Mass modification of D-meson at finite density in QCD sum rule

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Abstract

We evaluate the mass shift of isospin-averaged D-meson in the nuclear medium. Borel-transformed QCD sum rules are used to describe an interaction between the D-meson and a nucleon by taking into account all the lowest dimension-4 operators in the operator product expansion (OPE). We find at normal matter density the D-meson mass shift is about 10 times (~ 50 MeV) larger than that of J/ψ. This originates from the fact that the dominant contribution in the OPE for the D-meson is the nucleon matrix element of $m_c \bar{q} q$, where m_c is the charm-quark mass and q denotes light quarks. We also discuss that the mass shift of the D-meson in nuclear matter may cause the level crossings of the charmonium states and the $D\bar{D}$ threshold. This suggests an additional mechanism of the J/ψ suppression in high energy heavy-ion collisions.

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1. Introduction

Changes of hadron properties in the nuclear medium have recently attracted great interests in theoretical studies. They also have induced the on-going experiments and forthcoming experimental plans in heavy nuclei in GSI and in relativistic heavy-ion collisions in SPS at CERN and AGS, RHIC at BNL[1]. In particular the spectral changes of vector mesons are expected to be a possible signal to investigate such medium effect, because their leptonic decay in the medium can supply us information of the medium without disturbance of the strong interaction. Motivated by this experimental advantage, the in-medium effect of the light vector-mesons (ρ^0, ω and ϕ) has been mainly studied in effective hadronic models[2] and QCD sum rules (QSR's)[3–9]. Recently we applied the QSR method to low-lying heavy-quarkonium (J/ψ) in order to investigate its mass shift in nuclear matter and the J/ψ-nucleon interaction at low energy[10]. It was shown that the J/ψ mass at normal matter density drops by $0.1 \sim 0.2\%$ of its vacuum value. This negative mass-shift of J/ψ corresponds to a small decrease of the mass, $4 \sim 7$ MeV[10,11]. This is much smaller than the similar effect in the light mesons such as ρ^0 and ω[3–6]. This difference can be understood as follows: In QSR the magnitude of the spectral change is related to the in-medium change of quark and gluon condensates[12,13]. In the linear density approximation, the quark operators give a dominant contribution for ρ^0 and ω, and have significant change in the matter. On the other hand, the dominant in-medium gluon condensates at the normal matter density for J/ψ are only $5 \sim 10\%$ smaller than its vacuum value.

In this paper, we generalize the above calculations and apply the QSR analysis in Ref. 6 to light and heavy quark systems with unequal mass. We focus on in-medium properties of pseudoscalar mesons D's. Studying the mass modification of the D-mesons in the nuclear medium is important by the following physical reasons:

1. A part of total cross section of the D-meson production attributes a conversion of the $c\bar{c}$-states such as J/ψ and ψ' into open-charm pairs. The component is produced through the reaction, $\psi + N \rightarrow D + \bar{D}$, in high energy proton-nucleus collisions and

relativistic high energy heavy-ion collisions. It contains a heavy charm-quark and a light quark. The existence of a light quark in the D-meson causes much difference of the properties in the nuclear medium between the D-meson and J/ψ. The latter is dominated by the gluon condensates as discussed above, while the former has a large contribution from the light-quark condensates, multiplied by a charm-quark mass. Furthermore J/ψ predominantly interacts with light hadrons solely through gluonic content effects in matter, while the D-mesons will couple strongly through inelastic channels such as $DN \rightarrow \Lambda_c$ or Σ_c $(+\pi)$. Thus we can expect a large modification of the D-meson in the nuclear medium as well as ρ^0 and ω.

2. The existence of novel states, so-called the D-mesic nuclei, has been predicted by quark meson coupling (QMC) model[14]. The QMC model suggests that $D^-(\bar{c}d)$ may be bound in heavy atoms such as ^{208}Pb by an attractive scalar-meson exchange and an attractive Coulomb force. It also suggests that $D^0(c\bar{u})$ is deeply bound in the nuclei by an attractive ω-meson exchange. In view of QSR, it is of importance to investigate the D meson-nucleon (N) interaction.

