IN-MEDIUM MASS MODIFICATIONS OF THE $D_0$ AND $B_0$ MESONS WITH THE QCD SUM RULES

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Abstract

In this article, we calculate the in-medium mass modifications of the scalar mesons $D_0$ and $B_0$ using the QCD sum rules. In calculations, we observe that the $D_0N$ and $B_0N$ scattering lengths are about 1.1 fm and 4.1 fm respectively, the mass-shifts $\delta M_{D_0} = 69$ MeV and $\delta M_{B_0} = 217$ MeV, and the $D_0N$ and $B_0N$ interactions are repulsive. The positive mass-shifts indicate that the decays of the higher charmonium states into the $D_0\bar{D}_0$ pair are suppressed.

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1 Introduction

The in-medium properties of hadrons play an important role in understanding the strong interactions, the heavy ion collisions and the nuclear astrophysics. There have been many experiments devoted to study the in-medium hadron properties. The upcoming FAIR (Facility for Antiproton and Ion Research) project at GSI provides the opportunity to extend the experimental studies into the charm sector. The CBM (compressed baryonic matter) collaboration intends to study the in-medium properties of the hadrons, including charmed mesons\textsuperscript{[1]}, while the PANDA collaboration will focus on the charm spectroscopy, and mass and width modifications of the charmed hadrons in the nuclear matter\textsuperscript{[2]}. Therefore, the in-medium properties of the charmed mesons become excellent subjects in recent years.

The in-medium mass modifications of the $D$, $\bar{D}$ mesons have been studied with the QCD sum rules\textsuperscript{[8, 9]}, the quark-meson coupling model\textsuperscript{[10]}, the coupled-channel approach\textsuperscript{[11]}, etc. The works on the scalar $D_0$, $\bar{D}_0$ mesons are few\textsuperscript{[12, 13]}, it is interesting to study their mass modifications in the nuclear matter to see whether or not the decaying of the higher excited charmonium states to the $D_0\bar{D}_0$ pair is facilitated. Just like the $D$ meson, the $D_0$ meson contains a charm-quark and a light quark. The existence of a light quark in the charmed mesons results in much difference between the in-medium mass modifications of the charmed-mesons and charmonium states. The former have large contributions from the light-quark condensates, while the latter are dominated by the gluon condensates\textsuperscript{[8, 9, 12, 14]}. In this article, we study the mass modifications of the...
$D_0, \bar{D}_0$ mesons in the nuclear matter using the QCD sum rules \[15\], furthermore, we study the corresponding mass modifications of the $B_0, \bar{B}_0$ mesons considering the heavy quark symmetry. In Ref.\[12\], Hilger and Kampfer study the in-medium mass modifications of the scalar $D_0, \bar{D}_0$ mesons and quantify the $D_0 - \bar{D}_0$ mass splitting using the QCD sum rules, the hadronic spectral density they choose differs from that of the present work. In Ref.\[13\], Tolos et al study the in-medium properties of the scalar mesons $D_0, D_{a0}(2317)$ and $X(3700)$ in the coupled-channel approach and reproduce them as dynamically-generated resonances.

The article is arranged as follows: we study the in-medium mass modifications of the scalar mesons $D_0$ and $\bar{B}_0$ 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 In-medium mass modifications of the $D_0$ and $\bar{B}_0$ with QCD sum rules

We study the mass modification of the $D_0$ meson in nuclear matter with the two-point correlation function $\Pi(q)$. In the Fermi gas approximation for the nuclear matter, the $\Pi(q)$ can be divided into the vacuum part $\Pi_0(q)$ and the static one-nucleon part $\Pi_N(q)$, which is expected to be valid at relatively low nuclear density, and the $\Pi_N(q)$ can be approximated in the linear density of the nuclear matter $\[8, 10\]$, 

$$
\Pi(q) = i \int d^4x \, e^{iqx} \langle T \{ J(x) J^\dagger(0) \} \rangle_{\rho_N} = \Pi_0(q) + \Pi_N(q) \simeq \Pi_0(q) + \frac{\rho_N}{2M_N} T_N(q),
$$

(1)

where the $\rho_N$ denotes the density of the nuclear matter, and the forward scattering amplitude $T_N(q)$ is defined as

$$
T_N(\omega, q) = i \int d^4x e^{iqx} \langle N(p)|T \{ J(x) J^\dagger(0) \} |N(p)\rangle,
$$

