Evaluation of QCD sum rules for HADES

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QCD sum rules are evaluated at finite nucleon densities and temperatures to determine the change of pole mass parameters for the lightest vector mesons $\rho$, $\omega$ and $\phi$ in a strongly interacting medium at conditions relevant for the starting experiments at HADES. The role of the four-quark condensate is highlighted. A few estimates (within a fire ball model and BUU calculations) of dilepton spectra in heavy-ion collisions at 1 AGeV are presented.

1 Introduction

Hadrons are excitations of the QCD vacuum. There are many different species and masses of them. Their excitation spectra resemble, to some extent, to level schemes of atoms. The latter ones are known to be modified by external fields: The Stark and Zeeman effects manifest themselves as shifts of spectral lines. By analogy, one might expect a similar behavior of hadrons in strongly interacting matter. In a first approximation, hadrons should experience mass shifts, but considerable broadenings or strong modifications (e.g. mixing) of the spectral strengths are conceivable, too. A verification and understanding of such in-medium modifications of hadrons would allow to contribute to the ”mystery of particle masses”.

This topic was essentially triggered by the Brown-Rho scaling hypothesis \[1\] according to which the masses of the light vector mesons ($V = \rho, \omega, \phi$) behave as $m_V^2 \propto \langle \bar{q}q \rangle$, where $\langle \bar{q}q \rangle$ is the chiral condensate. In such a way fundamental and non-perturbative QCD quantities can be accessible. This idea influenced strongly the HADES project which aims at measuring the density and temperature dependence of $\langle \bar{q}q \rangle$ via vector meson ”spectral lines”. A strong impact on the field had QCD sum rules à la Shifman, Vainshtein, Zakharov \[2\] which, indeed, correlate condensates (and moments of the parton distributions in hadrons) and the pole masses of hadrons.

The purpose of the present contribution is two-fold. First, we revisit the QCD sum rule approach to estimate the change of the vector meson pole mass parameter at finite baryon density $n$ and temperature $T$. Second, we present dynamical calculations, basing on hydrodynamics (fire ball model) and transport calculations (BUU) to elucidate the prospects to verify in the starting experiments at HADES the predicted in-medium modifications of vector mesons.

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2 QCD sum rules revisited

The basic object is the retarded current-current correlation function

\[ \Pi_{\mu\nu}^R(q; \mu_N, T) = i \int d^4x \, e^{iqx} \Theta(x^0) \langle [J_\mu(x), J_\nu(0)] \rangle_{\mu_N,T}, \]  

(1)

with conserved vector currents \( J_\mu^\rho(\omega, \phi) = \frac{1}{2} (\bar{u} \gamma_\mu u \mp \bar{d} \gamma_\mu d), \) \( \bar{s} \gamma_\mu s \) carrying the quantum numbers of the \( \rho, \omega \) and \( \phi \) meson, respectively, but expressed by quark fields. \( \langle \cdot \cdot \rangle_{\mu_N,T} \) means grand canonical averaging with respect to chemical potential \( \mu_N \) and temperature \( T \). Here we focus on vector mesons at rest and constrain ourselves on the longitudinal part \( \Pi_L^R \) of the correlator \( \Pi_{\mu\nu}^R \).

The sum rule approach rests on a comparison of two different evaluations of \( \Pi_L^R \). One can use (i) the analytic properties to note a dispersion relation

\[ \Pi_L^R(q^0; \mu_N, T) = \frac{1}{\pi} \int_0^\infty ds \, \frac{\text{Im} \Pi_L^R(s; \mu_N, T)}{s - (q_0 + i\epsilon)^2} + \cdots, \]

(2)

where \( \cdots \) indicate subtractions, and (ii) the Operator Product Expansion at large \( Q^2 \equiv -q_0^2 > 0 \)

\[ \Pi_L^R(Q^2) = -C_0 \ln Q^2 + \sum_{n=1}^\infty \frac{C_n}{Q^{2n}}, \]

