



η bound states in nuclei: a probe of flavour-singlet dynamics

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Abstract

We argue that η bound states in nuclei are sensitive to the singlet component in the η . The bigger the singlet component, the more attraction and the greater the binding. Thus, measurements of η bound states will yield new information about axial U(1) dynamics and glue in mesons. η - η' mixing plays an important role in understanding the value of the η -nucleon scattering length.

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1. Introduction

Measurements of the pion, kaon and eta meson masses and their interactions in finite nuclei provide new constraints on our understanding of dynamical symmetry breaking in low energy QCD [1]. New experiments at the GSI will employ the recoilless (d, ³He) reaction to study the possible formation of η meson bound states inside the nucleus [2], following on from the successful studies of pionic atoms in these reactions [3]. The idea is to measure the excitation-energy spectrum and then, if a clear bound state is observed, to extract the in-medium effective mass, m_{η}^* , of the η in nuclei through performing a fit to this spectrum with the η -nucleus optical potential.

In this Letter we argue that m_{η}^* is sensitive to the flavour-singlet component in the η , and hence to non-perturbative glue [4,5] associated with axial U(1) dynamics. An important source of the in-medium mass modification comes from light-quarks coupling to the scalar σ mean-field in the nucleus. Increasing the flavour-singlet component in the η at the expense of the octet component gives more attraction, more binding and a larger value of the η -nucleon scattering length, $a_{\eta N}$. This result may explain why values of $a_{\eta N}$ extracted from phenomenological fits to experimental data where the η - η' mixing angle is

unconstrained give larger values than those predicted in theoretical models where the η is treated as a pure octet state.

We first introduce the basic physics. Next, in Section 2 we briefly review the QCD axial U(1) problem and its application to the η mass in nuclei. We motivate the *existence* of gluonic corrections to m_{η}^* which go beyond pure Goldstone boson dynamics. While QCD arguments imply information about the sign of the mass shift, a rigorous numerical calculation of m_{η}^* from QCD is presently not feasible. Hence, in Section 3, we consider QCD inspired model predictions for the η -nucleus and η' -nucleus systems and the vital role of flavour-singlet degrees of freedom in η bound-states. In Section 4 we summarize and conclude.

Meson masses in nuclei are determined from the scalar induced contribution to the meson propagator evaluated at zero three-momentum, $\vec{k} = 0$, in the nuclear medium. Let $k = (E, \vec{k})$ and m denote the four-momentum and mass of the meson in free space. Then, one solves the equation

$$k^2 - m^2 = \text{Re } \Pi(E, \vec{k}, \rho) \quad (1)$$

for $\vec{k} = 0$, where Π is the in-medium s-wave meson self-energy. Contributions to the in medium mass come from coupling to the scalar σ field in the nucleus in mean-field approximation, nucleon-hole and resonance-hole excitations in the medium. The s-wave self-energy can be written as [6]

$$\Pi(E, \vec{k}, \rho)|_{\vec{k}=0} = -4\pi\rho \left(\frac{b}{1 + b\langle \frac{1}{\tau} \rangle} \right). \quad (2)$$

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Here ρ is the nuclear density, $b = a(1 + \frac{m}{M})$, where a is the meson–nucleon scattering length, M is the nucleon mass and $\langle \frac{1}{r} \rangle$ is the inverse correlation length, $\langle \frac{1}{r} \rangle \simeq m_\pi$ for nuclear matter density [6]. (m_π is the pion mass.) Attraction corresponds to positive values of a . The denominator in Eq. (2) is the Ericson–Ericson–Lorentz–Lorenz double scattering correction.

What should we expect for the η and η' ?

2. QCD considerations

Spontaneous chiral symmetry breaking is associated with a non-vanishing chiral condensate

$$\langle \text{vac} | \bar{q}q | \text{vac} \rangle < 0. \quad (3)$$

The non-vanishing chiral condensate also spontaneously breaks the axial U(1) symmetry so, naively, in the two-flavour theory one expects an isosinglet pseudoscalar degenerate with the pion. The lightest mass isosinglet is the η meson, which has a mass of 547.75 MeV.

