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Abstract

The chromoelectric flux tube model is used to obtain a dynamical description of the evaporation of mesons from a quark-gluon plasma. The radiation pressure is computed to assess whether this process is an important mode for the disassembly of a compressed plasma. A new result for the creation rate of $q\bar{q}$ pairs in a constant color field is employed.

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Lattice gauge solutions of QCD provide fairly convincing theoretical support for the notion that hadrons will dissolve into a quark-gluon plasma at a sufficiently high energy density.[1] If such a plasma is formed in a high-energy hadron-hadron or nucleus-nucleus collision, the high internal pressure of the plasma will cause its rapid disassembly. Therefore, evidence for its prior existence must be found in its decay products. To assess whether signals of the plasma survive the evolution and what they are, a dynamical description of the disassembly is needed. One facet has been studied by Bjorken, by Kajantie et al., and by Baym et al., who propose a hydrodynamical model for its expansion.[2] As it expands it will cool, and the conditions for condensation back to the hadronic phase will be attained. In this note we address another facet of the disassembly, the formation and radiation of mesons at the surface of the hot plasma. The pre-freezeout radiation, which has also been emphasized in connection with the disassembly of a hadronic fireball,[3] plays two roles. It carries signals of the state of the plasma as it evolves. Second, it is coupled to the hydrodynamical expansion of the plasma, which may be inhibited by the backward radiation pressure. In the extreme, the plasma may evaporate mesons rather than expand collectively as a plasma. Such an extreme scenario was proposed recently by Danos and Rafelsky.[4] However, we shall see that their two-parameter model of evaporation is incompatible with the dynamics of a confining mechanism.

QCD is essentially different from all other theories of particle substructure because of the apparent permanent confinement of color. The conjectured difference between QED and QCD is depicted in Fig. 1. Two oppositely colored charges in QCD are connected by a confined tube of color flux. As a consequence, at distances for which the field lines are parallel, the force acting between them remains constant however great the distance. We adopt this picture quite literally. It provides an explicit mechanism for the formation of mesons at the plasma surface and explicit formula for their production rate and momentum distribution, as we shall show.

The plasma of almost free quarks, antiquarks, and gluons is viewed as a region of perturbative QCD vacuum embedded in the true nonperturbative vacuum that pervades most of space. Because of the thermal motion, a quark or antiquark may cross the boundary between the two vacua, but it will suffer a strong color interaction with the rest of the plasma (see Fig. 2) and cannot escape alone. However, there are two ways in which such a quark can initiate the evaporation of a meson. As the quark crosses the boundary, a flux tube of chromoelectric field connects it to the plasma, and it will experience an attractive potential that will pull it back unless the tube fissions. The tube can fission through the quantum tunneling to a real state of a virtual $q\bar{q}$ pair in the uniform color field of the tube. The energy and momentum of the meson so formed is then determined by the initial energy and momentum of the leading quark and the space-time position where the $q\bar{q}$ pair creation occurs. This mechanism of hadronization has been commonly used to describe particle production in high-energy e^+e^- annihilation.[5,6] Although in our problem the average momentum of the quarks is not so high ($T \sim 200$ MeV), we expect that the picture will qualitatively describe the fate of the high momentum component of the thermal distribution. The contribution of the low-momentum quarks to this hadronization process is suppressed as will be seen. This supports our idealization of a smooth plasma surface. A second mechanism for hadronization at the plasma surface is the coalescence of a quark-antiquark pair. We shall not further describe this process here, since at the expected plasma temperatures the thermal flux of quarks and antiquarks with specified color charge is small.

Consider a quark (or antiquark) of momentum $\underset{\sim}{k}_0$ that is outward directed with respect to the surface. As the quark passes through the surface, a tube of chromoelectric flux is built up behind it, out of its kinetic energy. The energy per unit length that is stored in the tube is $\sigma = \varepsilon^2 A/2$, where ε and A are the field strength and cross section of the tube. Gauss' law relates the flux εA to the quark charge, g/2 through $\varepsilon A = g/2$, yielding

$$\sigma = g\varepsilon/4 \tag{1}$$

This string constant can also be related to the Regge slope and so is essentially a known parameter, $\sigma = 0.177\ (\mathrm{GeV})^2$. We assume that the flux tube that shields its color will connect it to the plasma by the shortest path. The motion of the quark will be governed by the equations expressing the conservation of energy and of momentum parallel to the surface. The string dynamics can be solved analytically both for finite and zero quark mass. Its velocity parallel to the surface remains constant, and its motion perpendicular is illustrated in Fig. 3. It moves a distance E_{z0}/σ before being stopped at time (k_0/σ). Thereafter it is accelerated back into the plasma, unless the string fissions. The color flux tube can fission as the result of the tunneling of a virtual quark-antiquark pair to a real state inside the tube. The field connecting the leading quark to the plasma is then identically cancelled by that connecting the pair. If such a pair is created, say at a distance z' from the surface of the plasma and at a time t, a meson consisting of the original quark together with the antiquark of the created pair is thus formed. Its momentum perpendicular to the surface will be that possessed by the leading quark at the time t (see Fig. 3). Its energy will be E_0 less the energy carried back into the plasma by the fragment of string of length z' and the quark contained in it. Thus, the meson momentum and energy are given by the string dynamics (Fig. 3) and the space-time point where pair creation occurs (Fig. 2). For the detailed solution see Ref. 7.

