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Charmonium states in quark-gluon plasma

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Abstract. We discuss how the spectral changes of quarkonia at T_c can reflect the 'critical' behaviour of QCD phase transition. Starting from the temperature dependencies of the energy density and pressure from lattice QCD calculation, we extract the temperature dependencies of the scalar and spin-2 gluon condensates near T_c. We also parametrize these changes into the electric and magnetic condensate near T_c. While the magnetic condensate hardly changes across T_c, we find that the electric condensate increases abruptly above T_c. Similar abrupt change is also seen in the scalar condensate. Using the QCD second-order Stark effect and QCD sum rules, we show that these sudden changes induce equally abrupt changes in the mass and width of J/ψ, both of which are larger than 100 MeV at slightly above T_c.

Keywords. Quarkonium; quark-gluon plasma.

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1. Introduction

Heavy quark system emanating from heavy ion collision is considered to be an intriguing tool to investigate the properties of quark-gluon plasma (QGP) expected to form in the early stages of a heavy ion collision [1,2]. Indeed, measurements at past SPS and recent RHIC data show nontrivial J/ψ suppression patterns that could be consistent with the original prediction that J/ψ will dissolve in QGP. However, the subject has gained a new turn recently as lattice calculations show that J/ψ will survive past the critical temperature T_c up to about $1.6T_c$, while χ_c and ψ' will dissolve just above T_c [3,4]. The notion that J/ψ will not dissolve immediately above T_c was also considered before [5], as it was known that non-perturbative aspects of QCD persist in QGP [6,7]. These results suggest that the sudden disappearance of J/ψ at T_c is not the order parameter of QCD phase transition. On the other hand, the phase transition is accompanied by sudden changes in the chiral condensate, the heavy quark potential and the energy density, and should therefore have some effect on the properties of quarkonia at T_c.

Unfortunately, the lattice results based on the maximum entropy method have large error bars, and one cannot determine the detailed properties of quarkonium above T_c, except its existence. Solving the Schrödinger equation with potential extracted from the lattice is one possibility [8–11]. However, here there are some controversies on how to extract the correct potential from the free energy calculations. Even with a chosen prescription, one has to solve the Schrödinger equation at discrete temperature steps, and it is not clear how the sudden critical behaviour is translated into the discrete temperature steps. Recently there are progresses on extracting heavy quark potential at finite temperature using perturbative QCD approaches [12,13]. At the same time, thermal perturbation alone is also not sufficient to probe the temperature region from $T_{\rm c}$ to $2.5T_{\rm c}$ [14], which is now known to be strongly interacting.

Therefore, a systematic non-perturbative method is essential to treat the phase transition region. Here, we will summarize the recent developments [15–17] to attach the problem using QCD sum rules and the QCD second-order Stark effect. The inputs are the temperature dependencies of local gluonic operators, which undergo abrupt changes across the phase transition as does the energy density.

2. The gluon matter

The order parameters for QCD phase transition are the thermal Wilson line for pure gauge theory, and the quark condensates for QCD with massless quarks. For realistic QCD none of them are order parameters in the strict sense. But simulation shows that the sudden changes in both parameters take place at the same critical temperature $T_{\rm c}$, which is determined by the susceptibilities. At the same $T_{\rm c}$, the energy density also makes a drastic change, which is somehow universal for any quark flavour [18]. The sudden changes apparent in the energy density and pressure can be translated to temperature dependencies of local gluonic operators, which are expected to embody the non-perturbative nature of QCD. The link is obtained through the energy–momentum tensor, which can be written in terms of symmetric traceless part (twist-2 gluon) and the trace part (gluon condensate) [15].

$$T^{\alpha\beta} = -\mathcal{ST}(G^{a\alpha\lambda}G^{a\beta}_{\lambda}) + \frac{g^{\alpha\beta}}{4}\frac{\beta(g)}{2g}G^a_{\mu\nu}G^{a\mu\nu}, \tag{1}$$

where we will use the LO beta function $\beta(g) = -\frac{g^3}{(4\pi)^2}(11 - \frac{2}{3}N_{\rm f})$. The thermal expectation value of this operator can be related to the lattice measurements of energy–momentum tensor through the following relation:

$$\langle T^{\alpha\beta}\rangle_{\rm T} = (\varepsilon + p)\left(u^\alpha u^\beta - \frac{1}{4}g^{\alpha\beta}\right) + \frac{1}{4}(\varepsilon - 3p)g^{\alpha\beta}. \tag{2}$$