3. The nuclear absorption of the D-mesons created in π-A and p-A collisions (A denotes targets such as Be, Cu, Al, W and Au) has been measured via the decay $D \rightarrow K\pi$ at Fermilab[15]. The result indicates that D-mesons are not completely absorbed in the nuclear targets irrespective of the charge of the produced D-mesons. The theoretical analysis for the D-N interaction will give an important suggestion for the experimental data in this case too.

Motivated by these points, we investigate the properties of D-mesons in the nuclear matter through the D-N interactions, using an application of Borel-transformed QSR to the D-N forward scattering amplitude. Here we deal with isospin-averaged D-meson current for simplicity.

After a brief explanation of the QSR formulation, in the next section, we calculate the mass shift at finite density using the Borel QSR and compare the results with the mass

shift of J/ψ calculated in Ref. 10. On the basis of these results, we will discuss a possible mechanism of J/ψ suppression through the level crossings between the charmonium states and the $D\bar{D}$ threshold.

2. QCD sum rule analysis for D-meson mass shift

We start with a two-point in-medium correlation function Π_{PS}^{NM} to discuss hadron properties in the nuclear matter. In the Fermi gas approximation for the matter, Π_{PS}^{NM} is divided into two parts by applying the operator product expansion (OPE) to the correlators in the deep Euclidean region ($Q^2 = -q^2 > 0$). One is a vacuum part, Π_{PS}^0, and another is a static one-nucleon part, T_{PS}. This decomposition is expected to be valid at relatively low density[16]. Namely, in the framework of QSR, Π_{PS}^{NM} can be approximated reasonably well in the linear density of the nuclear matter that all nucleons are at rest[6,12,16]:

$$\Pi_{PS}^{\mathrm{NM}}(q) = i \int d^4x \; e^{iq \cdot x} \langle \mathrm{T} J_5(x) J_5^\dagger(0) \rangle_{\mathrm{NM}(\rho_N)} \simeq \Pi_{PS}^0(q) + \frac{\rho_N}{2M_N} T_{PS}(q), \qquad (1)$$

where ρ_N denotes the nuclear matter density, and the forward scattering amplitude T_{PS} of the pseudoscalar current-nucleon is

$$T_{PS}(\omega, \boldsymbol{q}) = i \int d^4x e^{iq \cdot x} \langle N(p) | \mathrm{T} J_5(x) J_5^\dagger(0) | N(p) \rangle. \qquad (2)$$

Here $q^\mu = (\omega, \boldsymbol{q})$ is the four-momentum carried by the D-meson current, $J_5(x) = J_5^\dagger(x) = (\bar{c}i\gamma_5 q(x) + \bar{q}\gamma_5 c(x))/2$, where q denotes u or d quark. $|N(p)\rangle$ represents the isospin, spin-averaged static nucleon-state with the four-momentum $p = (M_N, \boldsymbol{p} = \boldsymbol{0})$, where M_N is the nucleon mass, 0.94 GeV. The state is normalized covariantly as $\langle N(\boldsymbol{p}) | N(\boldsymbol{p}') \rangle = (2\pi)^3 2p^0 \delta^3(\boldsymbol{p} - \boldsymbol{p}')$. The QSR analysis on the forward scattering amplitude enables us to obtain the information for the D-N interaction. In Eq. (1), the second term means a slight deviation from in-vacuum properties of the D-meson determined by Π_{PS}^0. By applying QSR to T_{PS}, we get the D-N scattering length a_D in the limit of $\boldsymbol{q} \to \boldsymbol{0}$. In this limit, T_{PS} can be related to the T-matrix, $\mathcal{T}_{D\mathcal{N}}(m_D, \boldsymbol{q} = \boldsymbol{0}) = 8\pi(M_N + m_D)a_D$. Near the pole position of the

D-meson, the spectral function $\rho(\omega, \boldsymbol{q} = \boldsymbol{0})$ is given with three unknown phenomenological parameters a, b and c in terms of the T-matrix:

$$\rho(\omega, \boldsymbol{q} = \boldsymbol{0}) = -\frac{1}{\pi} \frac{f_D^2 m_D^4}{m_c^2} \mathrm{Im} \left[\frac{\mathcal{T}_{\mathcal{DN}}(\omega, \boldsymbol{0})}{(\omega^2 - m_D^2 + i\varepsilon)^2} \right] + \cdots \tag{3}$$

$$= a\, \delta'(\omega^2 - m_D^2) + b\, \delta(\omega^2 - m_D^2) + c\, \delta(\omega^2 - s_0). \tag{4}$$

