Analysis of mass modifications of the vector and axialvector heavy mesons in the nuclear matter with the QCD sum rules

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Abstract

In this article, we calculate the mass modifications of the vector and axialvector mesons $D^*$, $B^*$, $D_1$ and $B_1$ in the nuclear matter with the QCD sum rules, and obtain the mass-shifts $\delta M_{D^*} = -71$ MeV, $\delta M_{B^*} = -380$ MeV, $\delta M_{D_1} = 72$ MeV, $\delta M_{B_1} = 264$ MeV, and the scattering lengths $a_{D^*} = -1.07$ fm, $a_{B^*} = -7.17$ fm, $a_{D_1} = 1.15$ fm and $a_{B_1} = 5.03$ fm for the $D^*N$, $B^*N$, $D_1N$ and $B_1N$ interactions, respectively.

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1 Introduction

The modifications of the hadron properties in the nuclear matter can affect the productions of the open-charmed mesons and the $J/\psi$ in the relativistic heavy ion collisions, the higher charmonium states, such as the $\psi'$, $\chi_{c1}$, $\chi_{c2}$, etc, are considered as the major source of the $J/\psi$ \cite{1}. The charmed mesons can obtain mass augments or reductions in the nuclear matter, if the mass modifications are large enough, the decays of the higher charmonium states to the charmed meson pairs can be facilitated or suppressed remarkably due to the available phase-space, thus the decays to the lowest state $J/\psi$ are greatly modified \cite{2}. For example, the higher charmonium states can decay to the $D\bar{D}$ pairs instead of decaying to the lowest state $J/\psi$, if the mass reductions of the $D$ and $\bar{D}$ mesons are large enough. On the other hand, the suppression of the $J/\psi$ production in the relativistic heavy ion collisions is considered as an important signature to identify the possible phase transition to the quark-gluon plasma \cite{3}. We should be careful before making definite conclusions.

The QCD sum rules is a powerful theoretical tool in studying the in-medium hadronic properties \cite{4}, and has been applied extensively to study the light-flavor hadrons and charmonium states in the nuclear matter \cite{5, 6, 7}. The works on the heavy mesons and heavy baryons are few, only the $D$, $B$, $D_0$, $B_0$, $\Lambda_c$, $\Lambda_b$, $\Sigma_c$ and $\Sigma_b$ are studied with the QCD sum rules \cite{8, 9, 10, 11, 12, 13}. The heavy mesons contain a heavy quark and a light quark, the existence of a light quark in the heavy mesons leads to large difference between the mass-shifts of the heavy mesons and heavy quarkonia in the nuclear matter. The former have large contributions from the light-quark condensates, while the latter are dominated by the gluon condensates \cite{7, 8, 9, 10, 11}. In this article, we study the mass modifications of the vector mesons $D^*$, $B^*$ and axialvector mesons $D_1$, $B_1$ in the nuclear matter using the QCD sum rules. The present predictions can be confronted with the experimental data from the CBM and PANDA collaborations in the future \cite{14, 15}.

The article is arranged as follows: we study the mass modifications of the vector and axialvector mesons $D^*$, $B^*$, $D_1$ and $B_1$ in the nuclear matter with the QCD sum rules in Sec.2; in Sec.3, we present the numerical results and discussions; and Sec.4 is reserved for our conclusions.

2 Mass modifications of the $D^*$, $B^*$, $D_1$ and $B_1$ in the nuclear matter with QCD sum rules

We study the mass modifications of the $D^*$ and $D_1$ mesons in nuclear matter with the two-point correlation functions $\Pi_{\mu\nu}(q)$,

$$\Pi_{\mu\nu}(q) = i \int d^4x \, e^{iq\cdot x} \langle T \{ J_\mu(x) J_\nu^\dagger(0) \} \rangle_{\rho_N}, \quad (1)$$

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where the $J_\mu(x)$ denotes the isospin averaged currents $\eta_\mu(x)$ and $\eta_{5\mu}(x)$,

$$\eta_\mu(x) = \eta^{\dagger}_\mu(x) = \frac{\bar{c}(x)\gamma_\mu q(x) + \bar{q}(x)\gamma_\mu c(x)}{2},$$
$$\eta_{5\mu}(x) = \eta^{\dagger}_{5\mu}(x) = \frac{\bar{c}(x)\gamma_\mu \gamma_5 q(x) + \bar{q}(x)\gamma_\mu \gamma_5 c(x)}{2},$$

(2)

which interpolate the vector and axialvector mesons $D^*$ and $D_1$, respectively, the $q$ denotes the $u$ or $d$ quark. The $\bar{c}q$ and $\bar{q}c$ mesons maybe obtain different mass modifications in the nuclear matter, just like the $K^+$ and $K^-$ mesons $[10]$, for example, Hilger et al observe that there exist particle-antiparticle mass splittings for the scalar and pseudoscalar mesons $[10, 11]$. In this article, we intend to study whether or not the decays of the higher charmonium states to the $D^*D^*$ and $D_1\bar{D}_1$ states are facilitated in the phase-space, and prefer the average values as the particle-antiparticle mass splittings cannot modify the total mass of the particle-antiparticle pair.

