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Theory on Hadrons in Nuclear Medium

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After decades-long attempts to measure the mass shift and understand the origin of hadron mass, it became clear that one has to analyze hadrons with small vacuum width. Also, to identify the effect of chiral symmetry breaking, one has to start by looking at chiral partners. In this talk, I will review why such consideration inevitably led us to consider K^* and K_1 in nuclear matter [1]. With the kaon beam at JPARC, one could observe the mass shift of both particles in a nuclear target experiment. Once the masses and mass difference of K^* and K_1 mesons are measured, we will be closer to understanding the origin of the hadron masses and the effects of chiral symmetry breaking in them.

KEYWORDS: Hadron mass, strange vector meson, strange axial vector meson, JPARC, Kaon beam.

1. Introduction

The mass of a hadron can not be understood from the traditional picture of a composite particle, the mass of which can be expressed in terms of the masses of its constituents and a small binding energy. This is because the origin of a hadron mass is intricately related to the low energy property of QCD such as confinement and chiral symmetry breaking effects, which are analytically still not completely understood. It is known that chiral symmetry breaking is partly responsible for the masses [2–5]. Therefore, if one can identify the effect of chiral symmetry breaking to the masses of hadron, we will be able to build model and then a theory to understand the masses of the hadrons.

In the original in-medium QCD sum rules [4], the changes of the vector meson masses were dominantly due to the change of the four-quark condensate for the light quark system and the strange quark condensate for the strange quark system. Experiments have since then been performed world-wide to observe mass shift of hadrons at finite temperature or density [6]. The first hint of vector meson mass shift was first observe by the KEK-PS experiments from analyzing the invariant mass spectra of $e^+ e^-$ pairs from the nuclear targets. The experiment found excess signals at the lower end of the ω resonance peak that could be explained by the vector meson mass decrease of 9 % at normal nuclear density [7]. The possible mass shift of the ϕ meson has also been searched for in experiments. In fact, evidence of mass shift by 3.4 % at normal nuclear density was reported by the KEK-PS E325 collaboration [8]. Further measurements with higher statistics are planned at J-PARC by the E16 experiment [9].

Dilepton spectrum do not suffering from strong final state interaction with the medium when the signal emanates from inside the nucleus but is low in the yield and has many sources. Reactions involving hadronic decays have the opposite features. By analyzing the reaction $\gamma + A \rightarrow \omega + X \rightarrow \pi^0\gamma + X'$, the CBELSA/TAPS Collaboration observed that the ω meson mass decreased by about 60 MeV at an estimated average nuclear density of 0.6 ρ_0 [10]. However, the final state hadronic interaction could have contaminated the signal. Also, the QCD sum rule approach was found to have a large contribution from the scattering term [11–13], making the sum rule analysis less reliable. In the mean time, a successful set of experiments were performed by the CBELSA/TAPS collaboration to extract meson-nucleus optical potentials from near-threshold meson productions reactions. The success comes from focussing on mesons with small width in the vacuum $(K, \eta, \eta', \omega, \phi)$ [14]. Furthermore, the problem with final state interaction was found to be solvable by combining the measurement of excitation functions and the transparency ratios.

Motivated by such recent experimental progresses, the author has recently estimated the spectral shift of the f_1 meson, which is a chiral partner of ω in the limit where the disconnected diagrams are neglected, in the QCD sum rule approach [15]. Experimentally, the $f_1(1285)$ has been successfully identified by the CLAS collaboration in photoproduction from a proton target with a small width of 18±1.4 MeV [16], so that performing similar experiments on a nuclear target and comparing the result with that from the ω would be extremely useful for partial restoration of chiral symmetry in nuclei as suggested in Ref. [15]. In fact, the individual meson mass could behave differently depending on whether the hadron is in nuclear matter or at finite temperature, but the mass difference between chiral partners will only depend on the chiral order parameter and should be universal [17].