Here the leptonic decay constant f_D is defined by the relation $\langle 0|J_5|D(k)\rangle = f_D m_D^2/m_c$, where $|D(k)\rangle$ is the D-meson state with the four-momentum k. $m_D = 1.87$ GeV and $m_c = 1.35$ GeV are masses of the D-meson and charm quark respectively. The terms denoted by $\cdots$ in Eq.(3) represent the continuum contribution and δ' in Eq.(4) is the first derivative of δ function with respect to ω^2. The first term proportional to a is the double-pole term corresponding to the on-shell effect of the T-matrix and a is related to the scattering length a_D as $a = -8\pi(M_N + m_D)a_D f_D^2 m_D^4/m_c^2$. The second term proportional to b is the single-pole term corresponding to the off-shell effect of the T-matrix. The third term proportional to c is the continuum term corresponding to other remaining effects, where s_0 is the continuum threshold in the vacuum. Combining the single-pole term of Π_{PS}^0 in Eq. (1) with Eq. (4), we can relate the scattering length extracted from the QSR of T_{PS} with the mass shift of the D-meson,

$$\delta m_D = 2\pi \frac{M_N + m_D}{M_N m_D} \rho_N a_D. \tag{5}$$

We may determine these unknown parameters by using a dispersion relation and matching the phenomenological (ph) side with the OPE side. Before such analysis, we can impose a constraint from the low energy theorem among these parameters: In the low energy limit ($\omega \to 0$), $T_{PS}(\omega, \boldsymbol{0})$ becomes equivalent to the Born term $T_{PS}^{\mathrm{Born}}(\omega, \boldsymbol{0})$, $T_{PS}^{\mathrm{ph}}(0) = T_{PS}^{\mathrm{Born}}(0)$. We assume two cases to take the Born term into the "ph" side. The case (i) and the case (ii) are defined as follows[6]. At $q_\mu \neq 0$, we require the ω^2-dependence of the Born term explicitly in the case (i):

$$T_{PS}^{\mathrm{ph}}(\omega^2) = T_{PS}^{\mathrm{Born}}(\omega^2) + \frac{a}{(m_D^2 - \omega^2)^2} + \frac{b}{m_D^2 - \omega^2} + \frac{c}{s_0 - \omega^2}, \tag{6}$$

with the condition

$$\frac{a}{m_D^4} + \frac{b}{m_D^2} + \frac{c}{s_0} = 0. \tag{7}$$

In the case (ii), we do not require the ω^2-dependence of the Born term explicitly:

$$T_{PS}^{\text{ph}}(\omega^2) = \frac{a}{(m_D^2 - \omega^2)^2} + \frac{b}{m_D^2 - \omega^2} + \frac{c}{s_0 - \omega^2}, \tag{8}$$

with the condition

$$\frac{a}{m_D^4} + \frac{b}{m_D^2} + \frac{c}{s_0} = T_{PS}^{\text{Born}}(0). \tag{9}$$

If $T_{PS}^{\text{Born}}(0) = 0$, two cases coincide. We determine two unknown phenomenological param-eters a and b in the QSR after the parameter c is removed by the condition of Eq. (7) or (9). The analysis of the Born term is easily performed through a calculation of the Born diagrams at the tree level. The isospin states of the D-meson determine the contribution to the Born term as follows: For the $D^0(c\bar{u})$-N and $D^+(c\bar{d})$-N interactions, we need two reactions,

$$D^0(c\bar{u}) + p(uud) \text{ or } n(udd) \longrightarrow \Lambda_c^+, \Sigma_c^+(cud) \text{ or } \Sigma_c^0(cdd) \tag{10}$$

and

$$D^+(c\bar{d}) + p(uud) \text{ or } n(udd) \longrightarrow \Sigma_c^{++}(cuu) \text{ or } \Lambda_c^+, \Sigma_c^+(cud), \tag{11}$$

where the static masses of charmed baryons are 2.28 GeV for Λ_c^+ and 2.45 GeV for Σ_c^0, Σ_c^+ and Σ_c^{++}. Note that we use $M_B \sim 2.4$ GeV as the average value, where B means either Λ_c^+, Σ_c^+, Σ_c^{++} or Σ_c^0. After averaging over the nucleon spin, we obtain

$$T_{PS}^{\text{Born}}(\omega, \mathbf{0}) = \frac{1}{2} \frac{M_N(M_N + M_B)}{\omega^2 - (M_N + M_B)^2} \, g_p(\omega)^2. \tag{12}$$

