Analysis of the Heavy Mesons in the Nuclear Matter with the QCD Sum Rules

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Abstract

In this article, I calculate the contributions of the nuclear matter induced condensates up to dimension 5, take into account the next-to-leading order contributions of the nuclear matter induced quark condensate, and study the properties of the scalar, pseudoscalar, vector and axial-vector heavy mesons in the nuclear matter with the QCD sum rules in a systematic way. The present predictions for the shifts of the masses and decay constants can be confronted with the experimental data in the future. Furthermore, I study the heavy-meson-nucleon scattering lengths as a byproduct, and obtain the conclusion qualitatively about the possible existence of heavy-meson-nucleon bound states.

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1 Introduction

The suppression of \( J/\psi \) production in relativistic heavy ion collisions is considered as an important signature to identify the quark-gluon plasma \([1]\). The dissociation of \( J/\psi \) in the quark-gluon plasma due to color screening can result in a reduction of its production. The interpretation of suppression requires the detailed knowledge of the expected suppression due to the \( J/\psi \) dissociation in the hadronic environment. The in-medium hadron properties can affect the productions of the open-charmed mesons and the \( J/\psi \) in the relativistic heavy ion collisions, the higher charmonium states are considered as the major source of the \( J/\psi \) \([2]\). For example, the higher charmonium states can decay to the \( D\bar{D}, D^*\bar{D}^* \) pairs instead of decaying to the lowest state \( J/\psi \) in case of the mass reductions of the \( D, D^*, D, D^* \) mesons are large enough. We have to disentangle the color screening versus the recombination of off-diagonal \( c\bar{c} \) (or \( \bar{b}b \)) pairs in the hot dense medium versus cold nuclear matter effects, such as nuclear absorption, shadowing and anti-shadowing, so as to draw a definite conclusion on appearance of the quark-gluon plasma \([3, 4]\). The upcoming FAIR (Facility for Antiproton and Ion Research) project at GSI (Institute for Heavy Ion Research) in Darmstadt (Germany) provides the opportunity to study the in-medium properties of the charmoniums or charmed hadrons for the first time. The CBM (Compressed Baryonic Matter) collaboration intends to study the properties of the hadrons in the nuclear matter \([5]\), while the PANDA (anti-Proton Annihilation at Darmstadt) collaboration will focus on the charm spectroscopy, and mass and width modifications of the charmed hadrons in the nuclear matter \([6]\). However, the in-medium mass modifications are not easy to access experimentally despite the interesting physics involved, and they require more detailed theoretical studies. On the other hand, the bottomonium states are also sensitive to the color screening, the \( \Upsilon \) suppression in high energy heavy ion collisions can also be taken as a signature to identify the quark-gluon plasma \([7]\). The suppressions on the \( \Upsilon \) production in ultra-relativistic heavy ion collisions will be studied in details at the Relativistic Heavy Ion Collider (RHIC) and Large Hadron Collider (LHC).

Extensive theoretical and experimental studies are required to explore the hadron properties in nuclear matter. The connection between the condensates and the nuclear density dependence of the in-medium hadron masses is not straightforward. The QCD sum rules provides a powerful theoretical tool in studying the in-medium hadronic properties \([8, 9]\), and has been applied extensively to study the light-flavor hadrons and charmonium states in the nuclear matter \([10, 11, 12]\). The works on the heavy mesons and heavy baryons are few, only the \( D, B, D_0, B_0, D^*, B^*, D_1, D_2 \),

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B_1$, $\Lambda_Q$, $\Sigma_Q$, $\Xi_{QQ}$ and $\Omega_{QQ}$ are studied with the QCD sum rules. The heavy mesons (heavy baryons) contain a heavy quark and a light quark (two light quarks), the existence of a light quark (two light quarks) in the heavy mesons (heavy baryons) leads to large difference between the mass-shifts of the heavy mesons (heavy baryons) and heavy quarkonia in the nuclear matter. The former have large contributions from the light-quark condensates in the nuclear matter and the modifications of the masses originate mainly from the modifications of the quark condensates, while the latter are dominated by the gluon condensates, and the mass modifications are mild.

In previous works, the properties of the heavy mesons in the nuclear matter are studied with the QCD sum rules by taking the leading order approximation for the contributions of the quark condensates, and study the properties of the scalar, pseudoscalar, vector and axial-vector heavy mesons in the nuclear matter with the QCD sum rules in a systematic way, and make predictions for the modifications of the masses and decay constants of the heavy mesons in the nuclear matter. The present predictions can be confronted with the experimental data from the CBM and PANDA collaborations in the future, measuring those effects is a long term physics goal based on further theoretical studies on the reaction dynamics and on the exploration of the experimental ability to identify more complicated processes. Furthermore, I study the heavy-meson-nucleon scattering lengths as a byproduct. From the negative or positive sign of the scattering lengths, I can obtain the conclusion qualitatively that the interactions are attractive or repulsive, which favor or disfavor the formations of the heavy-meson-nucleon bound states. For example, the $\Sigma_c(2800)$ and $\Lambda_c(2940)$ can be assigned to be the $S$-wave $DN$ state with $J^P = \frac{1}{2}^-$ and the $S$-wave $D^*N$ state with $J^P = \frac{3}{2}^-$ respectively based on the QCD sum rules.

The article is arranged as follows: I study in-medium properties of the heavy mesons with the QCD sum rules in Sec.2; in Sec.3, I present the numerical results and discussions; and Sec.4 is reserved for my conclusions.

## 2 The properties of the heavy mesons in the nuclear matter with QCD sum rules

I study the scalar, pseudoscalar, vector and axial-vector heavy mesons in the nuclear matter with the two-point correlation functions $\Pi(q)$ and $\Pi_{\mu\nu}(q)$, respectively. In the Fermi gas approximation for the nuclear matter, I divide the $\Pi(q)$ and $\Pi_{\mu\nu}(q)$ into the vacuum part $\Pi(q)$ and $\Pi_{\mu\nu}(q)$ and the static one-nucleon part $\Pi_N(q)$ and $\Pi_{\mu\nu}^N(q)$, and approximate the $\Pi_N(q)$ and $\Pi_{\mu\nu}^N(q)$ up to the order $O(\rho_N)$ at relatively low nuclear density.

\[
\Pi(q) = \int d^4x e^{iq\cdot x} \left\langle T \left\{ J(x) J^\dagger(0) \right\} \right\rangle_{\rho_N} \\
= \Pi_0(q) + \frac{\rho_N}{2m_N} T_N(q),
\]

\[
\Pi_{\mu\nu}(q) = \int d^4x e^{iq\cdot x} \left\langle T \left\{ J_\mu(x) J_\nu^\dagger(0) \right\} \right\rangle_{\rho_N} \\
= \Pi_{\mu\nu}^0(q) + \frac{\rho_N}{2m_N} T_{\mu\nu}^N(q),
\]

where the $\rho_N$ is the density of the nuclear matter, and the forward scattering amplitudes $T_N(q)$ and $T_{\mu\nu}^N(q)$ are defined as

\[
T_N(\omega, q) = \int d^4x e^{iq\cdot x} \left\langle N(p) | T \left\{ J(x) J^\dagger(0) \right\} | N(p) \right\rangle, \\
T_{\mu\nu}^N(\omega, q) = \int d^4x e^{iq\cdot x} \left\langle N(p) | T \left\{ J_\mu(x) J_\nu^\dagger(0) \right\} | N(p) \right\rangle,
\]
where the $J(x)$ and $J_\mu(x)$ denote the isospin averaged currents $\eta(x)$, $\eta_5(x)$, $\eta_\mu(x)$ and $\eta_{5\mu}(x)$, respectively,