At the low nuclear density, the in-medium condensates $\langle O \rangle_{\rho_N}$,

$$\langle O \rangle_{\rho_N} = \langle O \rangle + \frac{\rho_N}{2M_N} \langle O \rangle_N,$$

(3)

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, the $\rho_N$ is the density of the nuclear matter $[6]$. Accordingly, the correlation functions $\Pi_{\mu\nu}(q)$ can be divided into a vacuum part $\Pi_{\mu\nu}^0(q)$ and a static one-nucleon part $T^N_{\mu\nu}(q)$ in the Fermi gas approximation for the nuclear matter $[6, 8]$,

$$\Pi_{\mu\nu}(q) = \Pi_{\mu\nu}^0(q) + \frac{\rho_N}{2M_N} T^N_{\mu\nu}(q),$$

(4)

where

$$T^N_{\mu\nu}(\omega, q) = i \int d^4xe^{i\omega x} \langle N(p)|T \{J_\mu(x)J^\dagger_\nu(0)\}|N(p)\rangle,$$

(5)

the $\langle N(p)\rangle$ denotes the isospin and spin averaged static nucleon state with the four-momentum $p = (M_N, 0)$, and normalized as $\langle N(p)|N(p')\rangle = (2\pi)^3 2\rho_0 \delta^4(p - p')$ $[8]$. The $T^N_{\mu\nu}(q)$ happen to be the current-nucleon forward scattering amplitudes. We can decompose the correlation functions $T^N_{\mu\nu}(\omega, q)$ as

$$T^N_{\mu\nu}(\omega, q) = T_N(\omega, q) \left( -g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2} \right) + \Pi^0_N(\omega, q) \frac{q_\mu q_\nu}{q^2},$$

(6)

according to Lorentz covariance, where the $T_N(\omega, q)$ denotes the contributions from the vector and axialvector mesons, and the $T^N(\omega, q)$ denotes the contributions from the scalar and pseudoscalar mesons. The interpolating currents $\eta_\mu(x)$ and $\eta_{5\mu}$ have non-vanishing couplings with the scalar and pseudoscalar mesons $D_0$ and $D$, respectively, i.e. $\langle 0|\eta_\mu(0)|D_0+\bar{D}_0\rangle = f_{D_0} q_\mu$ and $\langle 0|\eta_{5\mu}(0)|D+\bar{D}\rangle = if_D q_\mu$. We can exclude the contaminations by choosing the tensor structure $-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}$, and only take into account the vector and axialvector mesons through the definitions,

$$\langle 0|\eta_\mu(0)|D^*+\bar{D}^*\rangle = f_{D^*} M_{D^*} \epsilon_\mu,$$
$$\langle 0|\eta_{5\mu}(0)|D_1+\bar{D}_1\rangle = f_{D_1} M_{D_1} \epsilon_\mu,$$

(7)

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

In the limit of the 3-vector $q \to 0$, the correlation functions $T_N(\omega, q)$ can be related to the $D^*N$ and $D_1N$ scattering $T$-matrices, i.e. $T_{D^*N}(M_{D^*}, 0) = 8\pi(M_N + M_{D^*})a_{D^*}$ and $T_{D_1N}(M_{D_1}, 0) = 8\pi(M_N + M_{D_1})a_{D_1}$, where the $a_{D^*}$ and $a_{D_1}$ are the $D^*N$ and $D_1N$ scattering lengths, respectively.
Near the pole positions of the $D^*$ and $D_1$ mesons, the phenomenological spectral densities $\rho(\omega,0)$ can be parameterized with three unknown parameters $a, b$ and $c$ [8],
\[
\rho(\omega,0) = -\frac{f^2_{D^*/D_1} M^2_{D^*/D_1}}{\pi} \text{Im} \left[ \frac{T_{D^*/D_1,N}(\omega,0)}{(\omega^2 - M^2_{D^*/D_1} + i\varepsilon)^2} \right] + \cdots ,
\]
(8)
\[
= a \frac{d}{d\omega^2} \delta(\omega^2 - M^2_{D^*/D_1}) + b \delta(\omega^2 - M^2_{D^*/D_1}) + c \delta(\omega^2 - s_0),
\]
(9)
the terms denoted by $\cdots$ represent the continuum contributions. The first term denotes the double-pole term, and corresponds to the on-shell (i.e. $\omega^2 = M^2_{D^*/D_1}$) effects of the $T$-matrices,
\[
a = -8\pi (M_N + M_{D^*/D_1}) a_{D^*/D_1} f^2_{D^*/D_1} M^2_{D^*/D_1},
\]
(10)
and related with the mass-shifts of the $D^*$ and $D_1$ mesons through the relation
\[
\delta M^0_{D^*/D_1} = -\frac{\rho_N}{4M_N f^2_{D^*/D_1} M^3_{D^*/D_1}} a;
\]
(11)
the second term denotes the single-pole term, and corresponds to the off-shell (i.e. $\omega^2 \neq M^2_{D^*/D_1}$) effects of the $T$-matrices; and the third term denotes the continuum term or the remaining effects, where the $s_0$ is the continuum threshold.