However, while $\langle \bar{q}q \rangle$ is the order parameter of QCD, it is not clear how to directly relate them to the masses of the hadrons. On the other hand, one can show that the difference between the vector and axial-vector correlation functions in the open strange channel is also an order parameter of chiral symmetry [17]. This implies that their spectral densities will become degenerate if chiral symmetry is restored. In the vacuum, the low-lying modes that couple to the vector current are $K^*(892)$ and $K^*(1410)$ while for the axial vector current they are $K_1(1270)$ and $K_1(1400)$. There is a subtlety in the nature of the two K_1 states: They are assumed to be a mixture of the ${}^3P_1$ and ${}^1P_1$ quark-antiquark pair in the quark model [18]. However, if chiral symmetry is partially restored, the spectral density will tend to become degenerate so that the lowest distinctive poles in the respective current will approach each other. Therefore, in a recent work [1], we investigated the spectral modification of open strange meson K_1 through the axial-vector current in nuclear matter using QCD sum rules.

Measuring the open strange meson in the vector channel, namely the K^{*+} through the decay $K^{*+} \rightarrow K^+ + \gamma$ was suggested as a promising signal to measure the spectral change of the vector meson in Ref. [19]. Both the $K_1(1270)$ and the $K^*(892)$ have widths smaller than their non strange counter parts, namely 90 MeV and 47 MeV, respectively, compared to more than 250 MeV and 150 MeV for the a_1 and the ρ. At the same time, they are also chiral partners so that their mass difference is sensitive to the chiral order parameter.

We note that K_1^+ and K_1^- become non-degenerate in nuclear medium due to the presence of nucleons which break charge conjugation invariance in the medium. This is not a problem as long as we compare K^* and K^1 with the same charge states as the chiral partnership exits between them.

2. Importance of condensate

Let us first discuss the importance of condensate in discussing the hadron masses. We are still far from understanding the low energy physics directly from QCD variables. Yet, important physical effects can be summarized in terms of local operators. The most important operators are chiral order parameters: that is operators that characterizes the breaking of chiral symmetry. There are many chiral order parameters. Below are some of the most commonly cited order parameters.

$$\langle \bar{q}q \rangle, \ \ \langle (\bar{q}\gamma_\mu \tau^a q)^2 - (\bar{q}i\gamma_5\gamma_\mu \tau^a q)^2 \rangle, \cdots. \tag{1}$$

The method of constructing chiral order parameters are discussed in the next section.

On the other hand, there are no local operators that are directly related to confinement. Confinement is related to the Wilson loops. In particular, the area-law behavior in pure gauge theory can be related to the confinement physics. It is well known that while the space-time Wilson loops changes from the area law behaviour to the perimeter law behaviour as one increased the temperature across the the deconfinement temperature T_c, the area-law behaviour of space-space Wilson loop does not change [20]. There were several attempts to related the area-law behaviour to gluon condensates [21,22]. In Ref. [23], the author has therefore studied the non-perturbative part of the gluon condensate across T_c and found that while the condensate values changes, about 50% of the vacuum value survives above T_c, suggesting that non-perturbative nature survives as was seen in the persistent area-law behaviour of the space-space Wilson loop. However, when the total condensate values including the perturbative contributions are extracted [24], one finds that while the electric part of the gluon condensate values changes abruptly across T_c, the magnetic part does not change much. This behavior seems quite in line with what is seen in the space-time (space-space) Wilson loop where the leading condensate in their respective operator production expansion (OPE) is the $\langle \frac{\alpha_s}{\pi} E^2 \rangle$ ($\langle \frac{\alpha_s}{\pi} B^2 \rangle$).

