Pseudoscalar Mesons in Nuclear Medium

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Abstract

The behavior of pseudoscalar mesons in nuclear medium is reviewed with an emphasis on the possibility of their Bose-Einstein condensation in dense matter. In particular pion condensation is reexamined in detail, stimulated by recent theoretical and observational developments.

1 Introduction

Recently much attention has been paid for high density QCD: hadron matter at relatively low density and quark matter at high density are typical subjects there. As a key concept which goes through hadron and quark worlds chiral symmetry is realized in both matter, but in a different way. Parity-even and odd quantities interplay in this context. Since the vacuum is good parity state, there is no expectation value of parity-odd operators. However, they may have finite values in matter (parity violation); Bose-Einstein condensation of pseudoscalar mesons in hadronic matter or nonvanishing of the parity-odd mean-field, \( \langle \bar{q} \gamma_5 \Gamma^a q \rangle \), in quark matter. We consider the particle-hole operator here and discuss how the pseudoscalar quantity becomes nonvanishing in hadronic matter and what are its implications, by studying the behavior of pseudoscalar mesons in nuclear medium.

Pseudoscalar mesons (\( \pi, K \)) have some salient features; they are the lightest hadrons without and with strangeness and considered as the Nambu-Goldstone bosons as results of spontaneous breaking down of chiral symmetry. Since they are bosons, they may lead to the Bose-Einstein condensation in some situations.

To understand the behavior of these mesons in nuclear medium, we also take into account the effects of resonances. \( \Delta(1232) \) strongly couples with nucleon by the \( p \)-wave \( \pi N \) interaction and \( \Lambda(1405) \) gives rise to a peculiar feature in \( s \)-wave \( KN \) scattering near threshold. The mass difference between these resonances and nucleons is small \( (O(m_\pi)) \), so that they should play important roles even for low energy phenomena (energy-momentum scale \( \sim O((2-3)m_\pi) \)) we are interested in. Thus chiral symmetry and resonances may characterize the behavior of these mesons in nuclear medium.

There have been performed and planed many nuclear experiments to reveal the behavior of these mesons in nuclear medium. On the other hand, observations of compact stars have also provided information about it. Very recently appeared an interesting data about the surface temperature of young pulsar [1]. They reported that the pulsar inside

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Figure 1: Cooling curves for stars with medium FP EOS, taken from ref.2. Solid curve shows standard cooling of a $1.2M_\odot$ neutron star and dashed curve shows pion cooling of a $1.4M_\odot$ star. The data of $c, 2, 4, 5$ indicate cooler stars: $(c)3C58$, $(2)$ Vela, $(4)$ Geminga and $(5)$ RXJ 1856-3754.

The historical supernova 3C58 shows too low surface temperature to be explained by the standard cooling scenario, which essentially assumes usual neutron matter inside it (see Fig 1). The importance of this observation is in the age ($\sim 10^3$ yr) of this pulsar. As we shall see later if this fast cooling is attributed to the presence of pion condensation, the star cools very rapidly in the early neutrino-emitting phase and the difference of the surface temperature from the standard cooling scenario becomes remarkable there. So we can hope to see the evidence of pion condensation more clearly for young pulsars. In Fig. 1 we present our theoretical cooling curve with observational data. We can see that pion cooling can explain cooler stars including 3C58 [2].

## 2 Kaon Condensation

The low-energy $KN$ interaction is specified by three kinds of the $s$-wave interaction, the $KN\sigma$ term, the energy-dependent Tomozawa-Weinberg term and the resonant term with $\Lambda(1405)$. Empirically the scattering amplitude shows an interesting behavior in the $I = 0$ channel due to the existence of the resonance $\Lambda(1405)$ below the threshold. Recently some people propose a possibility of the deeply bound kaonic nuclei in relation to the property of this resonance [3]. On the other hand, the interaction becomes relatively weak in the $I = 1$ channel, and it is relevant when we consider kaon ($K^-$) condensation in neutron stars.

