Quarkonium in Hot Medium

Péter Petreczky
Department of Physics and RIKEN-BNL Research Center, Brookhaven National Laboratory, Upton, New York, 11973

Abstract. I review recent progress in studying quarkonium properties in hot medium as well as possible consequences for quarkonium production in heavy ion collisions.

PACS numbers: 11.15.Ha, 11.10.Wx, 12.38.Mh, 25.75.Nq

1. Introduction

There has been considerable interest in studying quarkonia in hot medium since the publication of the famous Matsui and Satz paper [1]. It has been argued that color screening in a deconfined QCD medium will suppress the existence of quarkonium states, signaling the formation of a quark-gluon plasma (QGP) in heavy-ion collisions. Although this idea was proposed a long time ago, first principle QCD calculations, which go beyond qualitative arguments, have been performed only recently. Such calculations include lattice QCD determinations of quarkonium correlators [2, 3, 4, 5, 6]; potential model calculations of the quarkonium spectral functions with potentials based on lattice QCD [7, 8, 9, 10, 11, 12, 13] (see also Ref. [14] for a review), as well as effective field theory approaches that justify potential models and reveal new medium effects [15, 16, 17, 18, 19]. Spectral properties of heavy quark bound states are important ingredients in modeling of heavy quarkonium production in hot medium as will be discussed later.

2. Color screening and deconfinement

At high temperatures, strongly-interacting matter undergoes a deconfining transition to a quark-gluon plasma (QGP). This transition is triggered by a rapid increase of the energy and entropy densities, as well as the disappearance of hadronic states. (For recent reviews, see Ref. [20, 21, 22]). According to current lattice calculations, at zero net baryon density deconfinement occurs at $T_c \sim (170 - 195)$ MeV [23, 24].

‡ In the view of crossover nature of the deconfinement transition it is difficult to define the corresponding temperature interval precisely. Moreover due to discretization errors there is a discrepancy in the value of the deconfinement transition temperature obtained in calculations using different discretization schemes. Calculations with so-called stout staggered fermion action give a deconfinement transition temperature around 170 MeV [25, 26]. The new calculation with so-called
The QGP is characterized by color screening: the range of interaction between heavy quarks becomes inversely proportional to the temperature. Thus at sufficiently high temperatures, forming a bound state with a heavy quark (c or b) and its anti-quark becomes impossible. Color screening is studied on the lattice by calculating the spatial correlation function of a static quark and anti-quark, which propagates in Euclidean time from $\tau = 0$ to $\tau = 1/T$ where $T$ is the temperature (see Ref. [28] for a recent review).

Two types of correlation functions are usually calculated on the lattice. The correlation function of Polyakov loops, which is also called the color averaged correlator

$$G(r, T) = \frac{1}{9}\langle \text{Tr} L(r) \text{Tr} L^\dagger(0) \rangle,$$

(1)

where the temporal Wilson line is defined in terms of link variable $U_0(x_0, r)$ as $L(r) = \prod_{x_0=0}^{N_t-1} U_0(x_0, r)$, and the color singlet correlator

$$G_1(r, T) = \frac{1}{3}\langle \text{Tr} L(r) L^\dagger(0) \rangle$$

(2)

HISQ action [27] seem to support this scenario.
The later is defined in Coulomb gauge [29, 30, 31, 32] or in a gauge invariant but path dependent manner [28]. The singlet correlators recently have been also calculated in perturbation theory [34]. The free energy of static quark anti-quark pair is defined through the logarithm of the color averaged correlator \( F(r, T) = -T \ln G(r, T) \).

Analogously, the so-called singlet free energy is defined as \( F_1(r, T)/T = -\ln G_1(r, T) \). The singlet free energy in 2+1 flavor QCD is shown in Fig. 1 (right). The numerical results were obtained on \( 16^3 \times 4 \) lattice using p4 action and quark masses which correspond to pion mass of about 220MeV [36]. As expected, in the zero temperature limit the singlet free energy coincides with the zero temperature potential calculated on the lattice [24]. Figure 1 also illustrates that, at sufficiently short distances, the singlet free energy is temperature independent and equal to the zero temperature potential. At large distances it approaches a constant value \( F_\infty(T) \), which is twice the free energy of an isolated static quark. The range of interaction decreases with increasing temperatures. For temperatures above the transition temperature, \( T_c \), the heavy quark interaction range becomes comparable to the charmonium radius. In particular, the combination \( r(F_1(r, T) - F_\infty(T)) \) is expected to decay exponentially at large distances, \( r m_D > 1 \), with \( m_D \) being the Debye mass. Indeed, for distances \( rT > 0.8\text{fm} \) the screening is exponential as can be seen from Fig. 1 (right). Based on this general observation, one would expect that the charmonium states, as well as the excited bottomonium states, do not exist above the deconfinement transition. (In the literature, this is often referred to as dissociation or melting).