The creation of a $q\bar{q}$ pair in the constant color field is similar to the QED process solved by Schwinger[8] more than 30 years ago. His result has been directly transcribed into QCD by several authors who describe hadron production in high-energy e^+e^- annihilation.[5,6]

However, Schwinger's calculation of the production rate of pairs neglects their mutual interaction and their mass in comparison with the energy in the macroscopic external field. These are essential difficulties in QED and the problem has never been solved exactly. While Schwinger's approximations are reasonable in QED, they are very serious in QCD where the field between the created pair is identical to the original field, and the field energy absorbed by the pair in tunneling to a real state is of the same order as the total field energy of the flux tube. We have therefore derived a new result for the rate per unit volume at which $q\bar{q}$ pairs will be created in a color field ε, which takes the mutual interaction and strict energy conservation into account. It is very interesting that a problem that has not been solved in 30 years in QED can be solved very easily in QCD in the model of confined color indicated in Fig. 1. The new result is[9]

$$p = \sum_{\text{flavor}} \frac{(g\varepsilon)^2}{64\pi^3} \sum_{n=1}^{\infty} \frac{1}{n^2} \exp\left(-\frac{4\pi^2 m_f^2 n}{g\varepsilon}\right) \xrightarrow[m_f=0]{} \frac{\sigma^2}{12\pi} \tag{2}$$

Having calculated the string dynamics (Fig. 3) and the rate per unit volume that a real pair will be formed and that hence the string will fission, we can now calculate the probability that the string will fission at the time t measured from the time that the outward-moving quark crosses the surface and in the interval dt. This is given by

$$dP(t) = pAz(t)\,dt\,[1 - P(t)] \tag{3}$$

where z(t) is the coordinate of the leading quark at time t and hence the length of the string. Note from Fig. 3 that the scale for both time and distance is controlled by the string tension, σ. Hence according to (3) the parameter that characterizes confinement is, where $\alpha_c = g^2/4\pi$,

$$k_c = \sqrt{\frac{\sigma^2}{pA}} = \sqrt{\frac{24\sigma}{\alpha_c}} \tag{4}$$

Equation (3) can be integrated analytically for both massive or massless quarks. For pedagogical reasons we exhibit the simpler massless limit in the form of the probability distribution in momentum k_z of the meson produced by a quark of initial momentum k_{zo}. The formula is remarkably simple and instructive

$$\frac{dP}{dk_z} = \frac{k_o}{k_{zo}} \frac{|k_z|-k_{zo}}{k_c^2} \exp\left[-\frac{1}{2}\frac{k_o}{k_{zo}}\left(\frac{k_z-k_{zo}}{k_c}\right)^2\right] \tag{5}$$

which shows that a quark emerging from the plasma with k_{zo} suffers a most probable loss of momentum k_c at the time of fission,

$$\hat{k}_z = k_{zo} - k_c \tag{6}$$

(see Fig. 4). The integrated probability that a quark of initial momentum k_{zo} will escape in a meson is

$$P_{esc} = 1 - e^{-k_{zo}k_o/2k_c^2} \tag{7}$$

The momentum scale k_c, set by confinement depends on the values of σ and α_c. These parameters must be related through $\mathcal{L}_{QCD}$, but this relationship is unknown. For $\sigma = 0.177\ (\text{GeV})^2$ and for $\alpha_c = 2$ or 0.55, corresponding respectively to the values used by Casher et al,[5] and by the MIT group,[10] we have

$$k_c = \begin{cases} 1.45\ \text{GeV}, & (\alpha_c = 2) \\ 2.78\ \text{GeV}, & (\alpha_c = 0.55) \end{cases}$$