The puzzle deepens when one considers SU(3). Spontaneous chiral symmetry breaking suggests an octet of would-be Goldstone bosons: the octet associated with chiral SU(3)_L ⊗ SU(3)_R plus a singlet boson associated with axial U(1)—each with mass squared $m_{\text{Goldstone}}^2 \sim m_q$. The physical η and η' masses are about 300–400 MeV too big to fit in this picture. One needs extra mass in the singlet channel associated with non-perturbative topological gluon configurations and the QCD axial anomaly [5]. The strange quark mass induces considerable η – η' mixing. For free mesons the η – η' mass matrix (at leading order in the chiral expansion) is

$$M^2 = \begin{pmatrix} \frac{4}{3}m_K^2 - \frac{1}{3}m_\pi^2 & -\frac{2}{3}\sqrt{2}(m_K^2 - m_\pi^2) \\ -\frac{2}{3}\sqrt{2}(m_K^2 - m_\pi^2) & [\frac{2}{3}m_K^2 + \frac{1}{3}m_\pi^2 + \tilde{m}_{\eta_0}^2] \end{pmatrix}. \quad (4)$$

Here $\tilde{m}_{\eta_0}^2$ is the gluonic mass term which has a rigorous interpretation through the Witten–Veneziano mass formula [7,8] and which is associated with non-perturbative gluon topology, related perhaps to confinement [9] or instantons [10]. The masses of the physical η and η' mesons are found by diagonalizing this matrix, viz.

$$\begin{aligned} |\eta\rangle &= \cos\theta|\eta_8\rangle - \sin\theta|\eta_0\rangle, \\ |\eta'\rangle &= \sin\theta|\eta_8\rangle + \cos\theta|\eta_0\rangle, \end{aligned} \quad (5)$$

where

$$\eta_0 = \frac{1}{\sqrt{3}}(u\bar{u} + d\bar{d} + s\bar{s}), \quad \eta_8 = \frac{1}{\sqrt{6}}(u\bar{u} + d\bar{d} - 2s\bar{s}). \quad (6)$$

One obtains values for the η and η' masses:

$$\begin{aligned} m_{\eta',\eta}^2 &= (m_K^2 + \tilde{m}_{\eta_0}^2/2) \\ &\pm \frac{1}{2}\sqrt{\left(2m_K^2 - 2m_\pi^2 - \frac{1}{3}\tilde{m}_{\eta_0}^2\right)^2 + \frac{8}{9}\tilde{m}_{\eta_0}^4}. \end{aligned} \quad (7)$$

The physical mass of the η and the octet mass $m_{\eta_8} = \sqrt{\frac{4}{3}m_K^2 - \frac{1}{3}m_\pi^2}$ are numerically close, within a few percent. However, to build a theory of the η on the octet approximation

risks losing essential physics associated with the singlet component. Turning off the gluonic term, one finds the expressions $m_{\eta'} \sim \sqrt{2m_K^2 - m_\pi^2}$ and $m_\eta \sim m_\pi$. That is, without extra input from glue, in the OZI limit, the η would be approximately an isosinglet light-quark state ($\frac{1}{\sqrt{2}}|\bar{u}u + \bar{d}d\rangle$) degenerate with the pion and the η' would be a strange-quark state $|\bar{s}s\rangle$ —mirroring the isoscalar vector ω and ϕ mesons.

Taking the value $\tilde{m}_{\eta_0}^2 = 0.73 \text{ GeV}^2$ in the leading-order mass formula, Eq. (7), gives agreement with the physical masses at the 10% level. This value is obtained by summing over the two eigenvalues in Eq. (7): $m_\eta^2 + m_{\eta'}^2 = 2m_K^2 + \tilde{m}_{\eta_0}^2$ and substituting the physical values of m_η , $m_{\eta'}$ and m_K [8]. The corresponding η – η' mixing angle $\theta \simeq -18^\circ$ is within the range from -17° to -20° obtained from a study of various decay processes in [11,12].¹ The key point of Eq. (7) is that mixing and gluon dynamics play a crucial role in both the η and η' masses and that treating the η as an octet pure would-be Goldstone boson risks losing essential physics.

2.1. η and η' interactions with the nuclear medium

What can QCD tell us about the behaviour of the gluonic mass contribution in the nuclear medium?

