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Quarkonium suppression

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Abstract. I discuss quarkonium suppression in equilibrated strongly interacting matter. After a brief review of basic features of quarkonium production I discuss the application of recent lattice data on the heavy quark potential to the problem of quarkonium dissociation as well as the problem of direct lattice determination of quarkonium properties in finite temperature lattice QCD.

Keywords. Quarkonium; heavy quark potential; finite temperature quantum chromodynamics.

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1. Introduction

The behavior of the heavy quarkonium states in hot strongly interacting matter was proposed as a test of its confining nature, since a sufficiently hot deconfined medium will dissolve any binding between the quark–antiquark pair [1]. Another possibility of dissociation of certain quarkonium states (subthreshold states at $T = 0$) is the decay into open charm (beauty) mesons due to in-medium modification of both the quarkonia and heavy–light meson masses [2,3].

The production of J/ψ and Υ mesons in hadronic reactions occurs in part through production of higher excited $c\bar{c}$ (or $b\bar{b}$) states and their decay into quarkonia ground state. Since the lifetime of different subthreshold quarkonium states is much larger than the typical lifetime of the medium which may be produced in nucleus–nucleus collisions, their decay occurs almost completely outside the produced medium. This means that the produced medium can be probed not only by the ground state quarkonium but also by different excited quarkonium states. Since different quarkonium states have different sizes (binding energies), one expects that higher excited states will dissolve at lower temperature than the smaller and more tightly bound ground states. These facts may lead to a sequential suppression pattern in J/ψ and Υ yield in nucleus–nucleus collision as a function of the energy density [4].

The rest of the paper is organized as follows: In §2 a brief review of the basic features of quarkonium production is presented. The possibility of quarkonium dissociation below deconfinement is discussed in §3. The use of lattice data on the heavy quark potential to determine the dissociation temperatures due to color screening is discussed in §4. Finally the problem of direct lattice determination of quarkonium properties at finite temperature is discussed in §5 followed by the conclusions presented in §6.

2. Quarkonium production and feed-down

It is well-known that J/ψ production in hadron–hadron collision is, to a considerable extent, due to the production and subsequent decay of higher excited $c\bar{c}$ states [5–7]. The feed-down from higher excited states was systematically studied in proton–nucleon and pion–nucleon interactions with 300 GeV incident proton (pion) beams [7]. In these studies the cross sections for direct production of different charmonium states (excluding feed-down) were measured. Then making use of the known branching ratios $B[\chi_1(1P) \to \psi(1S)] = 0.27 \pm 0.02$, $B[\chi_2(1P) \to \psi(1S)] = 0.14 \pm 0.01$, and $B[\psi(2S) \to \psi(1S)] = 0.55 \pm 0.05$, one obtains the fractional feed-down contributions f_i of the different charmonium states to the observed J/ψ production. These are shown in the second and third columns of table 1.

In the case of bottomonium the experiment provides only the inclusive (i.e. including also the feed-down from higher states) cross-section for different (nS) states [8]. The feed-down from (nP) states is known only for transverse momenta $p_{\mathrm{T}} \geq 8$ GeV/c [8]. To analyze the complete feed-down pattern, we thus have to find a way to extrapolate these data to $p_{\mathrm{T}} = 0$ as well as to determine the direct cross-section for different (nS) states. This can be done using the most simple and general model for quarkonium production, the color evaporation model [9]. In particular this model predicts that the ratios of cross-sections for production of different quarkonium states are energy independent. This prediction was verified for a considerable range of energies [10]. The ratios between the different $\chi_J(1P)$ states in this model are predicted to be governed essentially by the orbital angular momentum degeneracy [9]; we thus expect for the corresponding cross-sections

$$\chi_0(1P) : \chi_1(1P) : \chi_2(1P) = 1 : 3 : 5. \tag{1}$$

From table 1 we have for $\pi^- N$ collisions $\chi_2(1P)/\chi_1(1P) \simeq 1.44 \pm 0.38$ and thus reasonable agreement with the predicted ratio 1.67. Actually, for pN interactions, the experiment measures only the combined effect of χ_1 and χ_2 decay (30% of the overall J/ψ production); the listed values in table 1 are obtained by distributing this in the ratio 3 : 5.