Since we know the non-resonant $s$-wave terms can be well described in terms of $SU(3) \times SU(3)$ chiral symmetry, kaon condensation has been discussed on the basis of chiral Lagrangians [4, 5]. These $s$-wave terms cooperatively work to give a large decrease of the effective energy of kaons in nuclear medium. When the energy reaches the electron chemical potential as density increases, kaons begin to condense through the reaction,

$$nn \rightarrow npK^-.$$ (1)
This is very similar to the Bose-Einstein condensation of alkali atoms [4].

The most important consequence of kaon condensation is the large softening of the equation of state, which leads to an interesting phenomenon, delayed collapse of protoneutron stars to produce the low-mass black holes [6, 7, 8, 9, 10].

3 Pions in Nuclear Medium

Since pions couple with particle-hole and ∆-hole states with the same quantum number, we can study the properties of these states as well as pions themselves by considering the pion propagation in nuclear medium [11]. In Fig. 2 we show the longitudinal spin-isospin excitation (pionic) modes in the energy (ω)-momentum (k) plane.

![Figure 2: Schematic view of the spin-isospin modes in symmetric (N=Z) nuclear matter. NN^−1 and ΔN^−1 denote the continuum spectrum of the particle-hole and ∆-hole excitations, respectively. “ZS” means the zero-sound mode, which corresponds to the Gamow-Teller state in asymmetric matter.

3.1 Pionic Excitations within the Fermi Liquid Theory

Consider the pion propagator in nuclear medium at density ρ_B:

\[ D_π^{-1}(ω, k; ρ_B) = ω^2 - m_π^2 - k^2 - Π(k, ω; ρ_B), \]  

due to the self-energy Π. For the p-wave πN interaction \(^1\), it is simply given by the particle-hole polarization function \( U^{(0)}_N \),

\[ Π^{(0)}_p = -k^2 \left( U^{(0)}_N(ω, k; ρ_B) + U^{(0)}_Δ(ω, k; ρ_B) \right), \]  

in the lowest order, where we have taken into account the nucleon particle-hole (ph) and ∆-hole (Δh) states \(^2\). The polarization functions \( U^{(0)}_α \) are further given in terms of the

\(^1\)The s-wave interaction is negligible in symmetric (N = Z) nuclear matter.

\(^2\)We, hereafter, use the nonrelativistic approximation for nucleons. See ref. 12 for a relativistic treatment.
Lindhard functions $L_{\alpha}$ [13],

$$U^{(0)}_{\alpha} = \left( \frac{f_{\pi N\alpha} \Gamma(k)}{m_\pi} \right)^2 L_{\alpha},$$

(4)

where we introduced the form factor $\Gamma = (\Lambda^2 - m_{\pi}^2)/(\Lambda^2 + k^2)$ with the cut-off $\Lambda \sim O(1\text{GeV})$, and the $\pi NN$ and $\pi N\Delta$ coupling constants are $f_{\pi NN} \sim 1$ and $f_{\pi N\Delta} \sim 2$ (Chew – Low value), respectively. The Lindhard functions $L_{\alpha}$ are explicitly evaluated to be, e.g.

$$L_N = \text{Re}L_N + i\text{Im}L_N,$$

(5)

with

$$\text{Re}L_N = \frac{2m_N^*}{\pi^2} (p_F^p \phi_p(k, \omega) + p_F^n \phi_n(k, -\omega)),
\phi_i(k, \omega) = \frac{m_N^*}{2k^3 p_F^i} \left\{ -ab^i + \frac{a^2 - b^2}{2} \ln \frac{a + b}{a - b} \right\},
\begin{align*}
a &= \omega - k^2 / 2m_N^*,
b^i = kv^i_F,
\end{align*}$$

(6)

for $\pi^-$ propagation. Here we introduced the effective mass $m_N^*$ for nucleons.