The pseudoscalar form-factor $g_p(q^2)$ is introduced through the following relation[17],

$$\langle B(p')|\bar{Q}i\gamma_5 q|N(p)\rangle = g_p(q^2)\bar{u}(p')i\gamma_5 u(p), \tag{13}$$

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where $u(p)$ is a Dirac spinor and $Q(q)$ in the pseudoscalar current is the heavy(light)-quark. Furthermore, $g_p(q^2)$ is related to a coupling constant g_{NDB} such as

$$g_p(q^2) = \frac{2m_D^2 f_D}{m_q + m_c} \frac{g_{NDB}}{q^2 - m_D^2},$$ (14)

where m_q denotes a light-quark mass. $g_{ND\Lambda_c}$ has been estimated in Ref. 17 which gives $g_{DN\Lambda_c} \simeq 6.74$. However since $g_{DN\Sigma_c}$ has not been evaluated, we take an approximation $g_{DN\Sigma_c} \simeq g_{ND\Lambda_c}$. On the other hand, there are no inelastic channels like the above for $\bar{D}^0(\bar{c}u)$-N and $D^-(\bar{c}d)$-N interactions, i.e. $T_{PS}^{\text{Born}}(0) = 0$ in both case (i) and case (ii).

Eventually we determine parameters a and b simultaneously after applying the Borel transform to both the OPE side and the "ph" side.

3. Numerical results in the Borel sum rule

After performing the Borel transform $\hat{B}$ of the T_{PS} calculated in the OPE up to dimension-4[18], we obtain as a function of the Borel mass M

$$\hat{B}\left[T_{PS}^{\text{OPE}}\right] = \frac{1}{2}e^{-m_c^2/M^2}\left[-m_c\langle\bar{q}q\rangle_N + 2\langle q^\dagger i D_0 q\rangle_N\left(-1 + \frac{m_c^2}{M^2}\right)\right.$$
$$+ \frac{1}{2}m_q\langle\bar{q}q\rangle_N\left(1 + \frac{m_c^2}{M^2}\right) + \frac{1}{24}\langle\frac{\alpha_s}{\pi}G^2\rangle_N\left(2 - \frac{m_c^2}{M^2}\right)$$
$$+ \frac{\alpha_s}{\pi}\langle(u\cdot G)^2 - \frac{1}{4}G^2\rangle_N\frac{1}{3}\left\{\frac{4}{3} - \frac{1}{6}\frac{m_c^2}{M^2} + \frac{1}{2}\left(\frac{m_c^2}{M^2}\right)^3\right.$$
$$+ e^{m_c^2/M^2}\left(-2\gamma_E - \ln\left(\frac{m_c^2}{M^2}\right) + \int_0^{\frac{m_c^2}{M^2}}dt\frac{1 - e^{-t}}{t}\right)$$
$$\left.\left. + \left(1 - \frac{m_c^2}{M^2}\right)\ln\left(\frac{m_c^2}{4\pi\mu^2}\right)\right\}\right],$$ (15)

where $\langle\cdot\rangle_N$ denotes the nucleon matrix element. The renormarization scale is taken to be $\mu^2 = 1$ GeV^2 and the Euler constant is $\gamma_E = 0.5772\cdots$. Corresponding formula for the "ph" side, in the case (i) reads

$$\hat{B}\left[T_{PS}^{\text{ph}}\right] = a\left(\frac{1}{M^2}e^{-m_D^2/M^2} - \frac{s_0}{m_D^4}e^{-s_0/M^2}\right) + b\left(e^{-m_D^2/M^2} - \frac{s_0}{m_D^2}e^{-s_0/M^2}\right)$$

