function for $f(k_o/T_s)$ in Eq. (12). The result is

$$\int dk_n dE \frac{d^2(\text{Meson Flux})_n}{dk_n dE} k_n = \frac{10\gamma}{(2\pi)^2} T_s^4 \left(\frac{T_s}{k_c}\right)^2 + T_s^4 O((T_s/k_c)^4) \tag{16}$$

$$\int dk_n dE \frac{d^2(\text{Meson Flux})_n}{dk_n dE} E = \frac{80}{3} \frac{\gamma}{(2\pi)^2} T_s^4 \left(\frac{T_s}{k_c}\right)^2 + T_s^4 O((T_s/k_c)^4) \tag{17}$$

Then Eqs. (14) and (15), combined with Eqs. (7), (16) and (17), determine the surface temperature and the surface receding velocity v_s as functions of the bag constant B and k_c. Again we derive a power series expansion in $B^{1/2}/k_c^2$:

$$T_s = [0.705 + 0.0655(B^{1/2}/k_c^2) + O(B/k_c^4)]B^{1/4} \tag{18}$$

$$v_s = 0.50(T_s/k_c)^2 = 0.25(B^{1/2}/k_c^2) + O(B/k_c^4) \tag{19}$$

The leading order term in Eq. (18) corresponds to the surface temperature of the plasma in the absence of the leakage of quarks from the plasma surface. This temperature is also identical to the internal temperature of hadrons in the thermodynamic limit for the bag model. Recalling the relation (13) and choosing $B^{1/4}$ = 280 MeV which gives k_c = 1.62 GeV, corresponding to the flux tube radius R_{FT} = 0.43 fm, the first order correction due to the meson radiation is estimated to be 0.28% increase for T_s (thus T_s = 198 MeV) and v_s = 0.0075. These small numbers reassure us of the consistency of our calculations. For comparison note here that the critical temperature T_c of the first order phase transition from quark-gluon plasma to massless pion gas is determined by the Gibbs condition

$$p_{pl}(T_c) = B + p_{pion}(T_c) \tag{20}$$

where the equation of state for the noninteracting massless pion gas is given by

$$\varepsilon = bT^4 \quad , \quad p = \frac{1}{3}\,\varepsilon = \frac{1}{3}\,bT^4 \quad , \quad \sigma = \frac{4}{3}\,bT^3 \tag{21}$$

with $b = \pi^2/10$.

Using the same number for B, this critical temperature is estimated as $T_c = 0.72B^{1/4} = 202$ MeV. Thus we see that the space-like surface of the plasma is slightly supercooled. Also note here that although our bag constant is much larger than that determined phenomenologically by hadron spectroscopy ($B^{1/4} = 140 - 210$ MeV), it serves as a good parametrization of the result of a Monte Carlo calculation on the SU(3) lattice ($T_c = 190 \pm 20$ MeV).[12]

4. We now solve the hydrodynamic equations (Eqs. (3) and (6)) to determine the transverse flow of the plasma fluid, constrained by the surface boundary conditions (Eqs. (18) and (19)) at the transverse edge ($r = R_s$) of the plasma domain. As shown in ref. 6, for a cylindrically symmetric expansion accompanying the scaling longitudinal expansion, the hydrodynamic equations are reduced to a single characteristic equation

$$\frac{\partial}{\partial t} a + \frac{v_r + c_s}{1 + v_r c_s} \frac{\partial}{\partial r} a = \frac{-c_s}{1 + v_r c_s}\left(\frac{v_r}{r} + \frac{1}{t}\right) a \qquad (22)$$

where a and v_r are functions of r and t defined in the extended space interval $-R_s < r < R_s$. The local temperature and fluid velocity are given in terms of $a(r,t)$ by

