Mesons in Nuclei and Partial Restoration of Chiral Symmetry

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Recent topics on mesons in nuclei are discussed by especially emphasizing the role of the partial restoration of chiral symmetry in the nuclear medium. The spontaneously broken chiral symmetry in vacuum is considered to be incompletely restored in finite nuclear density systems with moderate reduction of the magnitude of the quark condensate. On the partial restoration of chiral symmetry, the wave function renormalization is important to be taken into account for the Nambu-Goldstone bosons. We also discuss the possible change of the meson properties in the nuclear medium and meson-nucleus systems for the $\bar{K}$, $\eta$, $K^+$ and $\eta'$ mesons.

KEYWORDS: chiral symmetry, partial restoration, nuclear medium, hadrons in nuclei, $\eta$ meson, $\bar{K}$ meson, $K$ meson, $\eta'$ meson

1. Introduction

The study of the hadron properties in nucleus attracts us in many different aspects. Nuclear physics investigates many-body systems governed by the strong force and aims at understanding the nature of the strong interaction. With this knowledge, hopefully, new bound systems of the strong interaction are desired to be discovered and predicted. Such exotic states themselves have attracted our continuous attentions. Hadrons in a nucleus change their properties due to the strong interaction with nucleons. These in-medium modifications can be described by hadronic many-body effective theories, since hadrons are fundamental degrees of the system in low energies thanks to the color confinement. Nevertheless, in hadron physics it is extremely important to have its more fundamental interpretation in terms of QCD. Its key concepts are spontaneous breaking of chiral symmetry and its partial restoration in the nuclear medium. The study of hadrons in nuclei also provides basic informations of high density physics and gives important constraints accessible in our reaction experiments.

Chiral symmetry is one of the fundamental symmetries in the underlying theory of the strong interaction, QCD, and is considered to be dynamically broken by the physical states. The dynamical chiral symmetry breaking determines vacuum properties and describes low-energy hadron dynamics. The dynamical breaking of chiral symmetry is a phase transition phenomenon and it can be restored at extreme conditions, such as high density and high temperature. The broken chiral symmetry is also considered to be partially, or incompletely, restored in nuclear medium. Accordingly hadrons change their properties as demended by the partial restoration. Actually experimental data of pionic atoms [1] and pion-nucleus elastic scatterings [2] together with theoretical consideration [3, 4] have suggested that chiral symmetry is partially restored in nuclear matter with 30% reduction of the magnitude of the quark condensate. This is consistent with the prior theoretical studies [5–9].

When chiral symmetry is partially restored, one expects the following phenomena in hadron properties: First of all, the reduction of the mass difference between parity parters, such as $\pi-\sigma$, $\rho-a_1$ and $N-N^*$, are expected, because these parity partners should degenerate when chiral symmetry is completely restored in the chiral limit. For example, chiral symmetry for the nucleon and its parity...
partners is discussed in Ref. [10]. The η mesonic nuclei, which are systems of a η meson and a nucleus governed by the strong interaction, are good to see the possible reduction of the mass difference among the nucleon chiral partners, since the η-nucleon systems strongly couples to N(1535), and N(1535) is a candidate of the parity partner of nucleon. Second of all, one expects a large wave function renormalization of the Nambu-Goldstone boson. Since, according to the low energy theorem of chiral symmetry, the amplitudes of the Nambu-Goldstone bosons are written in terms of energy expansion in low-energy, the in-medium self-energy of the Nambu-Goldstone bosons has strong energy dependence. This provides a large wave function renormalization. This can be seen in the $K^+A$ elastic scattering and the $\pi^0$ decay in medium. Finally, one expects the reduction of the hadron mass generated by chiral symmetry breaking. For instance, a part of nucleon mass is considered to be generated by spontaneous breaking of chiral symmetry. Such a mass should decrease when chiral symmetry is restored. Actually a part of the η' meson is also generated by chiral symmetry breaking [11]. Thus the η' mass is expected to be reduced in nuclear medium.