Using considerations based on color evaporation model, in particular eq. (1), the feed-down from higher excited $b\bar{b}$ states to Υ production can be predicted [11]; the feed-down fractions are summarized in table 1. Alternatively the feed-down fraction from higher excited $b\bar{b}$ states can be predicted using NRQCD factorization formula [11]. The results of this analysis are summarized in the last column of table 1.

Table 1. Feed-down fractions from higher excited states to the J/ψ and Υ states.

State	$f_i(\pi^- N)$ (%)	$f_i(p\,N)$ (%)	State	$f_i(p\bar{p})$ (%)	$f_i^{\mathrm{NRQCD}}(p\bar{p})$ (%)
$J/\psi(1S)$	57 ± 3	62 ± 4	$\Upsilon(1S)$	52 ± 9	52 ± 34
$\chi_1(1P)$	20 ± 5	16 ± 4	$\chi_b(1P)$	26 ± 7	24 ± 8
$\chi_2(1P)$	15 ± 4	14 ± 4	$\Upsilon(2S)$	10 ± 3	8 ± 7
$\psi(2S)$	8 ± 2	8 ± 2	$\chi_b(2P)$	10 ± 7	14 ± 4
			$\Upsilon(3S)$	2 ± 0.5	2 ± 2

3. Quarkonium dissociation below deconfinement

Recent lattice calculations of the heavy quark potential show evidence for the string breaking at finite temperature [12]. On the lattice the potential is calculated from the Polyakov loop correlator, to which it is related by

$$V(r,T) = -\ln\langle L(r)L^{\dagger}(0)\rangle + C, \tag{2}$$

where $L(r)$ is the Polyakov loop (see e.g. ref. [12] for definition). The normalization constant C contains both the cut-off dependent self-energy and the entropy contributions $-TS$ (for $T \neq 0$, $-\ln\langle L(r)L^{\dagger}(0)\rangle$ is actually the free energy of the static $Q\bar{Q}$ pair). For a properly chosen normalization constant C, $V(r,T)$ is the ground state energy of an infinitely heavy $Q\bar{Q}$ pair. In the absence of dynamical quarks (quenched QCD) $V(r,T)$ is linearly rising with r for large separations indicating the existence of a flux tube (string). If dynamical quarks are present, the flux tube can decay (the string can break) by creating a pair of light quarks $q\bar{q}$ from the vacuum once $V(r,T)$ is larger than twice the binding energy of a heavy–light $Q\bar{q}$ ($q\bar{Q}$) meson [12a] Thus the potential at very large distances is constant $V_{\infty}(T)$ and is equal to twice the binding energy of a heavy–light meson.

It is expected that medium effects are not important at very short distances. Therefore, at very short distances the potential $V(r,T)$ should be given by the Cornell potential [13]

$$V(r) = -\frac{e}{r} + \sigma r. \tag{3}$$

We use this fact to determine the normalization constant C and set the potential to be of the Cornell form at the the smallest distance $rT = 0.25$ available in lattice studies of [12], with $e = 0.4$ as expected for $(2+1)$-flavor QCD [14]. The resulting potential and $V_{\infty}(T)$ are shown in figure 1. Note the strong temperature dependence of $V_{\infty}(T)$. Since for sufficiently heavy quarks ($m_{\mathrm{Q}} \gg \Lambda_{\mathrm{QCD}}$) it does not matter whether the quark is infinitely heavy or just merely heavy, the open charm (beauty) meson masses are approximately given by $2M_{D,B}(T) = 2m_{c,b} + V_{\infty}(T)$.