$$T(r,t) = T_0(a(r,t)a(-r,t))^{c_s/2} \qquad (23)$$

$$v_r(r,t) = \frac{a(r,t) - a(-r,t)}{a(r,t) + a(-r,t)} \quad . \qquad (24)$$

This form automatically incorporates the boundary condition on the central axis of the cylindrical plasma domain, $v_r(r = 0, t) = 0$. The initial condition is now parametrized by $a(r, t = 0) = [1 + \alpha \exp((|r| - R)/\varepsilon)]^{-1/2}$ with α determined by $(1 + \alpha)^{-c_s/2} T_0 = T_s$. The initial temperature of the plasma drops from its interior value T_0 to the surface temperature T_s in the vicinity of the surface $R_s - |r| < \varepsilon$ while the initial transverse fluid velocity vanishes everywhere in the cylindrical plasma domain. The change of $a(r,t)$ along the characteristic line, whose slope is given by $dr/dt = (v_r(r,t) + c_s)/(1 + v_r(r,t)c_s)$, is then computed by $da(r,t) = -c_s(v_r/r + 1/t)/(1 + v_r c_s)a(r,t)dt$. In contrast to the free expansion without surface boundary conditions, each characteristic line terminates at $r = R_s(t)$, where the change of the surface location is determined by $dR_s(t) = (v_r(r = R_s(t),t) - v_s)/(1 - v_r(r = R_s(t),t)v_s)dt$, and a new characteristic line appears from $r = -R_s(t)$.[10] The value of $a(r,t)$ at $r = -R_s(t)$ on this new characteristic line should be computed so as to reproduce a constant surface temperature T_s. Therefore $a(-R_s(t),t) = (T_s/T_0)^{2/c_s}/a(R_s(t),t)$. One can continue to iterate this procedure as long as characteristic lines do not intersect.

A numerical solution constructed in the above method is presented in Fig. 2, where we start with the initial internal temperature $T_0 = 210$ MeV, which gives the initial energy density $\varepsilon = 3.1$ GeV/fm^3 (note that this energy density becomes $\varepsilon + B = 3.9$ GeV/fm^3 when it is measured from the non-perturbative QCD vacuum), and with the initial time $t_0 = 1$ fm/c and the initial radius of the system $R_s(t = t_0) = 7$ fm, corresponding to central U + U collisions. It is observed that, near the surface, a small rarefaction wave, generated by the initial pressure gradient, travels inward with the velocity of sound and a compression wave follows this. The rarefaction wave converts the heat energy of the fluid into the collective transverse flow energy, while the compression wave heats up the system again to reach the fixed surface temperature T_s. In the meantime, the temperature of the interior region falls off uniformly due to the longitudinal scaling expansion. This cooling is so rapid that, in the presence of the external bag pressure, significant compression occurs near the surface and, as a result, the system shrinks in the transverse direction. Eventually the compression wave would lead to a formation of a shock discontinuity, signaled by a crossing of characteristic lines, if a global hadronization, as described below, does not take place before.

5. As the time elapses, the interior region of the plasma continues to be diluted and, when it is sufficiently supercooled, a dynamical phase transition to a dense hadron gas may occur from the interior. On account of the stability argument on the longitudinal scaling solution[4,6] and the spatial uniformity in the transverse temperature distribution (see Fig. 2), we consider here an instantaneous uniform phase transition, neglecting the formation of hadron bubbles[8] triggered by irregularities in the system. Since the fluid is at rest in the interior (at $z = 0$) this happens in the space-time picture through a time-like surface characterized by a time-like four vector normal to this surface,

$$n_\mu = (1,0,0,0) \tag{25}$$

If the temperature is T_Q on the plasma side of this surface and T_H on the hadron gas side, the continuity of energy-momentum through this surface (Eq. (5)) imposes

$$\epsilon_{pl}(T_Q) = -B + \epsilon_{had}(T_H) \quad . \tag{26}$$

Also there is another constraint due to the second law of thermodynamics, which requires the increase of entropy density in this phase transition

$$\sigma_{pl}(T_Q) \leq \sigma_{had}(T_H) \tag{27}$$

To illustrate the consequence of these constraints we use here a massless pion gas approximation, neglecting effects of interactions, for the hadron gas equation of state expressed by Eq. (21). In Fig. 3 we illustrate the constraints (26) and (27) between T_Q and T_H. The shaded region is forbidden due to the entropy increase law. We see that the allowed combinations of T_Q and T_H are limited to a very narrow range on the solid curve. If the transition occurs through an isentropic path without generating extra entropy, one gets $T_Q = 0.5\ B^{1/4}$(=140 MeV) and $T_H = 2.31\ T_Q$(=323 MeV) where the numbers in the parentheses correspond to $B^{1/4} = 280$ MeV. Because in the scaling solution the plasma temperature decreases according to $T = T_o(t_o/t)^{1/3}$, this happens at the proper time $t_c = t_o(T_o/T_Q)^3 = 1.0$ fm/c(210 MeV/140 MeV)^3 = 3.4 fm/c. If the transition occurs Δt later than t_c, extra entropy is generated by $\Delta\sigma = \sigma_{had}(T_H) - \sigma_{pl}(T_Q) \simeq \sigma_{pl.}(1 - T_Q/T_H)\Delta t/t_c$ for small Δt. Note that, in any case, T_H is limited to a very narrow range, $1.00\ B^{1/4} < T_H < 1.16\ B^{1/4}$. Also the increase of pressure due to this bulk phase transition does not depend on the entropy increase:

$$\Delta p = B + p_{had}(T_H) - p_{pl}(T_Q) = \frac{4}{3} B \tag{28}$$

Thus the bulk phase transition generates high internal pressure which exceeds the external vacuum pressure and, therefore, the system starts to expand in the transverse direction. This stage of evolution of the system can be described by ordinary hydrodynamics without a surface boundary condition. Note that since this transverse expansion starts at $t_c (>> t_0)$, or perhaps later, the effect of the scaling longitudinal expansion is reduced and the internal fluid energy is more efficiently used to generate a transverse flow before the system breaks up into free streaming particles. This final stage of expansion is now under study, with a more realistic equation of state for the hot hadron gas. These results will be used to calculate the dilepton spectrum as well as the transverse momentum distribution of the secondary particles in order to assess the possible signatures of plasma formation.

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Figure captions

Fig. 1. Schematic space-time picture for the transverse evolution of a quark-gluon plasma produced in the central rapidity region. The plasma domain is surrounded by a space-like surface $(n^2 = \underset{\sim}{n}^2 - n_0^2 > 0)$ and a time-like surface $(n^2 < 0)$ as well.

Fig. 2. Temperature and fluid velocity distribution as functions of transverse distance r from the central axis of cylinder. The time which corresponds to each curve is indicated on the curve. Initial conditions at $t = t_0 = 1$ fm/c are set to $T_0 = 210$ MeV in the interior and $v(r) = 0$ for $0 \leq r \leq R_0 = 7$ fm. At the surface boundary, the plasma temperature is always maintained at $T_s = 200$ MeV according to the surface boundary condition (see Eq. (18)).

Fig. 3. Temperature of supercooled quark-gluon plasma (T_Q) versus a temperature of superheated hadron gas (T_H) at the global phase transition point. The allowed combinations are indicated by a solid curve. The shaded region is forbidden by the entropy increase law.