(2)

the $J(x)$ denotes the isospin averaged current,

$$
J(x) = J^\dagger(x) = \frac{\bar{\epsilon}(x)q(x) + \bar{q}(x)c(x)}{2},
$$

(3)

the $q$ denotes the $u$ or $d$ quark, the $q' = (\omega, q')$ is the four-momentum carried by the current $J(x)$, the $\langle N(p)\rangle$ 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)^32\rho_0\delta^3(p - p') \[8\]$. The terms proportional to $p^4_F, p^5_F, p^6_F, \cdots$ can be neglected at the normal nuclear matter with the saturation density $\rho_N = \rho_0 = \frac{2\rho_{sat}}{3}$, as the Fermi momentum $p_F = 0.27$ GeV is a small quantity $\[10\].$

In the limit $q \to 0$, the $T_N(\omega, q)$ can be related to the $D_0N$ scattering $T$-matrix, $T_{D_0N}(M_{D_0}, 0) = 8\pi(M_N + M_{D_0})a_{D_0}$, where the $a_{D_0}$ is the $D_0N$ scattering length. Near the pole position of the $D_0$-meson, the phenomenological spectral density $\rho(\omega, 0)$ can be parameterized with three unknown parameters $a, b$ and $c$ $\[8\],

$$
\rho(\omega, 0) = -\frac{f_{D_0}^2 M_{D_0}^2}{\pi} \text{Im} \left( \frac{T_{D_0N}(\omega, 0)}{(\omega^2 - M_{D_0}^2 + i\varepsilon)^2} \right) + \cdots,
$$

(4)

$$
= a \frac{d}{d\omega^2}\delta(\omega^2 - M_{D_0}^2) + b \delta(\omega^2 - M_{D_0}^2) + c \delta(\omega^2 - s_0),
$$

(5)

where the decay constant $f_{D_0}$ is defined by $\langle 0|J(0)|D_0(k) + \bar{D}_0(k)\rangle = f_{D_0} M_{D_0}$, the terms denoted by $\cdots$ represent the continuum contributions. The first term denotes the double-pole term, and corresponds to the on-shell effect of the $T$-matrix, $a = -8\pi(M_N + M_{D_0})a_{D_0}f_{D_0}^2 M_{D_0}^2$; the second term denotes the single-pole term, and corresponds to the off-shell effect of the $T$-matrix; and the
where the continuum term or the remaining effects, where the \( s_0 \) is the continuum threshold. Then the mass-shift of the \( D_0 \)-meson can be approximated as

\[
\delta M_{D_0} = 2\pi \frac{M_N + M_{D_0}}{M_N M_{D_0}} \rho_{aD_0} \, .
\]  

In the low energy limit \( \omega \to 0 \), the \( T_N(\omega, 0) \) is equivalent to the Born term \( T_{N}^\text{Born}(\omega, 0) \), i.e. \( T_N(0) = T_N^\text{Born}(0) \). We take into account the Born term at the phenomenological side,

\[
T_N(\omega^2) = T_N^\text{Born}(\omega^2) + \frac{a}{(M_{D_0}^2 - \omega^2)^2} + \frac{b}{M_{D_0}^2 - \omega^2} + \frac{c}{s_0 - \omega^2},
\]

with the constraint

\[
\frac{a}{M_{D_0}^2} + \frac{b}{M_{D_0}^2} + \frac{c}{s_0} = 0.
\]

The contributions from the intermediate spin-\( \frac{3}{2} \) charmed baryon states are zero in the soft-limit \( q_\mu \to 0 \), and we only take into account the intermediate spin-\( \frac{1}{2} \) charmed baryon states in calculating the Born term. The isospin states of the \( D_0 \)-mesons have the \( D_0^0N \) and \( \bar{D}_0^0N \) interactions,

\[
\begin{align*}
D_0^0(c\bar{u}) + p(uud) \text{ or } n(udd) & \rightarrow \Lambda_c^+ + \Sigma_c^+(cud) \text{ or } \Sigma_c^0(cud), \\
D_0^0(c\bar{d}) + p(uud) \text{ or } n(udd) & \rightarrow \Sigma_c^{++} + \Lambda_c^+ + \Sigma_c^0(cud),
\end{align*}
\]

where \( M_{\Lambda_c} = 2.286 \text{ GeV} \) and \( M_{\Sigma_c} = 2.454 \text{ GeV} \). We 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 \( \Sigma_c^0 \). It is straightforward to calculate the Born term by writing down the Feynman diagram, the result is