Now the temperature dependence of the different quarkonium states should be addressed. At zero temperature the heavy quark masses permit the application of potential theory for the description of quarkonium spectroscopy (see e.g. [15]). Furthermore, it turns out that the time scale of gluodynamics relevant for quarkonium spectroscopy is smaller than $(m_{\mathrm{Q}}v)^{-1}$ (v being the heavy quark velocity) [15]. For sufficiently heavy quarks this time scale is much smaller than the typical hadronic time scale $\Lambda_{\mathrm{QCD}}^{-1} \sim 1$ fm. The decay of the flux-tube like all other hadronic decays has time scale of order 1 fm. Therefore in the potential theory the potential must always be of Cornell form (i.e. linearly rising at large distances). These considerations have direct phenomenological support. Namely, simple potential models with linearly rising potential can describe reasonably well also the quarkonium states above the open charm (beauty) threshold. Many of these higher excited states have an effective radius of order of or even larger than 1 fm [14,13]. Contrary to this situation in the case of the potential becoming flat around 1 fm (the expected radius of string breaking at $T = 0$) the higher excited states above the open charm (beauty) threshold simply do not exist. Therefore, the temperature dependent quarkonium masses have been determined from the Schrödinger equation

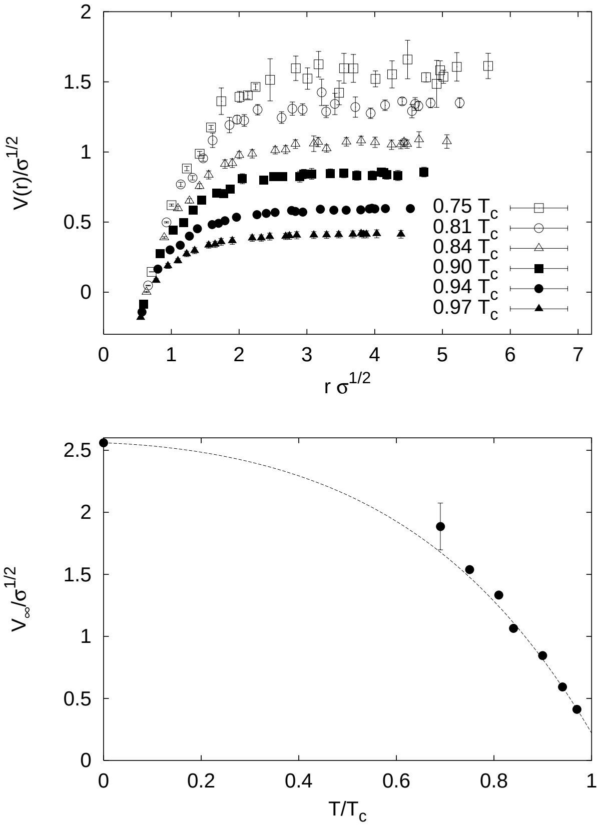


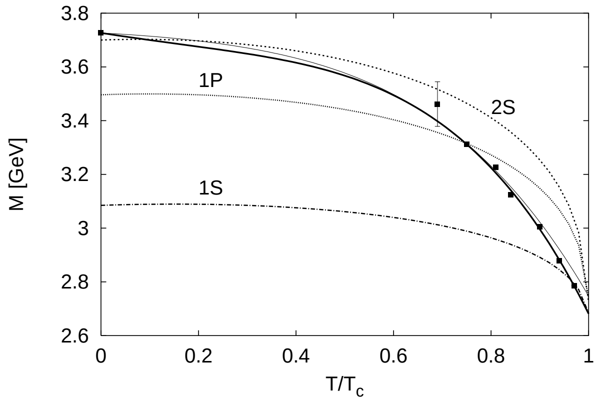
Figure 1. The heavy quark potential and its asymptotic value below deconfinement at different temperatures. The line on the bottom figure is a fit to the data points.

$$\left[2m_{\mathrm{Q}} - \frac{1}{m_{\mathrm{Q}}}\nabla^2 + V(r)\right]\Phi_i = M_i\Phi_i, \tag{4}$$

where the potential $V(r)$ is identified with the temperature dependent string potential [16]

$$\begin{aligned} V_{\mathrm{string}}(r,T) = &-\left(e - \frac{1}{6}\arctan(2rT)\right)\,\frac{1}{r} \\ &+ \left(\sigma(T) - \frac{\pi T^2}{3} - \frac{2T^2}{3}\arctan\frac{1}{rT}\right)\, r + \frac{1}{2}\ln(1+4r^2T^2). \end{aligned} \tag{5}$$