2. Mesons in nuclei

2.1 $\bar{K}$ meson

There are already a lot of studies on in-medium kaon [12]. It is known that the $\bar{K}N$ interaction is so attractive that the Λ(1405) resonance can be considered as a quasi-bound state of $\bar{K}N$ [13]. Thus, it is natural to expect that bound states of the $\bar{K}$ meson and a nucleus exist. One of the difficulties in observation of the bound states, even if they exist, is that $\bar{K}$ has large nuclear absorption. The one-body absorption into π and hyperon, $\bar{K}N \rightarrow \pi Y$, is substantially large, and also the nonmesonic two-body absorption, $\bar{K}NN \rightarrow YN$, is not small contribution as reported that the two-body absorption should be 30% of total absorption at saturation density [14]. This strong absorption comes through the Λ(1405) resonance. It may be hard to identify $\bar{K}$-nucleus bound states owing to the large decay width. To understand the in-medium properties of $\bar{K}$, it is extremely important to pin down the $\bar{K}N$ interaction as a fundamental interaction of the $\bar{K}$-nucleus system. For this purpose, the nature of the Λ(1405) resonance should be understood very well, since it is sitting at 30 MeV below the threshold of $\bar{K}N$. Chiral symmetry determines low-energy $\bar{K}N$ interaction, but the partial restoration of chiral symmetry may play a minor role on the in-medium $\bar{K}$ properties, because dynamics of $\bar{K}N$, or Λ(1405), is much more significant.

The model independent Tomozawa-Weinberg interaction of the chiral perturbation theory is attractive enough to form a bound state in the $\bar{K}N$ channel with isospin $I = 0$. Theoretically, the $\bar{K}N$ interaction with $I = 0$, however, is not so strong that the bound state is formed at 30 MeV below the $\bar{K}N$ threshold, which corresponds to 1405 MeV, but the theory predicts the bound state sitting at about 15 MeV below the threshold, which is 1420 MeV. To obtain the observed Λ(1405) spectrum, coupled channel analysis of $\bar{K}N$ and $\pi\Sigma$ is unavoidable, since the $\pi\Sigma$ channel is open at the Λ(1405) resonance and the $\bar{K}N-\pi\Sigma$ channel coupling is not negligible. The coupled-channel analyses based on chiral dynamics [15] reproduce the observed Λ(1405) spectrum by the two-pole structure for Λ(1405), in which there are two independent states with the same quantum number with Λ(1405) at around 1420 MeV with a small width as a bound state of $\bar{K}N$ and at around 1390 MeV with a large width as a $s$-wave resonance of $\pi\Sigma$, and the Λ(1405) spectrum is shown up as an interference of these two state [16, 17]. (This means that the nominal Λ(1405) at 1405 MeV with 50 MeV width is not an eigenstate of the Hamiltonian but the two states are.) Because the $\bar{K}N$ scattering state more strongly couples to the $\bar{K}N$ bound state rather than the $\pi\Sigma$ resonance state, the double pole scenario for Λ(1405) suggests that the peak position corresponding to Λ(1405) in the spectrum in the $\bar{K}N$ channel appears around 1420 MeV of the bound state energy not at 1405 MeV.

One of the ways to confirm this scenario is to observe Λ(1405) produced by the $\bar{K}N$ channel. Nevertheless, because Λ(1405) is located below the $\bar{K}N$ threshold, Λ(1405) cannot be produced
2.2 $K^+$ meson

It is known that the $K^+N$ interaction is repulsive and the $K^+N$ cross section is small compared to $K^-N$ [12,24]. The mean free path of the kaon in nuclear medium is estimated to be around 5 fm, which is larger than the size of medium-heavy nuclei. With these properties, kaon has been considered to be a clean hadronic probe to investigate nuclear matter. Here we would like to draw your attention to another characteristic of kaon. In contrast to baryon, the $\Lambda(1405)$ resonance is produced mainly by $\bar{K}N$, and the contribution of the $\pi\Sigma$ channel to create $\Lambda(1405)$ is small because multiple steps are necessary to pass the strangeness brought into the system by $K^-$ to the baryon. The J-PARC E31 experiment has already started and a preliminary result was reported in this conference [19]. The important thing is that, to extract the $\bar{K}N$ scattering amplitude, one has to make the production mechanism of $\Lambda(1405)$ well under control. A lot of theoretical progress were reported in this conference [20–23].