\[
T_N^\text{Born}(\omega, 0) = \frac{2f_{D_0} M_{D_0}^2 M_N(M_H - M_N)g_{D_0NH}}{[\omega^2 - (M_H - M_N)^2]^2},
\]

where the \( g_{D_0NH} \) denotes the strong coupling constants \( g_{D_0N\Lambda_c} \) and \( g_{D_0N\Sigma_c} \). On the other hand, there are no inelastic channels for the \( D_0^0N \) and \( \bar{D}_0^0N \) interactions, i.e. \( T_N^\text{Born}(0) = 0 \).

We carry out the operator product expansion to the condensates up to dimension-5 at the large space-like region in the nuclear matter. Once analytical results at the level of quark-gluon degree’s of freedom are obtained, then we set \( \omega^2 \approx q^2 \), and take the quark-hadron duality and perform the Borel transform with respect to the variable \( Q^2 = -\omega^2 \), finally obtain the following sum rule:

\[
\begin{align*}
a \left\{ \frac{1}{M^2} e^{-\frac{m_{D_0}^2}{M^2}} - \frac{s_0}{M_{D_0}^2} e^{-\frac{m_{D_0}^2}{M^2}} \right\} + b \left\{ \frac{e^{-\frac{m_{D_0}^2}{M^2}} - \frac{s_0}{M_{D_0}^2} e^{-\frac{m_{D_0}^2}{M^2}}} \right\} + \frac{2f_{D_0} M_{D_0}^2 M_N(M_H - M_N)g_{D_0NH}}{(M_H - M_N)^2 - M_{D_0}^2} \left\{ \frac{1}{(M_H - M_N)^2 - M_{D_0}^2} e^{-\frac{m_{D_0}^2}{M^2}} - \frac{1}{M_{D_0}^2} e^{-\frac{(M_H - M_N)^2}{M^2}} \right\} = \frac{m_c \langle \bar{q}q \rangle_N}{2} \\
-\langle q^4 iD_0 q \rangle_N + \frac{m_c^2(q^4 iD_0 q)N}{M^2} - \frac{2m_c \langle q^2 iD_0 q \rangle_N}{M^2} + \frac{m_c \langle q^2 iD_0 q \rangle_N}{M^4} \right\} e^{-\frac{m_{D_0}^2}{M^2}} + \frac{1}{16} \left\{ \frac{G_G}{\pi} \right\}_N \int_0^1 dx \left( 1 + \frac{\tilde{m}_G^2}{M^2} \right) e^{-\frac{\tilde{m}_G^2}{M^2}} - \frac{1}{48M^2} \left\{ \frac{G_G}{\pi} \right\}_N \int_0^1 \frac{x - x}{\tilde{m}_G^2} e^{-\frac{\tilde{m}_G^2}{M^2}},
\end{align*}
\]

where \( \tilde{m}_G^2 = \frac{m_G^2}{x} \).

Differentiate above equation with respect to \( \frac{1}{M^2} \), then eliminate the parameter \( b \), we can obtain the sum rule for the parameter \( a \). With the simple replacements \( m_c \to m_b \), \( D_0 \to B_0 \), \( \Lambda_c \to \Lambda_b \) and \( \Sigma_c \to \Sigma_b \), we can obtain the corresponding sum rule for the mass modification of the \( B_0 \) meson in the nuclear matter.
3 Numerical results and discussions

In calculations, we have assumed that the linear density approximation is valid at the low nuclear density, for a general condensate in the nuclear matter, \( \langle O \rangle_{\rho_N} = \langle O \rangle_0 + \frac{\rho_N}{2M_s} \langle N|O|N\rangle = \langle O \rangle_0 + \frac{\langle \bar{q}q \rangle_N}{2M_s} \). The input parameters are taken as \( \langle \bar{q}q \rangle_N = \frac{\sigma_N}{3}\langle N|O|N\rangle \), \( \langle \bar{q}q \rangle_N = -0.65\text{GeV}^2(2M_N) \), \( \langle \bar{q}q \rangle_N = 0.18\text{GeV}^2(2M_N) \), \( \langle \bar{q}q \rangle_N = 3.0\text{GeV}^2(2M_N) \), \( \langle \bar{q}q \rangle_N = 0.3\text{GeV}^2(2M_N) \), \( m_u + m_d = 12\text{MeV} \), \( \sigma_N = 45\text{MeV} \), \( M_N = 0.94\text{GeV} \), and \( \rho_N = (0.11\text{GeV})^3 \).