3. Chiral order parameter

The commonly known chiral order parameter can be rewritten in several forms.

$$\begin{aligned} \langle \bar{q}q \rangle &\equiv \lim_{x\to 0} -\langle \mathrm{Tr}[S(0,x)]] \rangle = \lim_{x\to 0} -\frac{1}{2}\langle \mathrm{Tr}[S(0,x) - i\gamma_5 S(0,x) i\gamma_5]] \rangle \\ &= -\pi \langle \rho(\lambda = 0) \rangle, \end{aligned} \tag{2}$$

where the second line shows the density of zero eigenvalue in the Euclidean formalism known as the Banks-Casher formula [25]. The gauge invariant part of the combination of propagators $(S(x,y) - i\gamma^5 S(x,y) i\gamma^5)$ can be identified as a convenient expression of a chiral order parameter. The second line, due to Banks-Casher, is useful to identify the origin or chiral symmetry breaking in terms of quark eigenvalues that can be usefully implemented in a lattice gauge theory calculations.

Using Eq. (2) one can identify many more order parameters of chiral symmetry breaking. A set of order parameters can be obtained by looking at the difference in the correlation functions of chiral partners. For example the difference between the scalar and pseudo scalar correlation function is an order parameter [26].

$$\begin{aligned} \Delta_{S-P}(q) &= \int d^4x e^{iqx} \langle \mathcal{T}\Big(\bar{q}\tau^a q(x) \bar{q}\tau^a q(0) - \bar{q}\tau^a i\gamma^5 q(x) \bar{q}\tau^a i\gamma^5 q(0) \Big) \rangle \\ &= \int d^4x e^{iqx} \frac{1}{2} \langle \mathrm{Tr}[\tau^a \tau^a \Big(i\gamma^5 S(x,0) i\gamma^5 - S(x,0) \Big)\Big(S(0,x) - i\gamma^5 S(0,x) i\gamma^5 \Big)] \rangle. \end{aligned} \tag{3}$$

As can be seen above, the combination of propagators appear, which will be proportional to the density of zero modes. Using these methods, one can construct chiral order parameters of different types of correlation functions. We have put in the flavor matrix in the scalar current. In principle, this is not necessary because the additional disconnected diagram in the scalar channel is also proportional to the chiral order parameter.

$$\begin{aligned} \Delta_S^{disconnected}(q) &= \int d^4x e^{iqx} \langle \mathcal{T}\Big(\bar{q}q(x) \bar{q}q(0) \Big) \rangle_{disconnected} \\ &= \int d^4x e^{iqx} \frac{1}{4} \langle \mathrm{Tr}[\Big(i\gamma^5 S(x,x) i\gamma^5 - S(x,x) \Big)] \mathrm{Tr}[\Big(S(0,0) - i\gamma^5 S(0,0) i\gamma^5 \Big)] \rangle. \end{aligned} \tag{4}$$

The difference between the vector and axial vector in the Kaon sector is also on order parameter. In this case, the difference is proportional to the the chiral symmetry breaking in both flavours.

$$\begin{aligned} \Delta^m_{V-A}(q) &= \int d^4x e^{iqx} \langle \mathcal{T}\Big(\bar{s}\gamma_\mu q(x) \bar{q}\gamma_\mu s(0) - \bar{s}\gamma^5 \gamma_\mu q(x) \bar{q}\gamma^5 \gamma_\mu s(0) \Big) \rangle \\ &= \int d^4x e^{iqx} \frac{1}{2} \langle \mathrm{Tr}[\gamma_\mu \Big(i\gamma^5 S_q(x,0) i\gamma^5 - S_q(x,0) \Big) \gamma_\mu \Big(S_s(0,x) - i\gamma_5 S_s(0,x) i\gamma_5 S(0,x) \Big)] \rangle, \end{aligned}$$

(5)

where the subscripts q, s in the propagator denotes the flavour. It should be noted that chiral partners in this case are between K^*, K_1 with the same flavour structure. That is between $K^*(\bar{s}q), K_1(\bar{s}q)$ or between $K^*(\bar{q}s), K_1(\bar{q}s)$.