In the limit $\omega \rightarrow 0$, the $T_N(\omega,0)$ is equivalent to the Born term $T^\text{Born}_{D^*/D_1,N}(\omega,0)$. We take into account the Born term at the phenomenological side,
\[
T_N(\omega^2) = T^\text{Born}_{D^*/D_1,N}(\omega^2) + \frac{a}{(M^2_{D^*/D_1} - \omega^2)^4} + \frac{b}{M^2_{D^*/D_1} - \omega^2} + \frac{c}{s_0 - \omega^2},
\]
(12)
with the constraint
\[
\frac{a}{M^2_{D^*/D_1}} + \frac{b}{M^2_{D^*/D_1}} + \frac{c}{s_0} = 0.
\]
(13)
The contributions from the intermediate spin-$\frac{3}{2}$ charmed baryons are zero in the soft-limit $q_\mu \rightarrow 0$ [17], and we only take into account the intermediate spin-$\frac{1}{2}$ charmed baryons in calculating the Born terms, and parameterize the hadronic matrix elements as
\[
\langle \Lambda_c/\Sigma_c(p-q)|D^*(-q)N(p)\rangle = \hat{U}_{\Lambda_c/\Sigma_c}(p-q) \left[ g_{\Lambda_c/\Sigma_c,D^*N} \not\!q + i g^T_{\Lambda_c/\Sigma_c,D^*N} \frac{\sigma^{\alpha\beta} q_\alpha q_\beta}{M_N + M_{\Lambda_c/\Sigma_c}} \right] U_N(p),
\]
\[
\langle \Lambda_c/\Sigma_c(p-q)|D_1(-q)N(p)\rangle = \hat{U}_{\Lambda_c/\Sigma_c}(p-q) \left[ g_{\Lambda_c/\Sigma_c,D_1N} \not\!q + i g^T_{\Lambda_c/\Sigma_c,D_1N} \frac{\sigma^{\alpha\beta} q_\alpha q_\beta}{M_N + M_{\Lambda_c/\Sigma_c}} \right] \gamma_5 U_N(p),
\]
(14)
where the $U_N$ and $\hat{U}_{\Lambda_c/\Sigma_c}$ are the Dirac spinors of the nucleon and the charmed baryons $\Lambda_c/\Sigma_c$, respectively; the $g_{\Lambda_c/\Sigma_c,D^*N}, g_{\Lambda_c/\Sigma_c,D_1N}, g^T_{\Lambda_c/\Sigma_c,D^*N}$ and $g^T_{\Lambda_c/\Sigma_c,D_1N}$ are the strong coupling constants in the vertexes. In the limit $q_\mu \rightarrow 0$, the strong coupling constants $g^T_{\Lambda_c/\Sigma_c,D^*N}$ and $g^T_{\Lambda_c/\Sigma_c,D_1N}$ have no contributions.

We draw the Feynman diagrams, calculate the Born terms and obtain the results
\[
T^\text{Born}_{D^*/N}(\omega,0) = \frac{2f^2_{D^*/D_1} M^2_{D^*/D_1} (M_H + M_N) g^2_{\bar{H}D^*/N}}{[\omega^2 - (M_H + M_N)^2] [\omega^2 - M^2_{D^*/D_1}]^2},
\]
\[
T^\text{Born}_{D_1,N}(\omega,0) = \frac{2f^2_{D_1} M^2_{D_1} (M_H - M_N) g^2_{\bar{H}D_1,N}}{[\omega^2 - (M_H - M_N)^2] [\omega^2 - M^2_{D_1}]^2},
\]
(15)
where the $H$ means either $\Lambda_c^+,$ $\Sigma_c^+$, $\Sigma_c^{++}$ or $\Xi_c^0$. The masses $M_{\Lambda_c} = 2.286$ GeV and $M_{\Sigma_c} = 2.454$ GeV from the Particle Data Group \cite{17}, we can take $M_H \approx 2.4$ GeV as the average value.

On the other hand, there are no inelastic channels for the $D^*N$ and $\bar{D}_1N$ interactions in the case of the charmed mesons $\bar{c}q$.

The scattering state $D^*N$ can translate to the scattering states $D^*N$, $\pi\Sigma_c$, $\eta\Lambda_c$, $DN$, $\pi\Lambda_c$, $\rho\Sigma_c$, $\rho\Lambda_c$, etc, we can take into account the infinite series of the intermediate baryon-meson loops with the Bethe-Salpeter equation to obtain the full $D^*N \to D^*N$ scattering amplitude, and the higher resonances, such as the $\Lambda_c(2595)$, $\Sigma_c(2800)$, etc, appear as dynamically generated baryon states \cite{19}. We can saturate the full $D^*N \to D^*N$ scattering amplitude with the tree-level Feynman diagrams of the exchanges of the higher-resonances including the negative spin-$\frac{1}{2}$ charmed baryons $\Lambda_c(2595)$ and $\Sigma_c(\frac{1}{2}^-)$ have the average value $M_{\Lambda_c} \approx 2.7$ GeV \cite{18, 20} and other excited spin-$\frac{1}{2}$ charmed baryons. The translations of the scattering state $D^*N$ to the ground states $\Lambda_c$ and $\Sigma_c$ are greatly facilitated in the phase-space, furthermore, the mass-shift $\delta M_{D^*}$ does not sensitive to contributions of the ground states $\Lambda_c$ and $\Sigma_c$, see Table 1; neglecting the loop-effects (or higher-resonance contributions) cannot impair the prediction power remarkably, in other words, the contaminations from the intermediate hadron states are small. As the scattering amplitudes contain both elastic and nonelastic parts, we use the scattering lengths and Born terms to parameterize the elastic and nonelastic scattering amplitudes respectively. The contributions of the scattering states (or continuum states) $\pi\Sigma_c$, $\eta\Lambda_c$, $DN$, $\pi\Lambda_c$, $\rho\Sigma_c$, $\rho\Lambda_c$, etc, are embodied in the Born terms.

