A QCD Sum Rule Study of Θ⁺ in Nuclear Matter

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We consider a [ud]\textsuperscript{s} current, in the finite-density QCD sum rule approach, to investigate the scalar and vector self-energies of the recently observed pentaquark state Θ\textsuperscript{+}(1540), propagating in nuclear matter. We find that, opposite to what was obtained for the nucleon, the vector self-energy is negative, and the scalar self-energy is positive. There is a substantial cancellation between them resulting in an attractive net self-energy of the same order as in the nucleon case.

The possible existence of a narrow exotic baryon Θ\textsuperscript{+} with strangeness +1 is, nowadays, possibly one of the most exciting topics in nuclear physics. There is a large number of experiments with evidence for the existence of this state \textsuperscript{1}, and a similar number of high-energy experiments that see no evidence \textsuperscript{2}. Even lattice QCD simulations of pentaquarks by several groups have not converged yet \textsuperscript{3, 4}. The apparent contradiction between evidences for and against the existence of Θ\textsuperscript{+} might be resolved if there is a production mechanism which is present in some experiments and absent in others. One such production mechanism was proposed in refs. \textsuperscript{3, 5} and is related with the formation of the mysterious \textsuperscript{N}(2400) resonance, which would decay into Θ\textsuperscript{+} Ki. The \textsuperscript{N}(2400) could be formed in the reactions γ + p → π\textsuperscript{+} + \textsuperscript{N}(2400) and π\textsuperscript{−} + p → \textsuperscript{N}(2400) \textsuperscript{6}, which are not present in e\textsuperscript{+}e\textsuperscript{−} collisions and could be difficult to happen at high-energy collisions. Other possible production mechanism is meson exchange between the nucleon and another hadron. Since at high energies all exchanges with the non-vacuum quantum numbers die out and only the gluonic pomerons survive, it may be difficult to excite the pentaquark by soft gluons on nucleon \textsuperscript{5, 7}.

If Θ\textsuperscript{+} is indeed produced through the decay of the \textsuperscript{N}(2400) resonance or through meson exchange, then one should expect that it could be seen in heavy ion collisions at RHIC, since these two processes could happen in the collision of nucleons and many “comoving” hadrons produced in the collisions at RHIC. However, if this is the case, the Θ\textsuperscript{+} would be formed in a nuclear medium which could change its mass and decay width. Therefore, to be able to identify positively the Θ\textsuperscript{+} signal in heavy ion collision at RHIC, it would be very important to know how the nuclear medium affects the Θ\textsuperscript{+} characteristics. In this work we will use the QCD sum rule (QCDSR) approach in nuclear matter \textsuperscript{8, 9, 10} to study the pentaquark Θ\textsuperscript{+} in a finite density medium.

The finite density QCDSR approach focuses on a correlation function evaluated in the ground state of nuclear matter, instead of the QCD vacuum (as in the usual sum rules). For spin-\frac{1}{2} baryons, this function can be decomposed into three invariant functions of two kinematic invariants. The appearance of an additional invariant function, compared with the vacuum case, is due to an additional four-vector: the four-velocity of the nuclear medium, u_\mu, which, together with the the nuclear density, ρ_\text{N}, characterizes the nuclear matter ground state. The quasibaryon excitations are characterized by scalar and vector self-energies. By introducing a simple ansatz for the spectral densities, one obtains a phenomenological representation of the correlation function.

The correlation function can be also evaluated at large spacelike momenta using an operator product expansion (OPE). This expansion is written in terms of the matrix elements of composite quark and gluon operators, evaluated in the nuclear matter ground state: the in-medium condensates. By equating the OPE and phenomenological representations of the correlation function, one obtains QCDSR sum rules that relate the baryon self-energies to the in-medium condensates.

Several zero-density QCDSR investigations of Θ\textsuperscript{+} already exist where the authors have used different interpolating fields \textsuperscript{11, 12, 13, 14, 15}. Here we use the interpolating field suggested in refs. \textsuperscript{1, 4, 16}:

$$\eta(x) = e_{abc} e_{def} e_{cfg} [u_a^T(x)C\delta_b(x)][u_c^T(x)C\gamma_5 d_e(x)]C\gamma_5 g_f(x),$$

where a, b, c, ... are color index and C = −CT is the charge conjugation operator. In Eq. \textsuperscript{11} each diquark pair has spin and isospin zero and is in the 3 representation of color SU(3). The total current has isospin zero, positive parity and spin 1/2. In ref. \textsuperscript{13} it was shown that the ground state described by this current has negative parity and a mass compatible with the experimental Θ\textsuperscript{+} mass.

