Analysis of the $\Lambda_Q$ baryons in the nuclear matter with the QCD sum rules

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Abstract

In this article, we study the $\Lambda_c$ and $\Lambda_b$ baryons in the nuclear matter using the QCD sum rules, and obtain the in-medium masses $M^*_\Lambda_c = 2.335$ GeV, $M^*_\Lambda_b = 5.678$ GeV, the in-medium vector self-energies $\Sigma^*_\Lambda_c = 34$ MeV, $\Sigma^*_\Lambda_b = 32$ MeV, and the in-medium pole residues $\lambda^*_\Lambda_c = 0.021 \text{GeV}^3$, $\lambda^*_\Lambda_b = 0.026 \text{GeV}^3$. The mass-shifts are $M^*_\Lambda_c - M_{\Lambda_c} = 51$ MeV and $M^*_\Lambda_b - M_{\Lambda_b} = 60$ MeV, respectively.

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1 Introduction

In the past years, there have been several progresses on the spectroscopy of the heavy baryon states. Several new charmed baryons, such as the $\Omega^*_c(2770)$, $\Lambda^+_c(2765)$, $\Lambda^+_c(2880)$, $\Sigma^+_c(2940)$, $\Sigma^+_c(2800)$, $\Xi^+_c(2980)$, $\Xi^+_c(3080)$, $\Xi^0_c(2980)$, $\Xi^0_c(3080)$, have been observed in recent years [1], and re-vivified the interest in the spectroscopy of the charmed baryons. On the other hand, the in-medium properties of hadrons play an important role in understanding the strong interactions, the relativistic heavy ion collisions and the nuclear astrophysics. The upcoming FAIR project at GSI provides the opportunity to extend the experimental studies of the in-medium hadron properties into the charm sector. The CBM collaboration intends to study the in-medium properties of the hadrons [2], while the PANDA collaboration will focus on the charm spectroscopy, and mass and width modifications of the charmed hadrons in the nuclear matter [3]. So it is interesting to study the in-medium properties of the charmed baryons.

The QCD sum rules is a powerful theoretical tool in studying the ground state hadrons [4]. In the QCD sum rules, the operator product expansion is used to expand the time-ordered currents into a series of quark and gluon condensates which parameterize the long distance properties. Based on the quark-hadron duality, we can obtain copious information about the hadronic parameters at the phenomenological side [4]. The in-medium properties of the light flavor hadrons have been studied extensively with the QCD sum rules [5,6], while the works on the in-medium properties of the heavy hadrons focus on the $J/\psi$, $\eta_c$, $D$, $B$, $D_0$ and $B_0$ [7,8]. In this article, we study the in-medium properties of the $\Lambda_c$ baryon in the nuclear matter using the QCD sum rules, furthermore, we study the corresponding properties of the $\Lambda_b$ baryon considering the heavy quark symmetry.

The article is arranged as follows: we study the in-medium properties of the heavy baryons $\Lambda_c$ and $\Lambda_b$ 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 properties of the $\Lambda_Q$ baryons with QCD sum rules

We study the $\Lambda_Q$ baryons in the nuclear matter with the two-point correlation functions $\Pi(q)$,

$$\Pi(q) = i \int d^4x e^{iq \cdot x} \langle \Psi_0 | T \left[ J(x) \bar{J}(0) \right] | \Psi_0 \rangle,$$

$$J(x) = \epsilon^{ijk} u_i(x) C \gamma_5 d_j(x) Q_k(x),$$  \hspace{1cm} (1)
where the \( i, j, k \) are color indexes, \( Q = c, b \), and the \( |\Psi_0\rangle \) is the nuclear matter ground state.

The correlation functions \( \Pi(q) \) can be decomposed as
\[
\Pi(q) = \Pi_s(q^2, q \cdot u) + \Pi_q(q^2, q \cdot u) \cdot q + \Pi_u(q^2, q \cdot u) \cdot \not{u},
\]
according to Lorentz covariance and parity, time reversal invariance. In the rest frame of the
\( J \)
the current operators
\[
\text{multiply both sides with the weight function (} \omega, \vec{q} \text{) into the correlation functions } \Pi(p) \text{ to obtain the hadronic representation}.
\]

We can insert a complete set of intermediate baryon states with the same quantum numbers as
the current operators \( J(x) \) into the correlation functions \( \Pi(p) \) to obtain the hadronic representation.

