Modification of the $\phi$-meson spectrum in nuclear matter

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Abstract

The vacuum spectrum of the $\phi$-meson is characterized by its decay into $K\bar{K}$. Modifications of the $K\bar{K}$-loops in baryonic matter change this spectrum. We calculate these in-medium modifications taking both $s$- and $p$-wave kaon-nucleon interactions into account. We use results of the in-medium $K$ and $\bar{K}$ spectra determined previously from a coupled channel approach based on a chiral effective Lagrangian. Altogether we find a very small shift of the $\phi$ meson mass, by less than 10 MeV at normal nuclear matter density $\rho_0$. The in-medium decay width of the $\phi$ meson increases such that its life time at $\rho = \rho_0$ is reduced to less than 5 fm/c. It should therefore be possible to observe medium effects in reactions such as $\pi^- p \rightarrow \phi n$ in heavy nuclei, where the $\phi$ meson can be produced with small momentum.

Introduction. The study of in-medium properties of hadrons is a topic of continuing interest. Experiments planned at GSI (HADES) and running at CERN (e.g. CERES) detect dilepton pairs in high energy collisions of nuclei. This opens the possibility of investigating vector mesons in hot and dense hadronic matter. In order to explore such medium effects it is necessary that the vector mesons decay inside the hot and dense region of the collision zone. The $\phi$ meson with its small width of 4.4 MeV has a lifetime of about 45 fm/c which is obviously too large to observe any medium effects. On the other hand, we demonstrate in this note that medium modifications are expected to increase the $\phi$ width and shorten its lifetime to less than 5 fm/c at the density of normal nuclear matter. One then enters the range in which medium effects of slowly moving $\phi$-mesons could become visible. A reaction of particular interest is $\pi^- p \rightarrow \phi n$ in heavy nuclei. This process violates the OZI rule but substantial $\omega \phi$ mixing makes the $\phi$ production rate large enough to be well detectable. Such an experiment could be performed at GSI where a pion beam can be used in combination with HADES to measure $\pi^- p \rightarrow \omega n$.

The purpose of this paper is to provide a systematic calculation of the in-medium $\phi$ meson self energy. We first review briefly the vacuum properties of the $\phi$ meson and discuss its self-energy. We then extend this for finite densities by including the in-medium

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interactions of the decay kaons, taking into account both s- and p-wave interactions of $K$ and $\overline{K}$ with nucleons in nuclear matter. In the final part a brief summary and discussion of the results will be given.

**The $\phi$ meson in vacuum.** The $\phi$ meson is observed as a pronounced resonance in the strange quark sector of the electromagnetic current-current correlation function [2]. The Fourier transform of the correlation function is

$$\Pi_{\mu\nu}(q) = i \int d^4 x \ e^{iq \cdot x} \langle 0 | T j_\mu(x) j_\nu(0) | 0 \rangle,$$

(1)

where in the present context $j_\mu$ represents the strange quark current,

$$j_\mu = -\frac{1}{3} (\bar{s} \gamma_\mu s).$$

(2)

Current conservation leads to a transverse tensor structure

$$\Pi_{\mu\nu}(q) = \left( g_{\mu\nu} - \frac{q_\mu q_\nu}{q^2} \right) \Pi(q^2),$$

(3)

which defines the scalar function $\Pi(q^2) = -\frac{1}{3} \Pi_{\mu\nu}(q)$. The low energy spectrum of the correlation function is well described by Vector Meson Dominance (VMD). We use our improved VMD approach of ref.[4] which gives

$$\text{Im} \Pi(q^2) = \frac{\text{Im} \Pi^\text{vac}(q^2)}{g_\phi^2} \left| \frac{(1 - a_\phi) q^2 - \hat{m}_\phi^2}{q^2 - \hat{m}_\phi^2 - \Pi^\text{vac}(q^2)} \right|^2.$$

(4)

Here we have introduced the bare mass $\hat{m}_\phi$ of the $\phi$ meson, $g_\phi = -3g/\sqrt{2} \simeq -14$ is its strong coupling constant, and $a_\phi$ a constant which describes deviations from universality of the $\phi \to e^+e^-$ and $\phi \to K\overline{K}$ couplings. This constant is close to unity, i.e. deviations from universality are small. For our present purpose extreme fine-tuning is not necessary and we can set $a_\phi = 1$.