This form of the potential describes quite well the temperature dependence of the heavy quark potential in quenched QCD for appropriately chosen $\sigma(T)$ [17]. In order to make contact to real QCD we set $e = 0.4$ and use $T_{\mathrm{c}}/\sqrt{\sigma} = 0.425$ from [12] for the deconfinement temperature (by σ we always denote the string tension at zero temperature). Furthermore we use the following values of the heavy quark masses, $m_c = 1.3$ GeV and $m_b = 4.72$ GeV as well as $\sqrt{\sigma} = 0.44$ GeV for the zero temperature string tension. This set of parameters



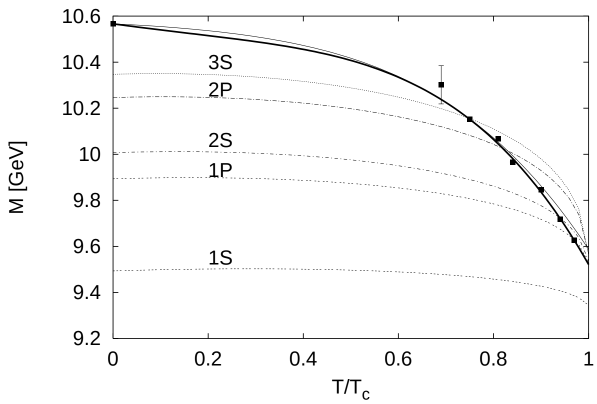


Figure 2. The masses of different quarkonia states and the open charm (beauty) threshold as a function of temperature. Shown are the charmonia masses and open charm threshold (top) and bottomonia masses and open beauty threshold (bottom) as function of the temperature. The thick solid line is the open charm (beauty) threshold obtained from normalization at $r = 1/(4T)$. The thin solid line is the open charm (beauty) threshold obtained from normalization at $r = \sqrt{2}/(4T)$ (see text).

gives a fairly good description of the observed quarkonium spectrum at zero temperature. The temperature dependence of the string tension was taken from [17]. The resulting quarkonia masses are shown in figure 2.

Since the smallest distance available on lattice is only $0.25T^{-1}$, one may worry about possible medium effects at this distance and their role in the determination of $V_\infty(T)$. Normalizing the Polyakov loop correlator (2) at $r = 1/(4T)$ with eq. (5), we thus obtain what might be a more reliable estimate of the plateau $V_\infty(T)$ than with the $T = 0$ form (3). It turns out, however, that the two forms of short distance behavior resulting from the zero temperature Cornell potentials (3) and (5) are practically identical, so that the normalization is in fact not affected by the in-medium modifications at larger distances. To consider further possible uncertainties of the normalization procedure, the Polyakov loop correlator has been normalized also at the next smallest distance $r = \sqrt{2}/(4T)$. The resulting two forms of $V_\infty(T)$ are shown in figure 2. The difference between the two curves of $V_\infty(T)$ provides an estimate of the normalization error. Except for the region very near $T = T_c$, the uncertainty is seen to be quite small.

From figure 2 one can see that ψ' and χ_c states become an open charm states well below T_c and can dissociate by decaying into $D\bar{D}$. The situation is similar for $\Upsilon(3S)$ and $\chi_b(2P)$ states which can decay into $B\bar{B}$ below T_c. For J/ψ, $\chi_b(1P)$ and $\Upsilon(2S)$ it is not possible to say whether they will dissociate above T_c or just below T_c. Finally, the $\Upsilon(1S)$ state will definitely dissociate above the deconfinement. We are going to estimate the dissociation temperatures of these states in the next section.

4. Quarkonium dissociation by color screening and the sequential suppression pattern

In the deconfined phase it is customary to choose the constant C in (2) to be the value of the correlator at infinite separation $C = T\ln\langle L(r)L^\dagger(0)\rangle|_{r\to\infty} \equiv \ln|\langle L\rangle|^2$. The resulting connected correlator defines the so-called color averaged potential [18]

$$V(r,T) = -T\ln\frac{\langle L(r)L^\dagger(0)\rangle}{|\langle L\rangle|^2}. \tag{6}$$

The color averaged potential written as the thermal average of the potentials in color singlet $V_1(T,r)$ and color octet $V_8(T,r)$ states

$$\exp(-V(r,T)/T) = \frac{1}{9}\exp(-V_1(r,T)/T) + \frac{8}{9}\exp(-V_8(r,T)/T). \tag{7}$$

