Charmonium states in the quark-gluon plasma

T. H. Hansson, Su H. Lee, and I. Zahed
Physics Department, State University of New York at Stony Brook, Stony Brook, New York 11794
(Received 11 January 1988)

The possibility that $\bar{c}c$ states remain bound even in the QCD plasma state is examined. Assuming that these states do exist, we calculate the thermal width due to the process $J/\psi\to\bar{D}D$ and show that low-lying states might be detected. The relevance of this to heavy-ion experiments is discussed.

At sufficiently high temperature or density, hadronic matter is believed to undergo a phase transition to a quark-gluon plasma[1] (QGP). While the nature of this transition is not well understood, its existence has been demonstrated by numerical simulations and the critical parameters are approximately known. Lattice calculations suggest $T_c\sim200$ MeV and $\rho\sim5-10\rho_0$, where ρ_0 is nuclear-matter density. It is an exciting possibility that the energy densities (~0.80 GeV fm^{-3}) believed to be needed to produce such a QGP can be reached by present and planned heavy-ion machines.[2] Experiments in this direction are at present being conducted at CERN SPS using oxygen on a fixed lead target at 200 GeV/nucleon. The planned relativistic heavy-ion collider (RHIC) at BNL will deliver ions as heavy as gold with energies up to 100 GeV/nucleon in the center-of-mass frame.

Various signatures have been proposed as a signal of QGP formation. Examples are strangeness enhancement, large dilepton yields, and enhanced photon emission.[3] Unfortunately, most of these signatures rely considerably on details of the hadronic "chemistry" and the (experimentally unknown) parameters governing the hydrodynamical evolution. Given this, it is clearly important to try to find model-independent signatures for QGP formation.

Recently, Matsui and Satz proposed that the suppression of J/ψ will provide a model-independent signature for quark-gluon plasma formation.[4] They argue that due to static deconfinement no $\bar{c}c$ resonance can exist above T_c, leading to a significant decrease of the $\mu^+\mu^-$ yield in the resonance region. Their estimate of the nonresonant production of $\mu^+\mu^-$ pairs from the heat bath does not make up for the disappearance of J/ψ.

In this paper we will present a different scenario for the behavior of heavy-quark resonances at high temperature. Following DeTar's original conjecture,[5] we assume that the quark-gluon plasma is dynamically confined at the magnetic mass scale; i.e., there are no colored states propagating over distances larger than $m_{\rm mag}^{-1}\sim(g^2T)^{-1}$. This assumption seems to be supported by recent numerical simulations[5-7] and perturbative calculations.[8,9] In particular, we will assume that J/ψ will remain bound at temperatures above the critical temperature T_c.

To put our assumption in perspective let us examine the arguments of Matsui and Satz in more detail. Their conclusion that J/ψ becomes unbound above T_c is based on lattice Monte Carlo measurements of the static potential between two infinitely heavy quarks. This potential changes from being "string"-like and confining at small T to being screened, Coulombic and nonconfining at large T. For sufficiently heavy quarks, there are always Coulomb bound states, and Matsui and Satz argue that if the radius of the naive Coulomb bound state is larger than the static screening length, it will disappear. This is certainly a possibility, but is not compatible with the assumption of dynamical confinement. A second possibility is that the charm quark is not to be considered heavy at temperatures at or above T_c. This would not be too surprising given the fact that charmonium at $T=0$ is not Coulombic even though $m_c\gg\Lambda_{\rm QCD}$. In this scenario the charmed quark would still be confined by momentum-dependent forces which are likely to be magnetic in origin. This kind of confining force cannot be seen in present lattice calculations of static correlation functions. Needless to say, we cannot give any strong argument for or against either possibility at this point, and we consider our picture as an alternative to the one discussed by Matsui and Satz.

For clarity, we shall first consider the case of charmed quarks in a pure gluon plasma, and discuss the effects of light quarks later. Since gluons cannot combine with a quark to form a color-singlet, dynamical confinement in the case of no light quarks (q) simply means that the valence structure of propagating states containing heavy quarks (Q) is $\bar{Q}Q$, QQQ, . . . just as at zero temperature. The masses of the resonances can of course depend on T. (There is even no reason to believe that the level ordering at high T is the same as the one at zero temperature. For a recent attempt to calculate the temperature dependence of resonance masses, see Refs. 10 and 11.) Note that we are not just saying that hadrons (or more precisely hadronlike excitations) change masses when in a heat bath, but that they remain bound [with radii $\simeq(g^2T)^{-1}$] even at temperatures above T_c.