The self-energy $\Pi^\text{vac}_\phi$ of the $\phi$ meson in the vacuum consists of three parts:

$$\Pi^\text{vac}_\phi = \Pi^\text{vac}_{\phi \to K^+K^-} + \Pi^\text{vac}_{\phi \to KK_0^0} + \Pi^\text{vac}_{\phi \to 3\pi},$$

(5)

describing the coupling of the $\phi$ to the $K\overline{K}$ and three-pion channels. The last term violates the OZI rule, but despite the small $\omega\phi$ mixing angle it contributes about 15 percent to the total $\phi$ meson decay width. We include its imaginary part as given in ref.[4] but focus here on the more important parts of the self-energy coming from the decay into $K\overline{K}$ channels. They are related by SU(3) to the $\rho - \pi\pi$ self energy and can be written in the form of a one-loop integral [4][3]:

$$\Pi^\text{vac}_{\phi \to K^+K^-}(q^2) = -\frac{ig^2}{6} \int \frac{d^4 l}{(2\pi)^4} \left[ \frac{(2l - q)^2}{(l^2 - m_K^2 + i\epsilon)(l^2 - m_K^2 + i\epsilon)^2} - \frac{8}{l^2 - m_K^2 + i\epsilon} \right].$$

(6)

The first term of the integrand involves a propagating $K^+K^-$ pair; the second (tadpole) term ensures gauge invariance at the level of the hadronic effective theory. Here we have
introduced the strong meson coupling $g = 6.5$ and the charged kaon mass $m_K = 493$ MeV. Evaluating this integral and applying regularization using a subtracted dispersion relation \[4\] we get

\[
\text{Re } \Pi_{\phi \to K^+ K^-}^{\text{vac}}(q^2) = c_0 q^2 - \frac{g^2}{48\pi^2} \left[ q^2 G(q^2, m_K^2) - 4m_K^2 \right],
\]

\[
\text{Im } \Pi_{\phi \to K^+ K^-}^{\text{vac}}(q^2) = -\frac{g^2}{96\pi} q^2 \left( 1 - \frac{4m_K^2}{q^2} \right)^\frac{3}{2} \Theta(q^2 - 4m_K^2),
\]

where the subtraction constant $c_0 = 0.11$ has been fixed to give a best fit to data as explained in ref. [4]. This leads to a bare mass $\hat{m}_\phi = 910$ MeV in eq.(4). In eq.(7) we insert [4]

\[
G(q^2, m^2) = \begin{cases} \left( \frac{4m^2}{q^2} - 1 \right)^\frac{3}{2} \arcsin \frac{\sqrt{q^2}}{2m} & : 0 < q^2 < 4m^2 \\ -\frac{1}{2} \left( 1 - \frac{4m^2}{q^2} \right)^\frac{3}{2} \ln \left[ \frac{1 + \sqrt{1 - \frac{4m^2}{q^2}}}{1 - \sqrt{1 - \frac{4m^2}{q^2}}} \right] & : 4m^2 < q^2 \text{ or } q^2 < 0. \end{cases}
\]

For $\Pi_{\phi \to K^0 K^0}^{\text{vac}}$ the same expressions as in eqs.[4][8] hold with the charged kaon mass $m_K$ replaced by $m_{K^0}$. Using these self-energies as input in eq.(4) we plot the spectrum of the vacuum correlation function $\Pi(q^2)$ in Fig.4a (dashed line). We also show the real part of the $\phi$ meson propagator $D_\phi(q^2) = [q^2 - \hat{m}_\phi^2 - \Pi_{\phi}^{\text{vac}}(q^2)]^{-1}$ in fig. 4b (dashed line). The zero of Re $D_\phi$ determines the physical mass of the free $\phi$ meson.