In potential models it is assumed that quarkonium is dominantly a singlet $Q\bar{Q}$ state. Furthermore the octet channel is repulsive (at least in perturbation theory) and therefore only a singlet $Q\bar{Q}$ pair can be bound in the deconfined phase. Thus we need to know the singlet potential. Lattice data in the relevant case of 3 flavor QCD exist only for the averaged potential [12,19]. The averaged potential in 3 flavor QCD is shown in figure 3 for three representative temperatures. Note that within the present accuracy of the lattice calculations the potential vanishes beyond some distance $r_0(T)$ denoted by vertical arrows in figure 3. In perturbation theory, the leading terms for both singlet and octet potentials are of Coulomb form at high temperature and small r ($r \ll T^{-1}$)

$$V_1(r,T) = -\frac{4}{3}\frac{\alpha(T)}{r}, \quad V_8(r,T) = +\frac{1}{6}\frac{\alpha(T)}{r}, \tag{8}$$

with $\alpha(T)$ for the temperature-dependent running coupling. In the region just above the deconfinement point $T = T_c$, there will certainly be significant non-perturbative effects of unknown form. We attempt to parameterize the existing non-perturbative effects by the following form

$$V_1(r,T) = -\frac{4}{3}\frac{\alpha(T)}{r}\exp\{-\mu(T)r\},\; V_8(r,T) = \frac{1}{6}c(T)\frac{\alpha(T)}{r}\exp\{-\mu(T)r\}, \tag{9}$$

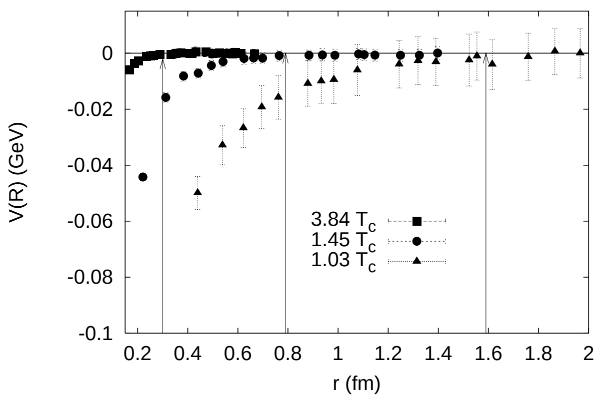


Figure 3. The color average potential calculated on lattice.

where $\mu(T)$ denotes the effective screening mass in the deconfined medium. We fit the lattice data of [12,19] by eqs (7) and (9) assuming $\alpha(T)$, $\mu(T)$ and $c(T)$ to be unknown functions of T. Since the present data are not precise enough to determine $\alpha(T)$, $\mu(T)$ and $c(T)$ simultaneously, some additional constraints coming from simulations of pure gauge theory [20–23] should be invoked in the fit procedure. The fit procedure is described in detail in ref. [11]. Here it is sufficient to mention that $\alpha(T)$ can be well described by the 1-loop formula for the QCD running coupling with $\Lambda_{\mathrm{QCD}} = (0.34 \pm 0.01)T_{\mathrm{c}}$ and the screening mass $\mu(T)$ is constant in units of the temperature $\mu(T) = (1.15 \pm 0.02)T$ [11].

Now we are in a position to discuss quarkonium dissociation due to color screening. It is natural to assume that the heavy $Q\bar{Q}$ pair cannot exist as a bound state if its effective binding radius (the mean distance between Q and $\bar{Q}$) is larger than the screening radius of the medium. The effective radii for different bound states are calculated from eq. (4) with $V(r) = V_1(r,T)$. The screening radius of the medium can be identified with $1/\mu(T)$. However, the value of $\mu(T)$ strongly depends on assumptions we have made to determine it. A less model dependent and more conservative approach would be to identify the screening radius with $r_0(T)$ defined above. We use the latter approach. In figure 4 we show the effective radius of J/ψ and Υ states and the screening radius as functions of temperature. The intersection of these curves defines the dissociation temperature of J/ψ and Υ states. Similar analysis was done for excited states which may survive above T_{c}.

