Charmonium in Nuclear Matter and Chiral Symmetry

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Changes in the masses of charmonium states in nuclear matter as a result of the modification of QCD vacuum in nuclear medium are studied in the perturbative QCD approach. For the leading-order effect due to change of gluon condensate, we use the leading-order QCD formula. The higher-twist effect that is related to the partial restoration of chiral symmetry, we use instead a hadronic model that includes the effect due to change of quark condensate in nuclear medium. It is found that although the mass of J/ψ decreases slightly in nuclear matter, those of $\psi(3686)$ and $\psi(3770)$ states are reduced appreciably. Experimental study of the mass shift of charmonium states in nuclear matter can thus provide valuable information on how QCD vacuum changes in nuclear medium.

§1. Introduction

Understanding hadron mass shifts in nuclear medium and/or at finite temperature can provide valuable information about the QCD vacuum.[1)–3)] It is also relevant phenomenologically to the interpretation of experimental results from relativistic heavy ion collisions,[4)] in which a hot dense matter is formed during the collisions. Previous studies have been largely concerned with hadrons that consist of only light quarks.[3)] Only recently were there studies of the in-medium masses of hadrons made also of heavy charm quarks. Using either the QCD sum rules[5),6)] or the quark-meson coupling model,[7)] it has been found that mass of D meson, which is made of a charm quark and a light quark, is reduced significantly in nuclear medium as a result of decrease of the light quark condensate. For J/ψ, which consists of a charm and anticharm quark pair, both the QCD sum rules analysis[8)] and the LO perturbative QCD calculation[9),10)] show that its mass is reduced slightly in nuclear matter mainly due to the reduction of the gluon condensate in nuclear medium.

The change of hadron masses at finite temperature is best studied using the lattice gauge theory, as it treats the non-perturbative aspect of QCD most reliably. Recent lattice gauge calculations at finite temperature with dynamical quarks have shown that even below critical temperature the interquark potential at large separation approaches an asymptotic value $V_\infty(T)$ that decreases with increasing temperature.[11)] This transition from a linearly rising interquark potential in free space to a saturated one at finite temperature is due to decrease in the strength of string tension and formation of $\bar{Q}q$ and $\bar{q}Q$ pairs, where q denotes a light quark,

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when the separation of two heavy quarks (Q) becomes large. Decrease in $V_\infty(T)$ can thus be interpreted as a reduction of the mass m_H of open heavy quark meson ($\bar{Q}q$ or $\bar{q}Q$), such as the D meson mass, at finite temperature.[12),13)] Furthermore, the decrease of m_H seems to be a consequence of the reduction in the constituent mass of light quark as the temperature dependence of $V_\infty(T)$ is similar to that of the chiral condensate $\langle\bar{q}q\rangle$.[14)] This relation between the mass m_H and the chiral order parameter also follows naturally from the heavy quark symmetry.[15)] From the solution of Schrödinger equation with the finite temperature interquark potential obtained from the lattice QCD calculation, it has also been found that masses of charmonium states are reduced at finite temperature.[16),17)]

At finite density, lattice gauge calculations are at present not feasible for studying the heavy quark potential or the mass m_H of open heavy quark meson. Masses of these heavy quark systems are, however, expected to change appreciably in nuclear medium. Model independent estimates have shown[18),19)] that condensates of the lowest dimensional operators $\langle\frac{\alpha_s}{\pi}G^2\rangle$ and $\langle\bar{q}q\rangle$ decrease, respectively, by 6% and 30% at normal nuclear matter, which are significant changes expected only near the phase transition at finite temperature.[20)] As in the case of finite temperature, the reduction of gluon condensate leads to a softening of the confining part of interquark potential,[21)] while the decrease of quark condensates implies a drop of the open heavy quark meson mass or the asymptotic value V_∞ of heavy quark potential. Both are expected to lead to nontrivial changes in the binding energies of charmonium states $\psi(3686)$ and $\psi(3770)$, as their wave functions are sensitive to both the confining part and the asymptotic value of interquark potential.

