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Dissociation kinetics and momentum-dependent J/ψ suppression in a quark-gluon plasma

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A quantum-statistical approach to two-particle states in a surrounding plasma is applied to heavy quarkonia. Thermal activation and dissociation of charmonium and bottomonium bound states is described by chemical reaction kinetics. The momentum dependence of the J/ψ suppression ratio is related to the temperature and the lifetime of the plasma phase. From recently observed data for high-energy nucleus-nucleus collisions, a plasma lifetime of the order of 20 fm/c is estimated.

Recent experiments with ultrarelativistic heavy-ion collisions of oxygen and sulfur at 200 GeV per nucleon on uranium and copper targets at CERN[1,2] show a decreasing ratio of J/ψ production, relative to the muon-pair continuum, when the energy density in the reaction increases. According to Matsui and Satz,[3] the suppression of charmonium production may serve as a possible signal of formation of a quark-gluon plasma. They proposed a scenario adopted from the concept of Debye screening in plasma physics where, due to the Mott effect, the bound states dissolve at high densities. In further works,[4–8] the observed momentum dependence of J/ψ suppression was interpreted by examining the formation time for the J/ψ in the quark-gluon plasma. However, alternate mechanisms, such as J/ψ absorption in a hadronic-resonance gas,[9] have also been discussed to find out whether suppression can be understood within more conventional physics (see also Refs. 10 and 11).

We will reexamine the suppression scenario of Matsui and Satz[3] from the point of view of a many-particle treatment of the quark-gluon plasma. Furthermore, we will show that the suppression of heavy quarkonia can be explained with the concepts of chemical reaction kinetics on a quark level.[12] In this way, our model allows also for a link with inelastic scattering in high-energy-density hadronic matter.[9]

Already in ordinary plasma physics, the simple concept of a statically screened potential gives no correct description of two-particle states in a dense plasma. In fact, as seen from a consequent quantum-statistical approach,[13] the continuum of scattering states is shifted downwards because of the effects of dynamic self-energy. Especially for Coulomb systems, the effects of dynamic screening and dynamic self-energy on the energy shift of bound states nearly compensate each other, and almost no shift with increasing plasma density occurs. This predicted behavior of the binding energy has been confirmed experimentally for the electron-hole plasma in strongly excited semiconductors.[14]

Similar to ordinary plasma physics, the application of the concept of a statically screened potential to the case of hadrons in a quark-gluon plasma environment leads to medium-dependent bound-state energies.[15–17] As an illustration, let us consider constituent quarks (masses in GeV: $m_u = m_d = 0.3$, $m_s = 0.5$, $m_c = 1.32$, and $m_b = 4.75$) interacting via an effective $q\bar{q}$ potential

$$V(r,\mu) = (\sigma/\mu)[1-\exp(-\mu r)] - (\alpha/r)\exp(-\mu r)\,, \quad (1)$$

where $\sigma = 0.192$ GeV^2 is the string tension and $\alpha = 0.471$ denotes the coupling constant. The screening parameter μ is supposed to be a function of the temperature.[15,16] The dependence of the bound-state energies on the screening parameter μ is shown in Fig. 1(a). However, the use of the static screening model gives incorrect results for the bound-state energy shifts already in ordinary plasma physics, so that this prediction of a mass shift for a quark-gluon plasma seems questionable. Two-particle bound-state energy shifts in a dense surrounding plasma should be derived from a systematic many-particle approach.

A quantum-statistical approach to a many-quark system has been worked out[18] introducing the concept of saturation of the quark-quark interaction within color-neutral clusters. The energy spectra of two-particle $c\bar{c}$ and $b\bar{b}$ states can immediately be obtained[12] and are shown in Fig. 1(b) as a function of the temperature. The bound-state energies are not shifted as long as there is no string flip, and Pauli blocking is not operative because of the low density of the heavy quarks. It is an advantage of the many-particle approach[18] to allow the formation of scattering states (quasifree quarks) although the two-particle interaction is of the confinement type. Supposing the usual Cornell form of the $q\bar{q}$ potential [$\mu = 0$ in Eq. (1)], the self-energy shift Δ^H is taken in the Hartree approximation according to Refs. 12, 18, and 19. A transition to the quark-gluon plasma where the light quarks move quasifreely is found at a temperature $T_c \approx 200$ MeV.[19] At this temperature, heavy-quark bound states (J/ψ and Υ) do exist, and their dissolution due to the Mott effect will take place at still higher values of temperature.[12]