Here, u^α is the four-velocity of the heat bath. With the following definitions,

$$\left\langle \frac{\alpha_{\rm s}}{\pi}G^a_{\mu\nu}G^{a\mu\nu}\right\rangle_{\rm T} = G_0(T) \tag{3}$$

$$\left\langle \mathcal{ST}\left(\frac{\alpha_{\rm s}}{\pi}G^{a\alpha\lambda}G^{a\beta}_{\lambda}\right)\right\rangle_{\rm T} = \left(u^\alpha u^\beta - \frac{1}{4}g^{\alpha\beta}\right)G_2(T), \tag{4}$$

we find for pure $SU(3)$ gauge theory,

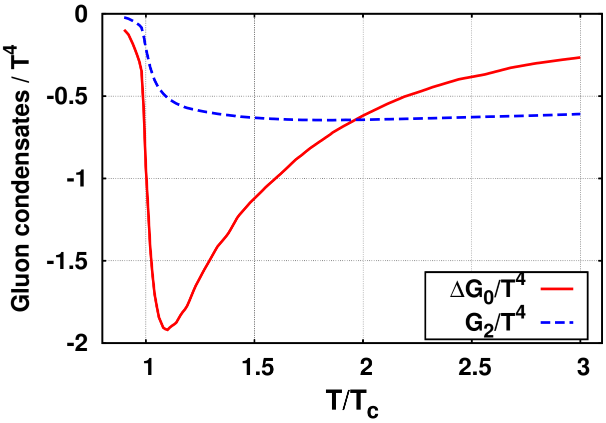

Figure 1. Gluon condensates near T_c.

$$G_0(T) = G_0^{\rm vac} - \frac{8}{11}(\varepsilon - 3p), \quad G_2(T) = -\frac{\alpha_s(T)}{\pi}(\varepsilon + p), \tag{5}$$

where $G_0^{\rm vac}$ is the value of the scalar gluon condensate in vacuum.

At low temperature, the scalar part is dominated and characterized by the non-perturbative contribution of the QCD vacuum. At high temperature, its behaviour will be dominated by the quasi-particle contribution. Just above T_c, it is still dominated by the non-perturbative contribution [7]. For the heat bath at rest, one can rewrite the thermal expectation values in terms of electric and magnetic condensate.

$$\left\langle \frac{\alpha_s}{\pi}\mathbf{E}^2 \right\rangle_{\rm T} = -\frac{1}{4}G_0(T) - \frac{3}{4}G_2(T) \tag{6}$$

$$\left\langle \frac{\alpha_s}{\pi}\mathbf{B}^2 \right\rangle_{\rm T} = \frac{1}{4}G_0(T) - \frac{3}{4}G_2(T). \tag{7}$$

Figure 2 shows the temperature dependence of G_0, G_2 or $\mathbf{E}^2$ and $\mathbf{B}^2$. One should first note that the sudden increase of energy density is translated to the anomalously large and sudden decrease in G_0, which greatly deviates from the asymptotic T^4 behaviour. In contrast, G_2 reaches the asymptotic value quickly. Similarly, one finds that there is a sudden increase in the electric condensate $\mathbf{E}^2$, while the magnetic condensate $\mathbf{B}^2$ hardly changes above T_c. This can be related to the fact that the area law behaviour of the space-time Wilson loop changes to the perimeter law above T_c, while that of the space-space Wilson loop does not [6]. The connection comes in through the operator product expansion (OPE) of rectangular Wilson loop, which was found to be expressible in terms of the electric condensates and the magnetic condensates for the space-time and space-space Wilson loops, respectively [19]. To investigate the consequences of the abrupt changes of condensates to the properties of J/ψ, we will use perturbative QCD and QCD sum rules.