The physics of axial U(1) degrees of freedom is described by the U(1)-extended low-energy effective Lagrangian [8]. In its simplest form this reads

$$\begin{aligned} \mathcal{L} &= \frac{F_\pi^2}{4} \text{Tr}(\partial^\mu U \partial_\mu U^\dagger) + \frac{F_\pi^2}{4} \text{Tr} M(U + U^\dagger) \\ &+ \frac{1}{2}iQ \text{Tr}[\log U - \log U^\dagger] + \frac{3}{\tilde{m}_{\eta_0}^2 F_0^2} Q^2. \end{aligned} \quad (8)$$

Here $U = \exp\{i(\phi/F_\pi + \sqrt{2/3}\eta_0/F_0)\}$ is the unitary meson matrix where $\phi = \sum \pi_a \lambda_a$ denotes the octet of would-be Goldstone bosons associated with spontaneous chiral SU(3)_L ⊗ SU(3)_R breaking and η_0 is the singlet boson. In Eq. (8) Q denotes the topological charge density ($Q = \frac{\alpha_s}{4\pi} G_{\mu\nu} \tilde{G}^{\mu\nu}$); $M = \text{diag}[m_\pi^2, m_\pi^2, 2m_K^2 - m_\pi^2]$ is the quark-mass induced meson mass matrix. The pion decay constant $F_\pi = 92.4 \text{ MeV}$ and F_0 is the flavour-singlet decay constant, $F_0 \sim F_\pi \sim 100 \text{ MeV}$ [11].

The flavour-singlet potential involving Q is introduced to generate the gluonic contribution to the η and η' masses and to reproduce the anomaly in the divergence of the gauge-invariantly renormalized flavour-singlet axial-vector current. The gluonic term Q is treated as a background field with no kinetic term. It may be eliminated through its equation of motion to generate a gluonic mass term for the singlet boson, viz.

$$\frac{1}{2}iQ \text{Tr}[\log U - \log U^\dagger] + \frac{3}{\tilde{m}_{\eta_0}^2 F_0^2} Q^2 \mapsto -\frac{1}{2}\tilde{m}_{\eta_0}^2 \eta_0^2. \quad (9)$$

The most general low-energy effective Lagrangian involves a $U_A(1)$ invariant polynomial in Q^2 . Higher-order terms in Q^2

¹ Closer agreement with the physical masses can be obtained by introducing the singlet decay constant $F_0 \neq F_\pi$ and including higher-order mass terms in the chiral expansion [13,14].

become important when we consider scattering processes involving more than one η' [15]. In general, couplings involving Q give OZI violation in physical observables.

To investigate what happens to $\tilde{m}_{\eta_0}^2$ in the medium we first couple the σ (correlated two-pion) mean-field in nuclei to the topological charge density Q . The interactions of the η and η' with other mesons and with nucleons can be studied by coupling the Lagrangian (8) to other particles. For example, the OZI violating interaction $\lambda Q^2 \partial_\mu \pi_a \partial^\mu \pi_a$ is needed to generate the leading (tree-level) contribution to the decay $\eta' \rightarrow \eta \pi \pi$ [15]. When iterated in the Bethe–Salpeter equation for meson–meson rescattering this interaction yields a dynamically generated exotic state with quantum numbers $J^{PC} = 1^{-+}$ and mass about 1400 MeV [16]. This suggests a dynamical interpretation of the lightest-mass 1^{-+} exotic observed at BNL and CERN.

Motivated by this two-pion coupling to Q^2 , we couple the topological charge density to the σ (two-pion) mean-field in the nucleus by adding the Lagrangian term

$$\mathcal{L}_{\sigma Q} = Q^2 g_\sigma^Q \sigma, \quad (10)$$

where g_σ^Q denotes coupling to the σ mean field—that is, we consider an in-medium renormalization of the coefficient of Q^2 in the effective chiral Lagrangian. Following the treatment in Eq. (9) we eliminate Q through its equation of motion. The gluonic mass term for the singlet boson then becomes

$$\tilde{m}_{\eta_0}^2 \mapsto \tilde{m}_{\eta_0}^{*2} = \tilde{m}_{\eta_0}^2 \frac{1+2x}{(1+x)^2} < \tilde{m}_{\eta_0}^2, \quad (11)$$

where

$$x = \frac{1}{3} g_\sigma^Q \sigma \tilde{m}_{\eta_0}^2 F_0^2. \quad (12)$$

That is, the gluonic mass term decreases in-medium independent of the sign of g_σ^Q and the medium acts to partially neutralize axial U(1) symmetry breaking by gluonic effects.

This scenario has possible support from recent lattice calculations [17] which suggest that non-trivial gluon topology configurations are suppressed inside hadrons. Further recent work at high chemical potential ($\mu > 500$ MeV) suggests that possible confinement and instanton contributions to $\tilde{m}_{\eta_0}^2$ are suppressed with increasing density in this domain [18]. We investigate the size of the η mass shift in Section 3 below.