(3)

where \( C_n \) are quantities containing Wilson coefficients and condensates and moments of parton distributions. Defining the spectral density \( \rho_{\text{had}}(s; \mu_N, T) = \frac{1}{\pi} \text{Im} \Pi_L^R(s; \mu_N, T) \) and performing a Borel transformation one arrives at the QCD sum rule (QSR)

\[ \int_0^\infty ds \, \rho_{\text{had}}(s) \, e^{-s/M^2} = M^2 \left( C_0 + \sum_{n=1}^\infty \frac{C_n}{(n-1)! M^{2n}} \right), \]

(4)

where \( M \) is the Borel mass. Assuming that the r.h.s. is calculable from first principles one can use the QSR (i) either as consistency check of hadronic models for \( \rho_{\text{had}}(s) \) (cf. [3]), or (ii) as one equation for a physically relevant parameter in a simple parameterization of \( \rho_{\text{had}}(s) \), or (iii) as constraint for a few parameters entering models of \( \rho_{\text{had}}(s) \) (e.g. to find a width-mass relation [4], or to interrelate various parameters with the chiral gap [4]).

We follow here the possibility (ii) as exercised by many previous authors like in [3],[4]. Inspired by the experimental results on the cross sections of the reactions \( e^+ e^- \rightarrow n\pi, K^+ K^- \), which is intimately related to \( \rho_{\text{had}} \) and which exhibits one low-lying distinctive resonance (either \( \rho \) or \( \omega \) or \( \phi \)) followed by a flat continuum, one parameterizes

\[ \rho_{\text{had}}^V(s; T, n) = F_V \delta(s - m_V^2) + C_0 \Theta(s - s_0^V) + (\rho_{\text{scatt}}^{V,\pi} + \rho_{\text{scatt}}^{V,N}) \delta(s), \]

(5)

where \( m_V \) is the wanted pole mass parameter, \( s_0^V \) stands for the continuum threshold, and \( \rho_{\text{scatt}}^{V,\pi(N)} \) denote the forward scattering amplitudes or Landau damping terms, which read for the \( \rho \) meson \( \rho_{\text{scatt}}^{\rho,\pi} = (48\pi^2)^{-1} \int_{4M_\rho^2}^\infty d\omega^2 \, \hat{n}_F \sqrt{1 - 4M_\rho^2/\omega^2} \), \( \rho_{\text{scatt}}^{\rho,N} = (24\pi^2)^{-1} \int_{4m_\rho^2}^\infty d\omega^2 \, \hat{n}_B \sqrt{1 - 4m_\rho^2/\omega^2} \) with \( \hat{n}_F = e^{(\omega - 2\mu_N)/2T} + 1 \)^{-1}, \( \hat{n}_B = e^{\omega/2T - 1} \)^{-1}, while for the \( \omega \) meson, \( \rho_{\text{scatt}}^{\omega,\pi} = 0 \) due to symmetry, and \( \rho_{\text{scatt}}^{\omega,N} = 9 \rho_{\text{scatt}}^{\rho,N} \).
due to the different isospin structure of the $\rho$ and $\omega$ mesons \[8\]; for the $\phi$ meson the Landau damping for the pion and nucleon gas is negligible \[10\].

Truncating the r.h.s of eq. \[7\] the sum rule might be cast in the form

$$m_V^2 = M^2 \left( \frac{C_0 \left( 1 - \left[ 1 + \frac{s_v^V}{M^2} \right] e^{-s_v^V/M^2} \right) - C_2 M^{-4} - C_3 M^{-6}}{C_0 \left( 1 - e^{-s_v^V/M^2} \right) + C_1 M^{-2} + C_2 M^{-4} + \frac{1}{2} C_3 M^{-6} - (\rho_{\text{scatt}} + \rho_{\text{scatt}})^V \pi M^{-2}} \right),$$

where the coefficients up to mass dimension 6 have been calculated by many authors in the low-temperature or low-density approximation. In the low-$T$ and low-$n_N$ approximation our result is

$$C_0 = \frac{1}{8\pi^2} \left( 1 + \frac{\alpha_s(\mu^2)}{\pi} \right), \quad C_1 = -\frac{3m_q^2}{4\pi^2}, \quad C_2 = q_2 + g_2 + a_2, \quad C_3 = q_4 + a_4, \quad q_2 = m_q \langle \bar{u}u \rangle_0 + 2M_N\sigma_N Y I_1^N + \frac{3}{4} m_\pi^2 \epsilon^{\rho(\omega)} I_1^\pi,$$