$$+ \frac{1}{2} \frac{M_N(M_N + M_B)}{(M_N + M_B)^2 - m_D^2} \frac{4 f_D^2 m_D^4}{(m_q + m_c)^2} g_{NDB}^2$$

$$\times \left[-\frac{e^{-(M_N + M_B)^2/M^2}}{(M_N + M_B)^2 - m_D^2} + \left(\frac{1}{(M_N + M_B)^2 - m_D^2} - \frac{1}{M^2} \right) e^{-m_D^2/M^2} \right], \qquad (16)$$

and that in the case (ii) reads

$$\hat{B} \left[T_{PS}^{\mathrm{ph}} \right] = a \left(\frac{1}{M^2} e^{-m_D^2/M^2} - \frac{s_0}{m_D^4} e^{-s_0/M^2} \right) + b \left(e^{-m_D^2/M^2} - \frac{s_0}{m_D^2} e^{-s_0/M^2} \right)$$

$$+ s_0 T_{PS}^{\mathrm{Born}}(0) \, e^{-s_0/M^2}. \qquad (17)$$

In Eq. (15), a convention, $(u \cdot G)^2 \equiv G_{\kappa\lambda}^a G_\rho^{a\,\lambda} \, u^\kappa u^\rho$ is used, where $u = (1, \mathbf{0})$ for the static nucleon. We equate Eq. (15) with Eq. (16) or Eq. (17) in the case (i) or the case (ii) respectively. Furthermore we take a first derivative of its equation with respect to M^2 in each case. We can derive the D-meson mass shift as a function of M^2 by removing the parameter b from two equations obtained thus. The Borel curve for the D-meson mass shift, δm_D, is shown in Fig.1. Here we adopt $f_D \simeq 0.18$ GeV for the decay constant and $s_0 \simeq 6.0$ GeV^2 for the continuum threshold. The values of these parameters were estimated by the analysis of the in-vacuum correlation function in Ref. 19,20. This f_D value is very close to the result of the lattice QCD calculation $(0.194 \pm 0.01$ GeV$)$[21] and is consistent with the upper bound of experimental data $(\leq 0.31$ GeV$)$[22]. Note that the parameters are fixed in our calculation. We use the nucleon matrix elements[13] for quark fields such as $\langle \bar{q}q \rangle_N \simeq 5.3$ GeV and $\langle q^\dagger i D_0 q \rangle_N \simeq 0.34$ GeV^2, and for gluon fields such as $\langle \frac{\alpha_s}{\pi} G^2 \rangle_N \simeq -1.2$ GeV^2 and $\frac{\alpha_s}{\pi} \langle (u \cdot G)^2 - \frac{1}{4} G^2 \rangle_N \simeq -0.1$ GeV^2. The light-quark mass is taken to be $m_q \simeq 0.008$ GeV. In Fig.1, the solid line ("Born = 0"[case (i)]) are calculated without the contribution of inelastic channels. In the doted line ("With Born term"[case (ii)]) and the dashed line ("Without Born term"), we allow for the contribution of the sub-threshold resonances Λ_c and Σ_c, lying very close to the D-N threshold. As shown in Fig.1, the sub-threshold effect is repulsive both in the case (i) and the case (ii) and its contribution in the case (ii) is stronger than that in the case (i). This is analogous to the case for the K^--p interaction[23,24]. In Fig.1, the remaining three lines are evaluated by extracting the contribution of $m_c \langle \bar{q}q \rangle_N$ term in the OPE from the solid, dashed and doted lines respectively. We find that the contribution

of $m_c\langle\bar{q}q\rangle_N$ term is more than 95% of total contribution corresponding to the solid, dashed and doted lines within the plateau regions.

The analysis of this graph is summarized in Table 1. We cannot determine th Borel windows, since there is only lowest-dimension term in the OPE side. Therefore we take the following procedure to determine a window of M^2: We focus on a plateau region of each line shown in Fig.1. First, we take the minimum point in the line as the smallest value of δm_D. Next, we determine two points in M^2 as the deviation from the minimum value of δm_D becomes less than 10% of the minimum value. We take the region between the two points as a window of M^2. The window of M^2 in each line is given in Table 1. As shown in Table 1, all the Borel curves are rather stable within the windows. Note that the windows discussed above are also close to that obtained by scaling up typical Borel-windows for the light-vector mesons[6]. An estimate from the light-vector mesons is $(\frac{m_D}{m_V})^2 \times 0.8 (= 4.7) <$ $M^2 < (\frac{m_D}{m_V})^2 \times 1.3 (= 7.7)$ for $V = \rho, \omega$ and $(\frac{m_D}{m_V})^2 \times 1.3 (= 4.4) < M^2 < (\frac{m_D}{m_V})^2 \times 1.8 (= 6.1)$ for $V = \phi$.