The QCDSR for Θ\textsuperscript{+} at finite density is based on the correlation function defined by

$$\Pi(q) \equiv i \int d^4x e^{iqx} \langle \Psi_0 | T \eta(x) \overline{\eta}(0) | \Psi_0 \rangle,$$

where |Ψ_0\rangle represents the nuclear matter ground state.

The correlator in Eq. \textsuperscript{2} can be decomposed in three distinct structures \textsuperscript{8}

$$\Pi(q) \equiv \Pi_s(q^2, q \cdot u) + \Pi_q(q^2, q \cdot u) g + \Pi_u(q^2, q \cdot u) h. \tag{3}$$

In vacuum \Pi_s and \Pi_q become function of q^2 only and \Pi_u vanishes. For simplicity we will work in the rest frame of
nuclear matter, which implies that $u_u = (1, 0)$. Therefore, $\Pi_i(q^2, q \cdot u) \rightarrow \Pi_i(q_0, |q|)$ ($i = \{s, q, u\})$.

In the phenomenological side, the analytic properties of $\Pi(q)$ can be studied through a Lehman representation, which leads to a dispersion relation in $q_0$, for each invariant function, of the form

$$\Pi_i(q_0, |q|) = \frac{1}{2\pi i} \int_{-\infty}^{\infty} d\omega \frac{\Delta \Pi_i(\omega, |q|)}{\omega - q_0},$$

where we have omitted polynomials arising from the contour at large $|q_0|$, which will be eliminated by the Borel transformation. The discontinuity, defined by $\Delta \Pi_i(\omega, |q|)$, contains the spectral information on the quasiparticle, quasihole, and higher-energy states.

In vacuum, the spectral weights for baryon and antibaryon are related by charge conjugation symmetry and one usually parametrizes the spectral density as a single sharp pole, representing the lowest resonance, plus a smooth continuum, representing higher-mass states. At finite density, the ground state is no longer invariant under charge conjugation and, therefore, the spectral densities for baryon and antibaryon are not simply related.

The width of $\Theta^+$ in free space is very small and can be ignored on hadronic scales. At finite density, the width of $\Theta^+$ can be broadened due to strong interactions. Since the introduction of one extra parameter in the spectral density would reduce the predictive power of the sum rule, here we assume that a sharp pole hypothesis is still valid at finite density. In the context of relativistic phenomenology, we assume that $\Theta^+$ couples to the same scalar and vector fields as the nucleon and the hyperons in nuclear matter. Therefore, we parametrize the discontinuities as

$$\Delta \Pi_s(\omega, |q|) = -2\pi i \frac{\pm \lambda_\Theta^2 m_\Theta^2}{2E_q^*} [\delta(\omega - E_q) - \delta(\omega - \bar{E}_q)],$$

$$\Delta \Pi_q(\omega, |q|) = -2\pi i \frac{\lambda_\Theta^2}{2E_q^*} [\delta(\omega - E_q) - \delta(\omega - \bar{E}_q)],$$

where $m_\Theta = m_\Theta + \Sigma_s$, $E_q^* = \sqrt{m_\Theta^2 + q^2}$, $E_q = \Sigma_v + E_q^*$ and $\bar{E}_q = \Sigma_u - E_q^*$. $\lambda_\Theta^2$ measures the coupling of the interpolating field with the physical $\Theta^+$ in the medium and has opposite signs in $\Delta \Pi_s$ depending on the parity of $\Theta^+$. The scalar and vector self-energies of $\Theta^+$ in nuclear matter are given by $\Sigma_s$ and $\Sigma_v$, respectively. The positive- and negative-energy poles are at $E_q$ and $\bar{E}_q$ respectively. In Eq. (4) we have omitted the contributions from higher-energy states, which will be included later in the usual form, i.e., being approximated by the OPE spectral density starting at an effective threshold.

At finite density, the OPE for the invariant functions takes the general form

$$\Pi_i(q_0, |q|) = \sum_n C_n^i(q_0, |q|) \langle \hat{O}_n \rangle_{\rho N},$$

where $\langle \hat{O}_n \rangle_{\rho N} = \langle \Psi | \hat{O}_n | \Psi \rangle$, are the in-medium condensates. The Wilson coefficients, $C_n^i(q_0, |q|)$ depend only on QCD Lagrangian parameters. Therefore, all the density dependence of the correlator is included in the in-medium condensates.

Following ref. [8] we separate the invariant functions into two pieces that are even and odd in $q_0$:

$$\Pi_i(q_0, |q|) = \Pi_i^{E}(q_0^2, |q|) + q_0 \Pi_i^{O}(q_0^2, |q|).$$

In ref. [8] it was shown that the choice $q_0 = \bar{E}_q$ completely suppresses sharp excitations at $\bar{E}_q$ and also strongly suppresses a broad excitation in this vicinity. Since we are interested in the positive energy pole, we will use $q_0 = \bar{E}_q$.