The path/gauge dependence of singlet free energy makes its interpretation somewhat complicated. Therefore, on the lattice one also calculates the Polyakov loop correlator. The results of the calculations in 2+1 flavor QCD with p4 action on \( 16^3 \times 4 \) lattices are shown in Fig. 2 in terms of the logarithm of the correlation function, i.e. the physical free energy of static quark anti-quark pair. Compared to the singlet free energy the physical free energy shows much stronger temperature dependence and is different from the zero temperature potential at all temperatures. In the low temperature region the strong temperature dependence of the free energy can be understood as resulting from significant contribution from higher excited states [33]. At high temperatures, the origin strong temperature dependence is more difficult to understand. In the short distance limit within the effective field theory framework it can be attributed to the contribution from octet degrees of freedom [35]. As one can also see from Fig. 2 the free energy of static \( QQ \) pair also shows exponential screening at large distances. Thus the strong color screening is a real effect. In summary, lattice calculations of static quark anti-quark correlators provide evidence that in the deconfined phase there are strong screening effect at distances relevant for quarkonium physics. In the next section the implication of these findings on quarkonium spectral functions will be discussed.

The identification of the correlators defined by Eqs. (1) and (2) with color averaged and color singlet channels is based on perturbative arguments, e.g. see discussion in Ref. [37]. Non-perturbatively such identification becomes problematic as was shown in Ref. [38]. At short distances effective field theory concepts can be invoked to define color singlet and color averaged channels [35].
3. Quarkonium spectral functions

In-medium quarkonium properties are encoded in the corresponding spectral functions, as are their dissolution at high temperatures. Spectral functions are defined as the imaginary part of the retarded correlation function of quarkonium operators. Bound states appear as peaks in the spectral functions. The peaks broaden and eventually disappear with increasing temperature. The disappearance of a peak signals the melting of the given quarkonium state.

In lattice QCD, the meson correlation functions, \( G(\tau, T) \), are calculated in Euclidean time. These correlation functions are related to the spectral functions \( \sigma(\omega, T) \) as

\[
G(\tau, T) = \int_0^\infty d\omega \sigma(\omega, T) \frac{\cosh(\omega(\tau - 1/(2T)))}{\sinh(\omega/(2T))}.
\]

Detailed information on \( G(\tau, T) \) could make it possible to reconstruct the spectral function from the lattice data. In practice, however, this turns out to be very difficult task because the time extent is limited to \( 1/T \), see the discussion in Ref. and references therein.

The quarkonium spectral functions can be calculated in potential models using the singlet free energy from Fig. or with different lattice-based potentials obtained using the singlet free energy as an input. The results of calculations in quenched QCD are shown in Fig. for S-wave charmonium and bottomonium spectral functions. All charmonium states are dissolved in the deconfined phase, while the bottomonium \( 1S \) state may persist up to temperature of about \( T_c \). The temperature dependence of the Euclidean correlators can be predicted using Eq. and the calculated spectral functions. On the lattice the temperature dependence of the correlation function is studied in terms of the ratio \( G(\tau, T)/G_{\text{rec}}(\tau, T) \), where

\[
G_{\text{rec}}(\tau, T) = \int_0^\infty d\omega \sigma(\omega, T = 0) \frac{\cosh(\omega(\tau - 1/(2T)))}{\sinh(\omega/(2T))}.
\]

If the spectral function is temperature independent \( G/G_{\text{rec}} = 1 \). It turns out, somewhat surprisingly, that the Euclidean correlation functions in the pseudo-scalar channel show very little temperature dependence irrespective of whether a state remains bound (the \( \eta_b(1S) \) ) or not (the \( \eta_c(1S) \) ). Note also that correlators from potential models are in accord with the lattice calculations (see insets in Fig.). Initially, the weak temperature dependence of the correlators was considered to be evidence for the survival of different quarkonium states. It is now clear that this conclusion was premature. There is a large enhancement in the threshold region of the spectral functions relative to the free spectral function, as shown in Fig. This threshold enhancement compensates for the absence of bound states and leads to Euclidean correlation functions with a very weak temperature dependence. It further indicates strong residual correlations between the quark and anti-quark, even in absence of bound states. Present lattice calculations of the spectral functions cannot discriminate between bound state peaks and threshold enhancement.
Figure 3. The $S$-wave charmonium (left) and bottomonium (right) spectral functions calculated in potential models [12]. Insets: correlators compared to lattice data [11]. The dotted curves are the free spectral functions.

