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Signals for hadronic modes in the quark plasma

Chu-Xin Chen and Carleton DeTar
Department of Physics, University of Utah, Salt Lake City, Utah 84112

Thomas DeGrand
Department of Physics, University of Colorado, Boulder, Colorado 80309
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We give a new argument for the occurrence of hadronic modes in the high-temperature quark plasma, present a crude phenomenological model of the plasma, and discuss possible experimental signatures for the low-lying modes.

I. INTRODUCTION

At high temperatures and zero baryon-number density, hadronic matter is expected to undergo a phase transition to a "quark-plasma" phase. The theoretically expected characteristics of this phase have only recently become apparent, although existence of this phase has been known for some time.[1] It has been traditional to assume that the phase is well described as a gas of relatively weakly interacting quarks and gluons. This "naive deconfinement" picture has been under attack for some time now, but recent evidence, which we outline, suggests that the simple picture is fundamentally wrong.

We focus on the low-lying modes of excitation of the plasma. They are associated with poles in frequency-momentum-space linear response functions of the type

$$S_{AB}(\mathbf{k},\omega)=\int d\mathbf{k}\,d\omega\, e^{(i\mathbf{k}\cdot\mathbf{x}-i\omega t)}\langle T[A(\mathbf{x},t)B(0,0)]\rangle \ , \quad (1)$$

where $A(\mathbf{x},t)$ and $B(\mathbf{x},t)$ are local operators and the average is taken over the Gibbs ensemble at temperature T. At zero temperature these modes are the familiar hadrons. At finite temperature they correspond to peaks in certain cross sections.

We present a new argument and cite recent evidence in support of the proposition that the plasma modes are vestiges of the familiar hadrons, introduce a crude model of the plasma, and discuss the prospects of detecting hadronic-plasma modes in heavy-ion collisions. Whether charmonium states survive the phase transition is a quantitative question of some recent controversy.[2] We are concerned here with lighter hadronic modes.

II. EVIDENCE FOR HADRONIC MODES

The phase boundary between the low- and high-temperature phases provides indirect information about the nature of the plasma excitations. In QCD several Lagrangian parameters can be varied to produce a multidimensional phase diagram. These parameters include the masses of the various quarks, the chemical potentials for each of the flavors, including baryon number and temperature. The key issue to be resolved is whether there is a break or gap anywhere in the phase boundary that permits a smooth passage between the phases. If so, there is a rigorous one-to-one correspondence between the low-temperature and high-temperature modes of excitation. That is, the poles in the linear response functions (1) can be tracked smoothly through the phase diagram from one phase to the other. Therefore, in particular, if there are no quark and gluon modes at low temperatures, there must be none at high temperatures as well.

It is sufficient to find any gap whatsoever in the phase boundary, even if the gap occurs at an unnatural value of the Lagrangian parameters. Of course, the rigorous argument does not rule out the possible practicality of a parton-model description at any temperature nor does it say anything about the stability of high-temperature modes. We argue below that some of them may indeed be sufficiently long lived at moderate temperatures to be observable.

What is known so far about the m vs T phase diagram, where m is a flavor-symmetric quark mass? For SU(2) color a gap almost certainly occurs, since the high-temperature phase transition is probably second order and so disappears when quarks of any large mass are introduced.[3] For SU(3) color several recent simulations on small lattices find a gap at small quark mass.[4,5] Even if the gap is found to be filled in as the lattice volume is increased, model studies suggest that it can be opened by reducing the number of flavors and increasing the flavor chemical potentials.[6]

The second indirect approach has been to study the behavior of the imaginary-time linear response functions at large spatial distances.[7-9] The imaginary-time functions (static susceptibilities) are analytically related to the real-time functions (1). In this case the poles are followed to zero frequency and imaginary wave number. The numerical simulations provide evidence for cleanly separated poles in color-singlet channels in imaginary momentum.[8,9] At temperatures about 25% above the critical temperature the degenerate pion and σ-meson modes are found to have screening masses of about $4T$ and the ρ-meson mode, about $7T$.[9]

Although there is a rigorous one-to-one correspondence between screening modes and real modes, the uniqueness of the equilibrium ensemble requires that all of the real-time modes be unstable to some degree, and the widths are not known from the simulations. We resort to models to illuminate this question. We shall assume here that for each mode the screening mass from the simulation is approximately equal to the gap frequency.