By using the quark propagators given in refs. [8, 10] and working to leading order in perturbation theory, up to dimension 5, we get (after Borel transforming both sides of the sum rules):

$$\lambda_\Theta^2 e^{-(E_q^2 - q^2)/M^2} \equiv \frac{M^{12} E_5}{2^{10} \pi^8 s^5 l^7} + \frac{m_s \langle \bar{s} s \rangle_{\rho N}}{2^{8} \pi^6 s^5 l^5} M^8 E_3 - \frac{M^8 E_3}{2^{6} \pi^6 s^5 l^3} \left( \langle s^4 i D_0 s \rangle_{\rho N} - \frac{1}{4} m_s \langle \bar{s} s \rangle_{\rho N} \right) + \frac{\langle \bar{q} \gamma_5 \bar{q} \rangle_{\rho N}}{2^{12} \pi^8 s^5 l^5} M^8 E_3$$

$$- \frac{2 \langle q^2 i D_0 q \rangle_{\rho N}}{\pi^6 s^4 l^6} (2 M^8 E_3 + 2 M^6 E_2) - \frac{\bar{E}_q \langle \bar{q} i D_0 q \rangle_{\rho N}}{2^{4} \pi^6 s^4 l^4} M^8 E_3 + \frac{\bar{E}_q M^6 E_2}{2^{2} \pi^6 s^4 l^2} \left( \langle \bar{s}^4 i D_0 i D_0 \bar{s} \rangle_{\rho N} \right)$$

$$+ \frac{1}{12} \langle g_s s^4 \sigma \cdot G s \rangle_{\rho N} \rangle_{\rho N} - \frac{\bar{E}_q \langle g_s s^4 \sigma \cdot G s \rangle_{\rho N}}{2^{9} \pi^6 s^4 l^4} M^6 E_2 - \frac{E_q (M^6 E_2 + 2 q^2 M^4 E_2)}{2^{9} \pi^6 s^4 l^2} \left( \langle q^4 i D_0 i D_0 q \rangle_{\rho N} \right)$$

$$+ \frac{1}{12} \langle g_s q^4 \sigma \cdot G q \rangle_{\rho N} \rangle_{\rho N} + \frac{\bar{E}_q \langle g_s q^4 \sigma \cdot G q \rangle_{\rho N}}{2^{5} \pi^6 s^4 l^2} M^6 E_2,$$
\[ 9 M^8 E_3 + 4 q^2 M^6 E_3 \left( \langle \bar{s} i D_0 i D_0 s \rangle_{\rho N} + \frac{g_s \bar{s} \sigma \cdot G_s}{2^9 \pi^6 4!} \right) \rho_N + m_s (g_s^2 G_s^2)_{\rho N} \rho_N \left( \langle q^4 \rangle_{\rho N} + \langle s^4 \rangle_{\rho N} \right) \]
\[ + \frac{m_s (g_s^2 G_s^2)_{\rho N}}{2^{12} \pi^6 4!} \rho_N \]
\[ + \frac{E_0 M^8 E_3}{2^{21} \pi^6 5!} \left( \langle q^4 \rangle_{\rho N} + \langle s^4 \rangle_{\rho N} \right) \]
\[ \lambda_{\rho N}^2 \sum_{\sigma} e^{-\left(\vec{q}^2 - \vec{q}_0^2\right)/M^2} = \frac{-M^{10} E_1}{2^{11} \pi^6 5!} \left( \langle q^4 \rangle_{\rho N} + 3 \langle s^4 \rangle_{\rho N} \right) - \frac{(5 M^8 E_3 + 2 q^2 M^6 E_3)}{2^{19} \pi^6 5!} \left( \langle i \bar{s} i D_0 i D_0 s \rangle_{\rho N} \right) \rho_N \]
\[ + \frac{1}{12} \left( g_s \bar{s} \sigma \cdot G_s \right)_{\rho N} + \frac{(g_s^2 G_s^2)_{\rho N} M^6 E_3 - (5 M^8 E_3 + 2 q^2 M^6 E_3)}{2^{22} \pi^6 5!} \left( \langle q^4 i D_0 i D_0 q \rangle_{\rho N} \right) \rho_N \]
\[ + \frac{1}{12} \left( g_s^2 G_s^2 \right)_{\rho N} \rho_N \left( \langle q^4 \sigma \cdot G q \rangle_{\rho N} + \langle q^4 \bar{q} \sigma \cdot G q \rangle_{\rho N} \right) \rho_N - \frac{E_0 (g_s^2 G_s^2)_{\rho N} M^6 E_3}{3^{15} \pi^6 5!} \]
\[ + \frac{E_0 M^8 E_3}{2^{11} \pi^6 5!} \left( \langle s^4 i D_0 s \rangle_{\rho N} - \frac{1}{4} \langle q^4 i D_0 q \rangle_{\rho N} \right) , \tag{10} \]