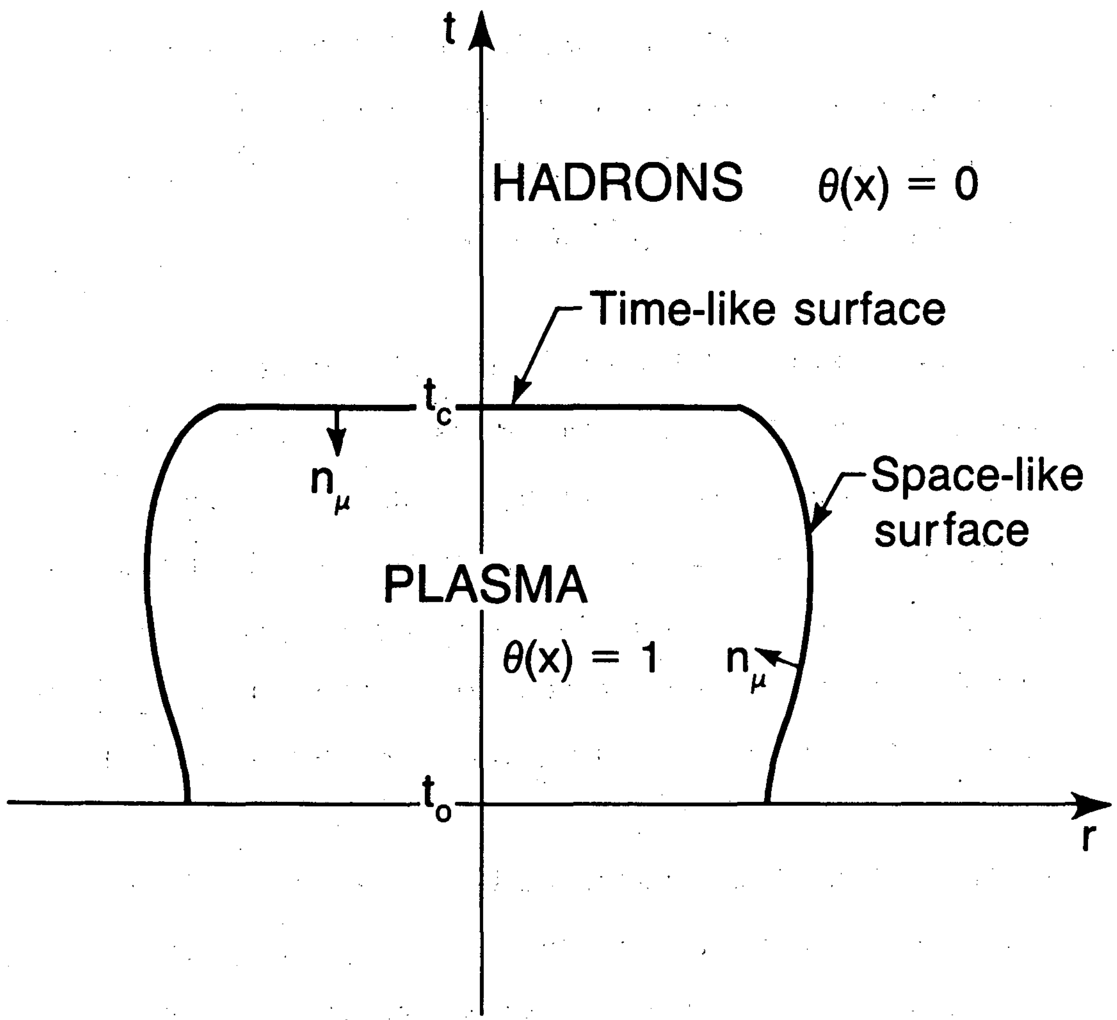

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Fig. 1

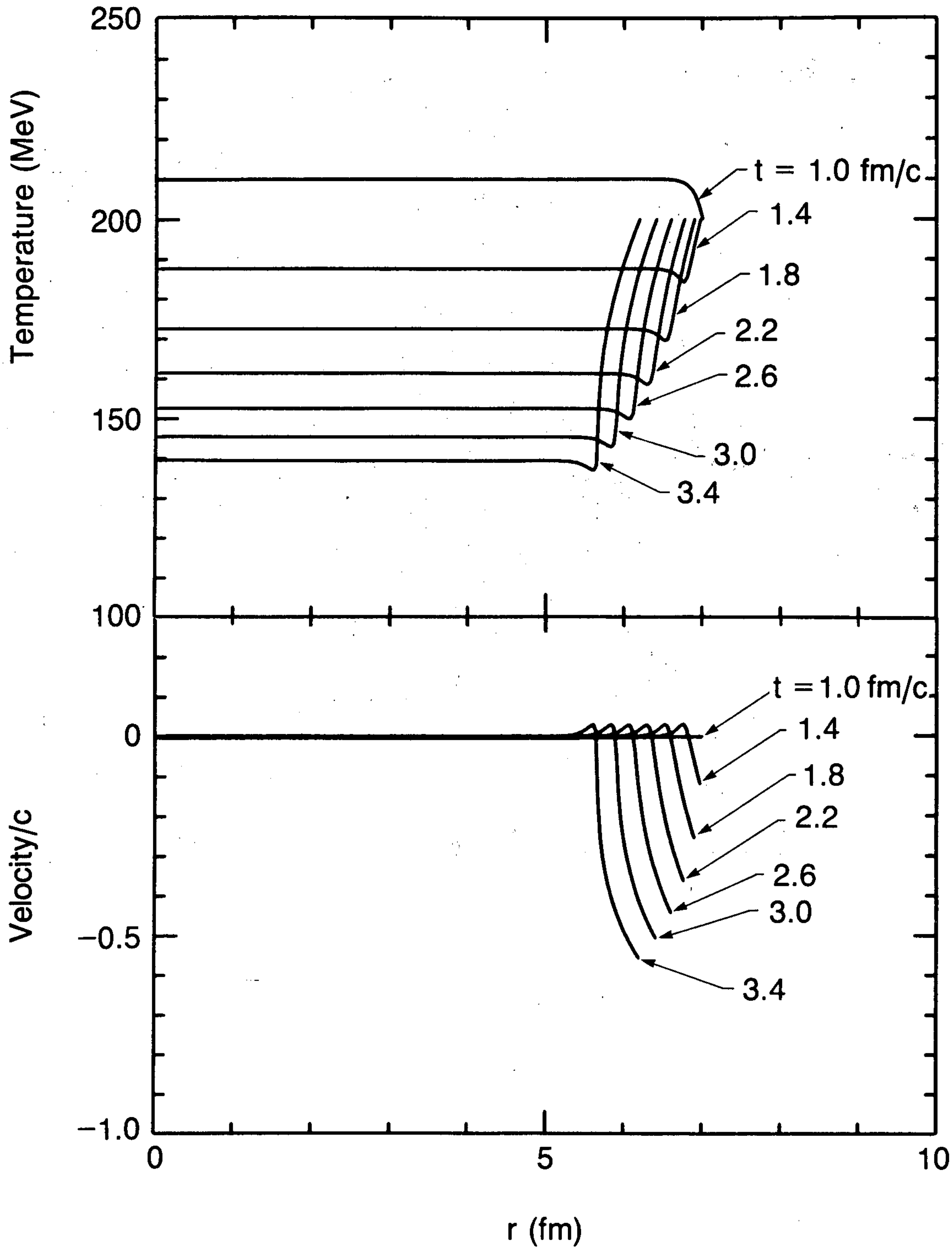

Fig. 2