Bearing in mind that the mean free path of kaon in nucleus is large, we expect that the $K^+$-nucleus scattering could be written well by a single step $K^+N$ interaction and multiple scatterings should be strongly suppressed. It is surprising, however, that the ratio of the total cross sections for $K^+$ elastic scattering off carbon and deuterium per nucleon is larger than unity in the range of laboratory momenta 450 to 900 MeV/c [25–27]. The theoretical study in the impulse approximation tells us that the ratio rather should be suppressed owing to the nuclear shadowing effect. This observation implies that the linear density approximation for the $K^+$ in nuclear matter is broken down in spite of the small interaction between $K^+$ and nucleon. It is also known that the so-called low density $\rho$ approximation for the optical potential is also broken down [12], which is a consequence of the linear density approximation. Having extracted the $K^+$-nucleus scattering amplitude by fitting the $K^+A$ elastic scattering data with the optical potential in the $\rho$ approximation, Ref. [12] has concluded that the $K^+$-nucleus scattering amplitude is repulsively enhanced in about 15% than the $K^+N$ scattering amplitude. This is also seen in pionic atoms as “missing repulsion” that the isospin-odd $\pi^-\Lambda$-nucleus scattering length is unexpectedly more repulsive than the $\pi^-N$ scattering length. Several possible explanations for the repulsive enhancement have been proposed: For instance, this is due to nucleon-nucleon correlation, “swelling” of nucleon [28] and mass reduction of vector mesons caused by the scale change in a nuclear medium [29].

Here we would like to explain this enhancement as the wave function renormalization [30]. If the in-medium self-energy has strong energy-dependence, the wave function renormalization plays an important role for the in-medium effect as one of the next-to-leading corrections of the linear density approximation [3]. Especially, for the Nambu-Goldstone bosons, the interactions should vanish in the soft limit ($p_\mu \to 0$) and the chiral limit (Adler zero), and thus the self-energy for the Nambu-Goldstone boson should have substantial energy dependence [4,31]. According to the argument by Kolomeitsev developed for pion [3], we obtain the energy-independent optical potential for the in-medium kaon as follows. The in-medium dispersion relation for $K^+$ at rest determines the in-medium $K^+$ mass as the solution of $m^2 - m^2 - \Sigma(m^2) = 0$. The optical potential for the in-medium kaon is given by the self-energy at the kaon energy $\omega = m^*$ as $2mV_{\text{opt}}(m^*) = \Sigma(m^*)$. Now assuming $V_{\text{opt}} \ll m$, that is, $m^* \approx m$, we expand the self-energy around $\omega = m$, set $\omega = m^*$ and obtain

$$2mV_{\text{opt}}(m^*) = \Sigma(m) + (m^2 - m^2) \left. \frac{\partial \Sigma}{\partial \omega^2} \right|_{\omega=m} + \cdots \approx \left( 1 + \frac{\partial \Sigma}{\partial \omega^2} \right) \Sigma(m) = Z \Sigma(m) , \quad (1)$$

where in the third equation we have neglected the higher order of $m^2 - m^2$ and assumed that the
difference of $\Sigma(m^*)$ and $\Sigma(m)$ gives a higher order contribution of the density expansion, and in the fourth equation we have introduced the wave function renormalization

$$Z = \left( 1 - \frac{\partial \Sigma}{\partial \omega^2} \right)^{-1} \bigg|_{\omega = m}.$$  

(2)

In this way, when the energy dependence of the self-energy is strong, which is the case for the Nambu-Goldstone boson, the wave function renormalization becomes one of the essential medium effects beyond the linear density approximation. Finally the optical potential is given by

$$2mV_{\text{opt}} = Z\rho T_{K^*N}.$$  

(3)

Let us now evaluate the $K^+N$ scattering amplitude in the chiral perturbation theory. The leading order contribution is calculated by the Tomozawa-Weinberg interaction as

$$T_{K^*N}^{I=0} = 0, \quad T_{K^*N}^{I=1} = \frac{k + k'}{2f_K^2},$$  

(4)

for the isospin 0 and 1 $KN$ channels, respectively. The averaged $K^+N$ scattering amplitude is given by

$$T_{K^*N} = \frac{1}{2}(T_{K^*p} + T_{K^*n}) = \frac{1}{4}(3T_{K^*N}^{I=1} + T_{K^*N}^{I=0}),$$  

(5)

and thus around the threshold the forward $K^+N$ scattering amplitude is written as

$$T_{K^*N}^{I=0}(\omega) = \frac{3 \omega}{4f_K^2}.$$  

(6)

With this amplitude, using the linear density approximation for the kaon self-energy in symmetric nuclear matter with density $\rho$, $\Sigma(\omega) = T_{K^*N}(\omega)\rho$, we obtain

$$Z = \left( 1 - \rho \frac{\partial T_{K^*N}}{\partial \omega^2} \bigg|_{\omega = m^*} \right)^{-1} = \left( 1 - \frac{1}{2m_K^*} \frac{3}{4f_K^2} \rho \right)^{-1} \approx 1 + 0.082 \frac{\rho}{\rho_c}$$  

(7)

where we have used $m_K = 493.7$ MeV, $f_K = 110$ MeV and $\rho_c = 0.17$ fm$^{-3}$. There around 10% enhancement of the optical potential is explained by the wave function renormalization.