Alternatively one can define the dissociation temperatures as the temperature where the effective bound state radius diverges [24] and the quark–antiquark pair is unbound. It is clear from figure 4 that such definition will lead to larger, though not very different value of the dissociation temperature. However, one should note that once the radius of the bound state is larger than the screening radius the present treatment based on Schrödinger equation is clearly not valid and the effect of the medium becomes so strong that quarkonium dissociation is very likely to happen.

Now the dissociation temperatures are known for all quarkonium states and summarized in table 2. Combining these dissociation temperature with the feed-down fractions determined in §2 we can predict the sequential suppression pattern of J/ψ and Υ states as functions of temperature. These are summarized in figure 5.

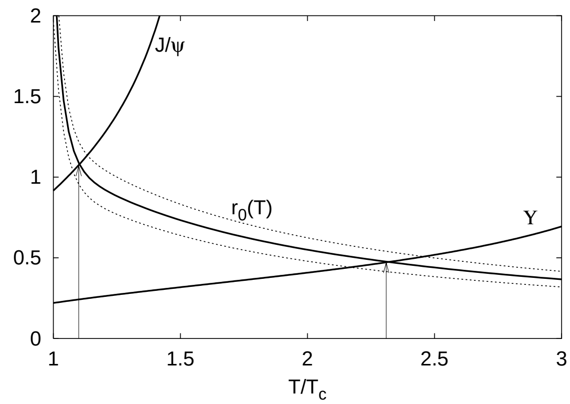


Figure 4. The effective radii of J/ψ and Υ states and the screening radius $r_0(T)$. The dotted lines denote the uncertainty in the value of $r_0(T)$.

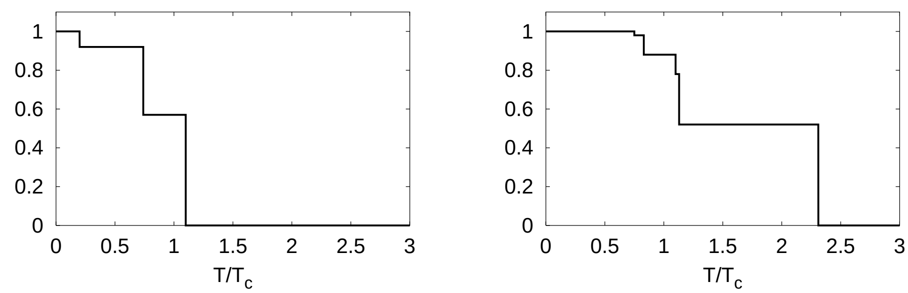


Figure 5. The suppression pattern of J/ψ (left) and Υ yield (right) as functions of temperature.

Table 2. Dissociation temperatures of different quarkonium states.

$q\bar{q}$	J/Ψ	χ_c	ψ'	$\Upsilon(1S)$	$\chi_b(1P)$	$\Upsilon(2S)$	$\chi_b(2P)$	$\Upsilon(3S)$
T/T_c	1.10	0.74	0.2	2.31	1.13	1.10	0.83	0.75

5. Lattice determination of quarkonium properties at finite temperature

At zero temperature, quarkonium properties (masses, decay constants etc.) can be obtained from the behavior of the mesonic correlators at large Euclidean time separations

$$\langle O(\tau,\vec{p}), O(0,-\vec{p})\rangle_{\tau\to\infty} = \sum_n A_n \mathrm{e}^{-m_n \tau}. \tag{10}$$

Here $O(\tau,\vec{p}) = \sum_{\vec{x}} \mathrm{e}^{i\vec{p}\vec{x}} O(\tau,x)$ is the mesonic operator bilinear in quark fields $O(\tau,\vec{x}) = \bar{q}(x)\Gamma q(x)$; $\Gamma = \gamma_5, \gamma_\mu$ for the pseudo-scalar and the vector channel, respectively. Direct application of eq. (10) for the determination of quarkonium masses is difficult because at large separations statistical errors in the correlation function become very large. To improve the situation one usually uses smeared (improved) operators $\tilde{O}$ instead of O such that only one coefficient A_n in eq. (10) is significantly different from zero. A possible choice of the operator $\tilde{O}$ which has an optimal projection onto a given quarkonium state is

$$\tilde{O}(\tau,\vec{z}) = \sum_{\vec{y}} \omega(\vec{y})\bar{q}(\tau,\vec{z})\Gamma q(\tau,\vec{z}+\vec{y}), \tag{11}$$

where $\omega(\vec{y})$ is the trial wave function [24a] (see [25] for further details).