where we have defined
\[ E_n \equiv 1 - e^{-s_0/M^2} \sum_{k=0}^{n} \left( \frac{s_0}{M^2} \right)^k \frac{1}{k!} , \tag{11} \]

which accounts for the continuum contribution with \( s_0 \) being the continuum threshold.

To extract the self-energies from the above finite density sum rules, one has to know the values of the intermediate condensates. To first order in the nucleon density, one can write \( \langle \bar{O} \rangle_{\rho N} \sim \langle \bar{O} \rangle_N + \langle \bar{O} \rangle_{\rho N} \) where \( \langle \bar{O} \rangle_N \) is the spin-averaged nucleon matrix element. The simplest in-medium condensates are \( \langle q^4 q \rangle_{\rho N} \) and \( \langle s^4 \rangle_{\rho N} \). Since the baryon current is conserved, \( \langle q^4 q \rangle_{\rho N} \) and \( \langle s^4 \rangle_{\rho N} \) are proportional to the nucleon and strangeness densities: \( \langle q^4 q \rangle_{\rho N} = \frac{3}{8} \rho_N \) and \( \langle s^4 \rangle_{\rho N} = 0 \). These are exact results. We extract the values of the other condensates from refs. [8, 11, 12]: \( \langle \bar{q} \bar{q} \rho N \rangle = \langle \bar{q} \bar{q} \rangle + \frac{m_s^{\perp} m_u + m_d}{2^{12} \pi^6 4!} \rho_N , \langle \bar{q} \rho N \rangle = \langle \bar{q} \rangle + \frac{m_s^{\perp} m_u + m_d}{2^{12} \pi^6 5!} \rho_N , \langle \bar{q} \sigma \cdot G \rangle_{\rho N} = \langle \bar{q} \sigma \cdot G \rangle + \frac{m_s^{\perp} m_u + m_d}{2^{12} \pi^6 5!} \rho_N \) and \( \langle \bar{q} \bar{q} \sigma \cdot G \rangle_{\rho N} = \langle \bar{q} \bar{q} \sigma \cdot G \rangle + \frac{m_s^{\perp} m_u + m_d}{2^{12} \pi^6 5!} \rho_N \). We see that there is a good agreement between them. In the lower part of Fig. 1 (solid line) we show the effective \( \Sigma^+ \) mass and continuum threshold are: \( m_{\Sigma^+} = 1.75 \text{ GeV} \), \( \Sigma^+ = -150 \text{ MeV} \) and \( s_0 = 3.6 \text{ GeV}^2 \). The first astonishing result is that, opposite to what was obtained in the nucleon and hyperon cases, the scalar self-energy is positive and the vector self-energy is negative. Therefore the effective \( \Sigma^+ \) mass in medium, \( m_{\Sigma^+} = m_{\Sigma^+} + \Sigma^+ \), is bigger than the free \( \Sigma^+ \) mass. However, since \( \Sigma^+ \approx 110 \text{ MeV} \), there is still a substantial cancelation between \( \Sigma^+ \) and \( \Sigma^+ \) in medium.

Looking at Eq. (10) it is easy to understand why we get \( \Sigma^+ < 0 \), since, in contrast to the nucleon and hyperon cases, all the terms in the right-hand side are negative.
The origin of this difference in the sign can be attributed to the existence of an antiquark in the Θ+ interpolating field. Since $\Sigma_s$ and $\Sigma_v$ are essentially the real parts of the optical potential, this qualitative result shows that there is an overall attractive $\Theta$–nucleon interaction.