The in-medium correction from the wave function renormalization is more significant in the $\pi^0$ decay [32, 33]. The in-medium amplitude of the $\pi^0$ decay into $\gamma\gamma$ is written as the one-particle irreducible vertex correction and the wave function renormalization, $M_{\gamma\gamma}^* = \sqrt{Z}M_{\gamma\gamma}$ [32]. It is known that there is no vertex correction in linear density approximation [34], and thus the in-medium change of the amplitude is given solely by the wave function renormalization, $M_{\gamma\gamma}^* = \sqrt{Z}M_{\gamma\gamma}$. If one neglects the change of the phase space of the $\pi^0$ decay in the medium and uses the result of the in-medium wave function renormalization obtained in Ref. [32], one finds

$$\frac{\Gamma_{\gamma\gamma}^*}{\Gamma_{\gamma\gamma}} = Z \approx 1 + 0.4 \frac{\rho}{\rho_c},$$  

(8)

and one expects the 40% enhancement of the $\pi^0$ decay into $\gamma\gamma$ in nuclear matter.
2.3 \( \eta \) meson

The \( \eta \) mesonic nuclei, which are bound systems of \( \eta \) in nuclei, were firstly predicted by Haider and Liu [35]. The hadron-nucleus bound systems are good “laboratories” to investigate hadrons in nuclei. The \( \eta \) meson is also one of the Nambu-Goldstone bosons associated with the spontaneous breaking of the SU(3) chiral symmetry. The Tomozawa-Weinberg interaction of the \( \eta N \) channel is, however, null, and thus the contribution from the leading-order term of the chiral perturbation theory vanishes. Instead, the \( \eta N \) system has a strong coupling to the \( N^\ast(1535) \) resonance. The excitation energy of \( \eta N \to N^\ast(1535) \) is so small that it is just about 50 MeV when we evaluate it in vacuum, in contrast to the \( \Delta \) resonance, which is located 150 MeV above \( \pi N \). Thus, thanks to the small excitation energy, for the in-medium \( \eta \) meson, the channel coupling to the \( N^\ast(1535)-N\)-hole \((N^\ast-h)\) mode should be unavoidably taken into account. In the context of chiral symmetry, the \( N^\ast(1535) \) resonance can be considered a chiral partner of the nucleon [10, 36, 37]. This implies that we can learn also the in-medium properties of \( N^\ast(1535) \) and investigate the chiral symmetry for the nucleons by studying the \( \eta \) meson in the nuclear medium where chiral symmetry is partially restored [38].

If the \( N^\ast(1535) \) resonance is the chiral partner of nucleon, the mass difference of \( N \) and \( N^\ast \) should be reduced as chiral symmetry is partially restored in the nuclear matter. The reduction is estimated as about 150 MeV at the saturation density in a chiral doublet model [37]. This implies that the energy of the \( N^\ast-h \) mode gets lower in nuclear matter at a certain density. The level crossing of the \( \eta \) and \( N^\ast-h \) modes takes place [39] as shown as the thin solid and dashed lines in Fig. 1. These two modes couple each other in the nuclear medium. The mode energies are calculated as a pole position of the in-medium \( \eta \) propagator \( G_\eta(\omega, k; \rho) = i/(\omega^2 - k^2 - m^2_\eta - \Sigma_\eta(\rho)) \) with the in-medium self-energy \( \Sigma_\eta(\rho) \). Here we take a \( N^\ast(1535) \) dominance model for the self-energy \( \Sigma_\eta \) as

\[
\Sigma_\eta(\omega, \rho) = g^2_\eta \frac{\rho}{\omega + m^*_N(\rho) - m^*_N^*(\rho) + i\Gamma_{N^\ast}(\omega, \rho)/2}
\]  