In this paper, we present results on the mass shifts of $\psi(3686)$ and $\psi(3770)$ due to changes in the gluon and quark condensates in nuclear medium.[22)] The effect of gluon condensate is determined using the leading-order QCD formula, which was developed in Refs. 9) and 23) and has been used to study the J/ψ mass in medium.[10)] The effect due to change in quark condensates is difficult to calculate using the quark and gluon degrees of freedom as they appear as higher twist effects in the operator product expansion.[24),25)] However, its dominant effect on a heavy quark system is to reduce V_∞ as a result of decrease of the D meson in-medium mass, about 50 MeV in normal nuclear matter due to the 30% reduction in the quark condensate.[5)–7),26)] Therefore, we can study the effect of changing quark condensate on charmonium states at finite density using a hadronic model to calculate their mass shifts due to the change of D meson in-medium mass. Combining effects from changes in the gluon condensate and in m_D, we find that both $\psi(3686)$ and $\psi(3770)$ masses are reduced appreciably at normal nuclear matter density.

§2. QCD vacuum in nuclear matter

The lowest dimensional QCD operators that characterize the non-perturbative nature of QCD vacuum are the quark and gluon condensates. Their expectation values in the vacuum are known to be large,[27)] i.e.,

$$\left\langle\frac{\alpha_s}{\pi}F^2_{\mu\nu}\right\rangle \sim 1.5\ \mathrm{GeV}\cdot\mathrm{fm}^{-3},$$

$$\langle \bar{q}q \rangle \sim 2 \text{ fm}^{-3}. \tag{2·1}$$

The gluon condensate can be written as the difference between the magnetic $B^2 = \frac{1}{2}F_{ij}^2$ and the electric $E^2 = F_{0i}^2$ condensate, which contribute equally to the gluon condensate in vacuum, i.e.,

$$\left\langle \frac{\alpha_s}{\pi} B^2 \right\rangle = -\left\langle \frac{\alpha_s}{\pi} E^2 \right\rangle = \frac{1}{2}\left\langle \frac{\alpha_s}{\pi} F_{\mu\nu}^2 \right\rangle. \tag{2·2}$$

To understand the above relation, we note that in the Euclidean formulation of QCD at zero temperature, such as in the lattice QCD, the electric condensate has the same expectation value as the magnetic one due to the space and time symmetry in the Euclidean space.[20),28)] The above relation thus follows naturally as the electric condensate in the Euclidean space is defined with a minus sign relative to its counterpart in the Minkowski space, while the magnetic condensate is defined with the same sign in both spaces.[27)]

In nuclear matter, both non-perturbative quark and gluon field configurations are expected to change appreciably. Model-independent studies have shown that the average gluon and quark condensate values decrease by 6% and 30%, respectively. These studies are based on linear density approximation and nucleon expectation values of the quark and gluon condensates, which are known, respectively, from the experimentally measured π-N sigma term and the nucleon expectation value of the trace anomaly relation.[29)]

Both the electric and magnetic parts of the gluon condensate in nuclear matter can be estimated using the twist-2 gluon operator,

$$\langle N(p)|\mathcal{ST}F_\mu^\alpha F_{\alpha\nu}|N(p)\rangle = \left(p_\mu p_\nu - \frac{1}{4}m_N^2 g_{\mu\nu}\right)2A_2(g), \tag{2·3}$$

where $A_2(g)$ is the second moment of the gluon distribution in a nucleon and has a value of about 0.45 at the renormalization scale of 1 to 2 GeV. In the linear density approximation, we then have

$$\begin{aligned} \left\langle \frac{\alpha_s}{\pi} E^2 \right\rangle_{\text{n.m.}} &= \left(\frac{4}{9}m_N m_N^0 + \frac{3}{2}m_N^2 \frac{\alpha_s}{\pi} A_2\right)\frac{\rho}{2m_N}, \\ \left\langle \frac{\alpha_s}{\pi} B^2 \right\rangle_{\text{n.m.}} &= -\left(\frac{4}{9}m_N m_N^0 - \frac{3}{2}m_N^2 \frac{\alpha_s}{\pi} A_2\right)\frac{\rho}{2m_N}. \end{aligned} \tag{2·4}$$

In the above, ρ and m_N are the nuclear density and nucleon mass, respectively. The mass $m_N^0 \sim 0.75$ GeV is the nucleon mass in the chiral limit[30)] and is obtained by taking the nucleon expectation value of the trace anomaly relation $T_\mu^\mu = -\frac{9}{8}\frac{\alpha_s}{\pi}F_{\mu\nu}^2$. Because of the small factor of $\frac{\alpha_s}{\pi}$ in the second terms in Eq. (2·4), change of the gluon condensate in nuclear medium is dominated by contributions from the first terms.