At finite nuclear density, the invariant functions \( \Pi \) can be decomposed as
\[
\Pi(q^2, \vec{q}) = \sum_n C_n(q^2, \vec{q}) \langle O_n \rangle_{\rho_N},
\]
where the \( C_n(q^2, \vec{q}) \) are the Wilson coefficients, the in-medium condensates \( \langle O_n \rangle_{\rho_N} = \langle |\Psi_0\rangle O_n |\Psi_0\rangle_{\rho_N} \) at the low nuclear density, the \( \langle O \rangle \) and \( \langle O \rangle_N \) denote the vacuum condensates and nuclear matter induced condensates, respectively. One can consult Refs.\[5, 6\] for the technical
details in the operator product expansion. We carry out the operator product expansion in the
nuclear matter at the large space-like region \( q^2 \ll 0 \) and obtain the QCD spectral densities, then
take the limit \( u_\mu = (1, 0) \) and obtain the imaginary parts,
\[
\Delta \Pi_i(\omega, \vec{q}) = \text{limit}_{\epsilon \to 0} \left[ \Pi_i(\omega + i\epsilon, \vec{q}) - \Pi_i(\omega - i\epsilon, \vec{q}) \right].
\]

We can match the phenomenological side with the QCD side of the spectral densities, and
multiply both sides with the weight function (\( \omega - \bar{E}_q \)) \( e^{-\frac{\omega}{\bar{E}_q}} \), which excludes the negative-energy pole contribution, then perform the integral \( \int_{-\omega_0}^{\omega_0} d\omega \),
\[
\int_{-\omega_0}^{\omega_0} d\omega \Delta \Pi_i(\omega, \vec{q}) (\omega - \bar{E}_q) e^{-\frac{\omega}{\bar{E}_q}},
\]
where the \(\omega_q\) is the threshold parameter, finally obtain the following QCD sum rules:

\[
\lambda_{\Lambda Q}^2 e^{-\frac{E_q^2}{2M^2}} = \int_{m_Q^2}^{\mu^2} ds \int_x^1 dx \left\{ \frac{3x(1-x)^2(s-\tilde{E}_Q^2)^2}{128\pi^4} - \frac{(1-x)^2m_Q^2}{384\pi^2x^2} \left( \frac{\alpha_s G}{\pi} \right)_{\rho N} \delta(s-\tilde{E}_Q^2) + \frac{x}{128\pi^2} \left( \frac{\alpha_s G}{\pi} \right)_{\rho N} - \frac{x(1-x)(q^1 i D_0 q)^{\rho N}}{3\pi^2} \left[ 1 + 2s\delta(s-\tilde{E}_Q^2) \right] \right. \\
\left. + \hat{E}_q \left[ \frac{x(1-x)(q^1 q)^{\rho N}}{4\pi^2} - \frac{x(1-x)(q^1 i D_0 i D_0 q)^{\rho N}}{2\pi^2} \right] \delta(s-\tilde{E}_Q^2) + \frac{x(1-x)(q^1 i D_0 i D_0 q)^{\rho N}}{\pi^2} \left( 1 + \frac{2s}{M^2} \right) \delta(s-\tilde{E}_Q^2) + \frac{x(1-x)(q^1 q)^{\rho N}}{6\pi^2} \right\} e^{-\frac{\hat{E}_q^2}{M^2}} + \frac{(\bar{q}q)^2_{\rho N} + (q^1 q)^2_{\rho N}}{6} e^{-\frac{E_q^2}{M^2}}, (8)
\]