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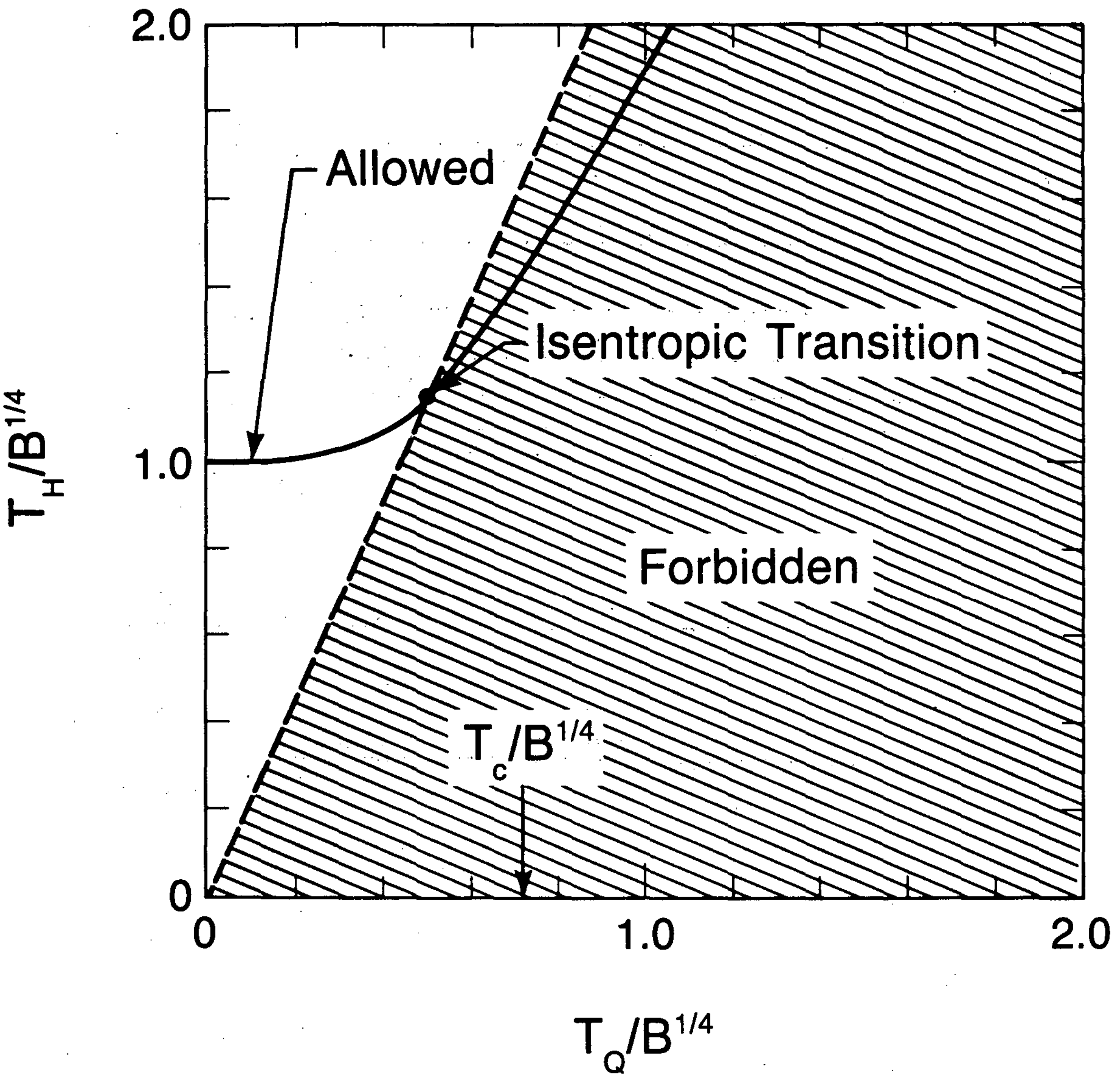

Fig. 3

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