The hadronic parameters \( M_{D_0}, M_{B_0}, f_{D_0}, f_{B_0} \) are determined by the conventional two-point correlation functions \( \Pi_0(q) \) using the QCD sum rules, \( M_{D_0} = 2.355\text{GeV} \), \( M_{B_0} = 5.74\text{GeV} \), \( f_{D_0} = 0.334\text{GeV} \), \( f_{B_0} = 0.28\text{GeV} \) with the threshold parameters \( s_{D_0}^0 = 8.0\text{GeV}^2 \) and \( s_B^0 = 39.0\text{GeV}^2 \), respectively. For the observed scalar meson \( D_0(2400) \), the value \( M_{D_0} = 2.355\text{GeV} \) is consistent with the average of the experimental data \( M_{D_0} = 2318\text{MeV} \) and \( M_{D_0} = 2403\text{MeV} \). The uncertainties come from the hadronic parameters \( f_{D_0} \) and \( f_{B_0} \) can be approximated as \( \frac{2\delta f_{D_0}}{f_{D_0}} \) and \( \frac{2\delta f_{B_0}}{f_{B_0}} \), respectively.

In Fig.1, we plot the mass-shifts \( \delta M \) versus the Borel parameter \( M^2 \). 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 \( M^2 = (6.1 - 7.4)\text{GeV}^2 \) and \( (33 - 39)\text{GeV}^2 \) in the charmed and bottom channels, respectively, the uncertainties originate from the Borel parameter \( M^2 \) are less than 1%. Furthermore, the exponential factor \( e^{-\frac{2g^2}{2M^2}} < e^{-1} \) at those intervals, the continuum contributions are greatly suppressed. The main contributions come from the terms \( m_c\langle \bar{q}q \rangle_N \) and \( m_b\langle \bar{q}q \rangle_N \), the operator product expansion is certainly convergent.

In the Borel windows, the mass-shifts \( \delta M_{D_0} = 61\text{MeV}, 65\text{MeV}, 68\text{MeV}, 72\text{MeV}, 75\text{MeV}, 79\text{MeV} \) and \( \delta M_{B_0} = 213\text{MeV}, 215\text{MeV}, 217\text{MeV}, 219\text{MeV}, 221\text{MeV}, 223\text{MeV} \) respectively at the values \( g^2 = 0, 20, 40, 60, 80, 100, \) where the \( g \) denotes the strong coupling constants \( g_{D_0N a}, g_{D_0N a}, g_{B_0N a}, g_{B_0N a}, g_{B_0N a}, \) the mass-shifts increase monotonously with increase of the squared strong coupling constants \( g^2 \).

The calculations based on the QCD sum rules indicate that the values of the strong coupling constants \( g_{N N a}(960) = 12 \pm 2 \) and \( g_{N N a}(980) = 11 \pm 2 \), which are of the same magnitude of the phenomenological value of the strong coupling constant \( g_{N N a} = 13.5 \). On the other hand, the value of the strong coupling constant \( g_{D_0N a} = 6.74 \) from the QCD sum rules is much smaller.

In this article, we take the approximation \( g_{D_0N a} \approx g_{D_0N a} \approx g_{B_0N a} \approx g_{B_0N a} \approx 6.74 \), and obtain the values \( \delta M_{D_0} = 69\text{MeV}, \delta M_{B_0} = 217\text{MeV}, \delta a_{D_0} = 1.1\text{fm} \) and \( a_{B_0} = 4.1\text{fm} \).

The positive scattering lengths indicate that the \( D_0N \) and \( B_0N \) interactions are repulsive, it is difficult to form the \( D_0N \) and \( B_0N \) bound states. Which are in contrast to the \( DN \) and \( BN \) interactions, where the negative scattering lengths indicate that the \( DN \) and \( BN \) interactions are attractive, it is possible to form the \( DN \) and \( BN \) bound states.

Due to positive mass-shift \( \delta M_{D_0} \), the decays of the high charmonium states to the \( D_0 \bar{D}_0 \) pair obtain additional suppression in the phase space, and the decay into the lowest state \( J/\psi \) is preferred. While the negative mass-shift \( \delta M_{D} \) indicate that the high charmonium states decays to the \( D \bar{D} \) pair are facilitated, and the production of the \( J/\psi \) is suppressed. The \( J/\psi \) production does not obtain additional suppression due to mass modification of the scalar meson \( D_0(2400) \) in the nuclear matter.