The vector mesons and the widths are given in the table 1. The ρ and a_1 are chiral partners, but with the large width, any previous attempts to measure their mass shift seem futile. The ω and the $f_1(1285)$ have small width, but are isospin singlets and are not chiral partners as discussed in Ref. [15]. Still, in the limit where we neglect the flavor changing disconnected diagrams, they form chiral partners and hence experimental measurements are extremely useful as they provide input for extracting the density dependence of condensates. On the other hand, K^*, K_1 are chiral partners but also have reasonably small widths. Furthermore, their first excited states are already degenerate so that the chiral symmetry breaking effects seems to be concentrated in the ground states. Therefore, we will concentrate on K^*, K_1.

Table I. Mass and width of vector mesons. Units are MeV.

$J^{PC} = 1^{--}$	Mass	Width	$J^{PC} = 1^{++}$	Mass	Width
ρ	770	150	a_1	1260	250 -600
ω	782	8.49	f_1	1285	24.2
ϕ	1020	4.266	f_1	1420	54.9
$K^*(1^-)$	892	50.3	$K_1(1^+)$	1270	90
$K^*(1^-)$	1410	236	$K_1(1^+)$	1400	174

4. $U_A(1)$ breaking effect

There is a cautionary point when taking the difference of correlation function. Consider taking the difference between the scalar and pseudo-scalar currents in SU(2) for simplicity.

$$\Delta^m_{s-\eta}(q) = \int d^4x e^{iqx} \langle \mathcal{T}\Big(\bar{q}q(x)\bar{q}q(0) - \bar{q}i\gamma^5 q(x)\bar{q}i\gamma^5 q(0)\Big)\rangle. \tag{6}$$

This correlation has contribution from both the connected and disconnected diagrams. Also, the differences in the connected diagram is related to the chiral order parameter and does not vanish when chiral symmetry is broken. However, even when chiral symmetry is restored, there is a contribution that is related to the topologically non-trivial configuration and does not vanish [27]. These are the $U_A(1)$ breaking effects and are related to the zero modes containing all flavours. The contribution can be well understood by looking whether the effective instanton vertex contributes [28]. The effective vertex in the quark form is of the following form in SU(2).

$$\det[\bar{q}_R q_L. + h.c] = \Big(\bar{u}_R u_L \times \bar{d}_R d_L - \bar{u}_R d_L \times \bar{d}_R u_L + .h.c.\Big). \tag{7}$$

It is clear that this term interpolates between the correlation function and gives a non-vanishing contribution to Eq. (6). The form in the meson degrees of freedom has the following form.

$$\det[\bar{q}_R q_L. + h.c] \propto \Big(\sigma^2 + \pi^2 - \eta^2 - \alpha^2\Big). \tag{8}$$

Hence, it is also clear that this equation will also give a non-vanishing contribution to Eq. (6).

5. K_1, K^* sum rules

Let us now present the sum rule calculation for the K^* and K_1 meson given in Ref [1]. The time-ordered current correlation function of the K^*, K_1 current is given by

$$\Pi_{\mu\nu}(\omega, \mathbf{q}) = i \int d^4x e^{iq\cdot x} \langle | T[j_\mu, \bar{j}_\nu] | \rangle, \tag{9}$$

where $q^\mu = (\omega, \mathbf{q})$ and

$$\begin{aligned} j_\mu^{K_1^+} &= \bar{s}\gamma_\mu\gamma_5 u \quad , \qquad j_\mu^{K_1^-} = \bar{u}\gamma_\mu\gamma_5 s \\ j_\mu^{K^{*+}} &= \bar{s}\gamma_\mu u \quad , \qquad j_\mu^{K^{*-}} = \bar{u}\gamma_\mu s. \end{aligned} \tag{10}$$

We will not consider currents with d quarks as we will only consider symmetric matter so that the results would be the same as those with the u quarks.