\[
\begin{align*}
\eta(x) &= \eta^I(x) = \frac{\bar{c}(x)q(x) + \bar{q}(x)c(x)}{2}, \\
\eta_5(x) &= \eta_5^I(x) = \frac{\bar{c}(x)i\gamma_5 q(x) + \bar{q}(x)i\gamma_5 c(x)}{2}, \\
\eta_\mu(x) &= \eta_\mu^I(x) = \frac{\bar{c}(x)\gamma_\mu q(x) + \bar{q}(x)\gamma_\mu c(x)}{2}, \\
\eta_{5\mu}(x) &= \eta_{5\mu}^I(x) = \frac{\bar{c}(x)\gamma_\mu\gamma_5 q(x) + \bar{q}(x)\gamma_\mu\gamma_5 c(x)}{2},
\end{align*}
\]

which interpolate the scalar, pseudoscalar, vector and axialvector mesons $D_0$, $D$, $D^*$ and $D_1$, respectively. I choose the isospin averaged currents since the $D_0$, $D$, $D^*$ and $D_1$ mesons are produced in pairs in the antiproton-nucleon annihilation processes. The $q$ denotes the $u$ or $d$ quark, the $q^\mu = (\omega, q)$ is the four-momentum carried by the currents $J(x)$ and $J_\mu(x)$, the $|N(p)|$ denotes the isospin and spin averaged static nucleon state with the four-momentum $p = (m_N, 0)$, and $\langle N(p) | N(p') \rangle = (2\pi)^3 2p_0 \delta^3(p-p')$ [13].

I can decompose the correlation functions $T_{\mu\nu}^N(\omega, q)$ as

\[
T_{\mu\nu}^N(\omega, q) = T_N(\omega, q) \left(-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \right) + T_N^0(\omega, q) q_\mu q_\nu + T_N^1(\omega, q) (q_\mu u_\nu + q_\nu u_\mu) + T_N^2(\omega, q) u_\mu u_\nu,
\]

according to Lorentz covariance, where the $T_N(\omega, q)$ denotes the contributions of the vector and axialvector charmed mesons, and the $T_N^{0,1/2}(\omega, q)$ are irrelevant in the present analysis.

In the limit $q \to 0$, the forward scattering amplitude $T_N(\omega, q)$ can be related to the $DN$ ($D_0N$, $D^*N$ and $D_1N$) scattering $T$-matrix,

\[
T_{D/D_0/D^*/D_1}^N(m_{D/D_0/D^*/D_1}, 0) = 8\pi(m_N + m_{D/D_0/D^*/D_1})a_{D/D_0/D^*/D_1},
\]

where the $a_{D/D_0/D^*/D_1}$ are the $D/D_0/D^*/D_1 N$ scattering lengths. I can parameterize the phenomenological spectral densities $\rho(0, x)$ with three unknown parameters $a$, $b$ and $c$ near the pole positions of the charmed mesons $D$, $D_0$, $D^*$ and $D_1$ according to Ref. [13],

\[
\rho(0, x) = -\frac{1}{\pi} \text{Im} \left[ \frac{T_{D/D_0}^N(\omega, 0)}{(\omega^2 - m_{D/D_0}^2 + i\varepsilon)^2} \right] \frac{f_{D/D_0}^2 m_D^4}{m_c^2} + \cdots,
\]

\[
= a \frac{d}{d\omega^2} \delta (\omega^2 - m_{D/D_0}^2) + b \delta (\omega^2 - m_{D/D_0}^2) + c \delta (\omega^2 - s_0),
\]

for the pseudoscalar and scalar currents $\eta_5(x)$ and $\eta(x)$,

\[
\rho(0, x) = -\frac{1}{\pi} \text{Im} \left[ \frac{T_{D^*/D_1}^N(\omega, 0)}{(\omega^2 - m_{D^*/D_1}^2 + i\varepsilon)^2} \right] \frac{f_{D^*/D_1}^2 m_{D^*/D_1}^2}{m_c^2} + \cdots,
\]

\[
= a \frac{d}{d\omega^2} \delta (\omega^2 - m_{D^*/D_1}^2) + b \delta (\omega^2 - m_{D^*/D_1}^2) + c \delta (\omega^2 - s_0),
\]

for the vector and axial-vector currents $\eta_\mu(x)$ and $\eta_{5\mu}(x)$.
Now the hadronic correlation functions $\Pi(\omega,0)$ and $\Pi_{\mu\nu}(\omega,0)$ at the phenomenological side can be written as

\[
\Pi(\omega,0) = \left( f_{D/D_0} + \delta f_{D/D_0} \right) \frac{2}{m_\pi^2} \frac{m_{D/D_0} + \delta m_{D/D_0}}{(m_{D/D_0} + \delta m_{D/D_0})^2 - \omega^2} \cdot \ldots
\]

\[
= f_{D/D_0}^2 \frac{m_{D/D_0}^4}{m_\pi^2} \frac{1}{m_{D/D_0}^2 - \omega^2} \cdot \ldots
\]

\[
+ \frac{\rho_B}{2m_B} \left[ \frac{a}{(m_{D/D_0}^2 - \omega^2)^2} + \frac{b}{m_{D/D_0}^2 - \omega^2} + \ldots \right], \tag{8}
\]

\[
\Pi_{\mu\nu}(\omega,0) = \left( f_{D^*/D_1} + \delta f_{D^*/D_1} \right) \frac{2}{m_\pi^2} \frac{m_{D^*/D_1} + \delta m_{D^*/D_1}}{(m_{D^*/D_1} + \delta m_{D^*/D_1})^2 - \omega^2} \cdot \ldots
\]

\[
= f_{D^*/D_1}^2 \frac{m_{D^*/D_1}^4}{m_\pi^2} \frac{1}{m_{D^*/D_1}^2 - \omega^2} \cdot \ldots
\]

\[
+ \frac{\rho_B}{2m_B} \left[ \frac{a}{(m_{D^*/D_1}^2 - \omega^2)^2} + \frac{b}{m_{D^*/D_1}^2 - \omega^2} + \ldots \right] \left( -g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \right) + \ldots, \tag{9}
\]

In Eqs. (6-7), the first term denotes the double-pole term, and corresponds to the on-shell effect of the $T$-matrix,

\[
a = -8\pi(m_N + m_{D/D_0})a_{D/D_0} f_{D/D_0}^2 \frac{m_{D/D_0}^4}{m_\pi^2}, \tag{10}
\]

for the currents $\eta_5(x)$ and $\eta(x)$ and

\[
a = -8\pi(m_N + m_{D^*/D_1})a_{D^*/D_1} f_{D^*/D_1}^2 \frac{m_{D^*/D_1}^4}{m_\pi^2}, \tag{11}
\]

for the currents $\eta_6(x)$ and $\eta_{6\mu}(x)$: the second term denotes the single-pole term, and corresponds to the off-shell effect of the $T$-matrix; the third term denotes the continuum term or the remaining effects, where the $s_0$ is the continuum threshold parameter. In general, the continuum contributions are approximated by $\rho_{QCD}(\omega,0)\theta(\omega^2 - s_0)$, where the $\rho_{QCD}(\omega,0)$ are the perturbative QCD spectral densities, and $\theta(x) = 1$ for $x \geq 0$, else $\theta(x) = 0$. In this article, the QCD spectral densities are of the type $\delta(\omega^2 - m^2_\pi)$, which include both the ground state and continuum state contributions, so the net continuum state contributions can be approximated as $c\delta(\omega^2 - s_0)$, then I obtain the result $c/ (s_0 - \omega^2)$ in the hadronic representation, see Eq.(15). The doublet $(D(2550), D(2600))$ or $(D_J(2580), D_J(2650))$ is assigned to be the first radial excited state of the doublet $(D, D^*)$ [20]. The single-pole contributions come from the doublet $(D(2550), D(2600))$ or $(D_J(2580), D_J(2650))$ are of the form $1/ (m_{D(2550)/D(2600)}^2 - \omega^2)$, so the approximation $c/ (s_0 - \omega^2)$ is reasonable.