III. THREE-COMPONENT MODEL OF THE QUARK PLASMA

For temperatures above the phase transition we propose a quark-matter version of the three-component plasma model of Ref. 10. The three components correspond to a crude division of distance scales of the plasma: at the longest distances, there are the hydrodynamic modes, namely, the phonon and thermal diffusion modes; at intermediate distances there are hadronic modes; at the shortest distances, we use quark and gluon degrees of freedom. Each of the components is treated as a noninteracting constituent in first approximation. This treatment is certainly incorrect in the vicinity of the phase transition where interactions produce the chiral condensate. The distance scales are enforced by introducing momentum cutoffs. Phonons have momentum $k < k_{\rm MF}$, a scale set by the mean free path. Hadronic constituents have momentum $k < k_H$, a scale set by their momentum-dependent lifetimes, as discussed below. The quark and gluon component is restricted to high momenta $k > k_M$, the "magnetic confinement" scale.

It is natural to introduce the lowest-lying hadronic modes explicitly. Judging from the screening masses,[9] they are likely to be the π, K, σ, ρ, and a_1 mesons. The quark and gluon component stands for all of the higher modes. They are not true plasma modes.

To obtain a rough estimate of the energy density, we put $k_{\rm MF} = k_H = k_M = \mu T$ for constant μ. Let the phonon dispersion relation be given by $\omega_{\rm ph} = v_s k$, the π, σ, and K dispersion relations by $\omega_s^2 = k^2 + (\mu_s T)^2$, the ρ and a_1 dispersion relations by $\omega_v^2 = k^2 + (\mu_v T)^2$, and the quark and gluon dispersion relations by $\omega_q = \omega_g = k$. The energy density is

$$\epsilon = B + \frac{T^4}{2\pi^2}(u_{\rm ph} + 8u_s + 18u_v + 16u_g + 36u_q) \ , \quad (2)$$

where the component energy densities are

$$u_{\rm ph}(\mu) = \int_0^\mu \frac{p^2 dp\, v_s p}{e^{v_s p} - 1} \ ,$$

$$u_h(\mu) = \int_0^\mu \frac{p^2 p d(p^2 + \mu_h^2)^{1/2}}{e^{(p^2+\mu_h^2)^{1/2}} - 1} \ ,$$

$$u_g(\mu) = \int_\mu^\infty \frac{p^3 dp}{e^p - 1}, \quad u_q(\mu) = \int_\mu^\infty \frac{p^3 dp}{e^p + 1} \ .$$

Here h stands for the label s or v for the scalar- and vector-meson modes. The coefficients 18 and 36 are appropriate for a color octet of gluons and three flavors of quarks. The other coefficients count the hadronic spin and flavor degrees of freedom. We have included a "latent" heat B to represent the energy required to liberate the chiral condensate.

To estimate the energy density in our model, we arbitrarily set $\mu_s = \mu_v = \mu$ and take $v_s = \sqrt{1/3}$. Then the components make the contributions shown in the figure to the total energy density. It is evident that for $\mu \lesssim 3$ the contributions are within 30% of the Stefan-Boltzmann value, consistent with numerical simulations,[4] even without the help of the latent heat B. We suggest that such a scale for strong interactions is warranted by the rather large screening mass ($\approx 7T$) obtained for the ρ-mesonic mode.[9]

Next, we turn to an estimate of the widths of the real-time modes. These widths arise from collisional and decay processes. Consider the ρ-meson mode. Judging from the temperature dependence of the screening masses,[9] we expect the $\pi\pi$ decay channel to close slightly above the phase transition and at higher temperatures only collisional processes remain. The collisional width of the ρ meson arises from the generic process

$$\rho + \phi \rightarrow \phi_1 + \phi_2 + \cdots \phi_m \ , \quad (3)$$

where ϕ is any constituent. Its contribution to the width of a zero-momentum ρ meson is[11]

$$\Gamma = \int \frac{d^3p}{(2\pi)^3} v_{\rm rel} n_\phi(p)(1 \pm n_{\phi 1}) \cdots (1 \pm n_{\phi m}) d\sigma \ , \quad (4)$$

where $v_{\rm rel}$ is the relative velocity between the constituent ϕ and the ρ meson, p is the momentum of the incident particle, $d\sigma$ is the differential cross section of the process, and n is the usual distribution function $n(E) = [\exp(E/T) \pm 1]^{-1}$. We are ignoring processes that produce the ρ meson. They decrease the width. If the momentum cutoff and masses of the hadronic constituents are much larger than the temperature, then $1 \pm n \simeq 1$. Summing over all final states gives

$$\Gamma = \int \frac{d^3p}{(2\pi)^3} v_{\rm rel} n(p) \sigma(p,T) \ . \quad (5)$$