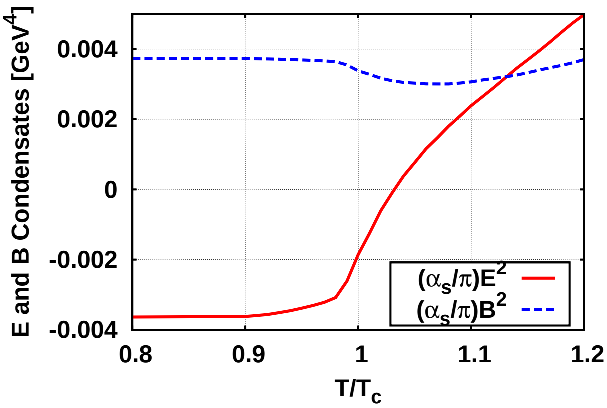

Figure 2. Electric and magnetic condensates near T_c.

Table 1. Physical processes involving two heavy quarks.

q^2	Process	Expansion parameter
0	Photoproduction of open charm	$\Lambda_{\rm QCD}^2/4m^2$
$-Q^2 < 0$	QCD sum rules for heavy quark system	$\Lambda_{\rm QCD}^2/(4m^2+Q^2)$
$m_{J/\psi}^2 > 0$	Dissociation cross-section of bound states	$\Lambda_{\rm QCD}^2/(4m^2-m_{J/\psi}^2)$

3. Quarkonium hadron interaction in QCD

Before discussing the main result, it would be useful to put the two main approaches in perspective with other QCD approaches. For that purpose, let us begin with some introduction on the propagation of heavy quarks in the QCD vacuum. The propagation of a heavy quark can be approximated by a perturbative quark propagation with a perturbative gluon insertion, which probes the non-perturbative gluon field configuration in the QCD vacuum. Hence, the full heavy quark propagator is,

$$iS^A(q) = iS(q) + iS(q)(-ig\not{A})iS(q)\ldots, \tag{8}$$

where $iS(q) = i/(\not{q} - m)$ and m is the heavy quark mass. The description in eq. (8) is valid even for $q \to 0$, because $m \gg A \sim \Lambda_{\rm QCD}$, where in the end only gauge invariant combination of the gauge field A will remain after taking the vacuum expectation value.

The starting point in discussing heavy quark–anti-quark system is the correlation function, typically defined as

$$\Pi(q) = i\int {\rm d}^4x e^{iqx}\langle T[J(x), J(0)]\rangle. \tag{9}$$

Here, the current $J = \bar{h}\Gamma h$ is the interpolating current having the quantum numbers of the meson that we want to study. The matrix element can be taken with respect to the vacuum, nucleon, nuclear matter and/or finite temperature, depending on the problem at hand. This correlation function can be thought of as an effective propagation of a system composed of a heavy quark and an antiquark and can also be approximated by combined perturbative heavy quark propagator with gluon insertions. However, since there a two heavy quarks involved, based on the operator product expansion, the propagation can typically be written in the following form:

$$\Pi(q) = \cdots + \int_0^1 \mathrm{d}x \frac{F(q^2, x)}{(4m^2 - q^2 - (2x-1)^2 q^2)^n} \langle G^n \rangle \cdots, \tag{10}$$

where $F(q^2, x)$ is a function depending on the structure of the two quark system and $\langle G^n \rangle \sim \Lambda_{\mathrm{QCD}}^{2n}$ denotes the typical gauge invariant expectation value of gluonic operator of dimension $2n$. The integration variable x can be thought of as the momentum fraction carried by one of the heavy quarks. Here, one notes that such perturbative expansion is valid when $4m^2 - q^2 \gg \Lambda_{\mathrm{QCD}}^2$. The cases where this condition is satisfied and perturbative QCD treatments are possible are summarized in table 1.

- $q^2 = 0$ corresponds to photoproduction of open heavy quarks. The expansion parameter is small, and thus perturbative description reliable.
- $q^2 < 0$ is used in the QCD sum rule analysis for heavy quark systems. The expansion parameter is small, and the results reliable. This is the basis for the successful description of heavy quark system in QCD sum rules, which predicted the masses of η_{c} to be smaller than that of J/ψ even before experiment.
- $q^2 > 0$ corresponds to the perturbative approaches for bound states. In the last line of table 1, $4m^2 - m_{J/\psi}^2 \approx (2m + m_{J/\psi})\epsilon_0$, where ϵ_0 is the binding energy of J/ψ. In QCD if $m \to \infty$, the bound state becomes Coulombic and $\epsilon_0 = m[N_{\mathrm{c}} g^2/(16\pi)]^2 \gg \Lambda_{\mathrm{QCD}}$ in the large N_{c} limit. Therefore, the expansion parameter becomes small, and perturbative description becomes possible. The QCD second-order Stark effect corresponds to this limit.