From (7) only quarks from the tail of the thermal distribution with $k \gtrsim k_c$ have an appreciable chance of forming a meson, and the momentum of the meson is shifted downward by k_c from the original quark momentum. Next we calculate the probability that the string will break at the point z' at the time t while the quark is outward moving. This probability is distributed uniformly along the string. Therefore,

$$d^2P(z',t) = pA\,dz'\,dt\,(1 - P(t)) \tag{8}$$

We fold d^2P with the flux of quarks of momentum k_o and sum over all such momenta to obtain the number of mesons radiated per unit time per unit surface area of the plasma having normal momentum k_z^M and energy E^M,

$$\frac{d^5N}{dS\ dt\ dk_z^M\ dE^M} = \frac{\gamma}{(2\pi)^3} \int d^3k_o\ e^{-E_o/T} \frac{k_{zo}}{E_o} \frac{pA}{\sigma^2}\left[1 - P(k_z^M)\right]\theta(E_z - m - (E_{zo}/E_o)E^M)$$
$$\times\ \theta(m + (E_{zo}/E_o)E^M - E_{zo})\ \theta(-k_z^M)\ \theta(k_z^M - k_{zo}) \qquad (9)$$

where θ is the step function $\theta(+) = 0$, $\theta(-) = 1$. The degeneracy factor γ is $\gamma = \gamma_C \times \gamma_S \times \gamma_F \times 2 = 24$ where 2 counts quarks and antiquarks. $P(k_M^z)$ means that it is evaluated at the time that yields the momentum k_z^M.

The radiation pressure exerted on the plasma by the radiation of mesons, and the energy flow per unit surface area per unit time carried in the radiation follow immediately:

$$P^M = \frac{d^3k}{dS\ dt} = \int_0 dk_z\ k_z \int_{E_\pi} dE^M \frac{d^5N}{dS\ dt\ dk_z\ dE^M} \qquad (10)$$

$$\frac{d^3E}{dS\ dt} = \int_0 dk_z \int_{E_\pi} dE^M\ E^M \frac{d^5N}{dS\ dt\ dk_z\ dE^M} \qquad (11)$$

where $E_\pi = \sqrt{k_z^2 + m_\pi^2}$.

Figure 5 shows, as a function of temperature, the ratio of radiation pressure and internal quark pressure. Since this ratio is less than 20% up to T = 500 MeV, we conclude that the radiation due to this process will be a minor perturbation on the collective hydrodynamical disassembly of the plasma.

This conclusion contradicts an assertion made by Danos and Rafelski.[4] Their conclusion, however, is based on a model that does not contain a dynamical description of color confinement but that instead is characterized by two adjustable parameters, a minimum momentum for emission and an efficiency factor defining the fraction of the quark energy carried off by the meson. Figure 3 contradicts the basis of their model, since they assert that a quark of initial momentum k_0 ($>p_M$) produces a meson of the same momentum.

The flux of mesons that are formed by the fission of the string is given by (9). Many of these, in their rest frame, are more massive than the pion. If the invariant mass M lies in the interval $n < M/m_\pi < n + 1$ it is assumed that the heavy meson will decay into n pions according to phase space distributions. These distributions can be calculated for up to three pion decay. For four or more pion decay modes, a high-energy scaling formula that agrees with the experimentally observed distributions[11] in $e^+e^- \rightarrow q\bar{q} \rightarrow n\pi$ annihilation is used with the irrelevant electromagnetic factors divided out. In terms of the above number distribution in rapidity, dn/dy', as seen in the rest frame of the heavy meson, the pion flux, rapidity distribution, dF/dy, in the rest frame of the plasma surface can be written

$$\frac{dF}{dy} = \int dy_M\ dM_M \frac{d^2F_M}{dy_M dM_M}\ dy'\ \frac{dn}{dy'}\ \delta(y - y_M - y')$$

where the flux distribution F_M of heavy mesons follows from (9).

In Fig. 6 this distribution is compared at a plasma temperature of 300 MeV with several other quantities. The flux of quarks and antiquarks at the surface is very large but relatively few of them lead to fission events. For those that do, the distribution of heavy mesons is shown (broken strings). Their decay as described above is shown by the pion curve. The number of emitted pions is, for the coupling constant $\alpha_c = 2$, much larger than the number in black-body radiation from a pion gas at the same temperatuare as the plasma. The peak in the momentum is also shifted down in rapidity as compared to the most probable rapidity of the quarks or of the pion gas. Because of the high density of radiated pions, we assume that they will form a thermalized expanding halo outside the plasma.

However, the results of Fig. 6 depend sensitively on the value of α_c. For the MIT value, deduced from hadron spectroscopy, the number of emitted pions is much smaller. Our prejudice is that the coupling appropriate to the hadronization process should be larger than the MIT value, as, for example, the one used in Ref. 5 and in Fig. 6. However, a resolution of this point requires a deeper understanding than we presently possess.

In summary, the hadronization at the surface of a quark-gluon plasma has been studied in the framework of a chromoelectric flux tube model. We find that the radiation pressure is sufficiently small compared to the internal pressure that it can be ignored to first approximation in the hydrodynamical expansion of the plasma. Thus our solution for meson radiation can be folded with a solution to the hydrodynamical expansion to obtain the spectrum of radiated mesons emitted over the history of the expansion. Of special interest is the distinction between strange and non-strange mesons, which is explicit in the theory through the dependences on the quark masses and the thermal populations in the plasma.