\[
\lambda_{\Lambda Q}^2 \Sigma e^{-\frac{E_q^2}{2M^2}} = \int_{m_Q^2}^{\mu^2} ds \int_x^1 dx \left\{ \frac{x(1-x)^2(s-\tilde{E}_Q^2)^2}{128\pi^4} - \frac{(1-x)^2s}{384\pi^2x^2} \left( \frac{\alpha_s G}{\pi} \right)_{\rho N} \delta(s-\tilde{E}_Q^2) + \frac{1}{192\pi^2x^2} \left( \frac{\alpha_s G}{\pi} \right)_{\rho N} - \frac{1}{128\pi^2} \left( \frac{\alpha_s G}{\pi} \right)_{\rho N} - \frac{(1-x)(q^1 i D_0 q)^{\rho N}}{6\pi^2} \left[ 1 + 2s\delta(s-\tilde{E}_Q^2) \right] + \hat{E}_q \left[ \frac{(1-x)(q^1 q)^{\rho N}}{4\pi^2} - \frac{(q^1 i D_0 i D_0 q)^{\rho N}}{2\pi^2} \right] \delta(s-\tilde{E}_Q^2) + \frac{(1-x)(q^1 i D_0 i D_0 q)^{\rho N}}{2\pi^2} \left( 1 + \frac{2s}{M^2} \right) \delta(s-\tilde{E}_Q^2) - \frac{(q^1 q)^{\rho N}}{16\pi^2} \delta(s-\tilde{E}_Q^2) + \frac{(q^1 i D_0 i D_0 q)^{\rho N}}{12\pi^2} \delta(s-\tilde{E}_Q^2) + \frac{(q^1 q)^{\rho N}}{6\pi^2} \left( 1 + \frac{2s}{M^2} \right) \delta(s-\tilde{E}_Q^2) + \frac{(q^1 q)^{\rho N}}{16\pi^2} \delta(s-\tilde{E}_Q^2) \right\} e^{-\frac{\hat{E}_q^2}{M^2}} + \frac{(\bar{q}q)^2_{\rho N} + (q^1 q)^2_{\rho N}}{6} e^{-\frac{E_q^2}{M^2}}, (9)
\]

where \(\tilde{E}_Q^2 = \frac{m_Q^4}{2} + \hat{q}^2\), \(E_q^2 = m_Q^2 + q^2\), \(x_i = \frac{m_i^2}{\tilde{E}_Q^2}\), \(s_0^2 = \omega_0^2 = s_0^2 - q^2\), the \(x_i\) in the spectral densities where the function \(\delta(s-\tilde{E}_Q^2)\) appears should be 0. We can obtain the in-medium masses \(M_{\Lambda Q}\), vector self-energies \(\Sigma\) and pole residues \(\lambda_{\Lambda Q}^2\) by solving above equations with simultaneous iterations.

### 3 Numerical results and discussions

In calculations, we have assumed that the linear density approximation is valid at the low nuclear density. The input parameters are taken as \(\langle q^1 q \rangle_{\rho N} = \frac{3}{2} \rho_N\), \(\langle \bar{q}q \rangle_{\rho N} = \langle \bar{q}q \rangle + \frac{\alpha_s G}{\pi} \rho_N\), \(\langle q^1 i D_0 q \rangle_{\rho N} = \langle q^1 i D_0 q \rangle - (0.65 \pm 0.15) \text{GeV} \rho_N\), \(\langle q^1 i D_0 q \rangle_{\rho N} = (0.18 \pm 0.01) \text{GeV} \rho_N\), \(\langle q^1 i D_0 q \rangle_{\rho N} + \frac{1}{12} \langle q^1 q \rangle_{\rho N}\).
Figure 1: The in-medium and vacuum masses versus the Borel parameter $M^2$, the A and B denote the $\Lambda_c$ and $\Lambda_b$ baryons, respectively.