As the currents are not conserved, the polarization function will have contributions from the scalar and pseudo-scalar mesons. Furthermore, there will be transverse and longitudinal polarization in the medium. Here, we will only consider the limit where we take the external three momentum to be zero $q^\mu = (\omega, \mathbf{0})$. In this work, we will extract the vector and axial-vector part by looking at the following projection.

$$\frac{1}{3}(q^\mu q^\nu / q^2 - g^{\mu\nu})\Pi_{\mu\nu}(q) \xrightarrow{\mathbf{q}\to 0} \Pi(\omega, 0). \tag{11}$$

5.1 OPE

5.1.1 Vacuum

The OPE in the vacuum has the following form up to dimension 6 operators [1].

$$\Pi(q^2) = B_0 Q^2 \ln\frac{Q^2}{\mu^2} + B_2 \ln\frac{Q^2}{\mu^2} - \frac{B_4}{Q^2} - \frac{B_6}{Q^4}, \tag{12}$$

where for the axial vector current,

$$\begin{aligned} B_0 &= \frac{1}{4\pi^2}\Big(1 + \frac{\alpha_s}{\pi}\Big) \\ B_2 &= \frac{3m_s^2}{8\pi^2} \\ B_4^A &= -m_s\langle\bar{u}u\rangle_0 + \frac{1}{12}\langle\frac{\alpha_s}{\pi}G^2\rangle_0 \\ B_6^A &= \frac{2\pi\alpha_s}{9}\Big(\langle(\bar{s}\gamma_\mu\lambda^a s + \bar{u}\gamma_\mu\lambda^a u)(\sum_q \bar{q}\gamma_\mu\lambda^a q)\rangle_0\Big) + 2\pi\alpha_s\Big(\langle(\bar{u}\gamma_\mu\lambda^a s)((\bar{s}\gamma_\mu\lambda^a u)\rangle_0\Big) \\ &= \frac{32\pi\alpha_s}{81}\Big(\langle\bar{u}u\rangle_0^2 + \langle\bar{s}s\rangle_0^2\Big) + \frac{32\pi\alpha_s}{9}\Big(\langle\bar{u}u\rangle_0\langle\bar{s}s\rangle_0\Big). \end{aligned}$$

Here, the superscript A in B_n refers to those specific for the axial vector current. For the vector current $J_\mu^{K^*} = \bar{u}\gamma_\mu s$, the OPE up to this order will be similar with $B_0^V = B_0$, $B_2^V = B_2$ and

$$\begin{aligned} B_4^V &= +m_s\langle\bar{u}u\rangle_0 + \frac{1}{12}\langle\frac{\alpha_s}{\pi}G^2\rangle_0 \\ B_6^V &= \frac{2\pi\alpha_s}{9}\Big(\langle(\bar{s}\gamma_\mu\lambda^a s + \bar{u}\gamma_\mu\lambda^a u)(\sum_q \bar{q}\gamma_\mu\lambda^a q)\rangle_0\Big) + 2\pi\alpha_s\Big(\langle(\bar{u}\gamma_\mu\gamma^5\lambda^a s)((\bar{s}\gamma_\mu\gamma^5\lambda^a u)\rangle_0\Big) \end{aligned}$$

$$= \frac{32\pi\alpha_s}{81}\Big(\langle\bar{u}u\rangle_0^2 + \langle\bar{s}s\rangle_0^2\Big) - \frac{32\pi\alpha_s}{9}\Big(\langle\bar{u}u\rangle_0\langle\bar{s}s\rangle_0\Big).$$

Here, $Q^2 \equiv -q^2$, $\mu = 1$ GeV is the renormalization scale, n in B_n indicates the canonical dimension of the operator, and $\langle O\rangle_0$ denotes the condensate of operator O in the vacuum. Here we take the value used recently [1].