Similar analysis has been done for the $P$-wave charmonium and bottomonium spectral functions [12, 13]. Here the contribution of the zero mode is important. The zero mode contribution is responsible for large enhancement and strong temperature dependence of $G/G_{\text{rec}}$ in the scalar and axial-vector channels [12, 39]. A smaller temperature dependence of $G/G_{\text{rec}}$ due to zero mode was observed in the vector channels as well [40]. Here the zero mode contribution is related to heavy quark diffusion. A comprehensive study of the zero mode contribution was presented in Ref. [41] where it was concluded that almost the entire temperature dependence of the correlation function is due to this contribution. Furthermore, the zero mode contributions which are related to generalized susceptibilities are well described by quasi-particle model [41]. This provides additional input for in-medium properties of heavy quarks, e.g. the temperature dependent heavy quark mass, which will be important for constructing more realistic potential model at finite temperature. While the above studies have been performed in quenched approximation, recently attempts to extend the quasi-particle model in the charm quark sector to full QCD have been presented in Ref. [42].

To estimate the dissociation temperatures, i.e. the temperature above which no resonance like peak are observed in the quarkonium spectral functions, it is not sufficient to consider only the effects due to color screening. One has to consider also the effect of thermal broadening of different states. Upper bound on the dissociation temperatures can be obtained from the analysis of the spectral functions calculated in potential model and using simple phenomenological estimates for the thermal width [13]. This analysis shows that most of quarkonium states dissolve at temperatures $< 1.2T_c$ and only the ground state bottomonium may survive up to temperatures as high as $2T_c$.

The application of potential models can be justified using an effective field theory approach. The existence of distinct energy scales related to the heavy quark mass $m$, the inverse size $mv$ (where $v$ is the heavy quark velocity), and the binding energy $mv^2$ makes it possible to construct a sequence of effective theories at zero temperature. The
effective theory which emerges after integrating out the scales $m$ and $mv$ is pNRQCD. The Lagrangian of this effective theory contains singlet and octet meson fields composed of heavy quarks, which are coupled to soft gluon fields at scale $mv^2$. pNRQCD is equivalent to the potential model at $T = 0$ at leading order, where the coupling to the octet channel can be neglected. It is possible to extend this approach to finite temperature where additional scales, the temperature $T$, the Debye mass $m_D \sim gT$, and the magnetic scale $g^2 T$ are present. In the weak coupling regime, $g \ll 1$, these scales are well separated. Depending on how the thermal scales are related to the zero temperature scales, the various hierarchies make it possible to derive different effective theories for quarkonium bound states at finite temperature [19]. In the weak-coupling QCD approach, thermal corrections to the potential are obtained when the temperature is larger than the binding energy. An important general result of these effective theories is that the potential acquires an imaginary part [15, 16, 17, 18, 19]. The imaginary part of the potential smears out the bound state peaks of the quarkonium spectral function, leading to their dissolution prior to the onset of Debye-screening in the real part of the potential (see e.g. the discussion in Ref. [18]).

4. Dynamical models for quarkonium production

While knowing the quarkonium spectral functions in equilibrium QCD is necessary, it is insufficient to predict effects on their production in heavy-ion collisions because, unlike the light degrees of freedom, heavy quarks are not fully thermalized in heavy-ion collisions. Therefore, it is nontrivial to relate the finite temperature quarkonium spectral functions to quarkonium production rates in heavy ion collisions without further model assumptions. The bridge between the two is provided by dynamical models of the matter produced in heavy-ion collisions. Some of the simple models currently available are based on statistical recombination [43]; statistical recombination and dissociation [44, 45]; or sequential melting [46]. Here we highlight a more recent model, which makes closer contact with both QCD and experimental observations [47].