In the present work and Ref.[8], the correlation functions \( \Pi(q) \) are divided into the vacuum part and the static one-nucleon part, and the nuclear matter induced effects are extracted explicitly, while in Refs.[9][12], the pole terms of the hadronic spectral densities are parameterized as \( \text{Im} \Pi(q,\omega = 0) = F_+\delta(\omega - M_+) - F_-\delta(\omega - M_-) \), where \( M_\pm = M \pm \Delta M \) and \( F_\pm = F \pm \Delta F \), and the sum rules for the mass center \( M \) and the mass splitting \( \Delta M \) are obtained. For the pseudoscalar \( D, \bar{D} \) mesons, Hayashigaki obtains the mass-shift \( \delta M_D = -50\text{MeV} \), while Hilger, Thomas and Kampfer obtain the mass-shift \( \delta M_D = +45\text{MeV} \). For scalar \( D_0, \bar{D}_0 \) mesons, the mass-shift
\[ \delta M_{D_0} = M - M_{D_0} < 0 \]

\[ \delta M_{D_0} = 69 \text{ MeV} \]

\[ \delta M_{D_0} = -50 \text{ MeV} \]

In those studies, irrespective of the parameterizations of the hadronic spectral densities, derivatives with respect to \( 1/M^2 \) are used to obtain additional QCD sum rules so as to take special superpositions to delete the unknown parameters and extract the explicit expressions of the mass-shifts. The QCD sum rules from the derivatives are not necessarily work well, as the QCD sum rules are just a QCD model, the predictions can vary with the special assumptions. Furthermore, the isospin currents and isospin-averaged currents can also lead to differences in the hadronic spectral densities. All those predictions can be confronted with the experimental data in the future.

In Ref. [13], Tolos et al study the in-medium properties of the scalar mesons \( D_0(2400), D_{s0}(2317) \) and \( X(3700) \) in the coupled-channel approach and reproduce them as dynamically-generated resonances, and observe that the \( D_{s0}(2317) \) and \( X(3700) \) enlarge their widths to the order of 100 and 200 MeV respectively at the normal nuclear matter, and the \( D_0(2400) \) meson obtains an extra widening from the already large width, due to the \( D \) meson absorptions in the nuclear matter via the \( DN \) and \( DNN \) inelastic reactions. In the vacuum, the mass and width of the \( D_0(2400) \) are \( M_{D_0} = (2318 \pm 29) \text{ MeV}, M_{D_0^\pm} = (2403 \pm 14 \pm 35) \text{ MeV}, \Gamma_{D_0} = (267 \pm 40) \text{ MeV}, \Gamma_{D_0^\pm} = (283 \pm 24 \pm 34) \text{ MeV}, \) respectively [18]. For \( D_0(2400), \) the spin-parity \( J^P = 0^+ \) assignment is favored, while the spin-parity \( J^P \) of the \( D_0^\pm(2400) \) still needs confirmation [18]. The mass-shifts \( \delta M_{D_0} = 69 \text{ MeV} \) and \( \delta M_{D} = -50 \text{ MeV} \) obtained in the present work and Ref. [8] respectively favor the decays \( D_0 \rightarrow D\pi, \) we can expect that the in-medium width of the \( D_0(2400) \) is larger, and our prediction is compatible with the observation of Ref. [13]. However, the large width disfavors the experimental observation of the relative small mass-shift.

### 4 Conclusion

In this article, we calculate the in-medium mass-shifts of the scalar mesons \( D_0 \) and \( B_0 \) using the QCD sum rules. At the low density of the nuclear matter, we can take the linear approximation, and extract the mass-shifts explicitly. Our numerical results indicate that the \( D_0N \) and \( B_0N \) scattering lengths are \( a_{D_0} = 1.1 \text{ fm} \) and \( a_{B_0} = 4.1 \text{ fm} \), respectively; the \( D_0N \) and \( B_0N \) interactions are repulsive, the mass-shifts are \( \delta M_{D_0} = 69 \text{ MeV} \) and \( \delta M_{B_0} = 217 \text{ MeV} \). The positive mass-shifts indicate that the \( J/\psi \) production does not obtain additional suppression due to the mass
modification of the scalar meson $D_0(2400)$ in the nuclear matter.

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