The optimized results for the ratio $m_{\Theta}^*/m_\Theta$ (solid line) and for the scalar (dashed line) and vector (dotted line) self-energies as functions of $M^2$ are plotted in Fig. 2 for $y = 0.3$. We see that the curves are quite flat, indicating a weak dependence of the predicted results on $M^2$. The results are also not sensitive to the value of $|\vec{q}|$. Using $|\vec{q}| = 0$ instead of $|\vec{q}| = 270\text{ MeV}$ does not alter our results. On the other hand, the results are very sensitive to the value of $y$. In the range $0 \leq y \leq 0.5$ we got $50\text{ MeV} \leq \Sigma_s \leq 150\text{ MeV}$ and $-90\text{ MeV} \leq \Sigma_v \geq -190\text{ MeV}$. What is very interesting is that the sum $\Sigma_s + \Sigma_v$, which can be associated with the depth of the $\Theta^+$ potential in nuclear matter, remains independent of $y$ and is about $-40\text{ MeV}$. However, this value is very sensitive to the value of the gluon condensate. If we change the value of the gluon condensate to $\langle g_s^2 G^2 \rangle = 0.24\text{ GeV}^4$, as used in ref. [19], we get $\Sigma_s \sim 90\text{ MeV}$ and $\Sigma_v \sim -180\text{ MeV}$ which would imply in a potential depth of about $-90\text{ MeV}$.

The depth of the $\Theta^+$ potential in nuclear matter, $U$, was studied in ref. [19] using a relativistic mean field framework, where $\Theta^+$ couples with scalar and isoscalar-vector mesons. They found $-90\text{ MeV} \leq U \leq -45\text{ MeV}$ depending on the values used for $m_N^*/m_N$ (the ratio of the nucleon mass), and the coupling constants. In our work we found that the potential depth depends strongly on the value of the gluon condensate. In the range $0.47\text{ GeV}^4 \leq \langle g_s^2 G^2 \rangle \leq 0.24\text{ GeV}^4$ we got $-40\text{ MeV} \geq U \geq -90\text{ MeV}$, which is compatible with the findings of ref. [19].

It is also interesting that, in spite of $m_{\Theta}^*$ being bigger than $m_\Theta$, the energy of the quasi-$\Theta$ in nuclear matter, $E_q$, is smaller than $m_\Theta$ and it is the most stable result of our calculation. It does not depend on $y$ neither on $\langle g_s^2 G^2 \rangle$. While $m_\Theta$, $m_{\Theta}^*$ and $\Sigma_v$ varies significantly with $y$ and $\langle g_s^2 G^2 \rangle$ we got $E_q \sim 1.60\text{ GeV}$ for all the values of the parameters (the smallest value obtained for the $\Theta^+$ mass was $m_{\Theta} \simeq 1.64\text{ GeV}$).

All the results given above were obtained at the nuclear matter saturation density $\rho_N$. In Fig. 3 we show the density dependence of the scalar and vector self-energies. We see that most of the variation happens for $\rho_N < 0.5\rho_0$. Since we are expanding the density dependence of the
condensates up to first order in the nucleon density, the obtained density dependence of our results is compatible with the approximations made.

We would like to point out that, since our sum rules do not depend on the four-quark condensates, our results are free from the uncertainties associated with the density dependence of these condensates, which were the biggest source of uncertainty in the case of the nucleon and hyperons studied in refs. \([7,8,9,10]\). In that case, the scalar self-energy was very sensitive to the density dependence of the four-quark condensates, and one could even get a repulsive net self-energy. The results were a bit more stable for the vector self-energy. Comparing the values of the vector self-energies one has from refs. \([8,9,10]\): \(|\Sigma_0^v/\Sigma_0^N| \approx 0.3 - 0.4, |\Sigma_v^v/\Sigma_v^N| \approx 0.8 - 1.1\). And here we got \(|\Sigma_0^\Theta/\Sigma_0^N| \approx 0.3 - 0.5\). In terms of relativistic hadronic model these results would imply that the coupling of the hyperon \(\Sigma\) to the Lorentz vector field is very similar to the corresponding nucleon coupling, while the coupling of the hyperon \(\Lambda\) is similar to the coupling of \(\Theta\) and both are much weaker than the corresponding nucleon coupling. These results can be understood in terms of the interpolating fields used to study these baryons. For \(\Lambda\) there is a \([ud]\) diquark with spin and isospin zero plus the strange quark, that carries the spin of \(\Lambda\). Assuming no admixture of strange quark content in the vector meson \(\omega\) (the Lorentz vector field), there will be no coupling between \(\Lambda\) and \(\omega\). The same happens in the case of \(\Sigma\), since in its interpolating field there is a \([qs]\) diquark with spin zero. Therefore, it is the light quark that carries the spin of \(\Sigma\), as in the nucleon case.

As a final remark, we would like to mention that in ref. \([20]\) it was suggested that if the parity of \(\Theta^+\) were positive, then there would be a strong attractive \(\Theta\)–nucleon interaction. In our calculation we got a negative parity for \(\Theta^+\). Its interaction with the nucleon is still attractive, and roughly of the same order of the nucleon-nucleon interaction. Bigger attraction is obtained with smaller values of the gluon condensate.

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