Then the shifts of the masses and decay constants of the charmed-mesons can be approximated as

\[
\delta m_{D/D_0} = 2\pi \frac{m_N + m_{D/D_0}/D^*/D_1}{m_N m_{D/D_0}/D^*/D_1} \rho_B a_{D/D_0}/D^*/D_1, \tag{12}
\]
\[ \delta f_{D/D_0} = \frac{m_c^2}{2f_{D/D_0}m_D^4} \left( \frac{b\rho_N}{2m_N} - \frac{4f_{D/D_0}^2m_D^2m_D^2\delta m_{D/D_0}}{m_c^2} \right), \]

\[ \delta f_{D^*/D_1} = \frac{1}{2f_{D^*/D_1}m_{D^*}^2} \left( \frac{b\rho_N}{2m_N} - \frac{2f_{D^*/D_1}^2m_D^2m_{D^*}^2\delta m_{D^*/D_1}}{m_c^2} \right). \]  \hspace{1cm} (13)

In calculations, I have used the following definitions for the decay constants of the heavy mesons,

\[ \langle 0|\bar{\eta}(0)|D_0 + \bar{D}_0 \rangle = \frac{f_{D_0}m_D^2}{m_c}, \]

\[ \langle 0|\bar{\eta}_5(0)|D + \bar{D} \rangle = \frac{f_{D}m_D^2}{m_c}, \]

\[ \langle 0|\bar{\eta}_\mu(0)|D^* + \bar{D}^* \rangle = f_D m_D \epsilon_\mu, \]

\[ \langle 0|\bar{\eta}_\mu(0)|D_1 + \bar{D}_1 \rangle = f_{D_1} m_{D_1} \epsilon_\mu, \] \hspace{1cm} (14)

with summations of the polarization vectors \( \sum_\lambda \epsilon_\lambda(\lambda, q)\epsilon_\lambda^*(\lambda, q) = -g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \).

In the low energy limit \( \omega \to 0 \), the \( T_N(\omega, 0) \) is equivalent to the Born term \( T_N^{\text{Born}}(\omega, 0) \). Now I take into account the Born terms at the phenomenological side,

\[ T_N(\omega^2) = T_N^{\text{Born}}(\omega^2) + \frac{a}{(m_D^2/m_{D_0}^2 + D_{D_0}/D_1 - \omega^2)^2} + \frac{b}{m_D^2/m_{D_0}^2 + D_{D_0}/D_1 - \omega^2} + \frac{c}{s_0 - \omega^2}, \] \hspace{1cm} (15)

with the constraint

\[ \frac{a}{m_D^2/m_{D_0}^2 + D_{D_0}/D_1} + \frac{b}{m_D^2/m_{D_0}^2 + D_{D_0}/D_1} + \frac{c}{s_0} = 0. \] \hspace{1cm} (16)

The contributions from the intermediate spin-\( \frac{3}{2} \) charmed baryon states are zero in the soft-limit \( q_\mu \to 0 \), and I only take into account the intermediate spin-\( \frac{1}{2} \) charmed baryon states in calculating the Born terms,

\[ (D/D_0/D^*+D_0)^0(cu) + p(uud) \text{ or } n(udd) \rightarrow \Lambda^+_c, \Sigma^+_c(cud) \text{ or } \Sigma^0_c(cdd), \]

\[ (D/D_0/D^*+D_1)^+ (cd) + p(uud) \text{ or } n(udd) \rightarrow \Sigma^{++}_c(cuu) \text{ or } \Lambda^+_c, \Sigma^+_c(cud), \] \hspace{1cm} (17)

where \( M_{\Lambda_c} = 2.286 \text{ GeV} \) and \( M_{\Sigma_c} = 2.454 \text{ GeV} \). I can take \( M_H \approx 2.4 \text{ GeV} \) as the average value, where the \( H \) means either \( \Lambda^+_c, \Sigma^+_c, \Sigma^{++}_c \) or \( \Sigma^0_c \). In the case of the bottom baryons, I take the approximation \( M_H = \frac{M_{B^+} + M_{B^0}}{2} \approx 5.7 \text{ GeV} \). I write down the Feynman diagrams and calculate the Born terms directly; and obtain the results,

\[ T_N^{\text{Born}}(\omega, 0) = \frac{2m_N(m_H + m_N)}{[\omega^2 - (m_H + m_N)^2] \left[ \omega^2 - m_D^2 \right]^2} \left( \frac{f_D m_D^2 g_{DNH}}{m_c} \right)^2, \] \hspace{1cm} (18)

for the current \( \eta_0(x) \),

\[ T_N^{\text{Born}}(\omega, 0) \to T_N^{\text{Born}}(\omega, 0) \text{ (with } m_N \to -m_N, \ D \to D_0 \), \hspace{1cm} (19)

for the current \( \eta(x) \),

\[ T_N^{\text{Born}}(\omega, 0) = \frac{2m_N(m_H + m_N)}{[\omega^2 - (m_H + m_N)^2] \left[ \omega^2 - m_{D^*}^2 \right]^2} (f_{D^*} m_{D^*} g_{D^*NH})^2, \] \hspace{1cm} (20)

for the current \( \eta_\mu(x) \),

\[ T_N^{\text{Born}}(\omega, 0) \to T_N^{\text{Born}}(\omega, 0) \text{ (with } m_N \to -m_N, \ D^* \to D_1 \), \hspace{1cm} (21)
for the current \( \rho_\mu(x) \), where the \( g_{D/D_0/D_r/D_1,N} \) denote the strong coupling constants \( g_{D/D_0/D_r/D_1,N} \), and \( g_{D/D_0/D_r/D_1,N} \). On the other hand, there are no inelastic channels for the \( (D/D_0/D_r/D_1,0,0) \) interactions, and \( T_N^{\text{box}}(0) = 0 \). In calculations, I have used the following definitions for the hadronic coupling constants,

\[
\begin{align*}
\langle \Lambda_c/\Sigma_c(p-q)|D(-q)N(p) \rangle &= g_{\Lambda_c/\Sigma_c,DN} \mathcal{T}_{\Lambda_c/\Sigma_c}(p-q) i\gamma_5 U_N(p), \\
\langle \Lambda_c/\Sigma_c(p-q)|D_0(-q)N(p) \rangle &= g_{\Lambda_c/\Sigma_c,D_0N} \mathcal{T}_{\Lambda_c/\Sigma_c}(p-q) U_N(p), \\
\langle \Lambda_c/\Sigma_c(p-q)|D^*(q)N(p) \rangle &= \mathcal{T}_{\Lambda_c/\Sigma_c}(p-q) \left( g_{\Lambda_c/\Sigma_c,D^*N} \gamma + i \frac{g_{\Lambda_c/\Sigma_c,D^*N}}{M_N + M_{\Lambda_c/\Sigma_c}} \epsilon_{\alpha\beta} q_{\beta} \right) \gamma_5 U_N(p), \\
\langle \Lambda_c/\Sigma_c(p-q)|D_1(-q)N(p) \rangle &= \mathcal{T}_{\Lambda_c/\Sigma_c}(p-q) \left( g_{\Lambda_c/\Sigma_c,D_1N} \gamma + i \frac{g_{\Lambda_c/\Sigma_c,D_1N}}{M_N + M_{\Lambda_c/\Sigma_c}} \epsilon_{\alpha\beta} q_{\beta} \right) \gamma_5 U_N(p),
\end{align*}
\]

(22)

where the \( U_N \) and \( \mathcal{T}_{\Lambda_c/\Sigma_c} \) are the Dirac spinors of the nucleon and the charmed baryons \( \Lambda_c/\Sigma_c \), respectively. In the limit \( q_\mu \to 0 \), the strong coupling constants \( g_{\Lambda_c/\Sigma_c,D^*N} \) and \( g_{\Lambda_c/\Sigma_c,D_1N} \) have no contributions.