$$g_2 = \frac{1}{24} \left( \frac{\alpha_G}{\pi} \right) G^2 \rho_{\text{scatt}} + \rho_{\text{scatt}}^V \pi M^{-2},$$

$$q_4 = -\frac{112}{81} \pi \alpha_s \kappa_0 \langle \bar{u}u \rangle_0^2 \left[ 1 - \frac{\kappa_N}{\kappa_0} \frac{4M_N\sigma_N}{m_q \langle \bar{u}u \rangle_0} Y I_1^N - \frac{36}{7f_\pi^2} \xi^{\rho} I_1^\pi \right],$$

$$a_2 = M_N^2 A_{2,N}^{u+d} I_1^N + \frac{4}{3} A_{2,N}^{u+d} I_2^N + \frac{3}{4} m_\pi^2 A_{2,\pi}^{u+d} I_1^\pi + A_{2,\pi}^{u+d} I_2^\pi,$$

$$a_4 = -\frac{5}{3} M_N^4 A_{4,N}^{u+d} I_1^N - \frac{20}{3} M_N^2 A_{4,N}^{u+d} I_2^N - \frac{16}{3} A_{4,N}^{u+d} I_3^N - \frac{5}{4} m_\pi^4 A_{4,\pi}^{u+d} I_1^\pi - 5m_\pi^2 A_{4,\pi}^{u+d} I_2^\pi - 4A_{4,\pi}^{u+d} I_3^\pi,$$

where $\xi^{\rho(\omega)} = 1$, $Y = 1$ for $\rho$ ($\omega$) mesons, and elsewhere $\xi^{-} = 0$. The sequence of replacements $m_q \rightarrow m_s$, $\langle \bar{u}u \rangle_0 \rightarrow \langle \bar{s}s \rangle_0$, $Y \rightarrow y m_s/m_q$, and $A_{u,d,N,\pi}^{\rho(\omega)} \rightarrow A_{u,d,N,\pi}^{\rho(\omega)}$ holds for the $\phi$ meson. The above integrals are $I_1^N = \int d^3k [2(\pi)^3 E_k]^{-1} k^{2n-2} n_F$, $I_2^N = \int d^3p [(2\pi)^3 E_p]^{-1} p^{2n-2} n_B$, where $n_B = |e^{E_p/T} - 1|^{-1}$, and $n_F = |e^{(E_k - \mu_N)/T} + 1|^{-1}$ are thermal Boson and Fermion distributions, and $\mu_N$ is the chemical potential related to the nucleon density $n_N$.

Specific for our approach is the treatment of the four-quark condensate for which we extend the widely used ground state saturation assumption for the in-medium four-quark condensate at $T = 0$ in the following way

$$\langle \bar{q} \gamma_\mu \gamma_5 \lambda^a q \rangle_{\mu N}^2 = -\langle \bar{q} \gamma_\mu \lambda^a q \rangle_{\mu N}^2 = \frac{16}{9} \kappa(n_N) \langle \bar{q} q \rangle_{\mu N}^2,$$

where the density dependent factor $\kappa(n_N)$ is introduced to control a deviation from the exact ground state saturation. As pointed out in \[\ref{8}\] $\kappa(n_N)$ reflects the contribution from the scalar low-energy excitations of the ground state and seems to be weekly dependent on $n_N$, so that one can use $\kappa(n_N) = \text{const}$ as first approximation. In this case the density dependence of the four-quark condensate appears only via the density dependence of the quark condensate squared. In linear density approximation it is given by $\langle \bar{q} q \rangle_{\mu N}^2 = \langle \bar{q} q \rangle_{\mu N}^2 + \langle \bar{q} q \rangle_{\mu N}^2 |N\bar{q} q| N_{\mu N}^2$. We go beyond this approximation
The limit poorly known parameter \( \kappa \) and of the gluon condensate employed also for evaluating the four-quark condensates in the general case with this matrix element does not depend on particle momenta and temperature it can be further described by a constant factor \( \kappa_0 \). Our parameterization of the four-quark condensate at \( T = 0 \) has then the form

\[
\langle (\bar{q}\gamma_\mu\gamma_5\lambda^a q)^2 \rangle_{\mu,N} = \frac{16}{9} \langle \bar{q}q \rangle_0^2 \kappa_0 \left( 1 + \frac{\kappa_N}{\kappa_0} \frac{\langle N|\bar{q}q|N \rangle}{\langle \bar{q}q \rangle_0} \frac{n_N}{M_N} \right). \tag{15}
\]