4. QCD second-order Stark effect

The perturbative QCD formalism for calculating the interaction between heavy quarkonium and partons was first developed by Peskin [20,21]. The counting scheme in this formalism is obtained in the non-relativistic limit. Because of this, the formula for the mass shift reduces to the second-order Stark effect in QCD, which was used to calculate the mass shift of charmonium in nuclear matter [22,23]. The information needed from the medium is the electric field square. Noting that the dominant change across the phase transition is that of the electric condensate, one finds that the QCD second-order Stark effect is the most natural formula to be used across the phase transition.

The QCD second-order Stark effect for the ground state charmonium with momentum space wave function normalized as $\int \frac{\mathrm{d}^3 p}{(2\pi)^3} |\psi(\mathbf{p})|^2 = 1$ is as follows:

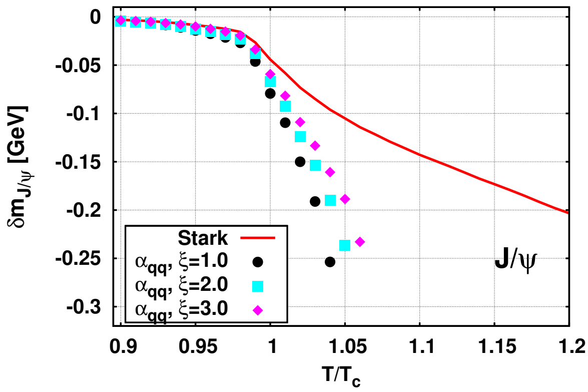

Figure 3. Mass shift from the second-order Stark effect (solid line) and the maximal mass shift obtained from QCD sum rules (points).

$$\delta m_{J/\psi} = -\frac{1}{18}\int_0^\infty \mathrm{d}k^2 \left|\frac{\partial\psi(k)}{\partial k}\right|^2 \frac{k}{k^2/m_\mathrm{c}+\epsilon}\left\langle\frac{\alpha_\mathrm{s}}{\pi}\Delta\mathbf{E}^2\right\rangle_\mathrm{T}$$
$$= -\frac{7\pi^2}{18}\frac{a^2}{\epsilon}\left\langle\frac{\alpha_\mathrm{s}}{\pi}\Delta\mathbf{E}^2\right\rangle_\mathrm{T}, \qquad (11)$$

where $k = |\mathbf{k}|$ and $\langle\frac{\alpha_\mathrm{s}}{\pi}\Delta\mathbf{E}^2\rangle_\mathrm{T}$ denotes the value of change of the electric condensate from its vacuum value so that $\delta m_{J/\psi} = 0$ in vacuum. The second line is obtained for the Coulomb wave function. Here, ϵ is the binding energy and m_c the charm quark mass. These parameters are fit to the size of the wave function obtained in the potential model [24], and to the mass of J/ψ assuming it to be a Coulombic bound state in the heavy quark limit [20]. The fit gives $m_\mathrm{c} = 1704$ MeV, $a = 0.271$ fm and $\alpha_\mathrm{s} = 0.57$. Few comments are in order. The minus sign in eq. (11) is a model independent result and follows from the fact that the second-order Stark effect is negative for the ground state. Second, the important parts of the formula are the Bohr radius a, the binding energy ϵ and the temperature dependence of the electric condensate squared. In the formula, the factor of a^2 follows from the dipole nature of the interaction, and the binding energy from the inverse propagator, characterizing the separation scale [20,25]. Therefore, the actual value of the mass shift does not depend much on the form of the wave function as long as the size of the wave function is fixed.