§3. LO QCD calculation

The mass shift of charmonium states in nuclear medium can be evaluated in the perturbative QCD when the charm quark mass is large, i.e., $m_c \to \infty$. In this limit, one can perform a systematic operator product expansion (OPE) of the charm quark-antiquark current-current correlation function between the quark bound states by taking the separation scale μ to be the binding energy of the charmonium.[9),23),31)] The forward scattering matrix element of the charm quark bound state with a nucleon then has the following form:

$$T(q^2 = m_\psi^2) = \sum_n \frac{C_n}{\mu^n} \langle \mathcal{O}_n \rangle_N. \tag{3·1}$$

Here, C_n is the Wilson coefficient evaluated with the charm quark bound state wave function and $\langle \mathcal{O}_n \rangle_N$ is the nucleon expectation value of local operators of dimension n.

For heavy quark systems, there are only two independent lowest dimension operators; the gluon condensate ($\langle \frac{\alpha_s}{\pi} G^2 \rangle$) and the condensate of twist-2 gluon operator multiplied by α_s ($\langle \frac{\alpha_s}{\pi} G_{\alpha\mu} G_\nu^\alpha \rangle$). These operators can be rewritten in terms of the color electric and magnetic fields: $\langle \frac{\alpha_s}{\pi} E^2 \rangle$ and $\langle \frac{\alpha_s}{\pi} B^2 \rangle$. Since the Wilson coefficient for $\langle \frac{\alpha_s}{\pi} B^2 \rangle$ vanishes in the non-relativistic limit, the only contribution is thus proportional to $\langle \frac{\alpha_s}{\pi} E^2 \rangle$, which is similar to the usual second-order Stark effect.

The QCD second-order Stark effect[10)] on charmonium masses can be evaluated by multiplying the leading term in Eq. (3·1) by the nuclear density ρ_N, and this gives

$$\Delta m_\psi(\epsilon) = -\frac{1}{9} \int dk^2 \left| \frac{\partial \psi(k)}{\partial \boldsymbol{k}} \right|^2 \frac{k}{k^2/m_c + \epsilon} \left\langle \frac{\alpha_s}{\pi} E^2 \right\rangle_N \frac{\rho_N}{2m_N}. \tag{3·2}$$

In the above, $\langle \frac{\alpha_s}{\pi} E^2 \rangle_N \sim 0.5$ GeV^2 is the nucleon expectation value of the color electric field obtained from Eq. (2·4), and $\epsilon = 2m_c - m_\psi$ is the binding energy of charmonium. In Ref. 9), the LO mass shift formula for J/ψ was derived in the large charm quark mass limit. As a result, the wave function $\psi(k)$ is Coulombic, and the mass shift is expressed in terms of the Bohr radius a_0 and the binding energy $\epsilon_0 = 2m_c - m_{J/\psi}$. This might be a good approximation for J/ψ but is not realistic for excited charmonium states as Eq. (3·2) involves the derivative of the wave function, which measures the dipole size of the system. We have thus taken the wave functions of charmonium states to be Gaussian with the oscillator constant β determined by their squared radii $\langle r^2 \rangle = 0.47^2$, 0.96^2 and 1 fm^2 for J/ψ, $\psi(3686)$ and $\psi(3770)$, respectively, as obtained from the potential models.[32)] This gives $\beta = 0.52$, 0.39 and 0.37 GeV if we assume that these charmonium states are in the $1S$, $2S$ and $1D$ states, respectively. Using these parameters and the QCD parameters $\alpha_s = 0.84$ and $m_c = 1.95$, which are fixed by the energy splitting between J/ψ and $\psi(3686)$ in free space,[9)] we find that the mass shifts at normal nuclear matter density obtained from the LO QCD formula Eq. (3·2) are -8, -100 and -140 MeV for J/ψ, $\psi(3686)$ and $\psi(3770)$, respectively.