We carry out the operator product expansion to the condensates up to dimension-5 at the large space-like region in the nuclear matter, and obtain the analytical expressions of the correlation functions at the level of quark-gluon degree’s of freedom,

\begin{equation}
\Pi_{\mu\nu}(q_0, \bar{q}) = (-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}) \sum_n C_n(q_0, \bar{q}) \langle O_n \rangle_{\rho N} + \cdots, \tag{16}
\end{equation}

where the $C_n(q_0, \bar{q})$ are the Wilson coefficients, the in-medium condensates $\langle O_n \rangle_{\rho N} = \langle O_n \rangle + \frac{\rho N}{2 M_N} \langle O_n \rangle_N$ at the low nuclear density, the $\langle O_n \rangle$ and $\langle O_n \rangle_N$ denote the vacuum condensates and nuclear matter induced condensates, respectively. One can consult Refs.\cite{5, 6} for the technical details in the operator product expansion. Then we collect the terms proportional to $\rho_N$ (or the nuclear matter induced condensates), take the quark-hadron duality,

\begin{equation}
T^N_{\mu\nu}(\omega, \bar{q}) = (-g_{\mu\nu} + \frac{q_\mu q_\nu}{q^2}) \sum_n C_n(\omega, \bar{q}) \langle O_n \rangle_N + \cdots, \tag{17}
\end{equation}

set $\omega^2 = q^2$, and perform the Borel transform with respect to the variable $Q^2 = -\omega^2$, finally obtain the following two QCD sum rules:

\begin{align*}
& a \left\{ \frac{1}{M^2} e^{-\frac{M_D^2}{M^2}} - \frac{s_0}{M_D^2} e^{-\frac{s_0^2}{M^2}} + b \left\{ \frac{1}{M_D^2} e^{-\frac{M_D^2}{M^2}} - \frac{s_0}{M_D^2} e^{-\frac{s_0^2}{M^2}} + \frac{2 f_D^2 M_D^2 (M_H + M_N) g_H^2}{(M_H + M_N)^2 - M_D^2} \right\} \right\} = \left\{ \frac{m_c(\bar{q}q)_N}{2} \right\} \\
& \frac{1}{3} \left\{ (M_H + M_N)^2 - M_D^2 \right\} e^{-\frac{M_D^2}{M^2}} - 1 M^2 - \frac{1}{M_D^2} e^{-\frac{M_D^2}{M^2}} \right\} \right\} e^{-\frac{m_c^2(q_i D_0 q_j)_N}{M^4}} \right\} \right\} e^{-\frac{m_c^2}{M^2}} \\
& \frac{-1}{24} \left( \frac{\alpha_s G_G}{\pi} \right)_N \int_0^1 dx \left( 1 + \frac{2 m_c^2}{2 M^2} \right) e^{-\frac{m_c^2}{M^2}} + \frac{1}{48 M^2} \left( \frac{\alpha_s G_G}{\pi} \right)_N \int_0^1 1 - x x \left( \frac{2 m_c^2}{M^2} \right) e^{-\frac{m_c^2}{M^2}}, \tag{18}
\end{align*}
mass-shifts of the external parameter $\rho$.

The approach developed for the light mesons still works for the heavy mesons in the nuclear matter, where we take the approximation $M_{H} = \frac{M_{c}+M_{h}}{2} \approx 5.7 \text{ GeV}$ \cite{13}.

The present approach was introduced by Koike and Hayashigaki to study the spin-isospin averaged meson-nucleon scattering lengths and the relevant mass-shifts for the $\rho$, $\omega$, $\phi$ mesons in the nuclear matter \cite{21}. The heavy mesons contain a heavy-quark and a light-quark. The existence of a heavy quark in the heavy mesons results in much difference between the in-medium properties of the heavy mesons and light mesons. The heavy quark interacts with the nuclear matter through the exchange of the intermediate gluons and the modifications of the gluon condensates in the nuclear matter are mild, while the modifications of the quark condensates in the nuclear matter are rather large. The relation in Eq.(3) is hold for the gluon condensates. We expect that the convergent behaviors of the heavy-light type interpolating currents are better than that of the light-light type interpolating currents, if the correlation functions $\Pi_{\mu\nu}(q)$ are expanded in terms of the external parameter $\rho_{N}$. The approach developed for the light mesons still works for the heavy mesons.

In Ref.\cite{22}, Klingl, Kaiser and Weise use an effective Lagrangian which combines chiral $SU(3)$ dynamics with vector meson dominance to calculate the forward vector-meson-nucleon scattering amplitudes, and take them as input parameters in the hadronic side of the QCD sum rules in nuclear matter, and observe a remarkable degree of consistency with the operator product expansion at the quark level. In Ref.\cite{23}, Leupold and Mosel study the electromagnetic current-current correlation functions in the nuclear matter, and expand the QCD sum rules in terms of the finite squared three-momentum $q^{2}$, and observe that the QCD sum rule can provide an interesting and non-trivial consistency check for the hadronic models, but cannot rule out the hadronic models which predict a different behavior of the vector mesons with different $q^{2}$. In this article, the hadronic side of the QCD sum rules of the order $\mathcal{O}(\rho_{N})$ consists of both elastic and inelastic scattering amplitudes, which are parameterized by the scattering lengths and the Born terms respectively, can lead to stable QCD sum rules with the suitable Borel parameters in a finite range, i.e. the predictions make sense.