The solid line in figure 3 shows the mass shift obtained from the second-order Stark effect with the extracted lattice input for the electric condensate shown in figure 2. As the formula is based on multipole expansion, it will break down if the higher-dimensional condensates become large. At this moment, there is no lattice calculation for the temperature dependence of higher-dimensional gluonic operators. However, as discussed before, up to $1.1T_\mathrm{c}$, the temperature dependence is dominated by the sudden increase in the electric condensate with little change in the magnetic counterpart. Moreover, the OPE in the QCD sum rules were found to be reliable up to $1.05T_\mathrm{c}$ [15]. Therefore, our result should at least be valid up to the same temperature. Above that temperature, the change of the electric condensate amounts to more than 100% of its vacuum value and higher-order corrections should be important. As can be seen in figure 1, the mass reduces abruptly above T_c and becomes smaller by about 100 MeV at $1.05T_c$, reflecting the critical behaviour of the QCD phase transition.

5. QCD sum rule result

5.1 *General remarks*

The QCD sum rule for quarkonium was found to be very reliable at zero temperature [26–28]. This is because the expansion parameter of the OPE for the correlation function for the heavy quark current–current correlation appearing in eq. (9) for the space-like momentum $q^2 = -Q^2 < 0$, can be typically written as

$$\Pi(q) = \sum_n \frac{C_n(q, m_c)}{(Q^2 + 4m_c^2)^n} \langle G^n \rangle. \tag{12}$$

C_n's are the Wilson coefficients and G^n the gluon operator of dimension $2n$. Assuming that the typical scale of the gluon operators is $\Lambda_{\rm QCD}$, one notes that the OPE is an asymptotic expansion in $\Lambda_{\rm QCD}^2/(Q^2 + 4m_c^2)$. Therefore, the OPE for the correlation function can be reliably estimated even at $Q^2 = 0$. Moreover, for the heavy quark system, only gluon operators appear, for which reliable estimates can be made at least for the lowest dimension 4 operator.

The generalization of the sum rule approach to finite temperature involves few additional considerations. As for the OPE side of the sum rule, the effects of the temperature can be put either into the Wilson coefficient C_n or into the temperature dependence of the operators $\langle G^n \rangle$. If the temperature is larger than the separation scale $Q^2 + 4m_c^2$, one calculates the temperature effect into the Wilson coefficient [29]. If the temperature is low, all the temperature effects can be put into the temperature-dependent operators [30]. In this case, the new expansion parameter in the OPE will be $(\Lambda_{\rm QCD} + cT)^2/(Q^2 + 4m_c^2)$, where c is some constant. Whether such approximation is valid or not can be directly checked by looking at the convergence of the OPE at finite temperature.

An additional ingredient at finite temperature is that unlike at zero temperature, where only scalar gluon operators appear, operators with Lorentz indices have to be added. Therefore, up to dimension 4, both the scalar gluon (G_0) and twist-2 gluon (G_2) operators contribute. Since we have extracted the temperature dependence of these operators from the lattice, no ambiguities exist for the OPE side of the sum rule for a consistent analysis. The generalization of the OPE for J/ψ to dimension 6 operators in medium was obtained before [31].

The OPE for the correlation function of heavy quark currents is related to the phenomenological side via the dispersion relation

$$\Pi(q^2) = \int \frac{\rho(s)}{s - q^2} {\rm d}s. \tag{13}$$

In the sum rule, one has to assume the form of the spectral density $\rho(s)$. Here, one notes that lattice calculations suggest that the peak structure in the spectral density persists above $T_{\rm c}$, although the resolution is not good enough for a critical study. Therefore, we can assume a relativistic Breit–Wigner form for the spectral density.

$$\rho(s) = \frac{1}{\pi}\frac{f\sqrt{s}\Gamma}{(s-m_{J/\psi}^2)^2+s\Gamma^2}. \tag{14}$$

The sum rule is obtained by calculating the moments of the correlation function,

$$M_n(Q^2) \equiv \frac{1}{n!}\left(\frac{\mathrm{d}}{\mathrm{d}q^2}\right)^n \Pi(q). \tag{15}$$

For the OPE, the moments will have contributions from scalar G_0 and spin-2 gluon operator G_2 up to dimension 4. For the phenomenological side, the moment is obtained through the dispersion relation,

$$M_n(Q^2)_{\rm phen} = \int \frac{\rho(s)}{(s+Q^2)^{n+1}}\mathrm{d}s. \tag{16}$$

The sum rule is obtained by taking the ratio between moments and comparing the OPE to the phenomenological side.