![Diagram](https://example.com/diagram.png)

Figure 3: Examples of the particle-hole and $\Delta$-hole interactions in the spin-isospin channel. They are written as sum of the one-pion exchange interaction and the phenomenological zero-range interaction with the Landau-Migdal parameters.

It is well known that the lowest order calculation is not sufficient to discuss the behavior of the pion in nuclear medium: we must take into account the correlations between ph and $\Delta h$ states. They can be easily incorporated in the spirit of Landau Fermi-liquid theory. Since we are interested in the region of $\omega, k \sim (2-3)m_\pi$, full ph-ph, ph-$\Delta h$ or $\Delta h-\Delta h$ interaction should be separated into two terms, depending on their length scales: we explicitly treat the long-range ($O(m_\pi^{-1})$) interactions by way of pion, ph and $\Delta h$ propagation, while the short-range ($O(m_N^{-1} \sim 0.2\text{fm})$) interactions are replaced by the momentum-independent local interactions parametrized by the Landau-Migdal parameters; e.g.

$$\mathcal{F}_{NN} = f_{NN} + f'_{NN} \mathbf{r}_1 \cdot \mathbf{r}_2 + g_{NN} \mathbf{\sigma}_1 \cdot \mathbf{\sigma}_2 + \left( \frac{f_{\pi NN}^2}{m_\pi^2} \right) g'_{NN} \mathbf{\sigma}_1 \cdot \mathbf{r}_1 \cdot \mathbf{\sigma}_2 \cdot \mathbf{r}_2,$$

(7)
where the strength in the spin-isospin channel is measured in the pion unit. It can be extended to include the isobar degrees of freedom:

\[
\mathcal{F}_{N\Delta} = \left( \frac{f_{\pi NN} f_{\pi N\Delta}}{m_\pi^2} \right) g'_{N\Delta} \sigma_1 \cdot S_2 \tau_1 \cdot T_2
\]

\[
\mathcal{F}_{\Delta\Delta} = \left( \frac{f^2_{\pi N\Delta}}{m_\pi^2} \right) g'_{\Delta\Delta} S_1 \cdot S_2 T_1 \cdot T_2
\]

with the transition spin and isospin operators, \(S\) and \(T\).

Then we have the \(p\)-wave self-energy of the pion by considering the one-line irreducible diagrams and the Dyson equations;

\[
\Pi_p = \Pi_N + \Pi_\Delta,
\]

where

\[
\Pi_N = -k^2 U_N = -k^2 U_N^{(0)} \left[ 1 + (g'_{\Delta\Delta} - g'_{N\Delta}) U_{\Delta}^{(0)} \right] / D,
\]

\[
\Pi_\Delta = -k^2 U_\Delta = -k^2 U_\Delta^{(0)} \left[ 1 + (g'_{NN} - g'_{N\Delta}) U_N^{(0)} \right] / D
\]

with

\[
D = 1 + g'_{NN} U_N^{(0)} + g'_{\Delta\Delta} U_\Delta^{(0)} + (g'_{NN} g'_{\Delta\Delta} - g'_{N\Delta}^2) U_N^{(0)} U_\Delta^{(0)}.
\]

We can study the spin-isospin modes in another way, starting from the \(\phi\) and \(\Delta h\) propagation within RPA. Both ways are equivalent with each other for the longitudinal modes. Considering the correlation function between the generalized spin-isospin density operator,

\[
\mathcal{O} = \psi^\dagger \tau \sigma \psi + \frac{f_{\pi NN}}{f_{\pi N\Delta}} \psi^\dagger_\Delta \tau \sigma \psi,
\]

we have the same excitation spectra of the spin-isospin modes as before [14, 15]. Indeed the nuclear response function in the longitudinal spin-isospin channel is defined as follows;

\[
R(\omega, k) = \frac{\text{Im} D_\pi(\omega, k; \rho_B)}{D_0(\omega, k)} = \frac{\text{Im} \Pi_p(\omega, k; \rho_B)}{D_0(\omega, k) [\omega^2 - m_\pi^2 - k^2 - \text{Re} \Pi_p(\omega, k; \rho_B)]^2 + [\text{Im} \Pi_p(\omega, k; \rho_B)]^2}
\]

with the free pion propagator \(D_0(\omega, k) = (\omega^2 - m_\pi^2 - k^2)^{-1}\).