For example, near the thresholds, the \( D^*N \) can translate to the \( DN, D^*N, \pi\Sigma_c, \eta\Lambda_c \), etc., we can take into account the intermediate baryon-meson loops or the re-scattering effects with the Bethe-Salpeter equation to obtain the full \( D^*N \to D^*N \) scattering amplitude, and generate higher baryon states dynamically \[29\]. We can saturate the \( D^*N \to D^*N \) scattering amplitude with the tree-level Feynman diagrams describing the exchanges of the higher resonances \( \Lambda_c(2595), \Sigma_c(\frac{1}{2}^-) \), etc. While in other coupled-channels analysis, the \( \Lambda_c(2595) \) emerges as a \( DN \) quasi-bound state rather than a \( D^*N \) quasi-bound state \[23\]. The translations \( D^*N \) to the ground states \( \Lambda_c \) and \( \Sigma_c \) are favored in the phase-space, as the \( \Lambda_c(2595) \) and \( \Sigma_c(\frac{1}{2}^-) \) with \( J^P = \frac{1}{2}^- \) have the average mass \( m_{\Lambda_c} \approx 2.7 \text{ GeV} \) \[22\] \[24\]. In fact, \( m_{\Lambda_c}^2 > s_0 \), I can absorb the high resonances into the continuum states in case the high resonances do not dominate the QCD sum rules. In calculations, I observe that the mass-shift \( \delta m_{D^*} \) does not sensitive to contributions of the ground states \( \Lambda_c \) and \( \Sigma_c \), the contributions from the spin-\( \frac{1}{2} \) higher resonances maybe even smaller. In this article, I neglect the intermediate baryon-meson loops, their effects are absorbed into continuum contributions.

At the low nuclear density, the condensates \( \langle O \rangle_{\rho_N} \) in the nuclear matter can be approximated as

\[
\langle O \rangle_{\rho_N} = \langle O \rangle + \frac{\rho_N}{2m_N} \langle O \rangle_N,
\]

(23)

based on the Fermi gas model, where the \( \langle O \rangle \) and \( \langle O \rangle_N \) denote the vacuum condensates and nuclear matter induced condensates, respectively \[11\]. I neglect the terms proportional to \( \rho_F \), \( \rho_F^2 \), \( \rho_F^3 \), \( \rho_F^4 \), \( \rho_F^5 \), \( \cdot \cdot \cdot \) at the normal nuclear matter with the saturation density \( \rho_N = \rho_0 = \frac{2\rho_F}{3\pi^2} \), as the Fermi momentum \( \rho_F = 0.27 \text{ GeV} \) is a small quantity \[11\].

I carry out the operator product expansion to the nuclear matter induced condensates \( \frac{\rho_N}{2m_N}\langle O \rangle_N \) up to dimension-5 at the large space-like region in the nuclear matter, and take into account the one-loop corrections to the quark condensate \( \langle \bar{q}q \rangle_N \). I insert the following term

\[
\frac{1}{2\pi^2} g_s \int d^D y \tilde{\psi}(y) \gamma^\mu \psi(y) \frac{\lambda^a}{2} G_\mu(x) \int d^D z \tilde{\psi}(z) \gamma^\nu \psi(z) \frac{\lambda^b}{2} G_\nu(z),
\]

(24)

with the dimension \( D = 4 - 2\epsilon \), into the correlation functions \( T_N(q) \) and \( T_N^N(q) \) firstly, where the \( \psi \) denotes the quark fields, the \( G_\mu \) denotes the gluon field, the \( \lambda^a \) denotes the Gell-Mann matrices, then contract the quark fields with Wick theorem, and extract the quark condensate \( \langle \bar{q}q \rangle_N \) according to the formula \( \langle N|\bar{q}_a \gamma_\mu q_\beta|N\rangle = -\frac{1}{12} \langle \bar{q}q \rangle_N \delta_{ij} \delta_{\alpha\beta} \) to obtain the perturbative corrections \( \alpha_s \langle \bar{q}q \rangle_N \), where
the $i$ and $j$ are color indexes and the $\alpha$ and $\beta$ are Dirac spinor indexes. There are six Feynman diagrams make contributions, see Fig.1. Now I calculate the first diagram explicitly for the current $\eta_5(x)$ in Fig.1,

$$2T_N^{(\alpha_1)}(q^2) = -\frac{\text{Tr}(\lambda^a \lambda^b)(\bar{q}q)_N g_s^2 \mu^{2\epsilon} i}{12(2\pi)^D} \int d^D k \text{Tr} \left\{ i\gamma_\alpha k^\alpha \gamma_\beta i\gamma_5 \frac{i}{g^+ k - m_c} \frac{-i\delta_{\alpha\beta} g_{a\beta}}{k^2} \right\}$$

$$= -\frac{4 \text{Dm}_c(\bar{q}q)_N g_s^2 \mu^{2\epsilon} i}{3(2\pi)^D} \frac{\partial (-2\pi i)^2}{\partial t} \int_{m_c^2}^{\infty} ds \frac{d^D k \delta \left( k^2 - t \right) \delta \left( (k + q)^2 - m_c^2 \right)}{s - q^2} \bigg|_{t=0}$$

$$= -\frac{4 \text{Dm}_c(\bar{q}q)_N g_s^2 \mu^{2\epsilon} [1 + \epsilon(\log 4\pi - \gamma_E)]}{12\pi^2} \int_{m_c^2}^{\infty} ds \frac{1}{s - q^2} \frac{s + m_c^2}{s^{1-\epsilon}(s - m_c^2)^{1+2\epsilon}}, \quad (25)$$

where I have used Cutkosky’s rule to obtain the QCD spectral density. There exists infrared divergence at the end point $s = m_c^2$. It is difficult to carry out the integral over $s$, I can perform the Borel transform $B_{M^2}$ firstly, then carry out the integral over $s$,

$$B_{M^2}2T_N^{(\alpha_1)}(q^2) = -\frac{\text{Dm}_c(\bar{q}q)_N g_s^2 \mu^{2\epsilon} [1 + \epsilon(\log 4\pi - \gamma_E)]}{12\pi^2 M^2} \int_{m_c^2}^{\infty} ds \frac{1}{s - q^2} \frac{s + m_c^2}{s^{1-\epsilon}(s - m_c^2)^{1+2\epsilon}} \exp\left(-\frac{s}{M^2}\right)$$

$$= \frac{m_c(\bar{q}q)_N g_s^2}{3\pi^2 M^2} \exp\left(-\frac{m_c^2}{M^2}\right) \left(1 - \log 4\pi + \gamma_E\right) + \frac{m_c(\bar{q}q)_N g_s^2}{3\pi^2 M^2} \Gamma(0, \frac{m_c^2}{M^2})$$

$$- \frac{m_c(\bar{q}q)_N g_s^2}{6\pi^2 M^2} \exp\left(-\frac{m_c^2}{M^2}\right) + \frac{m_c(\bar{q}q)_N g_s^2}{3\pi^2 M^2} \exp\left(-\frac{m_c^2}{M^2}\right) \log \frac{m_c^2}{M^2}. \quad (26)$$

where

$$\Gamma(0, x) = e^{-x} \int_0^\infty dt \frac{1}{t + x} e^{-t}. \quad (27)$$

Other diagrams are calculated analogously, I regularize the divergences in $D = 4 - 2\epsilon$ dimension, then remove the ultraviolet divergences through renormalization and absorb the infrared divergences into the quark condensate $\langle \bar{q}q \rangle_N$.