It should be stressed that the QSR for vacuum correlation function Π^0_{PS} up to dimension-6 operators cannot reproduce well the D-meson mass (m_D) in free space[25]. The Borel curve of m_D at $s_0 = 6.0$ GeV^2 does not have any stability and seems to give larger values than the experimental data ($m_D = 1.87$ GeV) in the plateau region ($M^2 = 3 \sim 8$ GeV^2) of δm_D discussed above. Furthermore, the curve of m_D has rather larger change than that of δm_D in the region. Therefore, if we perform the QSR analysis for the effective mass $m_D^* = m_D + \delta m_D$, the stability in the Borel curve of m_D^* will become obscure. This implies that the QSR analysis for m_D^* is not so valid in this case. However the application of the QSR to δm_D gives us a possible indication for the D-meson mass-shift as discussed above.

From the above analysis we obtain $\delta m_D = -48 \pm 8$ MeV by considering both cases ((i) and (ii)). Then the D-N scattering length a_D is -0.72 ± 0.12 fm. This result suggests that the D-N interaction is more attractive than the J/ψ-N interaction, where $\delta m_{J/\psi}$ is about -5 MeV and $a_{J/\psi}$ is about -0.1 fm[10,11]. It is noted the above results were performed at the nuclear matter density, $\rho_N = \rho_0 \sim 2\rho_0$, where ρ_0 is the normal matter density. Assuming

that the linear density approximation in the QSR is valid at the nuclear matter density, we find these results lead to a larger decrease of the D-meson mass than that of the charmonium. We also find that the large mass-shift of the D-meson originates from the contribution of $m_c\langle\bar{q}q\rangle_N$ term in the OPE. In the QMC model, the contribution from the $m_c\langle\bar{q}q\rangle_N$ term may correspond to a quark-σ meson coupling. In fact, the model predicts the mass shift of the D-meson becomes -60 MeV for the scalar potential at the normal matter density[14]. Their results are very close to our results.

4. Concluding remarks

We present an analysis of isospin-averaged D-meson mass-shift through a direct application of the Borel QSR to the forward scattering amplitude of the pseudoscalar current and nucleon. Here we perform the operator product expansion with all the terms up to dimension-4. The result predicts an attractive mass-shift of the isospin-averaged D-meson about 50 MeV (about 3% of the bare mass) at the normal matter density. The contribution of $m_c\langle\bar{q}q\rangle_N$ term in the OPE is more than 95% of the results evaluated up to all the dimension-4 operators. Our result is very close to that reported by the QMC model[14]. Then the D-N scattering length a_D is about -0.7 fm and indicates a strongly attractive force between the D-meson and nucleons. The mass modification of the D-meson in the nuclear matter corresponds to 10 times larger than that of J/ψ.

Recent NA50 Collaboration[26] has reported a strong (so called "anomalous") J/ψ suppression[27] in Pb-Pb collision at 158 GeV per nucleon. We suggest a following possibility to cause such a strong J/ψ suppression using the above large difference of the mass shifts between the D-meson and J/ψ: Suppose that the other $c\bar{c}$-states such as ψ' and χ_c's do not also have a large mass-shift in the nuclear matter as discussed above for J/ψ. Then on the basis of our calculations, the $D\bar{D}$ threshold (~ 3.74 GeV in free space) decreases by 100 MeV at $\rho_N = \rho_0$ and comes down between χ_{c2} (~ 3.55 GeV) and ψ' (~ 3.68 GeV). At $\rho_N = 2\rho_0$, it decreases by 200 MeV and comes down between χ_{c1} (~ 3.50 GeV) and