**The $\phi$ meson in medium.** We choose a Lorentz frame with nuclear matter at rest. In the following we consider the case with the $\phi$ meson at rest ($q = (\omega, \vec{q} = 0)$), so as to determine the in-medium mass of the $\phi$-meson. The tensor structure of the correlation function then reduces to a term proportional to the spacelike Kronecker symbol $\delta_{ij}$. All time components vanish, and one can single out a scalar function by taking the trace $\Pi = \frac{1}{3} \Pi_i$. The spectral function has a form analogous to that in the vacuum [3, 4]. One only needs to replace $q^2$ by $\omega^2$ and the vacuum self-energy by the in-medium self energy $\Pi_{\phi}(\omega^2, \rho)$ of the $\phi$ meson, with

\[
\text{Im} \Pi(\omega^2, \rho) = \frac{\text{Im} \Pi_{\phi}(\omega^2, \rho)}{g_\phi^2} \left| \frac{(1 - a_\phi) \omega^2 - \hat{m}_\phi^2}{\omega^2 - \hat{m}_\phi^2 - \Pi_{\phi}(\omega^2, \rho)} \right|^2, \quad (10)
\]

where we use $a_\phi = 1$ again as a good approximation. The difference between the vacuum and the in-medium self-energy defines the density dependent effective $\phi$-nucleon amplitude $T_{\phi N}$, as follows:

\[
\rho T_{\phi N}(\omega, \rho) = \Pi_{\phi}^{\text{vac}}(\omega^2) - \Pi_{\phi}(\omega^2, \rho).
\]

In ref.[7] we have shown that to leading order in density $\rho$ this quantity reduces to the free forward $\phi$-nucleon scattering amplitude $T_{\phi N}(\omega)$ (at $\vec{q} = 0$) and we write in this approximation:

\[
\Pi_{\phi}(\omega, \vec{q} = 0; \rho) = \Pi_{\phi}^{\text{vac}}(\omega^2) - \rho T_{\phi N}(\omega) + \ldots, \quad (12)
\]

where the dots represent terms of higher order in density.
The primary modification of the self-energy $\Pi_\phi$ comes from the interactions of the intermediate $K$ and $\bar{K}$ mesons with nucleons in the nuclear medium. The kaon propagators in eq. (13) are then to be replaced by the in-medium propagators,

$$\frac{1}{l^2 - m^2_{K^\pm} + i\epsilon} \to D_{K^\pm}(l_0, \vec{l}; \rho) = \frac{1}{l_0^2 - \vec{l}^2 - m^2_{K^\pm} - \Sigma_{K^\pm}(l_0, \vec{l}; \rho)},$$

where $\Sigma_{K^\pm}$ are the $K^+$ and $K^-$ self-energies in nuclear matter (or, correspondingly, those of $K_0$ and $\bar{K}_0$ where applicable). The kaon propagators have the following spectral representations at fixed kaon three-momentum $\vec{l}$

$$D_K(l_0, \vec{l}; \rho) = \int du^2 \frac{A_K(u, \vec{l}; \rho)}{l_0^2 - \vec{l}^2 - u^2 + i\epsilon}$$

with

$$A_K(u, \vec{l}; \rho) = \frac{-\Im \Sigma_K(u, \vec{l}; \rho)/\pi}{[u^2 - m^2_K - \Re \Sigma_K(u, \vec{l}; \rho)] + [\Im \Sigma_K(u, \vec{l}; \rho)]^2}. \quad (15)$$

Once these spectral functions are determined at a given density for both $K^+$, $K^0$ and $K^−, \bar{K}^0$, the $\phi$ meson self-energy is readily calculated. For example, the $\phi \to K^+K^−$ in-medium self-energy for a $\phi$ at rest becomes

$$\Pi_{\phi \to K^+K^−}(\omega, \vec{q} = 0; \rho) = -\frac{ig^2}{6} \int \frac{d^4l}{(2\pi)^4} \int du^2 \int du^2 \frac{A_{K^+}(u_+, \vec{l}; \rho) A_{K^−}(u_−, \vec{l}; \rho)}{(l_0^2 - \vec{l}^2 - u^2_+ + i\epsilon) ((l_0 - \omega)^2 - \vec{l}^2 - u^2_- + i\epsilon)} \left[(2l_0 - \omega)^2 - 4\vec{l}^2\right] + \text{tadpole}, \quad (16)$$

where the tadpole part, not written explicitly, contributes only to the real part of $\Pi_\phi$. Its imaginary part is now easily evaluated; for example:

$$\Im \Pi_{\phi \to K^+K^−}(\omega, \vec{q} = 0; \rho) = -\frac{g^2}{96\pi} \int du^2 \int du^2 A_{K^+}(u_+, \vec{k}; \rho) A_{K^−}(u_−, \vec{k}; \rho) \frac{\lambda^\phi(\omega^2, u^2_+, u^2_-)}{\omega^4} \quad (17)$$

with $|\vec{k}| = \lambda^\phi(\omega^2, u^2_+, u^2_-)$, where $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2ac - 2bc$ is the Källen function.