I calculate the contributions of other condensates at the tree level, the calculations are straight-
forward and cumbersome. In calculations, I use the following formulas,

\[
\langle q_\alpha(x) \bar{q}_\beta(0) \rangle_N = -\frac{1}{4} \left[ \left( \langle \bar{q} q \rangle_N + x^\mu \langle \bar{q} D_\mu q \rangle_N + \frac{1}{2} x^{\mu \nu} \langle \bar{q} D_\mu D_\nu q \rangle_N + \cdots \right) \delta_{\alpha \beta} 
+ \left( \langle \bar{q} \gamma^\lambda q \rangle_N + x^\mu \langle \bar{q} \gamma^\lambda D_\mu q \rangle_N + \frac{1}{2} x^{\mu \nu} \langle \bar{q} \gamma^\lambda D_\mu D_\nu q \rangle_N + \cdots \right) \gamma_{\lambda \alpha \beta} \right],
\]

and

\[
\langle g_s q_\alpha^i \bar{q}_\beta^j G^a_{\mu \nu} \rangle_N = -\frac{1}{96} \frac{\lambda_s^2}{2} \left\{ \langle g_s \bar{q} \sigma G q \rangle_N \left[ \sigma_{\mu \nu} + i (u_\mu \gamma_\nu - u_\nu \gamma_\mu) \right] \right\}_{\alpha \beta} + \langle g_s \bar{q} \gamma^5 \sigma G q \rangle_N 
\left[ \sigma_{\mu \nu} + i (u_\mu \gamma_\nu - u_\nu \gamma_\mu) \right]_{\alpha \beta} - 4 \langle \bar{q} u D_\mu D_\nu q \rangle_N \left[ \sigma_{\mu \nu} + 2i (u_\mu \gamma_\nu - u_\nu \gamma_\mu) \right]_{\alpha \beta},
\]

where \( D_\mu = \partial_\mu - ig_s \frac{\lambda_s^2}{2} G^a_{\mu \nu} \),

\[
\langle \bar{q} \gamma_\mu q \rangle_N = \langle \bar{q} q \rangle_N u_\mu , \\
\langle \bar{q} D_\mu q \rangle_N = \langle \bar{q} u D_\mu q \rangle_N u_\mu = 0 , \\
\langle \bar{q} \gamma_\mu D_\nu q \rangle_N = \frac{4}{3} \langle \bar{q} \bar{u} D_\mu q \rangle_N \left( u_\mu u_\nu - \frac{1}{4} g_{\mu \nu} \right), \\
\langle \bar{q} D_\mu D_\nu q \rangle_N = \frac{4}{3} \langle \bar{q} u D_\mu D_\nu q \rangle_N \left( u_\mu u_\nu - \frac{1}{4} g_{\mu \nu} \right) - \frac{1}{6} \langle g_s \bar{q} \gamma_5 \sigma G q \rangle_N \left( u_\mu u_\nu - g_{\mu \nu} \right), \\
\langle \bar{q} \gamma_\lambda D_\mu D_\nu q \rangle_N = 2 \langle \bar{q} \bar{u} D_\mu D_\nu q \rangle_N \left[ u_\lambda u_\mu u_\nu - \frac{1}{6} \left( u_\lambda g_{\mu \nu} + u_\mu g_{\lambda \nu} + u_\nu g_{\lambda \mu} \right) \right] \\
- \frac{1}{6} \langle g_s \bar{q} \gamma^5 \sigma G q \rangle_N \left( u_\lambda u_\mu u_\nu - u_\lambda g_{\mu \nu} \right),
\]

and

\[
\langle G^a_{\alpha \beta} G^b_{\mu \nu} \rangle_N = \delta^{ab}_{96} \langle G G \rangle_N \left( g_{3 \alpha \beta}, \bar{g}_{3 \beta \alpha} \right) - O ( (E^2 + B^2)^N ).
\]

Once analytical results at the level of quark-gluon degree’s of freedom are obtained, then I set \( \omega^2 = q^2 \), and take the quark-hadron duality below the threshold \( s_0 \), and perform the Borel transform with respect to the variable \( Q^2 = -\omega^2 \), finally obtain the following QCD sum rules:

\[
a C_a + b C_b = C_f ,
\]

\[
C_a = \frac{1}{M^2} \exp \left( -\frac{m_D^2}{M^2} \right) - \frac{s_0}{m_D^2} \exp \left( -\frac{s_0}{M^2} \right), \\
C_b = \exp \left( -\frac{m_D^2}{M^2} \right) - \frac{s_0}{m_D^2} \exp \left( -\frac{s_0}{M^2} \right),
\]

\[
C_f = \frac{2 m_N (m_H + m_N)}{(m_H + m_N)^2 - m_D^2} \left( \frac{f_D m_D^2 g_G N H}{m_c} \right)^2 \left[ \frac{1}{M^2} - \frac{1}{m_D^2 - (m_H + m_N)^2} \right] \exp \left( -\frac{m_D^2}{M^2} \right) \\
+ \frac{1}{(m_H + m_N)^2 - m_D^2} \exp \left( -\frac{(m_H + m_N)^2}{M^2} \right) - \frac{m_c \bar{\langle q q \rangle}_N}{2} \left[ 1 + \frac{\alpha_s}{\pi} \left( 6 - \frac{4 m_c^2}{3 M^2} \right) \right] \exp \left( -\frac{m_D^2}{M^2} \right) \\
- \frac{2}{3} \left( 1 - \frac{m_c^2}{M^2} \right) \log \frac{m_c^2}{\mu^2} - 2 \Gamma \left( 0, \frac{m_c^2}{M^2} \right) \exp \left( \frac{m_c^2}{M^2} \right) \right] \exp \left( -\frac{m_D^2}{M^2} \right) \\
+ \left\{ \frac{1}{2} \left[ - \frac{1}{2} \left( 1 - \frac{m_c^2}{M^2} \right) \langle q \gamma_\mu i D_\mu q \rangle_N + \frac{4 m_c}{M^2} \left( 1 - \frac{m_c^2}{2 M^2} \right) \langle \bar{q} i D_\mu i D_\mu q \rangle_N + \frac{1}{12} \frac{m_c^2}{\pi} \right] \right\} \exp \left( -\frac{m_D^2}{M^2} \right),
\]
for the current $\eta(x)$,
\[ C_i \rightarrow C_i (\text{with } m_N \rightarrow -m_N, \ m_c \rightarrow -m_c, \ D \rightarrow D_0), \quad (35) \]
for the current $\eta(x)$,
\[ C_a = \frac{1}{M^2} \exp \left( -\frac{m_D^2}{M^2} \right) - \frac{s_0}{m_D^2} \exp \left( -\frac{s_0}{M^2} \right), \]
\[ C_b = \exp \left( -\frac{m_D^2}{M^2} \right) - \frac{s_0}{m_D^2} \exp \left( -\frac{s_0}{M^2} \right), \quad (36) \]
\[ C_f = \frac{2m_N(m_H + m_N)}{(m_H + m_N)^2 - m_D^2} (f_D m_D g_{D^*} D_{NH})^2 \left\{ \frac{1}{M^2} - \frac{1}{m_D^2 - (m_H + m_N)^2} \right\} \exp \left( -\frac{m_D^2}{M^2} \right) \]
\[ + \frac{1}{(m_H + m_N)^2 - m_D^2} \exp \left( -\frac{(m_H + m_N)^2}{M^2} \right) \frac{m_c}{2} \left( m_{(\bar{q}q)} N \right) \exp \left( -\frac{m_D^2}{M^2} \right) \]
\[ + \frac{2}{3} \left( 2 + \frac{m_D^2}{M^2} \right) \log \frac{m_D^2}{\mu^2} - \frac{2m_D^2}{3M^2} \Gamma \left( 0, \frac{m_D^2}{M^2} \right) \exp \left( -\frac{m_D^2}{M^2} \right) \]
\[ + \frac{1}{2} \left( -\frac{4(q^i D_0 q) N}{3} + \frac{2m_N(q^i D_0 q) N}{M^2} + \frac{2m_c(q_{\overline{q}q}, \sigma G q) N}{3M^2} + \frac{16m_c(q i D_0 i D_0 q) N}{3M^2} \right) \]
\[ - \frac{2m^3(q i D_0 i D_0 q) N}{M^4} - \frac{1}{12} \left( \frac{\alpha_s G}{\pi} \right) N \exp \left( -\frac{m_D^2}{M^2} \right), \quad (37) \]
for the current $\eta_Q(x)$, where $i = a, b, f$. In this article, I neglect the contributions from the heavy quark condensates $\langle \bar{Q}Q \rangle, \langle \bar{Q}Q \rangle = -\frac{1}{12\pi m_Q} \left( \frac{\alpha_s G}{\pi} \right)$ up to the order $O(\alpha_s)$ (here I count the condensate $\langle \bar{Q}Q \rangle$ as of the order $O(\alpha_s)$), the heavy quark condensates have practically no effect on the polarization functions, for detailed discussions about this subject, one can consult Ref.\[25\]. In Ref.\[25\], Buchheim, Hilger and Kämpfer study the contributions of the condensates involve the heavy quarks in details, the results indicate that those condensates are either suppressed by the heavy quark mass $m_Q$ or by the additional factor $\frac{G}{\pi^2}$ (or $g_s^2/(4\pi)^2$). Neglecting the in-medium effects on the heavy quark condensates cannot affect the predictions remarkably, as the main contributions come from the terms $\langle \bar{q}q \rangle N$.