### 3.2 Pion Condensations

#### 3.2.1 Neutral Pion Condensation

First consider neutral pion condensation in symmetric (\(N=Z\)) nuclear matter, which require the following condition: the softening of the longitudinal spin-isospin mode

\[
R(0, k_c) \to \infty, \quad \text{Im} \Pi_p(0, k_c; \rho_c) \propto \omega \theta(\omega) \to 0,
\]

or equivalently

\[
D^{-1}(0, k_c; \rho_c) \to 0,
\]
in terms of the pion propagator. It is to be noted that pion condensation by no means implies a naive Bose-Einstein condensation of pions, but the softening of the longitudinal spin-isospin mode with the critical momentum \( k_c \). On the other hand, such mode is unstable and the Lindhard function has an imaginary part before the critical density \( \rho_c \).

We can see a peculiar enhancement of the strength function at small energy near the critical density.

The pion condensed phase can be represented in terms of chiral transformation as follows:

\[
|\pi^c\rangle = U(\theta_V(k_c \cdot r), \theta_A(k_c \cdot r))|\text{normal}\rangle = \exp \left[ -i \left( \int \theta_V \cdot V^0 d^3 x + \int \theta_A \cdot A^0 d^3 x \right) \right] |\text{normal}\rangle,
\]

with \( \langle \text{normal} | \pi | \text{normal} \rangle = 0 \). Then

\[
\langle \pi \rangle = \pi^c \neq 0.
\]

As an example,

\[
\pi^c = (0, 0, A \cos k_c z), \quad (\pi^0 \text{ condensation}).
\]

It would be worth mentioning that \( \pi^0 \) condensation gives rise to a magnetic ordering of nuclear matter: it exhibits a liquid-crystalline nature with one-dimensional antiferromagnetic ordering, called Alternating-Layer-Spin [ALS] structure (see Fig. 4) [17].

![Figure 4: Alternating-Layer-Spin [ALS] structure associated with \( \pi^0 \) condensate (\( \propto A \cos k_c z \)) in symmetric nuclear matter. Bold arrows denote proton spins and thin arrows neutron spins.](image)

Note that since \( U(\alpha)|\text{normal}\rangle = |\text{normal}\rangle, U \in SU(2) \), in the isospin space, the isospin rotated condensate, \( \tilde{\pi}^c = R \pi^c = (|\pi^c| \sin \theta \cos \phi, |\pi^c| \sin \theta \cos \phi, |\pi^c| \cos \theta), R \in O(3) \) is also possible with the same energy, which means \( \pi^\pm, \pi^0 \) condensation. Indeed all the propagators for \( \pi^0, \pi^\pm \) become identical in symmetric nuclear matter. Also note that \( g'_{NN} \) should be replaced by \( (g'_{NN} + g_{NN})/2 < g'_{NN} \) in neutron (\( Z = 0 \)) matter. Finally, since \( \omega = 0 \) in this case, a potential description is possible instead of the explicit introduction of pion field. This aspect has been emphasized in the study of the Alternating-Layer-Spin structure in the condensed phase [17].
3.2.2 Charged Pion Condensation

Next consider the charged pion condensation in neutron matter, which should have a direct relevance with neutron star phenomena. Consider, e.g., the $\pi^+$ propagator:

$$D_{\pi^+}^{-1}(\omega, k; \rho_B) = \frac{\omega^2 - m_{\pi}^2 - k^2 - \omega}{2f_{\pi}^2}\rho - \Pi_p(\omega, k; \rho_B).$$  \hfill (19)