Figure 2: The in-medium and vacuum pole residues versus the Borel parameter $M^2$, the A and B denote the $\Lambda_c$ and $\Lambda_b$ baryons, respectively.

\[
\langle \bar{q} q, \sigma G q \rangle + 3.0 \text{ GeV}^2 \rho_N, \quad \langle \bar{q} q' g_s G q \rangle_{\rho_N} = -0.33 \text{ GeV}^2 \rho_N, \quad \langle \bar{q} g_s G q \rangle = m_0^2 \langle \bar{q} q \rangle, \quad \langle \bar{q} q \rangle = -(0.23 \pm 0.01 \text{ GeV})^3, \quad m_0^2 = (0.8 \pm 0.2) \text{ GeV}^2, \quad \langle \pi^+ G \pi^- \rangle = (0.33 \text{ GeV})^4, \quad m_u + m_d = 12 \text{ MeV}, \quad \sigma_N = (45 \pm 10) \text{ MeV}, \quad q^2 = (0.27 \text{ GeV})^2, \quad \sigma_N = (0.11 \text{ GeV})^3, \quad \langle \bar{q} q \rangle_{\rho_N} = f \langle \bar{q} q \rangle_{\rho_N} + (1 - f) \langle \bar{q} q \rangle \times \langle \bar{q} q \rangle, \quad \rho_N = 0.77 \pm 0.8 \text{ GeV}, \quad \rho_N = 0.11 \text{ GeV}^3, \quad f = 0.5 \pm 0.5 \text{ GeV}.
\]

In the limit $\rho_N = 0$, we can recover the QCD sum rules in the vacuum, see Eqs.(8-9). We can take the threshold parameters $s_{\Lambda_c}^0 = 9.6 \text{ GeV}^2$ and $s_{\Lambda_b}^0 = 42.0 \text{ GeV}^2$ determined in Ref. to reproduce the experimental data. Taking the Borel parameters as $M^2 = (2.5 - 3.5) \text{ GeV}^2$ and $(4.3 - 5.3) \text{ GeV}^2$ for the $\Lambda_c$ and $\Lambda_b$ baryons, respectively, we obtain the hadronic parameters $M_{\Lambda_c} = 2.284_{-0.078}^{+0.049} \text{ GeV}$, $M_{\Lambda_b} = 5.618_{-0.104}^{+0.078} \text{ GeV}$, $\lambda_{\Lambda_c} = (0.022 \pm 0.002) \text{ GeV}^3$, $\lambda_{\Lambda_b} = (0.027 \pm 0.003) \text{ GeV}^3$, the uncertainties originate from the Borel parameters $M^2$. The central values of the masses are consistent with the experimental data $M_{\Lambda_c} = 2.2864 \text{ GeV}$ and $M_{\Lambda_b} = 5.6202 \text{ GeV}$.

Taking the same Borel parameters and threshold parameters as the QCD sum rules in the vacuum, we can obtain the in-medium hadronic parameters of the $\Lambda_Q$ baryons, $M_{\Lambda_c}^* = 2.335_{-0.072}^{+0.045} \text{ GeV}$, $M_{\Lambda_b}^* = 5.676_{-0.103}^{+0.077} \text{ GeV}$, $\Sigma_{\Lambda_c}^* = (34 \pm 1) \text{ MeV}$, $\Sigma_{\Lambda_b}^* = (32 \pm 1) \text{ MeV}$, $\lambda_{\Lambda_c}^* = (0.021 \pm 0.001) \text{ GeV}^3$, $\lambda_{\Lambda_b}^* = (0.026 \pm 0.003) \text{ GeV}^3$, again the uncertainties originate from the Borel parameters. In Figs.1-2, we present the numerical values of the in-medium masses and pole residues versus the Borel parameter $M^2$ compared with the QCD sum rules in the vacuum. We can eliminate some uncertainties by introducing the mass differences $\Delta M_{\Lambda_Q} = M_{\Lambda_Q}^* - M_{\Lambda_Q}$, and obtain the values $\Delta M_{\Lambda_c} = 51 \text{ MeV}$ and $\Delta M_{\Lambda_b} = 60 \text{ MeV}$. The ratios are $\frac{\Delta M_{\Lambda_c}}{M_{\Lambda_c}} = 2.2\%$ and $\frac{\Delta M_{\Lambda_b}}{M_{\Lambda_b}} = 1.1\%$, re-
spectively, the mass modifications are slight, on the other hand, \( \frac{\lambda_{\Lambda Q} - \lambda_{\Lambda_b Q}}{\lambda_{\Lambda Q}} \approx 4\% \), the pole residue modifications are also slight, we expect that the threshold parameters survive in the nuclear matter, and the in-medium effects cannot modify the pole contributions significantly. In calculations, we observe that the exponential factor \( e^{-\frac{M^2}{M^2}} \ll e^{-1} \) in the Borel windows, the contributions from the higher resonance and continuum states are greatly suppressed. There are uncertainties for the threshold parameters, in general, we expect that \( (M_{gr} + \frac{\Gamma_{gr}}{2})^2 \leq s_0 \leq (M_{re} - \frac{\Gamma_{re}}{2})^2 \), where the \( gr \) and \( re \) stand for the ground state and first radial excited state (or resonance), respectively. At this interval, there are some values which can satisfy the two criterions (pole dominance and convergence of the operator product expansion) of the QCD sum rules. In this article, we take the values \( \delta s_0 \) determined in our previous work \[9\]. In Fig.3, we plot the contributions from the perturbative term and \( \langle \bar{q}q \rangle^2_{\rho_N} + \langle q^1q \rangle^2_{\rho_N} \) term versus the Borel parameter \( M^2 \) in the operator product expansion. From the figure, we can see that the dominant (or main) contributions come from the perturbative term, while the contributions of the dimension-6 term \( \langle \bar{q}q \rangle^2_{\rho_N} + \langle q^1q \rangle^2_{\rho_N} \) are about \((6 - 19)\% \) (or \((14 - 29)\% \)) for the \( \Lambda_c \) (or \( \Lambda_b \)) baryon in the Borel window, the operator product expansion is well convergent. Compared with the corresponding QCD sum rules in the vacuum, the QCD sum rules in the nuclear matter have better convergent behavior in the operator product expansion, see Fig.3.