Let us now first discuss contributions to $\Im \Pi_\phi$ from $s$-wave $KN$ and $\bar{K}N$ interactions for which the leading process is illustrated in Fig.1a. The corresponding in-medium kaon spectral functions are determined using the coupled channels approach based on the chiral SU(3) effective Lagrangian as described in refs. [8, 9]. This approach successfully reproduces all available low-energy data of the $KN$ as well as the coupled $\bar{K}N$, $\pi\Sigma$, $\pi\Lambda$ and $\eta N$, $K\Lambda$, $K\Sigma$ systems. (For alternative approaches see refs. [11, 12, 13].)

For the $K^+$ and $K^0$ modes in matter one finds that the spectral functions can be well approximated by a $\delta$-function with the free kaon mass replaced by $m^*_K(\rho) = m^*_K(\rho)$:

$$A_{K^+,K^0}(u, \vec{l}; \rho) = -\delta(u^2 - m^2_{K^+}(\rho)). \quad (18)$$
At \( \rho = \rho_0 = 0.17 \text{ fm}^{-3} \) we find \( m_{K^+}^* \approx 535 \text{ MeV} \), nearly independent of \( \vec{l} \).

For the \( K^- \) and \( \overline{K}^0 \) modes the spectrum generated by \( s \)-wave \( KN \) interactions has more interesting features. We give examples in Fig.2 at \( \rho = \rho_0 \) for two energies \( \omega \) of the external \( \phi \) meson. The \( \overline{K}N \) system decays into \( \pi \Lambda \) and \( \pi \Sigma \). The spectral functions \( A_{K^-,\overline{K}^0} \) therefore have finite widths. The second interesting point is the two-mode structure that appears at finite \( \vec{k} \) three-momentum \(|\vec{k}|\). The low-energy \( \overline{K}N \) \( s \)-wave interaction is governed by the \( \Lambda(1405) \) resonance. The response of nuclear matter to \( S = -1 \) kaonic excitations therefore involves \( \Lambda(1405) \) particle-nucleon hole states. Whereas the \( \Lambda(1405) \) dissolves in matter for a \( K^- \) at rest \((\vec{k} = 0)\) due to the action of the Pauli principle \([9]\), it reappears at finite \( \vec{k} \). This is seen in the two-peak structure of the spectral function, Fig.2, under appropriate kinematical conditions. The lower mode can be identified with a \( \overline{K} \) (with its mass shifted downward) while the upper mode corresponds to a \( \Lambda(1405) \)-hole excitation. When implemented in the \( \phi \) meson self-energy, the upper \((\Lambda(1405))\) mode starts to become important at energies \( \omega \geq 1.1 \text{ GeV} \), above the free \( \phi \) resonance. At these energies the kaons move with relatively large momenta, and their scattering from nucleons in nuclear matter is not much influenced by Pauli blocking effects.

Next we incorporate \( p \)-wave \( KN \) and \( \overline{K}N \) interactions. They are generated by the axial vector coupling terms in the chiral \( SU(3) \) effective meson-baryon Lagrangian. The basic in-medium processes involve \( \Lambda \)- and \( \Sigma \)-hole excitations, and we also include excitations of the \( \text{spin-3/2 decuplet} \) \((\Sigma^*)\). The octet and decuplet couplings are connected by standard \( SU(3) \) relations. The \( p \)-wave \( KN \) and \( \overline{K}N \) contributions to the \( \phi N \) amplitude in leading order as sketched in Fig.1b and c. The additional terms, Fig.1d and e, involve \( \phi N \rightarrow KY \) contact interactions representing short-range processes such as \( K^* \) exchange. For baryon octet \((\Lambda \text{ and } \Sigma)\) intermediate states one finds in the heavy-baryon limit:

\[
\text{Im} T_{\phi N}^{p\text{-wave}}(\omega, \vec{q} = 0) = \frac{3}{4}(D-F)^2 \mathcal{H}(\omega, M_\Sigma-M_N, m_K) + \frac{1}{12}(D+3F)^2 \mathcal{H}(\omega, M_\Lambda-M_N, m_K),
\]