When light quarks are present, the situation is dramatically changed. Even if we were to assume dynamical confinement, there might still be no $\bar{Q}Q$ bound states whenever decay into color-singlet $(\bar{Q}q)(\bar{q}Q)$ pairs (here $D\bar{D}$) is possible. Because of the small total cross section for charm production recombination processes (here $D\bar{D}\to J/\psi$) can be neglected. This will certainly happen for all states above threshold. In fact, even states below threshold can become unstable because of the excitation process shown in Fig. 1. In this paper we demonstrate

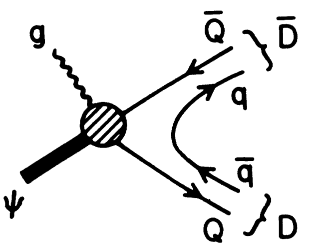

FIG. 1. Charmonium decays into a $D\bar{D}$ by a thermal gluon.

that this "gluoelectric effect" is not expected to make the low-lying $\bar{c}c$ states unstable. Thus, there are reasons to believe that charmonium resonances are present even in the high-temperature phase.

Because of the process shown in Fig. 1, the heavy-quark bound state acquires a temperature-dependent width. To estimate this width we need to know more about the bound-state wave function at temperatures that are above T_c. As already stressed, we assume that the QCD plasma still confines, but on a length scale $\sim(g^2T)^{-1}$ rather than $\sim\Lambda_{\rm QCD}^{-1}$. Clearly this qualitative statement is not sufficient for obtaining a wave function, and we need further assumptions. For a plasma temperature $T\sim250$ MeV and $g\sim2.5$ ($\alpha_s\sim0.5$), the electric screening length is $m_{\rm el}^{-1}=(gT)^{-1}\sim0.3$ fm to leading order. At this temperature, the lattice estimate is about 0.2–0.3 fm.[4] We shall assume that at temperatures relevant for heavy-ion experiments, the three length scales T^{-1}, $(gT)^{-1}$, and $(g^2T)^{-1}$ are not too far separated. Since $m_{\rm el}^{-1}\sim0.3$ fm, we therefore expect a confinement radius ≤1 fm inside the plasma.

At this point we would like to stress that we do not know much about the nature of the dynamically confining force, and therefore we will only make an assumption about its range. As a matter of fact we do not even know whether it can be described by a local potential (a velocity-dependent potential originating from magnetic effects would be more likely). Nevertheless, we shall calculate our wave functions using a family of very simple local potentials. This procedure will be justified *a posteriori* by showing that our final result is very insensitive to the shape of the wave function, and consequently to the nature of the potential. With this in mind, we describe heavy-quark bound states *above* T_c with a confining potential of the form

$$V(r)=\sigma r+\frac{4\alpha_s}{3}\frac{e^{-m_{\rm el}r}}{r}\ , \tag{1}$$

where the parameters are chosen so that the rms radius is ≤1 fm. The Debye screened Coulomb potential is chosen to reproduce the known short-distance large-T behavior for very heavy quarks, and the linearly confining potential to make contact to a known potential at $T=0$. It is well known that a potential of this type gives a good charmonium spectrum and for reference, the $T=0$ values are $\sigma(0)\sim0.1$ GeV^{-2}, $\alpha_s\sim0.6$, and $m_{\rm el}=0$. Again we stress that the particular form of $V(r)$ is of no importance for our conclusions.

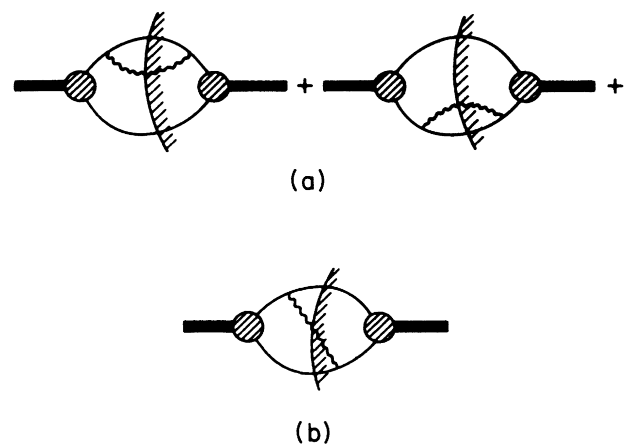

FIG. 2. First-order self-energy diagram of charmonium cut for the process of Fig. 1.

Given a potential and thus a bound-state spectrum, it is straightforward to calculate the rates by which the ground state is excited to higher-lying levels. In order for the process of Fig. 1 to take place, however, the excited state must lie above the $D\bar{D}$ threshold. Since neither the threshold energy nor the bound-state mass are known at finite temperature, we shall simply leave the difference $\Delta=2m_D-m_{j/\psi}$ as a free parameter, and assume that it will not change drastically from the zero-temperature value $\Delta\simeq0.69$ GeV. The sensitivity of our calculation to Δ will be discussed.