In this article, we have neglected the contributions of the perturbative \( \mathcal{O}(\alpha_s) \) corrections, which can be taken into account by introducing the formal coefficient \( 1 + \frac{\alpha_q}{\pi} f(m_Q, s) \) through the unknown function \( f(m_Q, s) \). As the dominant (or main) contributions come from the perturbative term, we
expect that the $O(\alpha_s)$ corrections to the perturbative term cannot change the predictions greatly. If the perturbative $O(\alpha_s)$ corrections have the typical value 30\%, i.e. \(1 + \frac{2\pi}{f(m_Q, s)} = 1.3\), where we have neglected the mass \(m_Q\) and energy \(s\) dependence to make a rough estimation, the resulting mass-shifts are \(\Delta M_{\Lambda_c} = 43\) MeV and \(\Delta M_{\Lambda_b} = 59\) MeV. In calculations, we observe that the mass-shifts \(\delta M\) decrease with the increase of the perturbative contributions.

The uncertainties originate from the \(\delta(\frac{\alpha_s}{\pi})_{\rho\rho_N}, \delta(q^\dagger iD_0 q)_{\rho\rho_N}, \delta m^2, \delta \sigma_N\) are tiny and can be neglected safely. The uncertainty of the quark condensate \(\delta\langle \bar{q}q\rangle\) leads to somewhat larger uncertainties for the masses \(M^*_{\Lambda_Q}\) and \(M_{\Lambda_Q}\), but the mass-shifts \(\Delta M\) change slightly and the uncertainties can be neglected. We can take into account the uncertainties of the heavy quark masses and threshold parameters, \(\delta m_c = 0.1\) GeV, \(\delta m_b = 0.1\) GeV, \(\delta s^0_{\Lambda_c} = 0.6\) GeV\(^2\), \(\delta s^0_{\Lambda_b} = 1.2\) GeV\(^2\) [9], as the dominant (or main) contributions come from the perturbative term, those uncertainties may result in considerable uncertainties for the mass differences \(\Delta M\). In calculations, we observe that the values of the mass differences \(\Delta M_{\Lambda_c} = 51\) MeV and \(\Delta M_{\Lambda_b} = 60\) MeV survive approximately, while other parameters obtain the uncertainties, \(\delta \lambda^*_{\Lambda_c} = (\pm 0.002 \pm 0.003)\) GeV\(^3\), \(\delta \lambda^*_{\Lambda_b} = (\pm 0.004 \pm 0.004)\) GeV\(^3\), \(\delta \Sigma_{\Lambda_c}^L = (\pm 1 \pm 1)\) MeV, and \(\delta \Sigma_{\Lambda_b}^L = (\pm 1 \pm 1)\) MeV.