(19)

where \( D = 0.75 \) and \( F = 0.51 \) are the \( SU(3) \) axial coupling constants with \( g_A = D + F = 1.26 \). The function \( \mathcal{H} \) (for \( \omega > 0 \)) is given by:

\[
\mathcal{H}(\omega, \Delta, m) = \frac{g^2}{24\pi f^2_\pi} \left[ \frac{\sqrt{(\omega-\Delta)^2-m^2}}{(\omega^2-2\omega\Delta)^2} \left( 3\omega^4 - 4\omega^2m^2 + 4m^4 + \Delta(8\omega m^2 - 12\omega^3) \right) \right.
\]

\[
\left. + \Delta^2(16\omega^2 - 8m^2) - 8\Delta^3\omega + 4\Delta^4 \right] - \frac{\Delta(\omega^2 - 4m^2)^2}{2\omega(\omega^2 - 4\Delta^2)^2} \left(3\omega^2 - 8m^2 - 4\Delta^2\right). \tag{20}
\]

We use \( M_N = 938 \text{ MeV} \), \( M_\Lambda = 1.116 \text{ GeV} \) and \( M_\Sigma = 1.191 \text{ GeV} \). The corresponding amplitude with a \( \Sigma^* \) intermediate state has its coupling scaled by the proper Clebsch-Gordan and \( SU(3) \) coefficients:

\[
\text{Im} T_{\phi N}^p(\omega) = \frac{3}{8}(D+F)^2 \mathcal{H}(\omega, M_\Sigma^*-M_N, m_K). \tag{21}
\]

with \( M_\Sigma^* = 1.383 \text{ GeV} \). In the actual calculations we have included form factors at the axial kaon-baryon vertices. For simplicity we have used the empirical nucleon axial form
factor $G_A(t) = G_A(0)[1 - q^2/\Lambda_A^2]^{-2}$ with $\Lambda_A = 1.05$ GeV. Note that the largest p-wave contribution comes from the intermediate $\Lambda(1116)$.

Summing all s- and p-wave contributions, we evaluate the imaginary part of the effective, density dependent $\phi N$ amplitude for a $\phi$ meson "at rest":

$$\rho \mathrm{Im} T_{\phi N}^{K^+K^-}(\omega, \vec{q} = 0; \rho) = \mathrm{Im} \Pi_{\phi N}^{\text{vac}}(\omega, \vec{q} = 0) - \mathrm{Im} \Pi_{\phi N}(\omega, \vec{q} = 0; \rho).$$

The real part of $T_{\phi N}$ is then determined by a subtracted dispersion relation at $\vec{q} = 0$:

$$\mathrm{Re} T_{\phi N}(\omega; \rho) = c_1 + \frac{\omega^2}{\pi} \mathcal{P} \int_0^\infty du \frac{\mathrm{Im} T_{\phi N}(u; \rho)}{u^2(u^2 - \omega^2)},$$

where the subtraction constant $c_1 = 0$ is fixed by the Thomson limit of the Compton amplitude involving the strange quark current (2).

Our results for the real and imaginary parts of $T_{\phi N}$ at $\rho = \rho_0$ are shown in Fig.3. Note that the lower plateau in $\mathrm{Im} T_{\phi N}$ comes mainly from the $\phi N \rightarrow K\Lambda$ channel, with an intermediate $\overline{K}N \rightarrow \Lambda$ p-wave coupling. The steep rise of $\mathrm{Im} T_{\phi N}$ just above 0.9 GeV has two major contributions. About half of it comes from $\phi N \rightarrow K\overline{K}$ with s-wave interactions of the $K$ and $\overline{K}$ in the nuclear medium. Roughly the other half originates from the $\phi N \rightarrow K\Sigma^*$ process. The complexity of the low energy $\phi N$ dynamics translates visibly into a highly structured pattern for $\mathrm{Re} T_{\phi N}$ below and around the $\phi$ resonance.