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χ_{c2} ($\sim$ 3.55 GeV). Such level crossings between the $D\bar{D}$ threshold and the charmonium spectrum are shown in Fig.2. It is well known from the p-A collision data that only 60% of J/ψ observed are directly produced and the remainder comes from excited states ($\chi_c(1P)$, $\psi'(2S)$) with a ratio of 3 to 1[28]. So if the disappearance of ψ' and χ_c due to their decays into $D\bar{D}$ takes place, it could lead to a decrease of the J/ψ yield. This J/ψ suppression will have stair-shaped form as a function of the energy density of the system. At low energy density, only $\psi' \to D\bar{D}$ takes place, which causes a slight suppression of J/ψ. At intermediate density, subsequent suppression occurs by the level crossing of (χ_{c2}, χ_{c1}) states with the $D\bar{D}$ threshold. At higher density, direct suppression of J/ψ by the decay $J/\psi \to D\bar{D}$ occurs. This is an alternative or supplementary effect to the deconfinement scenario in Ref. 29 and may have implications to the recent data of NA50 Collaboration[30]. Strictly speaking, as mentioned before, the formulation of QSR used here is applicable to dilute nuclear matter in comparison to highly dense nuclear matter produced in high energy heavy-ion collision. If we try to discuss the mass shifts at higher density, it will need an extension or some improvements in the above formulation. This is one of important future works for us. However we may say that our calculation in this paper can suggest a possibility of the level crossings even in only 1 or 2 times normal matter density and lets one expect the direct J/ψ suppression at sufficiently higher density.

We may also suggest another phenomenon caused by the level crossings discussed above. It is a change of the decay width of the charmonium. In the vacuum the resonances above $D\bar{D}$ threshold, for example the ψ'' state, have a width of order MeV because of the strong open-charm channel. On the other hand, the resonances below the threshold have a very sharp width of a few hundreds keV. So after the level crossings, the decay modes of the ψ' and χ_c states will change drastically at least one order of magnitude.

Needless to say, in the high energy heavy-ion collisions we must also perform the theoretical investigation of finite temperature effect[8,31,32] to the mass modification. Therefore, a future task will be to understand the medium effect in terms of the QSR when both temperature and density are finite. As for investigation of only the D-N interaction, an

inverse kinematics experiment will give useful information[33]. Since the projectile (target) is a heavy-ion beam (light-nuclei), the decays of the D-meson[15] and the charmonium inside a nucleus will be possible in such an experiment.

Before closing, it is stressed that we should take the above analysis for the isospin-averaged D-N interaction as a qualitative estimate. If we wish more quantitative discussion, the isospin-decomposition of the D-meson and higher correction terms beyond dimension-4 operators must be taken into account. Particularly the odd-components in the OPE play important roles for a difference of the mass shift between D and $\bar{D}$.