The bulk evolution of the matter produced in heavy-ion collisions is well modeled by hydrodynamics, see Ref. [48] for a recent review. The large heavy quark mass makes it possible to model its interaction with the medium by Langevin dynamics [49]. Such an approach successfully describes the anisotropic flow of charm quarks observed at RHIC [49, 50, 51] (see also Ref. [52] for a recent review). Potential models have shown that, in the absence of bound states, the $Q\bar{Q}$ pairs are correlated in space [12, 13]. This correlation can be modeled classically using Langevin dynamics, including a drag force and a random force between the $Q$ (or $\bar{Q}$) and the medium as well as the forces between the $Q$ and $\bar{Q}$ described by the potential. It has been shown that a model combining an ideal hydrodynamic expansion of the medium with a description of the correlated $Q\bar{Q}$ pair dynamics by the Langevin equation may explain why, despite the fact that a deconfined medium is created at RHIC, there is only a $40 - 50\%$ suppression in the charmonium yield. The attractive potential and the finite lifetime of the system prevents
Quarkonium in Hot Medium

the complete decorrelation of some of the $Q\overline{Q}$ pairs \cite{47}. Once the system has sufficiently cooled, these residual correlations make it possible for the $Q$ and $\overline{Q}$ to form a bound state. The details of the suppression pattern depend on the model parameters, such as $T_c$ and the way recombination is implemented. The later appears to be important in the central collisions \cite{53}. However, the results are not very sensitive to the choice of the potential \cite{54}.

The above approach, which neglects quantum effects, is applicable only if there are no bound states, as it is likely to be the case for the $J/\psi$. If heavy quark bound states are present, as is probable for the $\Upsilon(1S)$, the thermal dissociation rate will be most relevant for understanding the quarkonium yield. It is expected that the interaction of a color singlet quarkonium state with the medium is much smaller than that of heavy quarks. Thus, to first approximation, medium effects will only lead to quarkonium dissociation.

5. Concluding remarks

Lattice calculations of static quark anti-quark correlators provide evidence for strong screening effects in the deconfined phase. Potential model calculations based on lattice QCD, as well as re-summed perturbative QCD calculations indicate that all charmonium states and excited bottomonium states dissolve in the deconfined medium. This leads to the reduction of the quarkonium production yield in heavy-ion collisions compared to binary-scaling of $pp$ collisions. Survival of residual correlations in the deconfined phase as well as recombination effects at lower temperatures lead to non-zero quarkonium yield in heavy ion collisions.

It turns out, that despite melting of quarkonium states corresponding meson correlation function in Euclidean time show very little temperature dependence. Almost the entire temperature dependence of the Euclidean correlation functions is due to the zero mode contribution. Statements about the existence of heavy quark bound states based on lattice correlation functions and spectral functions need to be revisited. Yet, lattice results on correlation functions are still valuable tools for constraining potential models.

One of the great opportunities of the LHC and RHIC-II heavy-ion programs is measurement of bottomonium yields. From a theoretical perspective, bottomonium is an important and clean probe for at least two reasons. First, the effective field theory approach, which provides a link to first principle QCD, is more applicable for bottomonium due to better separation of scales and higher dissociation temperatures. Second, the heavier bottom quark mass reduces the importance of recombination effects, making bottomonium a good probe of dynamical models.