Differentiate above equation with respect to $\tau$, then eliminate the parameter $b (a)$, I can obtain the QCD sum rules for the parameter $a (b)$,
\[ a = \frac{C_f (-\frac{d}{d\tau}) C_b - C_b (-\frac{d}{d\tau}) C_f}{C_a (-\frac{d}{d\tau}) C_b - C_b (-\frac{d}{d\tau}) C_a}; \]
\[ b = \frac{C_f (-\frac{d}{d\tau}) C_a - C_a (-\frac{d}{d\tau}) C_f}{C_b (-\frac{d}{d\tau}) C_a - C_a (-\frac{d}{d\tau}) C_b}. \quad (39) \]
With the simple replacements $m_c \rightarrow m_b, \ D/D_0/D^*/D_1 \rightarrow B/B_0/B^*/B_1, \ \Lambda_c \rightarrow \Lambda_b$ and $\Sigma_c \rightarrow \Sigma_b$, I can obtain the corresponding the QCD sum rules for the bottom mesons in the nuclear matter.

3 Numerical results and discussions

At the normal nuclear matter with the saturation density $\rho_N = \rho_0 = \frac{2\rho^2}{3\pi^2},$ where the Fermi momentum $p_F = 0.27\text{GeV}$ is a small quantity, the condensates $\langle \mathcal{O} \rangle_{\rho_N}$ in the nuclear matter can
be approximated as $\langle O \rangle_{\rho N} = \langle O \rangle + \frac{\rho_{\pi}}{m_d} \langle O \rangle_N$, the terms proportional to $p_F^2$, $p_F^3$, $p_F^4$, ... can be neglected safely, where the $\langle O \rangle = \langle 0 | O | 0 \rangle$ and $\langle O \rangle_N = \langle N | O | N \rangle$ denote the vacuum condensates and nuclear matter induced condensates, respectively [11].

The input parameters at the QCD side are taken as $\rho_{\pi} = (0.11 \, \text{GeV})^3$, $\langle q \bar{q} \rangle_N = \frac{\sigma_\pi}{m_d + m_u} (2m_N)$, $\langle (\bar{c}cGq) \rangle_N = -0.65 \, \text{GeV}(2m_N)$, $\sigma_\pi = 45 \, \text{MeV}$, $m_u + m_d = 12 \, \text{MeV}$, $\langle q^i i D q \rangle_N = 0.18 \, \text{GeV}(2m_N)$, $\langle (\bar{q} \bar{p} \sigma Gq) \rangle_N = 3.0 \, \text{GeV}^2(2m_N)$, $\langle (\bar{q} \bar{q} i D q) \rangle_N + \frac{1}{8} \langle q \bar{q} \sigma Gq \rangle_N = 0.3 \, \text{GeV}^2(2m_N)$, $m_N = 0.94 \, \text{GeV}$ [11], $m_c = (1.3 \pm 0.1) \, \text{GeV}$, $m_b = (4.7 \pm 0.1) \, \text{GeV}$, $\alpha_s = 0.45$ and $\mu = 1 \, \text{GeV}$. If we take normalization $\langle N(p) | N(p') \rangle = (2\pi)^3 \delta^4(p - p')$, then $\langle O \rangle_{\rho N} = \langle O \rangle + \rho_{\pi} \langle O \rangle_N$, the unit $2m_N$ in the brackets in the values of the condensates $\langle q \bar{q} \rangle_N$, $\langle (\bar{c}cGq) \rangle_N$, ..., disappears.

The parameters at the hadronic side are taken as $m_D = 1.870 \, \text{GeV}$, $m_B = 5.280 \, \text{GeV}$, $m_{B_0} = 5.740 \, \text{GeV}$, $m_{D^*} = 2.010 \, \text{GeV}$, $m_{B^*} = 5.325 \, \text{GeV}$, $m_{D_{1s}} = 5.420 \, \text{GeV}$, $m_{B_{1s}} = 5.750 \, \text{GeV}$, $f_D = 0.210 \, \text{GeV}$, $f_B = 0.190 \, \text{GeV}$, $f_{D_0} = 0.334 \frac{m_c}{m_{D_{1s}}} \, \text{GeV}$, $f_{B_0} = 0.280 \frac{m_b}{m_{B_{1s}}} \, \text{GeV}$, $f_{D^*} = 0.195 \, \text{GeV}$, $f_{B^*} = 0.305 \, \text{GeV}$, $f_{B_{1s}} = 0.255 \, \text{GeV}$, $s_{D_{1s}}^0 = (6.2 \pm 0.5) \, \text{GeV}^2$, $s_{B_{1s}}^0 = (33.5 \pm 1.0) \, \text{GeV}^2$, $s_{D_{1s}}^1 = (6.5 \pm 0.5) \, \text{GeV}^2$, $s_{B_{1s}}^1 = (35.0 \pm 1.0) \, \text{GeV}^2$, $s_{D_{0s}}^0 = (8.0 \pm 0.5) \, \text{GeV}^2$, $s_{B_{0s}}^0 = (39.0 \pm 1.0) \, \text{GeV}^2$, $s_{D_{0s}}^1 = (8.5 \pm 0.5) \, \text{GeV}^2$ and $s_{B_{0s}}^1 = (39.0 \pm 1.0) \, \text{GeV}^2$, which are determined by the conventional two-point correlation functions using the QCD sum rules [14] [26]. I neglect the uncertainties of the decay constants to avoid double counting as the main uncertainties of the decay constants originate from the uncertainties of the continuum threshold parameters $s_0$.