$$\left.\frac{M_{n-1}}{M_n}\right|_{\rm OPE} = \left.\frac{M_{n-1}}{M_n}\right|_{\rm phen}. \tag{17}$$

In the $\Gamma \to 0$ limit, the right-hand side of eq. (17) becomes

$$\left.\frac{M_{n-1}}{M_n}\right|_{\rm phen} = m_{J/\psi}^2 + Q^2. \tag{18}$$

Therefore, after choosing some value for Q^2, typically 0 or some multiple of $4m_{\rm c}^2$, the temperature dependence of the OPE leads to the temperature dependence of the mass of J/ψ. If we allow for a finite value of the width Γ, the relation in eq. (17) reduces to a constraint for the temperature dependence of the mass and width of J/ψ due to the temperature dependence of the condensates.

5.2 *Results for J/ψ*

Using the moments for the correlation function, we can extract a constraint for the temperature dependence for the mass and width from the temperature dependence of the condensates. We find that if there is no change in the width, the mass critically decreases by few hundreds MeV slightly above $T_{\rm c}$. Similarly, if there is no change in mass, the width will increase by a similar amount [15]. The dots in figure 3 represent the maximum mass shift assuming no change in the width. This is a direct consequence of the critical change of the scalar condensate proportional

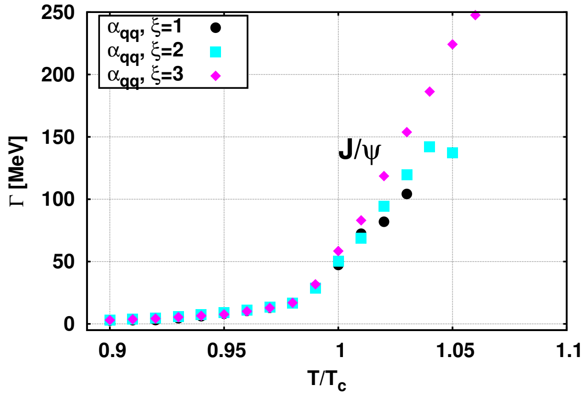

Figure 4. Thermal width of J/ψ obtained from the second-order Stark effect and QCD sum rule constraint.

to $\varepsilon - 3p$. As can be seen in the figure, the mass shift obtained from the second-order Stark effect is almost the same as the maximum mass shift obtained in the sum rule up to $T_{\rm c}$ and then becomes smaller. The mass shift at $T_{\rm c}$ is about -50 MeV. In the QCD sum rules, only a constraint for the combined mass shift and thermal width could be obtained. This constraint can be crudely summarized as follows:

$$-\Delta m + \Gamma_{\rm T} \simeq 80 + 17 \times (T - T_{\rm c}) \quad [\text{MeV}], \tag{19}$$

within the temperature range from $T_{\rm c}$ to $1.05T_{\rm c}$. Therefore, the difference between the Stark effect and the maximum mass shift obtained from QCD sum rules above $T_{\rm c}$ in figure 3 could be attributed to the non-perturbative thermal width at finite temperature. In figure 4, we plot the thermal width obtained from combining the constraint in eq. (19) with the mass shift obtained from the QCD second-order Stark effect. As can be seen in the figure, the thermal width at $1.05T_{\rm c}$ becomes larger than 100 MeV. Such a width slightly above $T_{\rm c}$ is larger than that estimated from a perturbative LO and NLO QCD matrix element calculation together with an assumption of weakly interacting quarks and gluons with thermal masses [32,33], but smaller than a recent phenomenological estimate [11].

The mass of quarkonium at finite temperature was also investigated in the potential models [9,10], where the mass was found to decrease at high temperature. However, the potential has to be extracted from the lattice at each temperature and hence a more detailed investigation is needed to identify the critical behaviour of J/ψ in the temperature region near $T_{\rm c}$.