Although higher twist effects on charmonium masses are expected to be non-trivial, the result for J/ψ is consistent with those from other non-perturbative QCD studies, such as the QCD sum rules[8),25)] and the effective potential model,[33)–35)] which are all based on the dipole interactions between quarks in the charmonium and those in the nuclear matter. To go beyond the leading order in OPE, we need to calculate the contributions from higher dimensional operators in Eq. (3·1), which include light quark operators. Explicit calculations from QCD sum rules for J/ψ up to dimension 6 operators[25] show that the effect due to change in the condensates of light quark operators at dimension 6, which include $\langle \bar{q}\Gamma q\bar{q}\Gamma q\rangle$ and $\langle \bar{q}DGq\rangle$,[24),25] is unimportant for the mass shift of J/ψ. However, such calculation cannot be easily generalized to the excited charmonium states $\psi(3686)$ and $\psi(3770)$, where the sum rules do not exist even in the vacuum. On the other hand, the higher twist effect due to light quark operators can be estimated by considering coupling of the charmonium to the $\bar{D}D$ states as in the potential model for charmonium states.[32] Therefore, instead of summing up the non-convergent contributions from the change in the light quark condensates in the OPE of Eq. (3·1), we estimate its contribution by evaluating the charmed meson one-loop effect on the mass of charmonium using the in-medium D meson mass predicted from the QCD sum rules[5),6] or the quark-meson coupling model.[7]

§4. Estimate of higher twist effect using an effective Lagrangian

Following the studies in Ref. 36) on ρ-π interactions and in Ref. 37) on ϕ-K interactions, we use the following Lagrangian for interacting charmonium ψ and D meson:

$$\mathcal{L} = \frac{1}{2}\left(|D_\mu \boldsymbol{D}|^2 - m_D^2|\boldsymbol{D}|^2\right) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_\psi^2\psi_\mu\psi^\mu, \tag{4·1}$$

where $F_{\mu\nu} = \partial_\mu\psi_\nu - \partial_\nu\psi_\mu$ is the charmonium field strength, $D_\mu = \partial_\mu - i2g_{\psi DD}\psi_\mu$ and $\boldsymbol{D} = (D^0, D^+)$.

The coupling constant $g_{\psi DD}$ can be determined using the $3P0$ model.[38] In this model, the coupling constant is given by the overlap integral between the relative quark wave function of the charmonium and that of the two outgoing charmed mesons, multiplied by a coupling parameter γ which characterizes the probability of producing a light quark-antiquark pair in the $^3P_0$ state. The result can be read off from Refs. 39) and 40) and is given by

$$g_{\psi DD}^2(q) = \gamma^2\pi^{3/2}\frac{m_\psi^3}{\beta_D^3}f_\psi(q^2,r)e^{-\frac{q^2}{2\beta_D^2(1+2r^2)}}, \tag{4·2}$$

where q is the three-momentum of D mesons in the ψ rest frame and $r = \beta/\beta_D$ with β_D (β) being the oscillator constant for D meson (ψ) wave function. The same values of β are used for charmonium states as in the LO QCD calculation. The values for γ and β_D are taken to be 0.35 and 0.31 GeV, respectively, to reproduce both the decay width of $\psi(3770)$ to $D\bar{D}$ and the partial decay width of $\psi(4040)$ to DD, DD^* and D^*D^*.[39),41] The function $f_\psi(q^2,r)$ denotes

$$f_{J/\psi}(q^2,r) = \frac{2^6r^3(1+r^2)^2}{(1+2r^2)^5},$$

$$f_{\psi(3686)}(q^2,r)=\frac{2^5(3+2r^2)^2(1-3r^2)^2}{3(1+2r^2)^7}\left(1-\frac{2r^2(1+r^2)}{(1+2r^2)(3+2r^2)(3r^2-1)}\frac{q^2}{\beta_D^2}\right)^2,$$

$$f_{\psi(3770)}(q^2,r)=\frac{2^9 5r^7}{3(1+2r^2)^7}\left(1-\frac{(1+r^2)}{5(1+2r^2)}\frac{q^2}{\beta_D^2}\right)^2, \qquad (4\cdot3)$$

for the three charmonium states.