The dissolution of heavy quarkonia also occurs below this Mott temperature because of thermal activation and

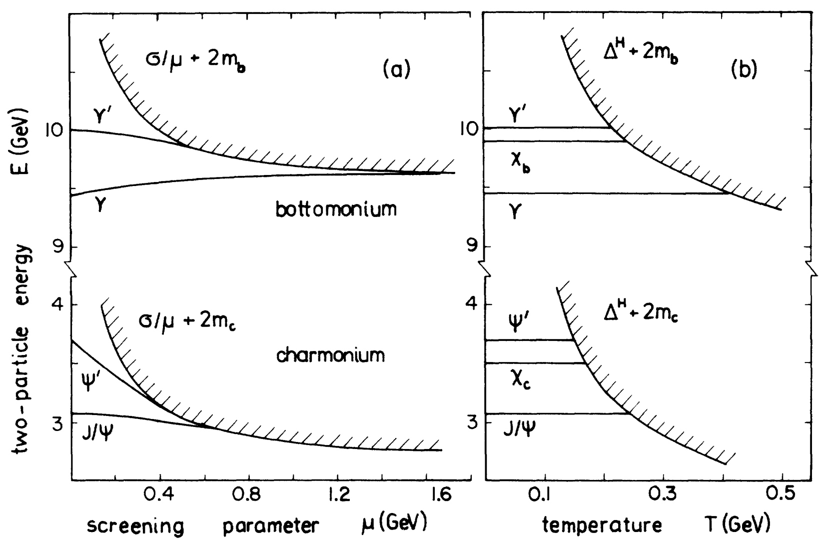

FIG. 1. Two-particle energies of heavy quarkonia ($b\bar{b}$ and $c\bar{c}$) in a quark-gluon plasma. (a) Static-screening model (Ref. 16). (b) Quark-potential model with color-saturated interaction (Ref. 12).

the formation of open-charm mesons. We suppose that after $c\bar{c}$ pairs have been formed due to hard collisions in the initial state of the reaction, they breakoff into light quarks of the surrounding plasma, due to string flips.

We apply reaction kinetics to estimate this suppression mechanism by calculating the reaction coefficient for the string-flip process from a $Q\bar{Q}$ bound state $(Q=c,b)$ to another quark q $(q=u,d,s)$ of the surrounding plasma of quasifree light quarks. This string flip may occur if a light quark is found at a distance smaller than the Q-$\bar{Q}$ distance of the bound-quarkonium state. For the string-flip cross sections we take the geometrical ones $\sigma_Q=\pi r^2_{Q\bar{Q}}$, with $r_{c\bar{c}}=0.45$ fm for J/ψ and $r_{b\bar{b}}=0.23$ fm for Υ. The frequency factor for the reactive collisions between bound heavy quarks and quasifree light ones is then given by

$$\tau_Q^{-1}=\sum_{\bar{q}}\int\frac{d^3p}{(2\pi\hbar)^3}\sigma^*(p)vf_{\bar{q}}(p)\,,$$

where v is the relative velocity, $f_{\bar{q}}(p)$ the quark-distribution function for $\bar{q}$, and the reactive cross section $\sigma^*(p)$ accounts only for those collisions that lead to a reaction process. As usual, we assume that for these reactive collisions the transferred kinetic energy exceeds a threshold value $\Delta E_Q(T)=E_{Q\bar{Q}}-2m_Q-\Delta^H(T)$ given by the dissociation energy of the bound meson into quasifree states. As is well known from chemical-reaction kinetics, the frequency factor of reactive collisions (in quasiclassical approximation) is given by

$$\tau_Q^{-1}=\pi r^2_{Q\bar{Q}}\sum_{\bar{q}}(8k_BT/\pi m_q)^{1/2}\tfrac{1}{3}n_{\bar{q}}(T)\times\exp[-\Delta E_Q(T)/k_BT]\,, \qquad (2)$$