The limit \( \kappa_N = \kappa_0 = 2.36 \) is used in \([3]\), while the parameterization \( \kappa_N = 1.4 \) and \( \kappa_0 = 3.3 \) with \( \langle \bar{q}q \rangle_0 = (-230\text{MeV})^3 \) corresponds to the choice in \([4]\). Below we vary the poorly known parameter \( \kappa_N \) to estimate the contributions of the four-quark condensates to the QSR. The needed ansatz for the nucleon matrix element of the scalar four-quark condensate can be extracted from the direct comparison of our parameterization in eq. (15) and the general expression for the condensates via the matrix elements (the latter ones to be taken at \( T = 0 \)) as \( \langle N|\bar{q}q|N \rangle = \frac{32}{9} \langle \bar{q}q \rangle_0 \langle N|\bar{q}q|N \rangle \kappa_N \). Since this matrix element does not depend on particle momenta and temperature it can be employed also for evaluating the four-quark condensates in the general case with \( T \neq 0 \) and \( \mu_N \neq 0 \).

The remaining parameters are specified by \( \alpha_s = 0.38 \), \( m_q = 0.0055 \text{GeV} \), \( m_s = 0.130 \text{GeV} \), \( f_\pi = 0.093 \text{GeV} \), \( M_N = 0.938 \text{GeV} \), \( M_0^2 = 0.770 \text{GeV} \), \( \sigma_N = 0.045 \text{GeV} \), \( y = 0.22 \), \( m_\pi = 0.138 \text{GeV} \), \( \langle \bar{u}u \rangle_0 = \langle \bar{d}d \rangle_0 = (-0.245 \text{GeV})^3 \), \( \langle \bar{s}s \rangle_0 = 0.8 \langle \bar{u}u \rangle_0 \), \( \langle \bar{q}q \rangle_0 = (0.33 \text{GeV})^4 \), \( A_{2,0}^{u+d} = 1.02 \), \( A_{4,0}^{u+d} = 0.12 \), \( A_{2,\pi}^{u+d} = 0.97 \), \( A_{4,\pi}^{u+d} = 0.255 \), \( A_{2,0}^{2,\pi} = 0.08 \), \( A_{4,\pi}^{2,\pi} = 0.93 \).

The density and temperature dependence of the chiral condensates \( \langle \bar{u}u \rangle \) and \( \langle \bar{s}s \rangle \) and of the gluon condensate \( \langle G^2 \rangle \) are exhibited in fig. 1. Note the linear density dependence of the condensates as a function of temperature and density (left: chiral condensate, top-right: strange chiral condensate, bottom-right: gluon condensate). Note that the validity range of the present calculations is constrained to small densities and temperatures. \( n_N \leq 3n_0 \) and \( T \leq 100 \text{MeV} \) are relevant for HADES.
dependence and the nearly independence of the temperature.

The parameter $s_0^{UV}$ in eq. (3) is fixed by requiring maximum flatness of $m_V(M)$ within a suitably chosen Borel window $M_{\text{min}} \cdots M_{\text{max}}$ determined by the "10% + 50%" rule [11]. If we follow [5] and choose $\kappa_N/\kappa_0 = 0.42$ and use $\kappa_0$ to adjust the vacuum values of $m_V$ to the physical values, we get the results displayed in fig. 2. Due to some cancellations of the coefficients entering the $\rho$ meson sum rule, the dropping mass parameter is determined by the four-quark condensate. For the $\omega$ meson the Landau damping is dominating [8]. The dropping $\phi$ meson mass parameter reflects directly the strange chiral condensate; a smaller value of $y$ diminishes strongly the mass shift.