Because of its momentum dependence, $g_{\psi DD}(q)$ takes into account the form factor at the ψDD vertex and allows also the coupling of the charmonium to off-shell D mesons. For $\psi(3770)$, it can decay to $D\bar{D}$ in free space, and its on-shell coupling constant is $g_{\psi(3770)DD}(q=(m_\psi^2/4-m_D^2)^{1/2})=15.4$. The coupling constants at $q=0$ are 15.3, 18.7 and 16.8 for J/ψ, $\psi(3686)$ and $\psi(3770)$, respectively. The value for J/ψ coupling to D mesons is slightly larger than that estimated by the vector meson dominance model[42),43)] and by the QCD sum rules.[44)] As the D meson momentum increases, its coupling constant to J/ψ has a simple exponential fall off due to the $1S$ quark wave function of J/ψ. In contrast, coupling constants of excited charmonium states $\psi(3686)$ and $\psi(3770)$ to D mesons fall off exponentially with D meson momentum but vanish at certain q^2 as a result of the nodes in the $2S$ or $1D$ wave function of the excited charmonium states.[39),41)]

Similar to the method introduced in Ref. 37), we have used the above Lagrangian to evaluate the self-energy $\Pi_{\mu\nu}(k)$ of charmonium due to the D meson loop. After performing the energy integral in the rest frame of ψ, i.e., $k=(m_\psi,0)$, the invariant part of $\Pi_{\mu\nu}(k)=(k_\mu k_\nu-k^2g_{\mu\nu})\Pi(k)$ then has the following form:

$$\Pi(k)=\frac{1}{6\pi^2}\mathcal{P}\int dq^2 g_{\psi DD}^2(q^2)\left[\frac{q}{\sqrt{m_D^{*2}+q^2}}\left(\frac{4q^2}{m_\psi^2-4m_D^{*2}-4q^2}+3\right)\right.$$
$$\left.-(m_D^*=m_D)\right], \qquad (4\cdot4)$$

where m_D^* is the in-medium D meson mass and $\mathcal{P}$ denotes that only the principle value of the integral is evaluated. The subtracted term in the above equation is a renormalization constant which is determined by requiring $\Pi(k^2=m_\psi^2)=0$ when $m_D^*=m_D$. This ensures that the D meson loop does not contribute to the real part of charmonium self energy in free space. The mass shift of charmonium at finite density is then given by $\Delta m_\psi=\Pi(k^2=m_\psi^2)$.

In Fig. 1, we show the mass shifts of charmonium states as functions of m_D^*. It is seen that the mass shift of $\psi(3770)$ is negative for small negative shift of D meson mass but becomes positive when the D meson mass drop is large. In contrast, the mass shift of $\psi(3686)$ is negative for all negative mass shifts of D meson. This difference can be understood from Eq. (4·4), where the integral is a convolution of the form factor $g_{\psi DD}^2(q^2)$ with the terms in the square bracket, which are singular when $q^2=m_\psi^2/4-m_D^{*2}$ and $q^2=m_\psi^2/4-m_D^2$. The integrand thus changes its sign whenever the D meson momentum q passes through these singularities and finally becomes negative when q^2 is larger than any of the singularities, which correspond to the energies of the virtual intermediate D meson states. As in second-order perturbation