(9)

where \( g_\eta \) is the coupling strength of the \( s\)-wave \( \eta NN^\ast(1535) \) coupling, \( m^*_N \) and \( m^*_N^* \) denote the in-medium masses of nucleon and \( N^\ast(1535) \), respectively, and the \( \Gamma_{N^\ast} \) is the \( N^\ast(1535) \) decay width, which depends on the energy and density. The real parts of the pole positions are plotted as the thick solid lines in the left panel of Fig. 1. These two modes are labeled 1 and 2 in Fig. 1. With the mode coupling, one can see level repulsion in which one of the modes is pushed above repulsively in higher densities. It is also interesting that, as a consequence of the level crossing, the strength of these modes changes drastically as the density increases. In the right panel of Fig. 1, we plot the residua of the
propagator $G_\eta$ at the pole positions, which are calculated by

$$Z = \left(1 - \frac{\partial \Sigma_\eta}{\partial \omega^2}\right)^{-1}.$$  \hspace{1cm} (10)

The residue of the propagator corresponds to the renormalization of the wave function of each mode. In vacuum, the wave function for the $\eta$ meson, or the residue of the $\eta$ mode, is normalized as unity, while the $N^*-h$ mode is absence in vacuum and its wave function is zero. Owing to the mode coupling, the wave functions of the two modes mix each other and the wave functions are renormalized. The mode originating from the $\eta$ meson in vacuum, labeled 2 in Fig. 1, has a larger residue in lower densities and the residue is reduced as the density increases, while the other mode labeled 1 in Fig. 1, which originates from the $N^*-h$ mode at $\rho = 0$, has a smaller residue in lower densities and the residue gets enhanced in higher densities. Thus, for the in-medium $\eta$ meson, which is an admixture of the $\eta$ and $N^*-h$ modes, looks attractive in lower densities and turns repulsive in higher densities. The spectral function $S_\eta(\omega, \rho) = -\text{Im}G_\eta(\omega, \rho)$ with $k = 0$ is plotted in Fig. 2. The level crossing is expected to be observed in nuclear reactions, such as $(\gamma, p)$, $(d, ^3\text{He})$ and $(\pi, N)$ with nuclear targets [39–43].

As another picture of $N^*(1535)$, it could be a dynamically generated resonance of the octet mesons and baryons. In such a case, it is reported that the $N^*(1535)$ resonance has large components of $K\Lambda$ and $K\Sigma$ [44] and is insensitive to the medium effects since the $\Lambda$ and $\Sigma$ are free from the Pauli blocking by nucleons, which is one of the main medium effects [45]. In this case the level crossing does not take place. These two pictures are distinguishable when one observes the missing mass spectroscopy of the formation reactions of the $\eta$ mesonic nuclei [39, 43]. One of the main differences can be seen in the formation spectrum around the quasi-free $\eta$ energy region. When the level crossing takes place, the $\eta$ mode feels repulsive interaction since the $N^*-h$ mode comes down below the $\eta$ mode. Accordingly some enhancement appears above the $\eta$ creation threshold. When the level crossing does not take place, such enhancement does not show up.

2.4 $\eta'$ meson

The $\eta'$ meson would be one of the Nambu-Goldstone bosons associated with spontaneous breaking of the three flavor chiral symmetry in classical theory of chromodynamics. In fact, quantum chromodynamics has no axial U(1) symmetry at the beginning and it is explicitly broken by quantum effect. Thus, the $\eta'$ meson does not have to be a Nambu-Goldstone boson of the chiral symmetry breaking and is not massless even in the chiral limit. The observed $\eta'$ meson mass is 958 MeV/c$^2$ as large as the proton mass. Actually the $\eta'$ mass has a strong connection also to the breaking of the SU(3) chiral symmetry [11, 46]. Since the $\eta'$ meson is a pseudoscalar boson and the $\eta'$ field is written in terms of both the left and right handed quark fields, it cannot couple to nonchiral gluonic field, which causes the axial U(1) anomaly, without breaking chiral symmetry in any sense, explicitly and/or spontaneously. Therefore, to generate the $\eta'$ mass, the axial U(1) anomaly is not the only
source, but also the SU(3) chiral symmetry is necessarily broken, as schematically shown in Fig. 3. This means that the mass gap of the $\eta$ and $\eta'$ is generated by the chiral symmetry breaking through the axial U(1) anomaly. Therefore, when the broken chiral symmetry is restored in the case of the chiral limit, the $\eta'$ meson should get degenerate with the octet pseudoscalar mesons as shown in Fig. 4.