It turns out that the actual $\rho$ and $\omega$ mass shifts depend on the poorly known parameter $\kappa_N$ which parameterizes the density dependence of the four-quark condensate. In contrast to the above ratio $\kappa_N/\kappa_0 = 0.42$ claimed in [5], the groundstate saturation hypothesis requires $\kappa_N/\kappa_0 = 1$. This gives a stronger dropping $m_\rho$. For the $\omega$ meson the rising $m_\omega$, advocated in [8], is changed into a drop, as seen in fig. 3. Thus, even within the present simple pole approximation, the qualitative behavior of the $\omega$ meson remains unsettled.

3 Estimates of dilepton spectra in heavy-ion collisions

Up to now we focused on the behavior of $m_V$ within the QSR approach in zero-width approximation. The width-mass relation derived in [4] is consistent with dropping $\rho$ mass as long as the width is constrained to $\Gamma < 600$ MeV. However, a unique prediction of the in-medium modification of the width and the mass separately seems not be possible within the QSR approach. Rather, one has to resort to hadronic models which predict a considerable in-medium broadening. To elucidate the consequences of
Figure 4: Dilepton emission rate at $T = 70$ MeV as a function of invariant mass. Left panel: integrated over all $M_\perp$, right panel: for $M_\perp = 1.3$ GeV (dashed curve: $v_r = 0$, solid curve: $v_r = 0.3$). A vector meson broadening factor of 3 is assumed.

these in-medium changes together we use the following parameterization of the emission rate of hadron matter \cite{12}

$$
\frac{dN}{d^4x d^2Q} = \sum_V \frac{2d_V}{(2\pi)^3} \exp \left\{ -\frac{u \cdot Q}{T} \right\} \frac{1}{\pi} \frac{(m_V \Gamma_V^{tot})(M_{V\rightarrow e^+e^-})}{(M^2 - m_V^2)^2 + (m_V \Gamma_V^{tot})^2},
$$

(16)

where $Q$ is the four-momentum vector of the lepton pair with invariant mass $M$, $u$ denotes the four-velocity of matter, and $d_V$ stands for the degeneracy factor of the vector meson $V$. The corresponding invariant mass spectrum is displayed in fig. 4 for $T = 70$ MeV, where we assume that all vector mesons experience the overall broadening $\Gamma_V^{tot} = b\Gamma_V^{tot,0}$ with $b = 3$. As seen in fig. 4, the $\rho$ meson is smeared out to such a degree that no distinct peak is visible. The $\omega$ and $\phi$ can be identified as they are still sufficiently narrow. It is the double differential spectrum $dN/d^4x dM dM_\perp$ which allows access to the $\rho$ peak. As seen in the right panel of fig. 4, for fixed $M_\perp = 1.3$ GeV there is a pronounced $\rho$ peak with the $\omega$ sitting on top. As known for a long time, such double differential spectra are sensitive also to the flow of matter. This is highlighted by a comparison of two flow velocities, $v = 0$ (lower curve in left panel of fig. 4) and 0.3 (upper curve). The full invariant mass spectra are insensitive to flow effects.

3.1 Hydrodynamics

In \cite{12} we have studied, within a scenario of an expanding fire ball, the resulting dilepton spectra. Instead of solving the hydrodynamical flow equations in full detail we have used a schematic expansion dynamics governed by $R(t) = R_i + v_r t$ with $v_r = 0.3$, entropy conservation, $s/n = const$, and baryon conservation ($N = 330$). The final (freeze-out) parameters are from an analysis of hadron multiplicities at SIS energies \cite{13} $n_f = 0.3n_0$, $T_f = 50$ MeV. Assuming an initial density $n_i = 3n_0$, as suggested by BUU calculations, one gets $T_i = 90$ MeV for the initial temperature on the isentropic line. Keeping the above assumed in-medium broadening by a factor $b = 3$ and using
Figure 5: Space-time integrated dilepton spectra from the expanding fire ball described in text. Left panel: invariant mass spectrum integrated over all $M_\perp$, right panel: $M_\perp = 0.9$ GeV. A vector meson broadening factor of 3 is assumed and a $\rho$ meson mass shift of $\Delta m_\rho = 300$ MeV.