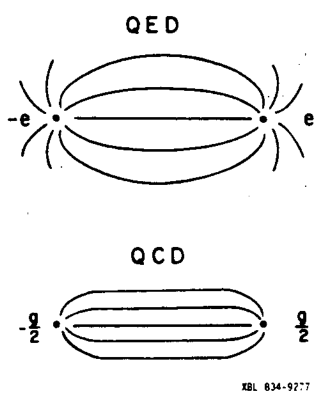

Fig. 1. Depicts field lines in QED and QCD.

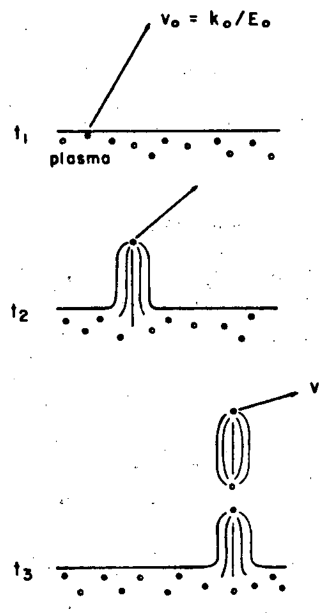

Fig. 2. A quark of velocity v crossing the surface is decelerated in the normal direction. At t_3 a $q\bar{q}$ pair is formed, which breaks the string. A meson of velocity v thereby emerges.

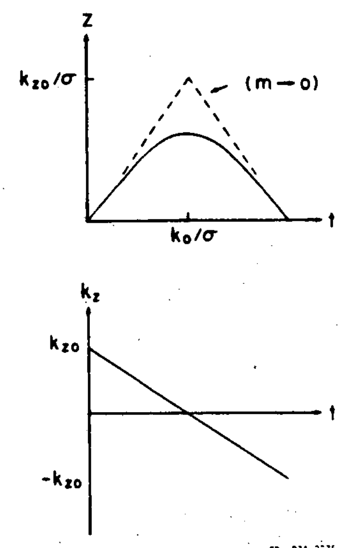

Fig. 3. The solution of the string dynamics showing quark position and momentum normal to the plasma surface as a function of time.

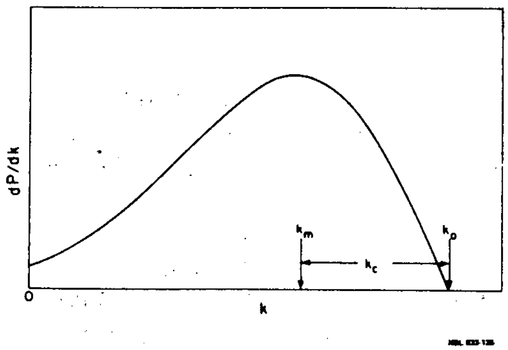

Fig. 4. Momentum distribution of radiated mesons for initial quark momentum $k_{zo} > k_c$.

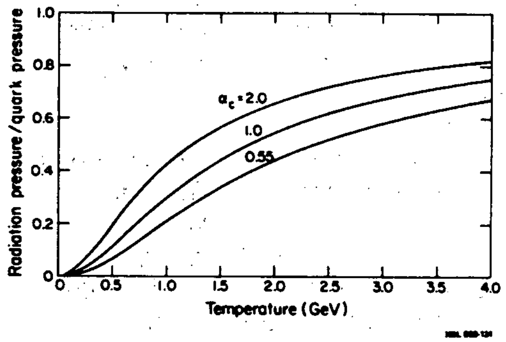

Fig. 5. Pressure acting on the surface of the quark-gluon plasma due to the radiation of mesons as a function of plasma temperature.

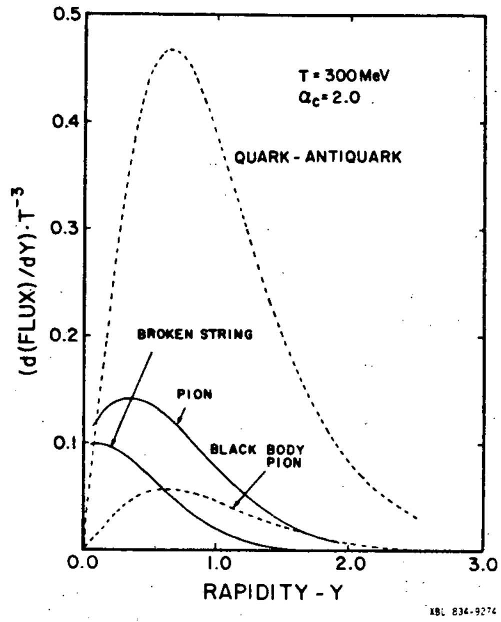

Fig. 6. Rapidity distributions as described in text.

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