The value of the strong coupling constant $g_{D_{1s}N}$ is $g_{D_{1s}N} = 6.74$ from the QCD sum rules [27], while the average value of the strong coupling constants $g_{D_{1s}N}$ and $g_{D_{1s}N}$ from the light-cone QCD sum rules is $\frac{g_{D_{1s}N} + g_{D_{1s}N}}{2} = 6.775$ [28], those values are consistent with each other. The average value of the strong coupling constants $g_{D_{1s}N}$ and $g_{D_{1s}N}$ from the light-cone QCD sum rules is $\frac{g_{D_{1s}N} + g_{D_{1s}N}}{2} = 3.86$ [28]. In this article, I take the approximation $g_{D_{1s}N} = g_{D_{1s}N} = g_{B_{1s}N} \approx g_{B_{1s}N} \approx g_{D_{1s}N} \approx g_{D_{1s}N} \approx g_{D_{1s}N} \approx 6.74$ and $g_{D_{1s}N} = g_{D_{1s}N} = g_{D_{1s}N} \approx 3.86$.

In Figs.2-3, I plot the shifts of the masses and decay constants of the heavy mesons in the nuclear matter versus the Borel parameter $M^2$, respectively. From the figures, I can see that there appear platforms. In this article, I choose the Borel parameters $M^2$ according to the criterion that the uncertainties originate from the Borel parameters $M^2$ are negligible. The values of the Borel parameters $M^2$ are shown explicitly in Table 1. I can obtain the shifts of the masses and decay constants of the heavy mesons in the nuclear matter in the Borel windows, which are shown explicitly in Table 2. From Table 2, I can obtain the fractions of the shifts $\frac{\delta m_{D^*/B^*/B_{1s}}}{{m}_{D^*/B^*/B_{1s}}}$, $\frac{\delta m_{D_{1s}N}}{m_{D_{1s}N}}$, $\frac{\delta m_{D_{1s}N}/m_{D_{1s}N}}{m_{D_{1s}N}}$ ($5 - 15\%$) and $\frac{\delta f_{D^*/B^*/B_{1s}}}{{f}_{D^*/B^*/B_{1s}}}$, $\frac{\delta f_{D^*/B^*/B_{1s}}}{f_{D^*/B^*/B_{1s}}}$ (25 - 55\%).

Table 1: The Borel parameters in the QCD sum rules for the shifts of the masses and decay constants of the heavy mesons in the nuclear matter, the unit is $\text{GeV}^2$.

| $M^2$ | $\delta m_{D_{1s}}$ | $\delta m_{D^*}$ | $\delta m_{B_{1s}}$ | $\delta m_{B^*}$ | $\delta m_{B_{1s}}$ | $\delta m_{B_{1s}}$ |
|-------|------------------|------------------|------------------|------------------|------------------|------------------|
| 0.44  | 4.4 - 5.4        | 4.6 - 5.6        | 6.0 - 7.0        | 6.6 - 7.6        | 29 - 33          | 30 - 34          |
| 0.35  | 3.5 - 4.5        | 4.3 - 5.3        | 5.3 - 6.3        | 25 - 29          | 27 - 31          | 30 - 34          |
| 0.29  | 1.9 - 2.9        | 3.5 - 4.5        | 4.3 - 5.3        | 5.3 - 6.3        | 25 - 29          | 27 - 31          |
| 0.25  | 3.5 - 4.5        | 4.3 - 5.3        | 5.3 - 6.3        | 25 - 29          | 27 - 31          | 30 - 34          |
| 0.23  | 1.9 - 2.9        | 3.5 - 4.5        | 4.3 - 5.3        | 5.3 - 6.3        | 25 - 29          | 27 - 31          |

The mass-shifts of the negative (positive) parity mesons are negative (positive), the decays...
Figure 2: (Color online) The shifts of the masses of the heavy mesons in the nuclear matter versus the Borel parameter $M^2$, the I (II) denotes contributions up to the next-to-leading order (leading order) are included.
Figure 3: (Color online) The shifts of the decay constants of the heavy mesons in the nuclear matter versus the Borel parameter $M^2$, the I (II) denotes contributions up to the next-to-leading order (leading order) are included.
Table 2: The shifts of the masses and decay constants of the heavy mesons in the nuclear matter, where the NLO (LO) denotes contributions up to the next-to-leading order (leading order) are included, the unit is MeV.

|        | $\delta m_D$ | $\delta m_{D^*}$ | $\delta m_{D_0}$ | $\delta m_{D_1}$ | $\delta m_B$ | $\delta m_{B^*}$ | $\delta m_{B_0}$ | $\delta m_{B_1}$ |
|--------|--------------|------------------|------------------|------------------|--------------|------------------|------------------|------------------|
| NLO    | $-72$        | $-102$           | $80$             | $97$             | $-473$       | $-687$           | $295$            | $522$            |
| LO     | $-47$        | $-70$            | $54$             | $66$             | $-329$       | $-340$           | $209$            | $260$            |
| [13]   |              | $-48$            |                 |                 |              |                  |                  |                  |
| [16]   | $+45$        |                  |                 | $+60$            |              |                  |                  |                  |
| [15]   | $-46$        |                  |                 |                  |              |                  |                  |                  |

Table 3: The fractions of the shifts of the masses and decay constants of the heavy mesons in the nuclear matter, where the NLO (LO) denotes contributions up to the next-to-leading order (leading order) are included.

|        | $\delta f_D$ | $\delta f_{D^*}$ | $\delta f_{D_0}$ | $\delta f_{D_1}$ | $\delta f_B$ | $\delta f_{B^*}$ | $\delta f_{B_0}$ | $\delta f_{B_1}$ |
|--------|--------------|------------------|------------------|------------------|--------------|------------------|------------------|------------------|
| NLO    | $-6$         | $-26$            | $11$             | $31$             | $-71$        | $-111$           | $56$             | $134$            |
| LO     | $-4$         | $-18$            | $7$              | $21$             | $-48$        | $-55$            | $39$             | $67$             |
| [15]   | $-2$         |                  |                 |                  |              |                  |                  |                  |

Table 4: The uncertainties of the shifts of the masses and decay constants of the heavy mesons in the nuclear matter originate from the uncertainties of the heavy quark masses and continuum threshold parameters, where the unit is MeV.

|        | $\delta (\delta m_D)$ | $\delta (\delta m_{D^*})$ | $\delta (\delta m_{D_0})$ | $\delta (\delta m_{D_1})$ | $\delta (\delta m_B)$ | $\delta (\delta m_{B^*})$ | $\delta (\delta m_{B_0})$ | $\delta (\delta m_{B_1})$ |
|--------|------------------------|-----------------------------|-----------------------------|------------------------|------------------------|-----------------------------|-----------------------------|-----------------------------|
| $\delta m_Q$ | $\pm 14$            | $\pm 14$                    | $\pm 26$                    | $\pm 6$                | $\pm 18$               | $\pm 14$                    | $\pm 13$                    | $\pm 65$                    |
| $\delta s_0$  | $\pm 9$             | $\pm 14$                    | $\pm 12$                    | $\pm 6$                | $\pm 1$                | $\pm 14$                    | $\pm 13$                    | $\pm 65$                    |

|        | $\delta (\delta f_D)$ | $\delta (\delta f_{D^*})$ | $\delta (\delta f_{D_0})$ | $\delta (\delta f_{D_1})$ | $\delta (\delta f_B)$ | $\delta (\delta f_{B^*})$ | $\delta (\delta f_{B_0})$ | $\delta (\delta f_{B_1})$ |
|--------|------------------------|-----------------------------|-----------------------------|------------------------|------------------------|-----------------------------|-----------------------------|-----------------------------|
| $\delta m_Q$ | $\pm 1$             | $\pm 1$                     | $\pm 3$                     | $\pm 2$                | $\pm 1$               | $\pm 1$                     | $\pm 3$                     | $\pm 0$                     |
| $\delta s_0$  | $\pm 1$             | $\pm 7$                     | $\pm 4$                     | $\pm 8$                | $\pm 21$              | $\pm 25$                    | $\pm 15$                    | $\pm 34$                    |
of the high charmonium states to the $D\bar{D}$ and $D^*\bar{D}^*$ ($D_0\bar{D}_0$ and $D_1\bar{D}_1$) pairs are enhanced (suppressed) in the phase space, and we should take into account those effects carefully in studying the production of the $J/\psi$ so as to identifying the quark-gluon plasmas.