At finite temperature further difficulties appear in direct lattice determination of quarkonium masses. First of all, eq. (10) is not applicable in principle because the time extent is limited by the inverse temperature. Since the temporal lattice size becomes smaller as the temperature is increased, fewer data points on the temporal correlators are available. Another problem which is also present at zero temperature but becomes more serious at finite temperature (just because one is enforced to consider mesonic correlators also at short distances) is the discretization errors of order $m_{\mathrm{Q}}a$ with m_{Q} being the heavy quark mass and a being the lattice spacing.

To study quarkonium properties at finite temperature in quenched QCD correlators of smeared operators were calculated on anisotropic lattices [26,27]. To deal with large discretization errors in [26] the NRQCD formalism was used (where the scale m_{Q} is integrated out). In this calculation the mesonic correlators corresponding to the ground state quarkonium show only small changes up to temperature $1.2T_{\mathrm{c}}$ compared to the zero temperature case while the first excited state shows very dramatic change with the temperature. The authors of [27] used a different formalism, namely they used the Fermilab action [28] on anisotropic lattices which has no discretization error of order $m_{\mathrm{Q}}a$ at tree level. They observed more dramatic changes of the mesonic correlators as the deconfinement point is crossed. However, the correlators are very far from the free ones even at temperatures as high as $1.5T_{\mathrm{c}}$ indicating possible existence of bound states at this temperature.

There are several problems with the approach presented in [26,27]. The most serious one is the use of optimized correlators which assumes existence of well-defined quarkonium states at finite temperature. Clearly, different particle states at zero temperature will appear as quasiparticles with finite width at non-zero temperature. The implementation of NRQCD formalism in [26] does not assume anti-periodic boundary condition in time direction for the quark propagators and therefore it is not clear to what extent the meson correlators calculated in [26] can be related to some finite temperature (retarded) correlation functions. In [27] the spatial lattice spacing was quite large and the effect of doublers cannot be completely neglected.

Another possibility to extract quarkonium properties in lattice QCD is to calculate usual point-to-point correlators $G(\tau,\vec{p}) = \langle O(\tau,\vec{p})O(0,-\vec{p})\rangle$ (i.e. correlators of point sources) in imaginary time and extract the spectral function. The information on the bound state are then encoded in the peaks of the spectral function.

The imaginary time correlator can be related to the retarded meson propagator $G_{\mathrm{R}}(\omega,p)$ by analytic continuation

$$G(i\omega_n,p) = \int_0^{1/T} \mathrm{e}^{i\omega_n\tau} G(\tau,p), \quad G(i\omega_n \to \omega + i\varepsilon, p) = G_{\mathrm{R}}(\omega,p). \tag{12}$$

which allows to write down the spectral representation for $G(\tau,p)$,

$$G(\tau,p)=\int_0^\infty \mathrm{d}\omega\sigma(\omega,p)\frac{\cosh(\omega(\tau-1/(2T))}{\sinh(\omega/(2T))}. \tag{13}$$