3 Numerical results and discussions

In calculations, we have assumed that the linear density approximation is valid at the low nuclear density, $\langle \mathcal{O}\rangle_{\rho_{N}} = \langle 0|\mathcal{O}|0\rangle + \frac{\rho_{N}}{2M_{N}}\langle N|\mathcal{O}|N\rangle = \langle \mathcal{O}\rangle + \frac{\rho_{N}}{2M_{N}}\langle \mathcal{O}\rangle_{N}$ for a general condensate $\langle \mathcal{O}\rangle_{\rho_{N}}$ in the nuclear matter. The input parameters are taken as $\langle \bar{q}q\rangle_{N} = \frac{<\bar{q}q>}{m_{c}+m_{d}}(2M_{N}) = (-0.65 \pm 0.15)\text{ GeV}(2M_{N})$, $\langle q^{i}D_{0}\bar{q}\rangle_{N} = (0.18 \pm 0.01)\text{ GeV}(2M_{N})$, $\langle \bar{q}g_{a}Gq\rangle = 3.0\text{ GeV}^{2}(2M_{N})$.\cite{23}.
\[
\langle \bar{q}qD_0 D_0 q \rangle_N + \frac{i}{\pi} \langle \bar{q}q \sigma G q \rangle_N = 0.3 \text{GeV}^2(2M_N), \quad m_u + m_d = 12 \text{MeV}, \quad \sigma_N = (45 \pm 10) \text{MeV}, \quad M_N = 0.94 \text{GeV}, \quad \rho_N = (0.11 \text{GeV})^3 \quad [5], \quad m_c = (1.35 \pm 0.1) \text{GeV} \quad \text{and} \quad m_b = (4.7 \pm 0.1) \text{GeV} \quad \text{at the energy scale } \mu = 1 \text{GeV}.
\]

For the well established vector mesons \(D^*\) and \(B^*\), we take the values from the Particle Data Group, \(M_{D^*} = 2.01 \text{GeV}\) and \(M_{B^*} = 5.325 \text{GeV} \quad [13]\), the decay constants \(f_{D^*}\) and \(f_{B^*}\) are determined by the QCD sum rules, \(f_{D^*} = 0.270 \text{GeV}\) and \(f_{B^*} = 0.195 \text{GeV} \quad [24]\), where the threshold parameters are taken as \(s^0_{D^*} = (5 - 7) \text{GeV}^2\) and \(s^0_{B^*} = (33 - 37) \text{GeV}^2 \quad [24, 25]\), here we have neglected the uncertainties of the decay constants. We can take the threshold parameters as \(s^0_{D^*} = 6.5 \text{GeV}^2\) and \(s^0_{B^*} = 35 \text{GeV}^2\) to reproduce the values \(M_{D^*} = 2.01 \text{GeV}\), \(M_{B^*} = 5.325 \text{GeV}\), \(f_{D^*} = 0.270 \text{GeV}\) and \(f_{B^*} = 0.195 \text{GeV}\) approximately for the QCD sum rules in the vacuum. The mass of the axialvector meson \(D_0^0\) is \(M_{D_0^0} = (2427 \pm 26 \pm 25) \text{MeV}\) from the Particle Data Group \[15\], and the axialvector meson \(B_1\) has not been observed yet. We calculate the hadronic parameters of the axialvector mesons \(D_1\) and \(B_1\) using the QCD sum rules in the vacuum, and obtain the values \(M_{D_1} = 2.42 \text{GeV}\), \(M_{B_1} = 5.75 \text{GeV}\), \(f_{D_1} = 0.305 \text{GeV}\) and \(f_{B_1} = 0.255 \text{GeV}\) with the threshold parameters \(s^0_{D_1} = 8.5 \text{GeV}^2\) and \(s^0_{B_1} = 39 \text{GeV}^2\). The value \(M_{D_1} = 2.42 \text{GeV}\) reproduces the experimental data \(M_{D_1} = (2427 \pm 26 \pm 25) \text{MeV}\) well \[13\]. The prediction of the mass \(M_{B_1}\) satisfies the relation \(M_{B_1} - M_{D_1} \approx M_{D_1} - M_{D^*}\). In this article, the threshold parameters are taken as \(s^0_{D_1} = (6.5 \pm 0.5) \text{GeV}^2\), \(s^0_{B_1} = (8.5 \pm 0.5) \text{GeV}^2\), \(s^0_{D_1} = (35 \pm 1) \text{GeV}^2\) and \(s^0_{B_1} = (39 \pm 1) \text{GeV}^2\), respectively, which satisfy the relations \(s^0_{D_1,D_1,B_1} = (M_{D_1,D_1,B_1} + 0.4 \sim 0.6 \text{GeV})^2\) and \(s^0_{B_1} = (M_{B_1} + 0.5 \sim 0.7 \text{GeV})^2\). In general, the energy gap between the ground state and the first radial excited state is about 0.5 GeV. For the explicit expressions of the QCD sum rules in the vacuum derived from the correlation functions \(\Pi_{\mu\nu}(q)\), one can consult Ref.\[26\] and the references therein.

The value of the strong coupling constant \(g_{D_1N,N}\) is \(g_{D_1N,N} = 6.74\) from the QCD sum rules \[27\], while the average value of the strong coupling constants \(g_{D_1N,N}\) and \(g_{D_1N,N}\) from the light-cone QCD sum rules is \(2g_{D_1N,N} = 6.775 \quad [28]\), those values are consistent with each other. The average value of the strong coupling constants \(g_{D_1N,N}\) and \(g_{D_1N,N}\) from the light-cone QCD sum rules is \(2g_{D_1N,N} = 6.775 \quad [28]\). In this article, we take the approximation \(g_{D_1N,N} \approx g_{D_1N,N} / 2\).