Acknowledgements

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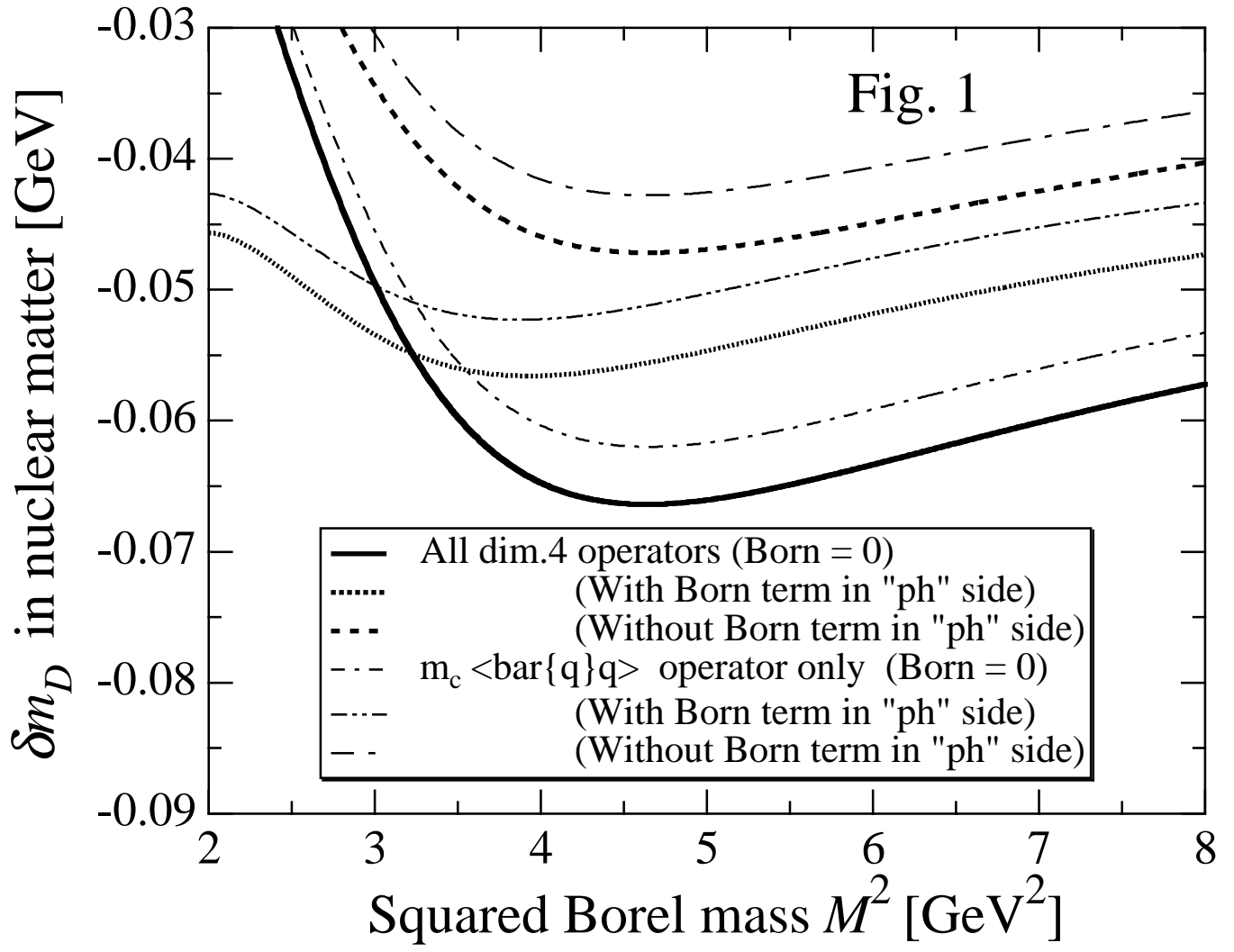
FIGURES

Fig. 1.   The D-meson mass-shift δm_D [GeV] at normal matter density ($\rho_0 = 0.17$ fm^{-3}) as a function of the squared Borel mass M^2 [GeV^2] : The solid, doted and dashed lines are calculated by taking into account all the dimension-4 operators in the OPE. The lines correspond to the cases of $T_{PS}^{\mathrm{Born}}(0) = 0$, with (case (i)) and without (case (ii)) an explicit Born term ($T_{PS}^{\mathrm{Born}}(0) \neq 0$) in the "ph"side respectively. The remaining three lines show the results obtained by extracting the contribution of $m_c \langle \bar{q}q \rangle_N$ term in the OPE from the above each line respectively.

Fig. 2.   Comparison between charmonim spectrum and $D\bar{D}$ threshold level: If the mass shifts of all the charmonium states are very small (0.1 ∼ 0.2% of the bare mass) in the nuclear medium ($\rho_N = \rho_0 \sim 2\rho_0$, $\rho_0 = 0.17$ fm^{-3}), we predict that the threshold lying just below $1D$ state in free space falls down below the ψ' state at normal matter density and the χ_{c2} state twice at the density.

Table 1. The result of an analysis for the solid, dashed and doted lines in Fig.1: The second raw shows the plateau region, which we determine as the deviation from a minimum value of δm_D in each line becomes less than 10% of the minimum value. The δm_D in the plateau region is given in the third raw.

Lines	Solid line	Doted line (case (i))	Dashed line (case (ii))
Plateau region [GeV^2]	$3.5 \leq M^2 \leq 7.0$	$2.6 \leq M^2 \leq 7.0$	$3.4 \leq M^2 \leq 8.0$
δm_D [GeV]	-63 ± 3	-53 ± 3	-44 ± 3

Fig. 1

Fig. 2