We mention that s-wave interactions of the $\overline{K}$ in the medium must be treated to all orders in the density, even at small $\rho$. In contrast, the iteration of p-wave interactions to higher orders in $\rho$ has a significant effect only at energies above the $\phi$ resonance, at least for densities $\rho \lesssim \rho_0$. We estimate the uncertainties in the high energy part of $\mathrm{Im} T_{\phi N}$ (at $\omega > m_\phi$) to be at the 20-30% level. However, this uncertainty has very little influence on the in-medium spectral function of the $\phi$ meson shown in Fig. 4a. In any case, the resulting in-medium mass shift of the $\phi$ stays within 1% of its free mass.

The predicted spectrum of the strange current-current correlation function, eq.(10), is shown at $\rho = \rho_0$ in Fig.4 together with the real part of the in-medium $\phi$ meson propagator,

$$D_\phi(\omega, \vec{q} = 0; \rho) = \frac{1}{\omega^2 - \overline{m}_\phi^2 - \Pi_\phi(\omega, \vec{q} = 0; \rho)}.$$  

One observes a very small and insignificant shift of the resonance position (by about 1% of its free mass at $\rho = \rho_0$). The primary in-medium effect is the broadening of the resonance due to inelastic $\phi N$ reactions. Its width at $\omega \approx m_\phi$,

$$\Gamma_\phi = -\frac{\mathrm{Im} \Pi_\phi}{m_\phi},$$

reaches $\Gamma_\phi \approx 45$ MeV at $\rho = \rho_0$ and exceeds the free (vacuum) width of 4 MeV by about an order of magnitude. The $\phi$ meson life time is reduced to

$$\tau_\phi = \frac{1}{\Gamma_\phi} \approx 4.4 \text{ fm/c} \text{ at } \rho = \rho_0.$$  

Consequently, $\phi$ mesons implanted with small velocity in a nucleus have a good chance to decay within nuclear dimensions.
Summary and conclusions. We have evaluated the in-medium spectrum of the $\phi$ meson taking the important s- and p-wave interactions of its $K\overline{K}$ components with surrounding nuclear matter fully into account. At normal nuclear matter density, $\rho = \rho_0$, we predict almost no mass shift of the $\phi$ meson, while its width is expected to increase by about an order of magnitude over its free width. The resulting short lifetime would make it possible to observe in-medium modifications of a slowly moving $\phi$ within the diameter of a typical medium-heavy nucleus. Such considerations are of interest for upcoming experiments using HADES at GSI. Finally we mention that our calculated $\phi$ meson spectrum is consistent with the in-medium QCD sum rule analysis of refs. [3, 7].

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Figure 1: Diagrams contributing to the leading s- and p-wave interactions of kaons in nuclear matter. The hyperon intermediate states $Y$ include those of the baryon octet and decuplet.

Figure 2: Spectral function $A_{K^-}$ of $K^-$ modes in nuclear matter at density $\rho = \rho_0 = 0.17 \text{ fm}^{-3}$ for two energies ($\omega = 1.0 \text{ GeV}$ and $\omega = 1.2 \text{ GeV}$) of the primary $\phi$ meson (taken at rest), as a function of the squared invariant mass $u^2 = k^2_0 - \vec{k}^2$. The calculation includes s-wave $KN$ interactions derived from chiral SU(3) dynamics with coupled channels [8, 9].
Figure 3: Real and imaginary parts of the effective $\phi N$ scattering amplitude $T_{\phi N}(\omega, \vec{q} = 0; \rho)$ in nuclear matter at density $\rho = \rho_0$. All s- and p-wave $KN$ and $\bar{K}N$ interactions are included.
Figure 4: a) Spectrum of the strange quark current-current correlation function in vacuum (dashed) and in nuclear matter at \( \rho = \rho_0 \) (solid curve). The normalization is chosen such that the vacuum result can be compared directly to the ratio \( \sigma(e^+e^- \rightarrow K^+K^-)/\sigma(e^+e^- \rightarrow \mu^+\mu^-) \) for which data are taken from ref. [10].

b) Real part of the \( \phi \) meson propagator in vacuum (dashed) and in nuclear matter at \( \rho = \rho_0 \) (solid curve).