Assuming that lowest-order perturbation theory holds (no final-state interactions), we can deduce the

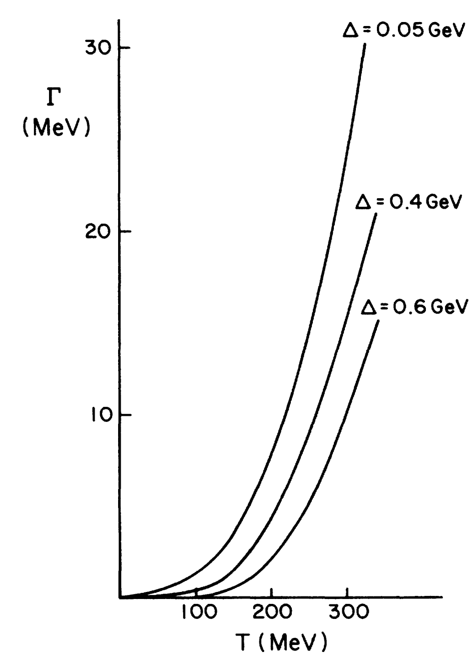

FIG. 3. Decay width as a function of temperature for $\Delta=0.05$ GeV, $\Delta=0.4$ GeV, and $\Delta=0.6$ GeV with charm-quark mass$=1.58$ GeV, $\frac{4}{3}\alpha=0.8$, and $\sigma=0.1$ GeV^{-2}.

temperature-dependent width of the bound state from the imaginary part of the diagrams shown in Fig. 2. Remember that in the frame of the heat bath, the energy of $\bar{Q}Q$ is above the $(\bar{Q}q)(\bar{q}Q)$ threshold. With this in mind, the bound-state wave function for J/ψ is

$$\psi(P)=\int \frac{dQ_0}{2\pi}\Delta\left[P+\frac{Q}{2}\right]\Gamma\left[P+\frac{Q}{2},P-\frac{Q}{2}\right]\times\Delta\left[P-\frac{Q}{2}\right], \quad (2)$$

where Δ is the standard quark propagator and Γ the usual covariant Bethe-Salpeter vertex:

$$\Gamma(p,-p')=i\int\frac{dk^4}{(2\pi)^4}V(k+p')\gamma^0\Delta(q+k)\Gamma\times(q+k,k)\Delta(k)\gamma^0 . \quad (3)$$

Here $2P=(p-p')$ and $Q=(p+p')$ are the relative and center-of-mass momenta of c and $\bar{c}$ in the bound state. Since the current-quark mass $m_c\sim1.58$ GeV is large, we can use the nonrelativistic version of (2),

$$\psi(p)=\tfrac{1}{2}(1+\gamma^0)\begin{bmatrix}0 & \Phi(p)\\ 0 & 0\end{bmatrix}\tfrac{1}{2}(1-\gamma^0) , \quad (4)$$

where $\Phi(p)$ is a triplet-spin state $(^3S_0)$ that satisfies the following Schrödinger equation:

$$\left[-E+\frac{p^2}{m}\right]\Phi(p)=\int\frac{d^3k}{(2\pi)^3}V(k-p)\Phi(k) . \quad (5)$$

Here V is the Fourier transform of the potential defined in (1). The discontinuity in the self-energy can be extracted using the generalized Cutkovsky's cutting rules in real time.[12] If we denote by $F(a)$ and $F(b)$ the contribution from cutting diagrams 2(a) and 2(b), respectively, then

$$F(a)=ig^2\int\int\frac{d^4p}{(2\pi)^4}\frac{d^4k}{(2\pi)^4}\mathrm{Tr}\left[\Delta(p)\tfrac{1}{2}(1+\gamma^0)\Phi(p)\Delta^-(p-q)\frac{\lambda^a}{2}\gamma^\mu\Delta(k-q)\tfrac{1}{2}(1-\gamma^0)\Phi^\dagger(k)\Delta^+(k)(p)\frac{\lambda^a}{2}\gamma^\mu\right]\times\left[-E+\frac{p^2}{m}\right]\left[-E+\frac{k^2}{m}\right]G^-(p-k)n_g , \quad (6)$$

$$F(b)=ig^2\int\int\frac{d^4p}{(2\pi)^4}\frac{d^4k}{(2\pi)^4}\mathrm{Tr}\left[\Delta(p)\tfrac{1}{2}(1+\gamma^0)\Phi(p)\Delta^-(p-q)\tfrac{1}{2}(1-\gamma^0)\Phi^\dagger(k)\Delta(p)\frac{\lambda^a}{2}\gamma^\mu\Delta^+(p-k)\frac{\lambda^a}{2}\gamma^\mu\right]\times\left[-E+\frac{p^2}{m}\right]^2G^-(p-k)n_g , \quad (7)$$