Note that there appears the isovector s-wave coupling term $\propto (\rho_n - \rho_p)$ besides the p-wave term. Poles of $D_{\pi^+}^{-1}$ include the energies of $\pi^\pm$ mesons and the $pn^{-1}$ collective mode with the same quantum number of $\pi^+$, called $\pi^+_s$, besides single ph and $\Delta h$ excitations. The threshold condition for the charged pion condensation is

$$\omega_{\pi^+_s} + \omega_{\pi^-} = 0,$$  \hfill (20)

which implies the $\pi^+_s\pi^-$ pair condensation. In terms of the propagator we have

$$D_{\pi^+}^{-1}(\omega = \mu^\pi_s, k_c; \rho_c) = 0, \quad \frac{\partial D_{\pi^+}^{-1}}{\partial k}\bigg|_{k=k_c} = 0, \quad \frac{\partial D_{\pi^+}^{-1}}{\partial \omega}\bigg|_{\omega=\mu^\pi_s} = 0,$$  \hfill (21)

which are called the double-pole condition.

The condensed phase can be represented as

$$|\pi^c\rangle = \exp(i \int V_\pi k_c \cdot r d^3x) \exp(iQ^5_\pi \theta)|\text{normal}\rangle,$$  \hfill (22)

and we see

$$\langle\pi\rangle = \pi^c = (\sin \theta \cos k_c \cdot r, \sin \theta \sin k_c \cdot r, 0).$$  \hfill (23)

Accordingly nucleons form the quasi-particles $\eta, \zeta$ in the condensed state, which are superposition of $p, n$. Then the pion cooling works through the process;

$$\eta \rightarrow \eta' + e^- + \bar{\nu}_e$$  \hfill (24)

and

$$\eta' + e^- \rightarrow \eta + \nu_e.$$  \hfill (25)

It provide an efficient cooling mechanism for young neutron stars.

3.3 Gamow-Teller Resonance and Critical Densities of Pion Condensation

3.3.1 New Information on Landau-Migdal Parameters

The Landau-Migdal parameters play a crucial role in the study of the spin-isospin excitation modes as well as the spin-dependent structure of nuclei. Experimentally their values are not well known yet. Since experimental information was very limited, especially for $g'_{N\Delta}$ and $g'_{\Delta\Delta}$, one analyzed the spin dependent phenomena by using the universality ansatz,

$$g_{NN}' = g_{N\Delta}' = g_{\Delta\Delta}' \equiv g'.$$  \hfill (26)
e.g., based on the $SU(6)$ quark model [18]. Then it has been found that many experimental data can be reproduced by the value of $g' = 0.6 - 0.8$. Furthermore they concluded that pion condensation does not occur at several times nuclear density.

The giant Gamow-Teller (GT) states correspond to the spin-isospin dependent particle-hole states with small momentum, and thereby their coupling with pions should be small. Hence their excitation energy and strength should provide us with important information about the Landau-Migdal parameters. The strength, however, was not fully determined. At least about 50% of the GT sum-rule value was observed, but it was not clear why the rest of strength is missing [19]. Nevertheless, if one assumes that the strength is quenched owing to the coupling of the particle-hole states with the $\Delta$-hole states, the universality ansatz with $g' = 0.6 - 0.9$ can explain well both the excitation energy and strength.

Recent analysis of the giant GT resonance observed by $^{90}$Zr($p, n$) $^{90}$Nb at 295MeV suggests the quenching factor is small [20],

$$Q = \left[ 1 - \frac{g'_{N\Delta}U_{\Delta}^{(0)}}{1 + g'_{\Delta\Delta}U_{\Delta}^{(0)}} \right] = 90 \pm 5\%$$

(c.f. $Q=0.5 - 0.7$ given by previous results). Using also the excitation energy $\omega_{GT}$, we obtain [21]

$$g'_{N\Delta} = 0.18 + 0.05g'_{\Delta\Delta}$$

$$g'_{NN} \simeq 0.59.$$  \hspace{1cm} (28)

Thus two Landau-Migdal parameters have been experimentally fixed, while $g'_{\Delta\Delta}$ is still left as an unknown parameter.