The $\rho$ meson mass shift of $\Delta m_\rho = 300$ MeV (a maximum value with respect to the above QSR studies) one gets the dilepton spectra exhibited in fig. 5. According to the above QSR results for $\kappa_N \approx 2$ we employ here no mass shift of the $\omega$ meson. While in the invariant mass spectrum again no $\rho$ peak is visible, the double-differential spectrum with $M_\perp = 0.9$ GeV shows clearly the shifted peak.

In [12] also such scenarios are studied in detail, where the density depend broadening and mass shifts are incorporated. These studies support the need to measure double-differential cross sections to have access to the $\rho$ peak.

3.2 BUU calculations

The above hydrodynamical calculations are based on local thermal and chemical equilibrium. These stringent assumptions are not longer needed in transport model calculations. We employ here a BUU code to calculate the dilepton spectra for central collisions Au(1 AGeV) + Au. The results are displayed in fig. 6 for $b = 3$ and either no mass shift (left column) or a $\rho$ mass shift of 150 MeV (right column). We have included the dileptons from $\pi^+\pi^-$ annihilations via the intermediate $\rho$ meson, decays of $\rho$ mesons from baryon resonance decays, $pn$ bremsstrahlung, $\Delta$ and $\eta$ Dalitz decays and $\omega$ decays. In contrast to the above hydrodynamical calculations the $\rho$ and $\omega$ mesons are not allowed to undergo equilibration in the BUU code, but rather are propagated in mean field approximation along their original trajectories. The present BUU calculations confirm the results of the hydrodynamical model: Only when selecting a suitable $M_\perp$ bin one gets access to the $\rho$ peak. However, even in case of no mass shift, due to phase space weighting the $\rho$ peak seems shifted (see left column in fig. 6). The role of the phase space weight is highlighted in the right column of fig. 6, where the $\rho$ peak appears down-shifted by somewhat more than 150 MeV.

In addition, the BUU calculations show that the Dalitz and bremsstrahlung yields
Figure 6: Dilepton spectra as a function of invariant mass for several $M_\perp$ bins. Left column: no mass shift, right column: density independent mass shift $\Delta m_\rho = 150$ MeV. An overall vector meson broadening by a factor 3 is assumed. $\omega$ and $\phi$ contributions are not shown. $d\sigma^{\text{scal}}/dM = \pi R^2 (A_{\text{projectile}} \ast A_{\text{target}})^{-1} \int_{\Delta M_\perp} dM_\perp dN/(dM dM_\perp)$. 
do not disturb the wanted $\rho$ signal.

Various estimates demonstrate that the $\omega$ meson, even for some broadening, remains visible as distinctive sharp peak. Therefore, also small shifts should be visible. Here, however, a subtle interplay of $\omega$ decays in medium and after freeze-out occur. To elucidate the chances to verify possible in-medium $\omega$ modifications we have assumed $m_\omega = m_\omega^0 - \delta m_\omega (n/n_0)$ and $\Gamma_\omega = \Gamma_\omega^0 - \delta \Gamma_\omega (n/n_0)$ with $\delta m_\omega = 70$ MeV and $\delta \Gamma_\omega = 50$ MeV and obtain the results exhibited in fig. 7. One can see clearly a double-peak structure from vacuum decays and in-medium decays. As expected, selecting low-$Q_\perp$ dileptons the vacuum contribution is strongly reduced since the parent $\omega$ is nearly at rest. The smearing of the $\omega$ due to the density dependence can mimic a peak structure at the vacuum $\rho$ position.

4 Summary

In summary we have revisited the QCD sum rule approach to the in-medium behavior of the light vector mesons in zero-width approximation. Due to cancellations the $\rho$ meson mass shift depends on the four-quark condensate, but is always decreasing with increasing density. The large Landau damping and the uncertain four-quark condensate
do not allow a definite statement on the $\omega$ meson mass; for the ground state saturation the mass drops. The actual mass shift of the $\phi$ meson depends sensitively on the uncertain strangeness fraction in the nucleon; the $\phi$ mass is directly related to the strange chiral condensate.

Our hydrodynamical fire ball calculations show that a double-differential cross section needs to be measured to get access to the $\rho$ peak in case of strong in-medium broadening. This is supported by transport calculations relying on a BUU code. To amplify a possible in-medium $\omega$ shift the selection of low-momentum pairs is very helpful.

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