At the interval \( f = 0 \sim 1\), the masses \(M^*_{\Lambda_Q}\) increase monotonously with the increase of the parameter \(f\), the uncertainty \(\delta f = \pm 0.5\) result in the uncertainties \(\delta M^*_{\Lambda_c} = +68\) MeV and \(\delta M^*_{\Lambda_b} = +113\) MeV. The corresponding uncertainties of other hadronic parameters are rather small, \(|\delta \lambda^*_{\Lambda_c}| < 0.001\) GeV\(^3\), \(|\delta \lambda^*_{\Lambda_b}| < 0.002\) GeV\(^3\), \(|\delta \Sigma_{\Lambda_c}^L| \leq 1\) MeV, \(|\delta \Sigma_{\Lambda_b}^L| \leq 2\) MeV, and can be neglected.

For the nucleons, \(m_u \approx m_d \approx 0\), the in-medium hadronic parameters \(\Sigma^L_{\Lambda_Q}\) and \(\Sigma^L_{\Lambda_b}\) can be approximated as \(M^*_{\Lambda_Q} = -\frac{s^0_{\Lambda_Q}}{3\pi^2} \langle \bar{q}q\rangle_{\rho\rho_N}\) and \(\Sigma^L_{\Lambda_b} = \frac{s^0_{\Lambda_b}}{3\pi^2} \langle q^\dagger q\rangle_{\rho\rho_N}\), respectively, and the in-medium mass modification is large, about 40\%. In the present case, the quark masses \(m_c\) and \(m_b\) are large, no such simple relations can be obtained, and the mass modifications are rather mild. In the QCD sum rules, the modifications of the phenomenological spectral densities are related to the in-medium quark and gluon condensates. In the linear density approximation, the in-medium modification of the quark condensate is large while the modification of the gluon condensate is mild, \(\langle \bar{q}q\rangle_{\rho\rho_N} \approx 0.60 \langle \bar{q}q\rangle\) and \(\langle \frac{\alpha_s}{\pi} GG\rangle_{\rho\rho_N} \approx 0.93 \langle \frac{\alpha_s}{\pi} GG\rangle\) [9]. Furthermore, there appear additional quark condensates associated with the light flavors, such as \(\langle q^\dagger q\rangle_{\rho\rho_N}, \langle q^\dagger D_0 q\rangle_{\rho\rho_N}, \langle q^\dagger D_0 D_0 q\rangle_{\rho\rho_N}, \langle \bar{q}iD_0D_0q\rangle_{\rho\rho_N}\), etc, which also play an important role. The \(\Lambda_Q\) baryons have a heavy quark besides two light quarks, the heavy quark interacts with the nuclear matter through the exchange of the intermediate gluons, we expect the mass modifications are smaller than that of the nucleons, which have three light quarks, \(uud\) or \(uud\).

### 4 Conclusion

In this article, we study the in-medium properties of the heavy baryons \(\Lambda_c\) and \(\Lambda_b\) using the QCD sum rules, and obtain the analytical expressions for the in-medium masses \(M^*_{\Lambda_Q}\), vector self-energies \(\Sigma^e\), and pole residues \(\lambda^*_{\Lambda_Q}\), then get the values \(M^*_{\Lambda_c} = 2.335\) GeV, \(M^*_{\Lambda_b} = 5.678\) GeV, \(\Sigma^L_{\Lambda_c} = 34\) MeV, \(\Sigma^L_{\Lambda_b} = 32\) MeV, \(\lambda^*_{\Lambda_c} = 0.021\) GeV\(^3\), \(\lambda^*_{\Lambda_b} = 0.026\) GeV\(^3\), \(M^*_{\Lambda_c} - M_{\Lambda_c} = 51\) MeV, \(M^*_{\Lambda_b} - M_{\Lambda_b} = 60\) MeV, and discuss the relevant uncertainties. The present predictions can be confronted with the experimental data in the future.

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