In Fig.1, we plot the mass-shifts \(\delta M\) versus the Borel parameter \(M^2\) at large intervals. From the figure, we can see that the values of the mass-shifts are rather stable with variations of the Borel parameter at the intervals \(M^2 = (4.5 - 5.4) \text{GeV}^2\), \((6.5 - 7.6) \text{GeV}^2\), \((22 - 24) \text{GeV}^2\) and \((34 - 37) \text{GeV}^2\) for the \(D^*, D_1, B^*\) and \(B_1\) mesons, respectively; in other words, the uncertainties originate from the Borel parameter \(M^2\) are less than 1%. The main contributions come from the terms \(\pm m_{c,(\bar{q}q)} N\) and \(\pm m_{b,(\bar{q}q)} N\), see Eqs.\(18-19\) and Fig.2, the operator product expansion is well convergent. The spectral densities at the level of the quark-gluon degrees of freedom consist of the medium-induced condensates, and have the form \(A_1 \delta(s - m_{c/(b)}^2) + A_2 \delta(s - \bar{m}_{c/(b)}^2)\), where the \(A_1\) and \(A_2\) denote the coefficients, we carry out the integrals,

\[
\int_{m_{c/(b)}^2}^{s_0} ds \left[ A_1 \delta(s - m_{c/(b)}^2) + A_2 \delta(s - \bar{m}_{c/(b)}^2) \right] e^{-\frac{s}{M^2}},
\]

(20)

to obtain the right side of the QCD sum rules in Eqs.\(18-19\), there are no perturbative terms to approximate the continuum contributions at the regions \(s > s_0\). At the phenomenological side, the exponential factors

\[
e^{-\frac{s_{c/(b)}}{M^2}} = e^{-(1.20-1.44)}, e^{-(1.12-1.31)}, e^{-(1.46-1.59)}, e^{-(1.05-1.15)},
\]

(21)

at the intervals \(M^2 = (4.5 - 5.4) \text{GeV}^2\), \((6.5 - 7.6) \text{GeV}^2\), \((22 - 24) \text{GeV}^2\) and \((34 - 37) \text{GeV}^2\) for the \(D^*, D_1, B^*, B_1\) mesons respectively, where we take the central values of the threshold parameters, the corresponding exponential factors of the ground states are

\[
e^{-\frac{s_{c/(b)}}{M^2}} = e^{-(0.75-0.90)}, e^{-(0.77-0.90)}, e^{-(1.18-1.29)}, e^{-(0.89-0.97)},
\]

(22)
where the \( m \) stands for the \( D^*, D_1, B^*, B_1 \) mesons respectively; the continuum contributions are suppressed more efficiently. Furthermore, we expect that the couplings of an special interpolating current to the excited states are more weak than that to the ground state mesons. For example, the decay constants of the pseudoscalar mesons \( \pi(140) \) and \( \pi(1800) \) have the hierarchy: \( f_{\pi(1300)} \ll f_{\pi(140)} \) from the Dyson-Schwinger equation [29], the lattice QCD [30], the QCD sum rules [31], etc, or from the experimental data [32].

We can take the Borel windows as \( M^2 = (4.5 - 5.4) \) GeV\(^2\), \( (6.5 - 7.6) \) GeV\(^2\), \( (22 - 24) \) GeV\(^2\) and \( (34 - 37) \) GeV\(^2\) for the \( D^*, D_1, B^* \) and \( B_1 \) mesons, respectively, and obtain mass-shifts \( \delta M_{D^*} = -72^{+22}_{-23} \) MeV, \( \delta M_{D_1} = 72^{+22}_{-23} \) MeV, \( \delta M_{B_1} = 264^{+70}_{-66} \) MeV, respectively; and the scattering lengths \( a_{D^*} = -1.07^{+0.30}_{-0.34} \) fm, \( a_{B^*} = -7.17^{+1.53}_{-1.71} \) fm, \( a_{D_1} = 1.15^{+0.35}_{-0.32} \) fm and \( a_{B_1} = 5.03^{+1.46}_{-1.31} \) fm for the \( D^* \), \( B^* \), \( D_1 \) and \( B_1 \) interactions, respectively. For the technical details in analyzing the uncertainties, one can consult Ref. [33].

In Fig. 3, we plot the mass-shifts \( \delta M \) versus the Borel parameter \( M^2 \) and the strong coupling constants \( g^2 \). From the figure, we can see that the mass-shifts decrease (increase) monotonously with increase of the squared strong coupling constants \( g^2 \) for the vector mesons \( D^* \) and \( B^* \) (axial-vector mesons \( D_1 \) and \( B_1 \)) in the Borel windows. The precise values of the mass-shifts and scattering lengths are presented in Table 1.

Although the present QCD sum rules are not stable with variations of the Borel parameters, the uncertainties originate from the Borel parameters are less than 1% in the Borel windows, i.e. we choose suitable platforms to avoid large uncertainties. On the other hand, we can take moments of the correlation functions and derive QCD sum rules to study the mass-shifts as in Ref. [10], the pole terms of the hadronic spectral densities are parameterized as \( \frac{\text{Im}\Pi(\omega,0)}{\pi} = F_\pm \delta(\omega - M_\pm) - F_\mp \delta(\omega + M_\mp) \), where \( M_\pm = M \pm \Delta M \) and \( F_\pm = F \pm \Delta F \), and QCD sum rules for the mass center \( M \) and the mass splitting \( \Delta M \) are obtained. For the pseudoscalar \( D, D \) mesons, Hayashigaki obtains the mass-shift \( \delta M_D = -50 \) MeV [8], while Hilger, Thomas and Kampfer obtain the mass-shift \( \delta M_D = +45 \) MeV [10]. For scalar \( D_0, D_0 \) mesons, the mass-shift \( \delta M_{D_0} = M - M_{D_0} < 0 \) obtained by Hilger and Kampfer [11] differs from the result \( \delta M_{D_0} = +69 \) MeV obtained by Wang and Huang [9]. In Ref. [12], Hilger, Kampfer and Leupold study the chiral partners of charmed mesons in the nuclear matter, and focus on the differences between the pseudoscalar and scalar as well as vector and axialvector D mesons and derive the corresponding Weinberg type sum rules, while the mass-shifts are not presented.