In principle this equation allows to determine $\sigma(\omega)$ from $G(\tau,p)$. However, in practice the value of G is available only for $N_\tau/2\sim 8$ different τ-values. For reasonably fine discretization in ω-space one has $N_\omega\sim 700$ degrees of freedom to be reconstructed. This problem can be solved only using the maximum entropy method (MEM) (see [29] for a review). Mesonic spectral functions in different channels were successfully reconstructed using this method [30]. Very recently it has been demonstrated that spectral function also at finite temperature can be reconstructed using this method and isotropic lattices with $N_\tau=12,\ 16$ sites in temporal direction [31,32]. Thus a natural alternative to the approach presented in [26,27] could be the calculation of point-to-point mesonic correlators using isotropic lattice with non-perturbatively improved clover action [32]. Compared to [27] this approach has the advantage that all lattice artifacts of order Ta pa are completely removed but also the disadvantage that $m_\mathrm{Q}a$ effects are present even at the tree level. The latter, however, can be controlled if a is small enough. This approach is currently being investigated [33]. Calculations so far have been performed on $48^3\times16$, $48^3\times12$ and $64^3\times16$ lattices and gauge coupling $\beta=6.499,\ 6.640$ and 6.872. These parameters correspond to temperatures $T=0.9T_\mathrm{c},\ 1.2T_\mathrm{c}$ and $1.5T_\mathrm{c}$ and lattice spacing from $a^{-1}=6.5$ GeV, 4.9 GeV and 4.0 GeV. The quark mass in these calculations was about 1.5 GeV. The discretization errors in the correlation function turns out to be large, e.g the value of $G(\tau=1/(2T))$ at $1.5T_\mathrm{c}$ calculated with $a^{-1}=4.9$ GeV and $a^{-1}=6.5$ GeV differs by 16% while the same difference in the light quark sector is about 1%. Because of this MEM cannot be applied to extract the spectral function unless extrapolation to the continuum limit is performed. Nevertheless, correlation functions themselves can provide some information about the existence of quarkonium bound states. Figure 6 shows the ratio of the meson correlator calculated at $1.2T_\mathrm{c}$ ($48^3\times12$, $\beta=6.499$) to the one calculated at $0.9T_\mathrm{c}$ ($48^3\times16$, $\beta=6.499$). Except the last point at $\tau\simeq0.3$ fm this ratio stays very close to unity at 1σ level. The deviation at $\tau\simeq0.3$ fm is simply due to periodic boundary condition ($\tau\simeq0.3$ fm corresponds to $\tau=1/(2T)$ on $48^3\times12$ lattice). While this fact does not necessarily imply that the

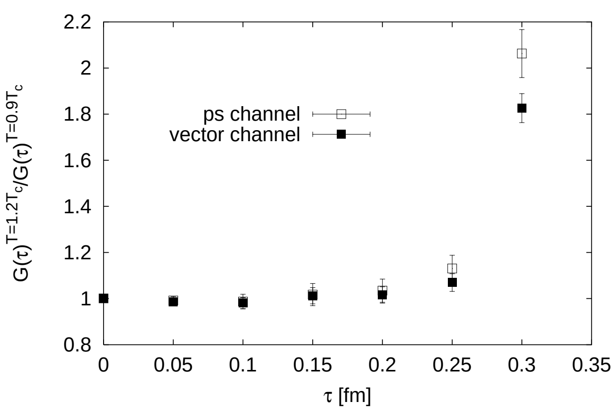


Figure 6. The ratio of the correlators in pseudo-scalar and vector channels at $1.2T_\mathrm{c}$ to the one at $0.9T_\mathrm{c}$. Calculations were done at $\beta=6.499$ on $48^3\times12$ and $48^3\times16$ correspondingly.

quarkonium spectral function does not change as the deconfinement point is crossed, it does imply that the propagation of heavy quarks in the deconfined phase is far from free propagation of $Q\bar{Q}$ pair in agreement with studies performed in [26,27].

6. Conclusions

I have considered quarkonium dissociation in hot strongly interacting matter below as well as above the deconfinement. In a confined medium, dissociation of certain quarkonium states occurs due to in-medium modification of the open charm (beauty) threshold as well as the quarkonia masses. In the deconfined medium quarkonium, dissociation is due to color screening. In this analysis quarkonium masses were extracted from Schrödinger's equation. The singlet potential used in Schrödinger's equation was extracted from lattice data on the Polyakov loop correlator using some additional assumptions. For a more accurate determination of the quarkonium suppression patterns, it would be desirable to carry out direct lattice studies of the color singlet potential and of its quark mass dependence, which may become important near the critical temperature. Furthermore, to make contact with nuclear collision experiments, a more precise determination of the energy density via lattice simulations is clearly needed, as is a clarification of the role of a finite baryochemical potential. For the latter problem, lattice studies are so far very difficult; nevertheless, a recent new approach [34] could make such studies feasible.

In §5 the problem of direct lattice determination of quarkonium properties was discussed. An unbiased determination of these properties free of lattice artifacts seems to be very difficult and not available so far even in the quenched approximation. Such an analysis will be a very important check of the results obtained from the simple potential picture discussed in §§3 and 4.

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