where

$$n_{\bar{q}}(T)=\sum_p(\exp\{[(p^2c^2+m_q^2c^4)^{1/2}+\Delta^H]/k_BT\}+1)^{-1}$$

is the density of the corresponding light quark. Since the color is fixed (color neutrality of the clusters), we have to consider only two spin directions and the different flavors $(q=u,d,s)$. For example, using the values $r_{c\bar{c}}$, $r_{b\bar{b}}$ given above and taking the dissociation energy $\Delta E_Q(T)$ according to Fig. 1(b) (see also Ref. 12), we find for $T=180$ MeV: $\tau_c=14$ fm/c for J/ψ, and $\tau_b=208$ fm/c for Υ. The small value for the reactive collision frequency of b quarks results from the relative smaller value of $r_{b\bar{b}}$ as well as from the larger value of the dissociation energy.

Before applying this model to high-energy nucleus-nucleus collisions, one has to take into account the space-time evolution of the plasma. For small x_F, i.e., small values of the longitudinal momentum, the extension of the plasma $R(T_c,t)$ is characterized by the critical isotherm T_c according to[7] $R(T_c,t)=R_0[1-(t/t_m)^2]^{1/2}$. Here R_0 is taken as the radius of the projectile ($R_0=1.2\,A^{1/3}$ fm) and t_m is the plasma lifetime. Taking the r dependence of the initial heavy-quark density in the $z=0$ plane[7] as $\rho_Q(r)=\rho_Q^0(1-r^2/R_0^2)^{1/2}$, one obtains the suppression ratio (hadronic suppression in the mixed-phase regime being disregarded)

$$\mathcal{R}_Q(p_T)=\frac{\int_0^{R_0}r\,dr\int_0^{2\pi}d\theta\,\rho_Q(r)\exp(-t_Q/\tau_Q)}{\int_0^{R_0}2\pi r\rho_Q(r)dr}\,, \qquad (3)$$

where $t_Q(r,\theta;p_T)$ is the time of flight across the plasma for a heavy quarkonium (Q) with momentum p_T, which is emitted at distance r in the direction θ relative to the axis 0-r. A simple geometric analysis gives for the determination of t_Q the relation

$$R_0^2(1-t_Q^2/t_m^2)=r^2+\frac{p_T^2c^2}{m_Q^2c^2+p_T^2}t_Q^2+2r\cos\theta\frac{p_Tc}{(m_Q^2c^2+p_T^2)^{1/2}}t_Q\,.$$

The parameters that determine the behavior of $\mathcal{R}_Q(p_T)$ [Eq. (3)] are the frequency factor of reactive collisions

τ_Q^{-1} [Eq. (2)], the plasma lifetime t_m, and the plasma radius R_0. In particular, for $p_T=0$ the value $\mathcal{R}_Q(0)$ is solely determined by the ratio of the plasma lifetime to the collision frequency according to

$$\mathcal{R}_Q(p_T=0)=6(\tau_Q/t_m)^3\left\{1-\exp(-t_m/\tau_Q)\left[\frac{1}{2}\left[\frac{t_m}{\tau_Q}\right]^2+\frac{t_m}{\tau_Q}+1\right]\right\}.$$

In Fig. 2, the suppression ratio $\mathcal{R}_Q(p_T)$ is shown for J/ψ as a function of the transverse momentum p_T. The experimental values of the J/π abundance (normalized to the continuum background) at high transverse energy ($E_T>85$ GeV) to the abundance at low energy ($E_T<34$ GeV) where no suppression is found are taken according to the most recently available data.[2] The theoretical values according to Eq. (3) are shown for different parameter values R_0/t_m, but for the same value $t_m/\tau_Q=1.2$ corresponding to a value $\mathcal{R}_c(0)=0.418$. A good fit between experimental values and theoretical results is found for $R_0/t_m=0.15c$.