This scenario suggests that the mass gap of $\eta$ and $\eta'$ gets reduced in the nuclear medium where the spontaneously broken chiral symmetry is to be partially restored. Assuming that the $\eta$ meson mass is insensitive to the partial restoration of chiral symmetry because the $\eta$ mass is generated by the explicit chiral symmetry breaking through the quark mass, we expect that the $\eta'$ mass gets reduced in the nuclear medium. According to a simple estimation done in Ref. [47], 100 MeV reduction of the $\eta'$ mass is expected at the nuclear saturation density, where one expects 35% reduction of the magnitude of the quark condensate. The linear sigma model analysis also suggests 80 MeV reduction [47], and the NJL models predict 150 MeV reduction [48, 49]. With such enough attraction for the $\eta'$ meson in the nuclear medium, one expects $\eta'$-nucleus bound systems. The first calculation of the formation spectra was done in Ref. [51] for a ($\gamma$, $p$) reaction with $^{12}$C targets. Later ($p$, $d$) reactions have been studied [50] and recently, the feasibility study of observing $\eta'$ mesonic nuclei has been done in Ref. [52].

In the view of the linear sigma model [47], (at least a part of) the nucleon mass is generated by the spontaneous breaking of chiral symmetry where the chiral field $\sigma$ gets condensed in vacuum, and the nucleon is expressed by $m_N = -g(\sigma_0)$. This implies the presence of the strong $\sigma NN$ coupling.
This is the origin of the scalar attraction in the $NN$ interaction. In the same way, the chiral symmetry breaking generates a part of the $\eta'$ mass with the help of the axial anomaly. This implies again that there exists a strong coupling of $\sigma\eta'\eta'$. This leads to a strong attraction in the $\eta'N$ interaction in the scalar channel by exchanging the $\sigma$ meson. It is also known that the Tomozawa-Weinberg interaction, which is the leading term of the chiral perturbation theory, is zero for the $\eta'N$ channel. Thus, there may be no (or very small) repulsion in the vector channel, which plays an important role for the $NN$ interaction. With this knowledge, a two body bound state of $\eta'N$ is evaluated in the linear sigma model using the same machinery in which $\Lambda(1405)$ is calculated as a quasi-bound state of $KN$, and it is found that two-body $\eta'N$ bound state is obtained with 6 MeV binding energy [47] and that with the coupled channel effect of $\eta'N$ and $\eta N$ the bound state is found with 12 MeV binding energy and 6 MeV width [53].

3. Summary

We have discussed the pseudoscalar mesons in the nuclear medium under the situation that chiral symmetry is partially (30%) restored. Such in-medium changes of the meson properties are observed in meson and nucleus systems. Our expectations of the physical consequences when the partial restoration of chiral symmetry takes place in the nuclear medium are as follows:

**Reduction of the mass difference among the chiral partners:** The chiral symmetry breaking resolves parity degeneracy in the hadron mass spectrum and thus, it is responsible for the mass splitting of the chiral partners. When the broken chiral symmetry is restored in the nuclear medium, the mass gap should get smaller than that in vacuum. The $N^*(1535)$ nucleon resonance can be regarded as the chiral partner of nucleon, and the mass difference of $N^*(1535)$ and $N$ gets reduced in this case. The reduction of the mass gap could be observed in the formation spectrum of the $\eta$ mesonic nuclei, since the $\eta$ meson couples to the $N^*\text{h}$ mode in the nuclear medium.

**Substantial effect from the wave function renormalization of the Nambu-Goldstone bosons:** The low energy theorem of chiral symmetry tells us that the interactions of the Nambu-Goldstone bosons are described in terms of energy expansion. Thus, the in-medium self-energy of the Nambu-Goldstone boson necessarily has energy dependence, and the wave function renormalization can be a substantial medium effect as one of the corrections beyond the linear density approximation. We have seen that the enhancement of the $K^+$-nucleus scattering can be partially explained by the $K^+$ wave function renormalization. We have also seen that one expects a large enhancement of $\pi^0 \to \gamma\gamma$ in nuclear medium due to the large wave function renormalization for $\pi^0$.

**Reduction of hadron mass:** The masses of some hadrons are generated by the spontaneous breaking of chiral symmetry, such as nucleon. A part of the $\eta'$ mass is generated by the SU(3) chiral symmetry breaking through the axial U(1) anomaly. A 100 MeV reduction of the $\eta'$ mass at the saturation density is expected, and this provides an attractive optical potential for $\eta'$ meson and nucleus systems. Thus, $\eta'$ bound states are expected with this attractive potential. In addition, the mass generation by the spontaneous chiral symmetry breaking implies a strong coupling to the $\sigma$ field. This leads to strong attraction of the $\eta'-N$ interaction from the isoscalar-scalar $\sigma$ exchange. With this attraction, a two-body $\eta'N$ bound state can be formed with a several MeV binding energy and a few MeV width.

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