5.1.2 Medium

In general, the correlation function in medium will have contribution from terms odd in q^μ because of the presence of the nucleons that breaks the charge conjugation invariance in the medium. Therefore,

$$\Pi_1(q^2) = \Pi^e(q^2) + q_0\Pi^o(q^2), \tag{13}$$

where

$$\begin{aligned}\Pi^e(q^2) &= B_0Q^2\ln\frac{Q^2}{\mu^2} + B_2\ln\frac{Q^2}{\mu^2} - \frac{B_4^{*V,A}}{Q^2} - \frac{B_6^{*V,A}}{Q^4},\\ \Pi^o(q^2) &= \frac{1}{3Q^2}(A_1^u - A_1^s)\rho - \frac{2m_N^2}{3Q^4}(A_3^u - A_3^s)\rho,\end{aligned} \tag{14}$$

with

$$\begin{aligned}B_4^{*A,V} &= B_4^{A,V} + \frac{m_N}{2}(A_2^u + A_2^s)\rho\\ B_6^{*A} &= B_6^A + \frac{8\pi}{3q^2}q^\mu q^\nu\alpha_s\langle(\bar{u}\gamma_\mu\lambda^a s)((\bar{s}\gamma_\nu\lambda^a u)\rangle_{\mathcal{ST}} - \frac{5}{6}m_N^3(A_4^u + A_4^s)\rho,\\ B_6^{*V} &= B_6^V + \frac{8\pi}{3q^2}q^\mu q^\nu\alpha_s\langle(\bar{u}\gamma_\mu\gamma^5\lambda^a s)((\bar{s}\gamma_\nu\gamma^5\lambda^a u)\rangle_{\mathcal{ST}} - \frac{5}{6}m_N^3(A_4^u + A_4^s)\rho,\end{aligned} \tag{15}$$

where we have included the dimension 6 spin 2 operators, which vanishes in the vacuum. Furthermore, while these operators have non-vanishing expectation values in general, we will take the medium expectation values to be zero based on the vacuum saturation hypothesis together with the chiral order parameter arguments given in the next section.

$\Pi^{e(o)}$ denotes even(odd) dimensional terms of correlation function. $\langle O\rangle_\rho$ is the condensate of operator O in nuclear matter, where we use the linear density approximation: $\langle O\rangle_\rho = \langle O\rangle_0 + \langle O\rangle_N\rho$ with ρ being the baryon density. $\langle\bar{u}u\rangle_N$, $\langle\bar{s}s\rangle_N$, and $\langle G^2\rangle_N$ are the nucleon matrix elements. $A_n^q(\mu^2) = 2\int_0^1 x^{n-1}\{q(x,\mu^2)+(-1)^n\bar{q}(x,\mu^2)\}dx$, where $q(x,\mu^2)$ and $\bar{q}(x,\mu^2)$ are, respectively, quark and antiquark distribution functions in the nucleon at scale μ^2, and are defined through the twist-two operators $\langle\mathcal{ST}(\bar{q}\gamma_{\mu_1}D_{\mu_2}...D_{\mu_n}q(\mu^2))\rangle_N = (-i)^{n-1}A_n^q(\mu^2)\frac{T_{\mu_1\cdots\mu_n}}{2m_N}$, where $\mathcal{ST}$ means 'symmetric and traceless', and expressed on the right side with the tensor $T_{\mu_1\cdots\mu_n}$.

5.1.3 Order parameter

The difference in the correlation function between the vector and axial vector current has the following form up to dimension 6 operators

$$\Pi^{AA} - \Pi^{VV} = \frac{2}{Q^2}m_s\langle\bar{u}u\rangle_0 - \frac{2\pi}{Q^4}\alpha_s\Big(\langle(\bar{u}\gamma_\mu\lambda^a s)((\bar{s}\gamma_\mu\lambda^a u)\rangle - \langle(\bar{u}\gamma_\mu\gamma^5\lambda^a s)((\bar{s}\gamma_\mu\gamma^5\lambda^a u)\rangle\Big)$$

$$+\frac{8\pi}{3Q^6}q^\mu q^\nu \alpha_s\Big(\langle(\bar{u}\gamma_\mu\lambda^a s)((\bar{s}\gamma_\nu\lambda^a u)\rangle_{\mathcal{ST}} - \langle(\bar{u}\gamma_\mu\gamma^5\lambda^a s)((\bar{s}\gamma_\nu\gamma^5\lambda^a u)\rangle_{\mathcal{ST}}\Big). \quad (16)$$