5.3 *Corrections from dynamical quarks*

To consider the quark effects, first we consider the quark operators appearing in the OPE. We can neglect the light quark contribution to the OPE, because the light quark operators appear in the OPE at order $\alpha_s^2(q^2)$: This is why the light quark condensate can be neglected in the sum rules for heavy quark system. On the other hand, thermal heavy quarks that directly couple to the heavy quark current contribute to the OPE at leading order. This is different from the heavy quark condensates that are perturbatively generated from the gluon condensates, and contribute to the OPE through gluon condensates, whose Wilson coefficients are calculated in the momentum representation. The direct thermal quark contributions are called the scattering terms. However, similar terms also appear in the phenomenological spectral density, which also has free charm quark mode that is not coupled with a light quark in the form of a D meson above T_c as recently studied in [34]. Therefore, the scattering term will cancel out between the OPE and the phenomenological spectral density in the QCD sum rule analysis for the deconfined medium.

Second, the gluon condensates themselves can have a different temperature dependence in the presence of dynamical quarks. As discussed before, the important input for the mass and width change is the temperature dependence of gluon condensates in figure 1; in particular the dominant contribution comes from the temperature dependence of G_0. For that purpose, we note that the trace of the energy–momentum tensor to leading order is given as

$$T^{\mu}_{\mu} = -\left(\frac{11-2/3N_f}{8}\right)\left\langle\frac{\alpha_s}{\pi}G^a_{\mu\nu}G^{a\mu\nu}\right\rangle + \sum_q m_q\langle\bar{q}q\rangle. \tag{20}$$

Therefore, we start from the lattice calculation of the trace of the energy–momentum tensor for the full QCD with realistic quark masses as given in [35]. Then, we subtract the fermionic part of the trace anomaly, which was also shown in the literature, from the total. Then we divide the result for the relevant prefactor with $N_f = 3$ multiplying the gluon condensate as given in eq. (20). Since the critical temperature T_c differs, we compared it as a function of T/T_c in which $T_c = 196$ MeV for the full QCD case [35]. As can be seen in figure 5, the magnitude of the resulting change near the critical temperature are remarkably similar between the full and pure gluon QCD; although the slope at T_c is milder for full QCD as a consequence of rapid cross-over transition instead of a first-order phase transition. Since the change of the condensate sets in at a lower T/T_c in the full QCD case, the mass and width of charmonia might start varying at a lower temperature in the realistic case than in the pure glue theory. Therefore, we believe our main argument and the quantitative result will not be altered even in the realistic situation.

5.4 *Additional comments*

We have shown that the sudden changes in the energy density lead to sudden changes in the spectral properties of J/ψ across the phase transition. This suggests that the sudden mass shift could effectively be an 'order parameter' of QCD phase transition. The expected mass shift might be too small for the present experimental resolution. However, with expected upgrades at RHIC and expected resolutions at LHC, such shifts from heavy ion collisions could be systematically studied.

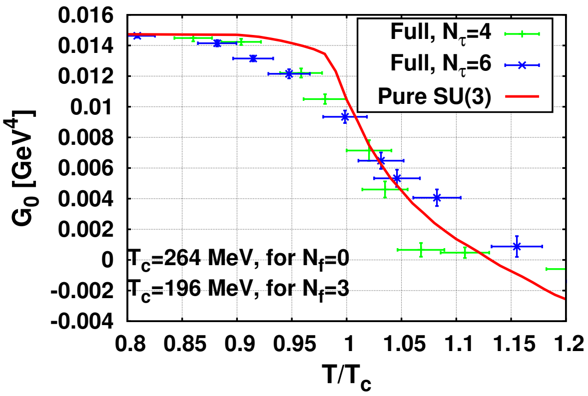

Figure 5. Comparison of the scalar gluon condensate in the pure gauge theory with the one extracted from full lattice QCD. Horizontal error bars in the full QCD case are drawn by assuming 2% uncertainty in the conversion from the lattice units to physical temperature.

Moreover, the shift originates from the dipole-type of interaction, the quarkonium with the medium. Therefore, for larger size system the expected mass shift is larger.

A precursor phenomena will also be observable at nuclear medium [16,23,36]. Such effects could be observed in the anti-proton project at FAIR GSI.

6. Summary

We have shown that the critical mass shift and width increase of quarkonia slightly above T_c could effectively be an 'order parameter' of QCD phase transition. The expected shift for J/ψ is small and will be a challenge for LHC. However, larger shift is expected for χ_c at its formation temperature slightly above T_c. Therefore, direct measurements and confirmation are possible. The changes will also lead to changes in production ratios in the statistical model and J/ψ suppression effects, which need further detailed studies.

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