### 3.3.2 Critical Densities

Using the new information about the Landau-Migdal parameters, we reexamine the possibility of pion condensations [22]. In Figs 5, 6 the critical densities of $\pi^0$ condensation are presented with those under the universality ansatz for comparison. In the case of neutron matter we use $g_{NN} = g'_{NN}$ for simplicity, since the value of $g_{NN}$ is not well-known yet. We can see that critical density moderately increases as $g'_{\Delta\Delta}$ does, while the universality ansatz gives a sharply increasing function; the critical densities result in low densities: $1 < \rho_c/\rho_0 < 2.5$ for $g'_{\Delta\Delta} < 1$.

Recently sophisticated variational calculations have been done for symmetric nuclear matter and pure neutron matter, using modern potentials [23, 24]. They also found that there are phase transitions to pion condensation at low densities, $2\rho_0$ and $1.3\rho_0$ for symmetric nuclear matter and pure neutron matter, respectively. It would be interesting to compare these values with our results.

The critical density for charged pion condensation is presented in Fig. 6. The behavior is almost the same as that for $\pi^0$ condensation and the critical density is low; $\rho_0 < \rho_c < 2\rho_0$ for $g'_{\Delta\Delta} < 1$ and $m'_N = 0.8m_N$. It would be interesting to refer the works by Tsuruta et al. [2, 25] in this context: they set charged pion condensation at $\rho_c = 2.5\rho_0$ in their calculation of neutron star cooling.
4 Summary and Concluding Remarks

We have seen three recent results, which may support possible existence of pion condensation at low densities. The new experiment on the Gamow-Teller resonance tells us the universality ansatz about the Landau-Migdal parameters by no means hold; $g'_{NN}$ should be much less than $g'_{NN}$ or $g'_{\Delta\Delta}$. The critical densities of pion condensations are $\rho_c \sim 1.5\rho_0(N = Z, m^*_N/m_N = 0.8, g'_{\Delta\Delta} = 1)$ and $\rho_c \sim 2.4\rho_0(Z = 0, m^*_N/m_N = 0.8, g'_{\Delta\Delta} = 1)$ for neutral pion condensation, while $\rho_c \sim 2.2\rho_0(Z = 0, m^*_N/m_N = 0.8, g'_{\Delta\Delta} = 1)$ for charged pion condensation.

As another theoretical work, a new calculation of nuclear matter with a modern potential has also suggested the phase transition at $\rho_c \sim 2\rho_0(N = Z)$ and $\rho_c \sim 1.3\rho_0(Z = 0)$, which corresponds to neutral pion condensation [23].

Besides these theoretical developments, the current observation about the surface temperature of a neutron star inside 3C58 suggests we need exotic cooling mechanisms beyond the standard cooling scenario. We have seen that a consistent calculation about pion cooling have been done by taking into account nucleon superfluidity in a proper way, and it can explain the data [2].

Unfortunately these are indirect evidences for pion condensation, and we hope for a direct evidence by heavy-ion collision experiments in near future.

Finally I would like give a comment about another theoretical aspect of pion condensation. It means a spontaneous violation of parity in nuclear matter and the condensed phase can be described as a chirally rotated state. We may also consider its analog in quark matter: nonvanishing of the parity-odd mean-filed $\langle \bar{q}\gamma_5\Gamma^a q \rangle$. It would be interesting in this context to refer to recent studies about ferromagnetism in quark matter, where a magnetic ordering is realized under the axial-vector mean-field [26].
Figure 6: Critical density for charged pion condensation in pure neutron matter. $m_N^* = 0.8m_N$ and $Q = 0.9$ are used here. The symbol “U” means that by the universality ansatz.

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