The second line of Eq. (16) corresponds to the twist-4 matrix element given in Ref. [29–31]. Since the difference in Eq. (16) is a chiral order parameter, the operators at each dimension are proportional to the chiral symmetry breaking effects; $m_s\langle\bar{u}u\rangle$ at dimension 4 and $\langle\bar{s}s\rangle\langle\bar{u}u\rangle$ at dimension 6 after vacuum saturation. This is so because the difference are proportional to the chiral symmetry breaking in the strangeness sector multiplied by those in the light quark sector. Therefore, vacuum saturation hypothesis should be a good approximation for these combination of 4-quark operators, for which the twist-4 matrix do not contribute. Therefore, even in the medium, one can estimate the change of the combination of 4-quark operators by just using the vacuum saturation hypothesis and change the light and strange quark condensate as is typically done in the factorization formula of the 4-quark condensate in medium.

5.1.4 Results

In a numerical analysis recently performed in Ref. [1], we find that the upper limit of the mass shift of K_1^- (K_1^+) in nuclear matter is as large as -249 (-35) MeV.

6. Phenomenological observations

The dominant hadronic decay channels are given in Table II. This means that with the kaon beam, the K^* (K_1) can be produced by the π (ρ) exchange with the nucleon.

Table II. Dominant hadronic decay channels of K^* and K_1 meson.

$J^{PC} = 1^-$	Decay Mode	Fraction	$J^{PC} = 1^+$	Decay mode	Fraction
$K^*(892)$	$K\pi$	100%	$K_1(1270)$	$K\rho$	42%
				$K^*\pi$	16%

As discussed before, the chiral partnership between the K^* and K_1 exists between the same charge states. If a K^- beam is used to produce these particles, the expected final states are given for K_1

$$K_1^- \to \left(\begin{array}{l} \rho^0 K^- \\ \rho^- \bar{K}^0 \end{array} \right. \left(\begin{array}{l} \pi^0 K^{*-} \\ \pi^- \bar{K}^{*0} \end{array} \right. , \quad \bar{K}_1^0 \to \left(\begin{array}{l} \rho^+ K^- \\ \rho^0 \bar{K}^0 \end{array} \right. \left(\begin{array}{l} \pi^+ K^{*-} \\ \pi^0 \bar{K}^{*0} \end{array} \right. ,$$

and for K^*

$$K^{*-} \to \left(\begin{array}{l} \pi^0 K^- \\ \pi^- \bar{K}^0 \end{array} \right. , \quad \bar{K}^{*0} \to \left(\begin{array}{l} \pi^+ K^- \\ \pi^0 \bar{K}^0 \end{array} \right. ,$$

7. Conclusion

After decades-long attempts to measure the mass shift and understand the origin of hadron mass, it became clear that one has to analyze hadrons with small vacuum width. Also, to identify the effect of chiral symmetry breaking effects, one has to start by looking at chiral partners. Such consideration inevitably led us to consider K^* and K_1 in matter [1]. In particular, with the kaon beam at JPARC, the possibility of observing both of these particles in a nuclear target experiment is quite feasible. The problem of final state interaction from hadronic decays can be overcome by combining the study of excitation energy [14]. Once the masses and mass difference of K^* and K_1 mesons are measured, one will be a closer to understanding the origin of the hadron masses and the effects of chiral symmetry breaking in them.

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