[1] T. Matsui and H. Satz, Phys. Lett. B 178 (1986) 416.
[2] T. Umeda, K. Nomura and H. Matsufuru, Eur. Phys. J. C 39 S1 (2005) 9
[3] M. Asakawa and T. Hatsuda, Phys. Rev. Lett. 92, 012001 (2004);
[4] S. Datta, F. Karsch, P. Petreczky and I. Wetzorke, Phys. Rev. D 69, 094507 (2004)
[5] A. Jakovac, P. Petreczky, K. Petrov and A. Velytsky, Phys. Rev. D 75 (2007) 014506
[6] G. Aarts et al., Phys. Rev. D 76 (2007) 094513
[7] S. Digal, P. Petreczky and H. Satz, Phys. Rev. D 64 (2001) 094015
[8] C. Y. Wong, Phys. Rev. C 72 (2005) 034906.
[9] Á. Mócsy and P. Petreczky, Phys. Rev. D 73 (2006) 074007; Eur. Phys. J. C 43 (2005) 77
[10] W. M. Alberico, A. Beraudo, A. De Pace and A. Molinari, Phys. Rev. D 75 (2007) 074009.
[11] D. Cabrera and R. Rapp, Phys. Rev. D 76 (2007) 114506
[12] Á. Mócsy and P. Petreczky, Phys. Rev. D 77 (2008) 014501
[13] Á. Mócsy and P. Petreczky, Phys. Rev. Lett. 99 (2007) 211602
[14] H. Satz, J. Phys. G 32 (2006) R25
[15] M. Laine, O. Philipsen, P. Romatschke and M. Tassler, JHEP 0703, 054 (2007)
[16] M. Laine, O. Philipsen and M. Tassler, JHEP 0709 (2007) 066
[17] M. Laine, JHEP 0705 (2007) 028
[18] M. Laine, Nucl. Phys. A 820 (2009) 25C
[19] N. Brambilla, J. Ghiglieri, A. Vairo and P. Petreczky, Phys. Rev. D 78 (2008) 014017
[20] C. DeTar and U. M. Heller, Eur. Phys. J. A 41, 405 (2009)
[21] P. Petreczky, Nucl. Phys. Proc. Suppl. 140 (2005) 78
[22] Z. Fodor and S. D. Katz, arXiv:0908.3341 [hep-ph].
[23] A. Bazavov et al., Phys. Rev. D 80, 014504 (2009)
[24] M. Cheng et al., Phys. Rev. D 77, 014511 (2008); Phys. Rev. D 81, 054504 (2010)
[25] Y. Aoki, Z. Fodor, S. D. Katz and K. K. Szabo, Phys. Lett. B 643, 46 (2006)
[26] Y. Aoki et al., JHEP 0906, 088 (2009)
[27] A. Bazavov and P. Petreczky, PoS LAT2009, 163 (2009)
[28] A. Bazavov, P. Petreczky and A. Velytsky, arXiv:0904.1748 [hep-ph].
[29] O. Kaczmarek, F. Karsch, P. Petreczky and F. Zantow, Phys. Lett. B 543, 41 (2002)
[30] O. Kaczmarek, F. Karsch, F. Zantow and P. Petreczky, Phys. Rev. D 70 (2004) 074505 [Erratum-ibid. D 72 (2005) 059903]
[31] P. Petreczky and K. Petrov, Phys. Rev. D 70, 054503 (2004)
[32] O. Kaczmarek and F. Zantow, Phys. Rev. D 71, 114510 (2005)
[33] A. Bazavov, P. Petreczky and A. Velytsky, Phys. Rev. D 78, 114026 (2008)
[34] Y. Burnier, M. Laine and M. Vepsalainen, JHEP 1001, 054 (2010)
[35] N. Brambilla, J. Ghiglieri, P. Petreczky and A. Vairo, TUM-EFT 2/09
[36] RBC-Bielefeld Collaboration, work in progress
[37] P. Petreczky, Eur. Phys. J. C 43, 51 (2005)
[38] O. Jahn and O. Philipsen, Phys. Rev. D 70, 074504 (2004)
[39] T. Umeda, Phys. Rev. D 75, 094502 (2007)
[40] P. Petreczky and D. Teaney, Phys. Rev. D 73, 014508 (2006)
[41] P. Petreczky, Eur. Phys. J. C 62 (2009) 85
[42] P. Petreczky, P. Hegde and A. Velytsky, PoS LAT2009, 159 (2009)
[43] A. Andronic, P. Braun-Munzinger, K. Redlich and J. Stachel, Nucl. Phys. A 789 (2007) 334
[44] R. L. Thews, M. Schroedter and J. Rafelski, Phys. Rev. C 63, 054905 (2001)
[45] X. Zhao and R. Rapp, Phys. Lett. B 664 (2008) 253
[46] F. Karsch, D. Karzsew and H. Satz, Phys. Lett. B 637 (2006) 75
[47] C. Young and E. Shuryak, Phys. Rev. C 79 (2009) 034907
[48] D. Teaney, arXiv:0905.2433 [nucl-th].
[49] G. D. Moore and D. Teaney, Phys. Rev. C 71 (2005) 064904
[50] H. van Hees and R. Rapp, Phys. Rev. C 71, 034907 (2005)
[51] P. B. Gossiaux, V. Guiho and J. Aichelin, J. Phys. G 32, S359 (2006).
[52] R. Rapp and H. van Hees, arXiv:0903.1096 [hep-ph].
[53] C. Young and E. Shuryak, arXiv:0911.3080 [nucl-th].
[54] C. Young, private communication 2009
$F_1(r,T) \ [\text{GeV}]$

$r \ [\text{fm}]$

$T/T_c = \begin{align*}
0.82 \\
0.89 \\
0.97 \\
1.02 \\
1.09 \\
1.58 \\
1.95 \\
2.54 \\
3.29
\end{align*}$