where

$$i\Delta^\pm(p)=2\pi(\not{p}+m)\Theta(\pm p_0)\delta(p^2-m^2) , \quad (8)$$

$$iG^\pm(p)=2\pi\Theta(\pm p_0)\delta(p^2) , \quad (9)$$

and n_g is the finite-temperature gluon distribution. The width of J/ψ is $\Gamma=2[F(a)+F(b)]$. Straightforward algebra gives

$$\Gamma(J+g\to c\bar{c})=\frac{4\alpha_s}{3\pi^2}\int\int dw_1dw_0n_gw_1w_2[\,|\Phi(p_1)|^2-\Phi(p_1)\Phi(p_2)] , \quad (10)$$

where $w_0=|p_1+p_2|$ is the gluon energy and $w_{1,2}=m+p_{1,2}^2/2m$ are the c and $\bar{c}$ kinetic energies.

In Fig. 3 we show the behavior of Γ vs T for different values of the binding energy $\Delta=2m_D-m_{J/\psi}$. For fixed values of T and Δ we can study the dependence on the shape of the wave function. Table I shows the variation in Γ with the size of the bound state for $\Delta=0.4$ GeV at $T=250$ MeV. Although smaller Δ increases the phase space (larger contribution from n_g) and larger sizes yield wider wave functions (peaked at low momentum), the destructive interference in (10) between diagrams 2(a) and 2(b) for $p_1\sim p_2$ will tend to cut down these two effects. The thermal width of J/ψ is in the range 5–10 MeV and is insensitive to rather large variations in the binding energy ($\Delta\simeq0.1$–0.6 GeV). We therefore conclude that thermally induced $D\bar{D}$ decays do not substantially suppress the presence of J/ψ in a dynamically confined quark-gluon plasma. This also implies that the possible

TABLE I. Decay width for different potentials of the type $V(r)=\sigma r+\frac{4}{3}\alpha_s/r$. At $T=250$ MeV, $\Delta=0.4$ GeV and charm-quark mass = 1.58 GeV.

$\frac{4}{3}\alpha$	σ (GeV^{-2})	Γ (MeV)	$\langle r^2\rangle^{1/2}$ (fm)
0.80	0.10	9.3	0.15
0.80	0.03	10.9	0.22
0.80	0	12.4	0.57
0.65	0	6.5	0.67
0.50	0	2.3	0.71

plasma formation in heavy-ion collisions does not necessarily exclude J/ψ production. This conclusion follows from the assumption that at high temperature the QGP is dynamically confined on length scales that are comparable to the magnetic mass scale. The sizable increase in the J/ψ width (0.06 MeV$\rightarrow$10 MeV) due to thermal effects, might be detectable in the $\mu^+\mu^-$ yield around the resonance region. This effect could be used as an implicit signal for QGP formation. The increase in the total width of J/ψ might also be accompanied by an increase in the $D\bar{D}$ yield, as suggested above. It is therefore interesting to investigate whether these two signals can be detected in present and future heavy-ion experiments.

This work was supported in part by the U.S. Department of Energy under Contract No. DE-AC02-76ER13001.

[1]J. Kuti, J. Polonyi, and K. Szlachanyi, Phys. Lett. **98B**, 199 (1981).

[2]H. Satz, in *Quark Matter '84*, proceedings of the Fourth International Conference on Ultrarelativistic Nucleus-Nucleus Collisions, Helsinki, Finland, 1985, edited by K. Kajantie (Springer, New York, 1985).

[3]J. Cleymans, R. V. Gavai, and E. Suhonen, Phys. Rep. **130**, 218 (1986).

[4]T. Matsui and H. Satz, Phys. Lett. B **178**, 416 (1986).

[5]C. DeTar, Phys. Rev. D **32**, 276 (1985).

[6]T. A. DeGrand and C. E. DeTar, in *Nuclear Chromodynamics: Quarks and Gluons in Particles and Nuclei*, proceedings of the ITP Workshop, Santa Barbara, California, 1985, edited by S. Brodsky and E. Moniz (World Scientific, Singapore, 1986).

[7]E. Manousakis and J. Polonyi, Phys. Rev. Lett. **58**, 847 (1987).

[8]T. H. Hansson and I. Zahed, Nucl. Phys. **B292**, 725 (1987).

[9]T. H. Hansson and I. Zahed, Phys. Rev. Lett. **58**, 2397 (1987).

[10]K. Hirose, T. Hashimoto, O. Miyamura, and T. Kanki, Phys. Rev. Lett. **57**, 2123 (1986).

[11]R. Vogt and A. Jackson, State University of New York at Stony Brook report (in preparation).

[12]R. L. Kobes and G. W. Semenoff, University of British Colombia report (unpublished).