For a plasma radius $R_0=3$ fm (corresponding to the radius of ^{16}O), $t_m=20$ fm/c results for the plasma lifetime, and for the collision frequency $\tau_c=16.6$ fm/c follows. According to Eq. (2), one obtains $T_c=178$ MeV for the plasma temperature. The value of the plasma lifetime $t_m\approx 20$ fm/c is on the line with the expected long duration of the mixed phase during the plasma rehadronization (see Refs. 20–23), but is longer than the values assumed within the model of static screening[4–8] which are of the order of 1 fm/c. This smaller value results from the assumption that the suppression ratio is determined by the formation time of the bound-quarkonium state. Our estimation for the plasma temperature T_c is also in agreement with other theoretical results and seems consistent with the experimental findings for particle spectra.[24]

If the radius R_0 of the plasma is increased so that R_0/t_m is also increased, the ratio $\mathcal{R}_c(p_T)$ as a function of p_T is flattened (as seen in Fig. 2). Therefore, for experiments with ^{32}S lower values of $\mathcal{R}_c(p_T)$ than for experiments with ^{16}O are expected. In general, the value of the plasma radius R_0 should depend on the energy density in the reaction.

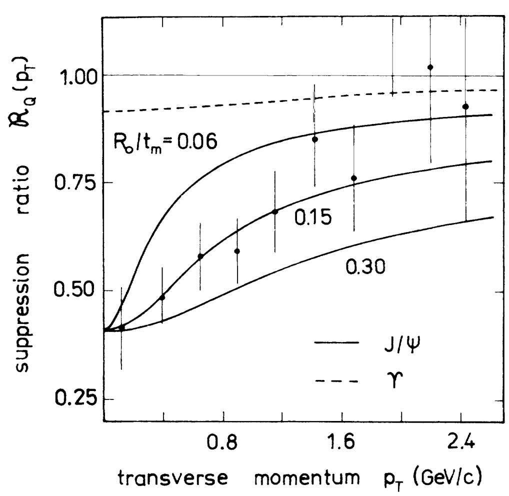

FIG. 2. Suppression ratio $\mathcal{R}_Q(p_T)$ of heavy quarkonia as a function of the transverse momentum p_T. Filled circles with error bars: NA38 data (Ref. 2) for ^{16}O+U collisions. Solid lines: model predictions for J/ψ suppression (3) with $t_m/\tau_c=1.2$ and different values of R_0/t_m (in units of c). Dashed line: model prediction for Υ suppression with the parameter values $t_m/\tau_b=0.1$ and $R_0/t_m=0.15c$.

Our simple model also allows us to predict the suppression of other heavy quarkonia. As an example, the suppression ratio $\mathcal{R}_b(p_T)$ for the plasma parameters $R_0=3$ fm, $t_m=20$ fm/c, is also shown in Fig. 2. It is expected that the suppression of Υ is small [$\mathcal{R}_b(p_T=0)\sim 0.93$].

Of course, the estimation of the suppression ratio in high-energy nucleus-nucleus collisions given here could be improved by taking into account the contribution of excited-quarkonia states and the analysis of the x_F dependence. Furthermore, the hydrodynamical model to describe the spacetime evolution of the plasma and the evaluation of the collision frequency can be refined using more detailed descriptions which will be the subject of further works.

Our model will be further improved if additionally reactions within the hadronic state of matter are taken into account. Hadronic J/ψ suppression due to inelastic J/ψ nucleon scattering has been considered recently by Gavin, Gyulassy, and Jackson[9] (see also Refs. 10 and 11). The concept of reactive scattering processes is also used in our model, but on a quark level where the medium-dependent cross sections are estimated in the plasma phase from simple geometry ($r_{Q\bar{Q}}$) and the activation energy $\Delta E_Q(T)$. It is expected that the possible transition from dense hadronic matter to a quark-gluon plasma phase may be connected with an abrupt change of the reactive cross section which should be seen in the E_T dependence of the heavy-quarkonia suppression.

In conclusion, we presented a model calculation of the suppression of heavy-quarkonia states based on a quantum-statistical treatment of the many-quark system,[18] which is an improved version of the more phenomenological static-screening model introduced by Matsui and Satz.[3,16] Instead of the quantum-mechanically not-well-founded concept of the formation time of a bound state, we describe the suppression of heavy-quark bound states within the framework of chemical reaction kinetics which is also more coherent to other conventional approaches.[9] Our simple model allows to extract the parameters of the possible plasma phase immediately from the experimental suppression patterns which are observed in different reactions[1,2] (see also Ref. 25). Estimations given here for the plasma lifetime and the plasma temperature differ from those of the static screening model,[4–8] but correspond to expectations derived within other approaches.[20–23]

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