In the limit \( m_q \to 0 \), the quark condensate \( \langle \bar{q}q \rangle \) serves as the order parameter and indicates that the chiral symmetry is broken. The quark condensate undergoes reduction in the nuclear matter, \( \langle \bar{q}q \rangle_{\rho N} = \langle \bar{q}q \rangle + \frac{\rho}{2M_N} \langle \bar{q}q \rangle_{\rho N} \), the chiral symmetry is partially restored, for example, the in-medium nucleon mass \( M_N^m \) can be approximated as \( M_N^m = -\frac{8\pi^2}{3M_N^3} \langle \bar{q}q \rangle_{\rho N} \) and the mass reduction is rather large: on the other hand, there appear new medium-induced condensates, for example, the \( \langle \bar{q}iD_0D_0 \rangle, \langle \bar{q}g_\sigma Gq \rangle, \) etc, which also break the chiral symmetry. In the present case, the medium-induced condensates are associated with the large heavy quark masses \( m_Q, m_Q^2, m_Q^3, m_Q^4 \), the net effects do not always warrant that the chiral symmetry is monotonously restored with the increase of the density of the nuclear matter. The light vector current \( \bar{q}\gamma_\mu q \) and axialvector current \( \bar{q}g_\mu^\gamma q \) are invariant under the chiral transformation \( q \to e^{i\alpha_5}q \), however, the heavy vector current \( Q_\gamma q \) and axialvector current \( Q_\gamma^\gamma q \) are mixed with each other under the transformation, the heavy quark currents \( Q_\gamma^\gamma q \) and \( Q_\gamma^\gamma q \) are not conserved in the limit \( m_q \to 0 \), it is better to take the doublets \( (D^*, D_1), (B^*, B_1) \) as the parity-doublets rather than the chiral-doublets. If we take into account the flavor \( SU(3) \) symmetry of the light quarks, the chiral \( SU(3)_L \times SU(3)_R \) transformations require that the ground states \( (D^{*0}, D^{*-}, D^{*+}) \) and \( (B^{*+}, B^{*0}, B^{*0}) \) have their chiral partners \( (D^{*0}_L, D^{*0}_R, D^{*0}_T) \) and \( (B^{*0}_L, B^{*0}_R, B^{*0}_T) \), respectively, those parity-doublets are chiral-
doublets. When the density of the nuclear matter is large enough, the order parameter \( \langle \bar{q}q \rangle_{NS} \rightarrow 0 \), the chiral symmetry is restored, the Fermi gas approximation for the nuclear matter does not survive, there are free of the non-perturbative contributions from the condensates, and the parity-doublets (or chiral-doublets) maybe have degenerated masses approximately. In the present case, we study the parity-doublets (or chiral-doublets) in the low nuclear density, the mass breaking effects of the parity-doublets (or chiral-doublets) maybe even larger.

The axialvector current \( \bar{u}(x)\gamma_\mu\gamma_5 c(x) \) interpolates the axialvector meson \( D_1(2430) \) has non-vanishing coupling with the scattering state \( D^*\pi \), in the soft \( \pi \) limit, the coupling constant \( \lambda_{D^*\pi} \) can be estimated as

\[
\langle 0 | \bar{u}(x)\gamma_\mu\gamma_5 c(x) | D^*\pi \rangle = -\frac{i}{f_\pi} \langle 0 | [Q_5, \bar{u}(x)\gamma_\mu\gamma_5 c(x)] | D^* \rangle
\]

\[
= -\frac{i}{f_\pi} \langle 0 | \bar{d}(x)\gamma_\mu c(x) | D^* \rangle = -\frac{i f_{D^*} M_{D^*}\epsilon_\mu}{f_\pi} = i\lambda_{D^*\pi} \epsilon_\mu, \tag{23}
\]

\[
= f_{D_1} M_{D_1} \epsilon_\mu \frac{i}{p^2 - M_{D_1}^2} \langle D_1 | D^* \rangle, \tag{24}
\]

where the axial-charge \( Q_5 = \int d^3yd^3(\gamma_5 u(y) \right) \) and the \( \epsilon_\mu \) and \( \epsilon_\mu \) are the polarization vectors of the vector and axialvector mesons \( D^* \) and \( D_1 \), respectively \([34]\), the formula in Eq.(24) survives beyond the soft \( \pi \) limit. The coupling constant \( \lambda_{D^*\pi} \) is a large quantity and cannot be neglected. The rescatterings

\[
D^*\pi \rightarrow D^*\pi,
\]

\[
D^*\pi \rightarrow D^*\pi, D_1^* K, D^*\eta \rightarrow D^*\pi,
\]

\[
D^*\pi \rightarrow D^*\pi, D_1^* K, D^*\eta \rightarrow D^*\pi, D_1^* K, D^*\eta \rightarrow D^*\pi,
\]

\[
D^*\pi \rightarrow D^*\pi, D_1^* K, D^*\eta \rightarrow D^*\pi, D_1^* K, D^*\eta \rightarrow D^*\pi, \tag{25}
\]

\[ \cdots \]

also have contributions to the hadronic spectral densities. In the heavy meson chiral unitary approach, we can use the Bethe-Salpeter equation to perform the summation of the infinite series of the intermediate meson-loops (such as the \( D^*\pi, D_1^* K, D^*\eta \)) to obtain the full \( D^*\pi \rightarrow D^*\pi \) scattering amplitude, and generate the axialvector meson \( D_1(2430) \) dynamically \([15,35]\). If we saturate the full \( D^*\pi \rightarrow D^*\pi \) scattering amplitude with the exchanges of the intermediate axialvector meson \( D_1(2430) \), the \( D^*\pi \) rescattering effects lead to the renormalization