For want of a better estimate take $\sigma(p,T) = \sigma_0/T^2$. Then,

$$\Gamma \simeq \frac{\sigma_0 T}{2\pi^2} \int_A^B v_{\rm rel} n(p) p^2 dp \ , \quad (6)$$

where the limits of integration are $(A,B) = (0, 2\mu)$ for hadronic constituents, and $(A,B) = (2\mu, \infty)$ for quarks and gluons. Using the same choices for cutoffs and mass as in the estimate of the energy density above, and summing over all constituents of the plasma, we get the total collisional width shown in Fig. 1. For $\mu = 3$ and $\sigma_0 = 1$, corresponding to 10 mb at $T \approx 200$ MeV, we have $\Gamma(T)$

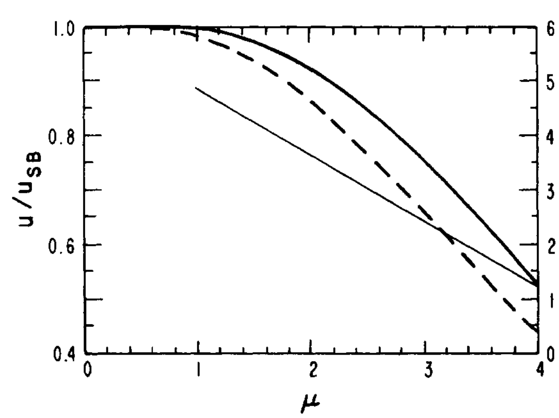

FIG. 1. Left scale: contribution to the energy density (2) as a function of cutoff parameter μ normalized to the Stefan-Boltzmann value. Dashed line: quark and gluon contribution; heavy line: quark, gluon, phonon, and hadron contributions. Right scale and thin line: hadronic width from all components from Eq. (6).

$\simeq 2.4T$, smaller than the measured screening masses. Similar results are expected for the other hadronic modes. Since with these plausible assumptions the widths of the plasma modes are smaller than the gap frequencies, our treatment of the hadronic modes as discrete plasma excitations is justifiable *a posteriori.*

Typical hadronic cross sections rise rapidly as resonances are excited. Therefore, as the momentum of the excitation increases relative to the medium, the widths may also increase so as to overtake the frequency of excitation. An ideal-gas treatment would then be inappropriate.

IV. EXPERIMENTAL SIGNATURES

We turn now to the prospects of observing some of the plasma modes in heavy-ion collisions. The model proposed in the previous section assumed an equilibrium ensemble. This condition is certainly not met on the large scale in heavy-ion collisions, but it has been customary to assume that local thermal equilibrium is achievable. Thus the longest-range excitations, the phonon and other hydrodynamic modes, must propagate in a distorted medium due to longitudinal expansion and rapid cooling, but the shorter-range hadronic modes and the partons may enjoy local equilibrium and survive to be observable.

To observe the excitations requires a nonhadronic probe. Perhaps the best prospects are found in the charged-dilepton channels, which probe the vector mesons. The cross section as a function of the invariant mass of the lepton pair may reveal a distortion in the ρ-meson peak. Two factors contribute to an observable distortion.

(1) A thermal shift in the gap frequency $\omega_\rho(0,T)$ due to formation of the ρ-a_1 chiral multiplet and changes in the confinement scale.[12]

(2) A moderately small width as discussed in the previous section.

Two factors obscure the observable distortion.

(1) Since the produced lepton pair may come from a wide variety of environments at any of a range of temperatures, the observed peak position will be smeared out, according to the variation of the gap frequency with temperature and the thermal histories of the events. Some other event-selection criterion, e.g., high multiplicity, may be helpful to reduce the contribution from events that do not ignite the plasma.

(2) Since the dispersion relation for the hadronic modes is not Minkowskian, there will be Doppler broadening of the peaks due to the motion of the excitation relative to a local plasma rest frame.

Because of theoretical uncertainties in knowing the plasma-mode cross sections, the excitation frequencies and their temperature dependence, and, to some extent, the thermal histories of plasma events, it is not possible at present to predict definitely whether a clear signal will survive. Only experiment will tell.

The two-photon channel offers but faint hope of observing the shifts in the π^0, σ, and η mode frequencies.[13] When the photon polarizations are not distinguished, all three peaks should be present. Compounding the usual combinatoric difficulties associated with pairing of the observed photons, the strong temperature dependence of the pion and σ screening masses[9] suggests that the peaks will be broadly smeared. However, the rapid crossover of the pion gap frequency at the critical temperature and the radical change in the σ mode might introduce some structure in the invariant-mass spectrum just above the low-temperature pion peak. Precisely what happens depends on how the plasma evolves at the crossover temperature. More detailed model calculations could help with this question. Collisional broadening will obscure the effect. Again, whether these modes are in fact observable must be left for experiment to decide.

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