The currents $\bar{Q}q$ and $\bar{Q}i\gamma_5q$ (also $Q\gamma_\mu q$ and $\bar{Q}\gamma_\mu \gamma_5q$) are mixed with each other under the chiral transformation $q \to e^{i\alpha_\mu}q$, the currents $\bar{Q}q$, $\bar{Q}i\gamma_5q$, $\bar{Q}\gamma_\mu q$, $\bar{Q}\gamma_\mu \gamma_5q$ are not conserved in the limit $m_q \to 0$, it is better to take the doublets ($D, D_0$) and ($D^*, D_1$) as the parity-doublets rather than the chiral-doublets. The quark condensate $\langle \bar{q}q \rangle_{\rho_N}$ serves as the order parameter, and undergoes reduction in the nuclear matter, the chiral symmetry is partially restored; however, there appear new medium-induced condensates, which also break the chiral symmetry. In this article, the $\langle O \rangle_N$ are companied by the heavy quark masses $m_Q$, $m^2_Q$ or $m^3_Q$, the net effects cannot warrant that the chiral symmetry is monotonously restored with the increase of the $\rho_N$. When the $\rho_N$ is large enough, the order parameter $\langle \bar{q}q \rangle_{\rho_N} \to 0$, the chiral symmetry is restored, the Fermi gas approximation for the nuclear matter breaks down, and the parity-doublets maybe have degenerated masses approximately. In this article, we study the parity-doublets at the low $\rho_N$, the mass breaking effects of the parity-doublets maybe even larger, see Table 2.

In Table 5, I show the scattering lengths $a_D$, $a_{D^*}$, $a_{D_0}$, $a_{D_1}$, $a_B$, $a_{B^*}$, $a_{B_0}$, $a_{B_1}$ explicitly, the $a_D$, $a_{D^*}$, $a_B$, $a_{B^*}$ are negative, which indicate the interactions $D\bar{N}$, $D^*\bar{N}$, $B\bar{N}$, $B^*\bar{N}$ are attractive, the $a_{D_0}$, $a_{D_1}$, $a_{B_0}$, $a_{B_1}$ are positive, which indicate the interactions $D_0\bar{N}$, $D_1\bar{N}$, $B_0\bar{N}$, $B_0\bar{N}$ are repulsive. It is difficult (possible) to form the $D\bar{N}$ state with $J^P = \frac{3}{2}^-$ and the $D^*\bar{N}$ state with $J^P = \frac{3}{2}^-$ respectively.

In the present work and Refs. [13, 14, 15], the correlation functions are divided into a vacuum part and a static one-nucleon part, and the nuclear matter induced effects are extracted explicitly; while in Refs. [16, 17], the pole terms of the total hadronic spectral densities are parameterized as $\frac{\text{Im}\Pi(\omega, \mathbf{p})}{\delta}\omega = F_+\delta(\omega - m_+) - F_-\delta(\omega + m_-)$, where $m_{\pm} = m \pm \Delta m$ and $F_\pm = F \pm \Delta F$, and QCD sum rules for the mass center $\mathbf{m}$ and the mass splitting $\Delta m$ are obtained. In the leading order approximation, the present predictions for the $\delta m_D$ and $\delta m_B$ are compatible with that of Refs. [16, 17] and differ greatly from that of Refs. [16, 17], see Table 2. In Refs. [16, 17], the perturbative terms up to the order $O(\alpha_s)$ are taken into account, then the QCD sum rules are dominated by the perturbative terms $20$, which are not affected by the nuclear matter, the perturbative $O(\alpha_s)$ corrections play an important rule. Then the continuum contributions are well approximated by $\rho_{QCD}(\omega^2)\theta(\omega^2 - s_0)$, this is one of the advantage of Refs. [16, 17], however, we have to extract the mass modifications from the QCD sum rules where the condensates play a less important role, which impairs the predictive ability. In the present work and Refs. [13, 14, 15], the nuclear matter induced effects are extracted explicitly, this is the advantage, however, the QCD spectral densities are of the form $\delta(\omega^2 - m_Q^2)$, we have to model the continuum contributions with some functions, which have unknown uncertainties. More importantly, in the present work and Refs. [13, 14, 15], the Borel parameters $M^2$ are much larger than that chosen in Refs. [16, 17], different Borel windows lead to different predictions. On the other hand, the coupled-channel approach exhibits complicated structures of the $D$ mesons, the two parameters denoting in-medium mass and width are insufficient. The two approaches based on the QCD sum rules both have shortcomings, and the corresponding predictions can be confronted with the experimental data in

|     | $a_D$ | $a_{D^*}$ | $a_{D_0}$ | $a_{D_1}$ | $a_B$ | $a_{B^*}$ | $a_{B_0}$ | $a_{B_1}$ |
|-----|-------|-----------|-----------|-----------|-------|-----------|-----------|-----------|
| NLO | -1.1  | -1.5      | 1.3       | 1.6       | -8.9  | -12.9     | 5.6       | 9.9       |
| LO  | -0.7  | -1.1      | 0.9       | 1.1       | -6.2  | -6.4      | 4.0       | 5.0       |

Table 5: The heavy-meson-nucleon scattering lengths, where the NLO (LO) denotes contributions up to the next-to-leading order (the leading order) are included, the unit is fm.
the future.

The upcoming FAIR project at GSI provides the opportunity to study the in-medium properties of the charmoniums or charmed hadrons for the first time, however, the high mass of charmed hadrons requires a high momentum in the antiproton beam to produce them, the conditions for observing in-medium effects seem unfavorable, as the hadrons sensitive to the in-medium effects are either at rest or have a small momentum relative to the nuclear medium. We have to find processes that would slow down the charmed hadrons inside the nuclear matter, but this requires more detailed theoretical studies. Further theoretical studies on the reaction dynamics and on the exploration of the experimental ability to identify more complicated processes are still needed.

4 Conclusion

In this article, I calculate the contributions of the nuclear matter induced condensates up to dimension 5, take into account the next-to-leading order contributions of the in-medium quark condensate, and study the properties of the scalar, pseudoscalar, vector and axial-vector heavy mesons in the nuclear matter with the QCD sum rules in a systematic way. The present predictions for the shifts of the masses and decay constants can be confronted with the experimental data in the future, and we should take into account those effects carefully in studying the production of the $J/\psi$ (and $\Upsilon$) so as to identifying the quark-gluon plasmas. Furthermore, I study the heavy-meson-nucleon scattering lengths as a byproduct, and obtain the conclusion qualitatively about the possible existence of the heavy-meson-nucleon bound states.

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