\[
\gamma(p^2 - M_{D_1}^2) \rightarrow \gamma(p^2 - M_{D_1}^2 - \Pi(p) + i\epsilon)
\]

in the hadronic representation of the correlation functions, where the \( \Pi(p) \) denotes the renormalized self-energy of the intermediate \( D^*\pi \) loops, and contributes a finite imaginary part to modify the dispersion relation. In fact, the contributions of the intermediate meson-loops are very large, we have to take the mass \( M_{D_1} \) as the bare mass \( M_{D_1} \) to absorb the real part of the un-renormalized self-energy to reproduce the physical mass, the net effects are embodied in the finite imaginary part. We can take into account those meson-loops effectively by taking the following replacement for the hadronic spectral density,

\[
\delta(s - M_{D_1}^2) \rightarrow \frac{1}{\pi} \frac{\sqrt{\Gamma_{D_1}^*}}{(s - M_{D_1}^2)^2 + s\Gamma_{D_1}^*}, \tag{26}
\]

here we neglect the complicated renormalization procedure for simplicity \([36]\). In Ref.[36], we observe that a width about (or less than) 400 MeV cannot change the prediction significantly, the \( \delta \) function approximation for the spectral densities still survives. In the present case, \( \Gamma_{D_1} = 384^{+107}_{-75} \pm 75 \) MeV from the Particle Data Group \([18]\), the contaminations from the intermediate state \( D^*\pi \) are expected to be small. Analogical discussions can be applied to the contaminations from the intermediate state \( B^*\pi \).

The negative scattering lengths \( a_{D^*} = -1.07 \text{ fm} \) and \( a_{B^*} = -7.17 \text{ fm} \) indicate that the \( D^* N \) and \( B^* N \) interactions are attractive, it is possible to form the \( D^* N \) and \( B^* N \) bound states; while
Table 1: The mass-shifts $\delta M$ and the scattering lengths $a$ versus the strong coupling constants $g^2$.

| $g^2$ | 0   | 10  | 20  | 30  | 40  | 50  |
|-------|-----|-----|-----|-----|-----|-----|
| $\delta M_{D^*}$ (MeV) | -75 | -72 | -70 | -67 | -65 | -62 |
| $\delta M_{B^*}$ (MeV) | -382 | -381 | -380 | -380 | -379 | -378 |
| $\delta M_{D_1}$ (MeV) | 70  | 71  | 73  | 74  | 76  | 78  |
| $\delta M_{B_1}$ (MeV) | 262 | 263 | 264 | 265 | 266 | 267 |
| $a_{D^*}$ (fm) | -1.13 | -1.09 | -1.05 | -1.02 | -0.98 | -0.94 |
| $a_{B^*}$ (fm) | -7.20 | -7.18 | -7.17 | -7.15 | -7.14 | -7.13 |
| $a_{D_1}$ (fm) | 1.11 | 1.14 | 1.16 | 1.19 | 1.21 | 1.24 |
| $a_{B_1}$ (fm) | 5.00 | 5.02 | 5.04 | 5.05 | 5.07 | 5.09 |

4 Conclusion

In this article, we calculate the mass-shifts of the vector and axialvector mesons $D^*$, $B^*$, $D_1$ and $B_1$ in the nuclear matter using the QCD sum rules. We take the linear approximation at the low density of the nuclear matter, and extract the mass-shifts and scattering lengths explicitly, $\delta M_{D^*} = -71$ MeV, $\delta M_{B^*} = -380$ MeV, $\delta M_{D_1} = 72$ MeV, $\delta M_{B_1} = 264$ MeV, $a_{D^*} = -1.07$ fm, $a_{B^*} = -7.17$ fm, $a_{D_1} = 1.15$ fm and $a_{B_1} = 5.03$ fm. Our numerical results indicate that the $D^*N$ and $B^*N$ interactions are attractive while the $D_1N$ and $B_1N$ interactions are repulsive; it is possible (difficult) to form the $D^*N$ and $B^*N$ ($D_1N$ and $B_1N$) bound states. The $J/\psi$ production can obtain additional suppressions due to mass modifications of the negative parity charmed mesons $D$ and $D^*$ in the nuclear matter.

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Figure 1: The mass-shifts $\delta M$ versus the Borel parameter $M^2$, the $A$, $B$, $C$ and $D$ denote the $D^*$, $D_1$, $B^*$ and $B_1$ mesons, respectively.
Figure 2: The contributions from different terms versus the Borel parameter $M^2$ in the operator product expansion, where the $A$, $B$, $C$ and $D$ denote the $D^*$, $D_1$, $B^*$ and $B_1$ mesons, respectively; the $\alpha$, $\beta$, $\lambda$ and $\tau$ denote the $\langle \bar{q}q \rangle_N$, $\langle \bar{q}iD_0q \rangle_N$, $\langle \bar{q}iD_0iD_0q \rangle_N + \langle \bar{q}g_\sigma Gq \rangle_N$, and $\langle \bar{q}qG \rangle_N$ terms, respectively.
Figure 3: The mass-shifts $\delta M$ versus the Borel parameter $M^2$, the A, B, C and D denote the $D^*, D_1, B^*$ and $B_1$ mesons, respectively; the $\alpha, \beta, \gamma, \lambda, \tau$ and $\rho$ correspond the strong coupling constants $g^2 = 0, 10, 20, 30, 40$ and 50, respectively.
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