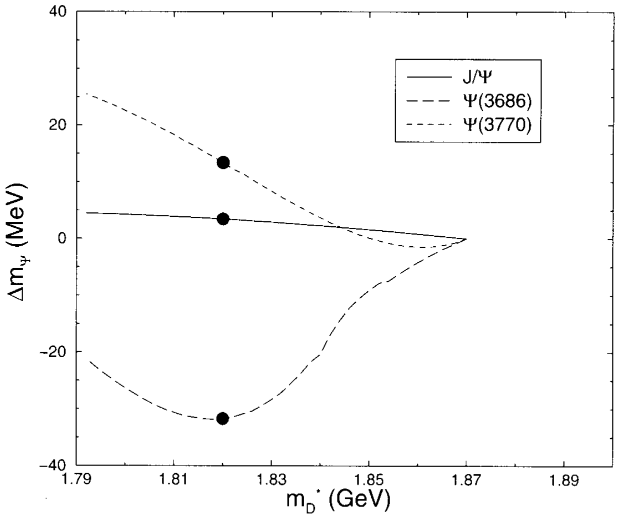

Fig. 1. Mass shifts of charmonium states J/ψ (solid curve), $\psi(3686)$ (long dashed curve) and $\psi(3770)$ (short dashed curve) as functions of D meson in-medium mass m_D^*. The circle indicates the expected mass shifts at normal nuclear matter density.

theory, the contribution is attractive when the energy of the intermediate state is larger than the charmonium mass. Since the form factor decreases exponentially with q^2 and can even be zero, the large negative contribution expected for a constant form factor is suppressed, leading thus to an increase of the $\psi(3770)$ mass when $m_D - m_D^* \geq 10$ MeV. On the other hand, the singularity of the integrand in Eq. (4·4) for the case of $\psi(3683)$ occurs only when $2m_D^*$ falls below its mass and therefore has only a small positive contribution when q^2 is very small, leading to a reduction of its mass for any D meson mass shift. For J/ψ, we find that its mass only increases slightly with dropping D meson mass and depends weakly on m_D^*. For $m_D - m_D^* = 50$ MeV, which is the expected mass shift of D meson at normal nuclear matter density, the mass shift of J/ψ is about 3 MeV. This result is consistent with that from the QCD sum rules[25] and is also expected from the potential model,[32] where the J/ψ wave function has only a small $D\bar{D}$ component. We note that the density dependence of the mass shifts of charmonium states, particularly the excited ones, is nonlinear if we use a linearly density-dependent D in-medium meson $m_D^* = m_D - 50\ \rho/\rho_0$ MeV in the denominator of Eq. (4·4). On the other hand, the mass shift obtained from the LO QCD formula in Eq. (3·2) depends linearly on the nuclear density.

Adding the mass shift from the D meson loop effect to the result from the LO QCD calculation, we find that masses of charmonium states are changed by the following amount at normal nuclear matter density:

$$\begin{aligned}\Delta m_{J/\psi} &= -8 + 3 \text{ MeV},\\ \Delta m_{\psi(3686)} &= -100 - 30 \text{ MeV},\\ \Delta m_{\psi(3770)} &= -140 + 15 \text{ MeV},\end{aligned} \tag{4·5}$$

where the first number represents the shift from the LO QCD while the second number is from the D meson loop. The above results thus show that masses of the excited charmonium states are reduced significantly in nuclear matter, largely due to the non-trivial decrease of the gluon condensate in nuclear medium.

§5. Summary

In summary, we have studied the mass shifts of charmonium states in nuclear matter in the perturbative QCD approach. The leading-order QCD formula is used to evaluate the effect due to the change of gluon condensate in nuclear matter, while a hadronic model that takes into account the interaction of charmonium states with D mesons is used to calculate the higher-twist effect due to the partial restoration of chiral symmetry. We find that while the mass of J/ψ in normal nuclear matter decreases by less than 10 MeV, those of $\psi(3686)$ and $\psi(3770)$ states are reduced by more than 100 MeV. Our results thus demonstrate that the masses of $\psi(3686)$ and $\psi(3770)$ are sensitive to both the confining part of the interquark potential and the $D\bar{D}$ threshold, which, according to recent lattice gauge calculations, are related to the chiral symmetry breaking of the QCD vacuum. The mass shifts we find for both $\psi(3686)$ and $\psi(3770)$ in nuclear medium are large enough to be observed in experiments involving $\bar{p}$-A annihilation as proposed in the future accelerator facility at the German Heavy Ion Accelerator Center (GSI).[45] In these experiments, $\psi(3770)$ and $\psi(3686)$ produced inside a heavy nucleus will be studied via the dilepton spectrum emitted from their decays. The observation of shifts in their masses in these experiments would give us valuable information on non-trivial changes of the QCD vacuum in nuclear medium and on the origin of masses in QCD.

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