2.2. The η nucleon scattering length and anomalous glue

Further insight is provided from looking at the scattering length. When the U(1)-extended chiral Lagrangian is coupled to nucleons one finds new OZI violating couplings in the flavour-singlet sector [19]. An example is the gluonic contribution to the singlet Goldberger–Treiman relation [20] which connects axial U(1) dynamics and the spin structure of the proton studied in polarized deep inelastic scattering and high-energy polarized proton–proton collisions—for a recent review see [21]. In the chiral limit the singlet analogy to the Weinberg–Tomozawa term does not vanish because of the anomalous glue terms. Starting from the simple Born term one finds anomalous

gluonic contributions to the singlet-meson–nucleon scattering length proportional to $\tilde{m}_{\eta_0}^2$ and $\tilde{m}_{\eta_0}^4$ [22].

We briefly summarize this section.

The masses of the η and η' receive contributions from terms associated with both explicit chiral symmetry breaking and with anomalous glue through the Witten–Veneziano term. Mixing is important and, ideally, one would like to consider the medium dependence of the different basic physics inputs. At the QCD level, OZI-violating gluonic couplings have the potential to affect the effective η and η' masses in nuclei and, through Eq. (2), the η –nucleon and η' –nucleon scattering lengths. It is interesting to also mention the observation of Brodsky et al. [23] that attractive gluonic van der Waals type exchanges have the potential to produce flavour-singlet η_c bound-states in the (d, ^3He) reaction close to threshold.

The above discussion is intended to motivate the *existence* of medium modifications to $\tilde{m}_{\eta_0}^2$ in QCD. However, a rigorous calculation of m_η^* from QCD is beyond present theoretical technology. Hence, one has to look to QCD motivated models and phenomenology for guidance about the numerical size of the effect. The physics described in Eqs. (4)–(7) tells us that the simple octet approximation may not suffice.

3. Models

We now discuss the size of flavour-singlet effects in m_η^* , $m_{\eta'}^*$ (the η' mass in-medium) and the scattering lengths $a_{\eta N}$ and $a_{\eta' N}$. First we consider the values of $a_{\eta N}$ and $a_{\eta' N}$ extracted from phenomenological fits to experimental data. There are several model predictions for the η mass in nuclear matter, starting from different assumptions. We collect and compare these approaches and predictions with particular emphasis on the contribution of η – η' mixing. We also compare model predictions for the internal structure of the $S_{11}(1535)$ nucleon resonance and its in-medium excitation energy.

3.1. Phenomenological determinations of $a_{\eta N}$ and $a_{\eta' N}$

Green and Wycech [24] have performed phenomenological K -matrix fits to a variety of near-threshold processes ($\pi N \rightarrow \pi N$, $\pi N \rightarrow \eta N$, $\gamma N \rightarrow \pi N$ and $\gamma N \rightarrow \eta N$) to extract a value for the η –nucleon scattering. In these fits the $S_{11}(1535)$ is introduced as an explicit degree of freedom—that is, it is treated like a 3-quark state—and the η – η' mixing angle is taken as a free parameter. The real part of $a_{\eta N}$ extracted from these fits is 0.91(6) fm for the on-shell scattering amplitude.

From measurements of η production in proton–proton collisions close to threshold, COSY-11 have extracted a scattering length $a_{\eta N} \simeq 0.7 + i0.4$ fm from the final state interaction (FSI) based on the effective range approximation [25]. For the η' , COSY-11 have deduced a conservative upper bound on the η' –nucleon scattering length $|\text{Re } a_{\eta' N}| < 0.8$ fm [26] with a preferred a value between 0 and 0.1 fm [27] obtained by comparing the FSI in π^0 and η' production in proton–proton collisions close to threshold.

3.2. Chiral models

Chiral models involve performing a coupled channels analysis of η production after multiple rescattering in the nucleus which is calculated using the Lippmann–Schwinger [28] or Bethe–Salpeter [29] equations with potentials taken from the SU(3) chiral Lagrangian for low-energy QCD. In these chiral model calculations the η is taken as pure octet state ($\eta = \eta_8$) with no mixing and the singlet sector turned off. These calculations yield a small mass shift in nuclear matter

$$\frac{m_\eta^*}{m_\eta} \simeq 1 - 0.05\rho/\rho_0. \quad (13)$$

The values of the η –nucleon scattering length extracted from these chiral model calculations are $0.2 + i0.26$ fm [28] and $0.26 + i0.24$ fm [29] with slightly different treatment of the intermediate state mesons.

3.3. The quark–meson coupling model

The third approach we consider is the Quark–Meson Coupling model (QMC) [30]. Here one uses the large η mass (which in QCD is induced by mixing and the gluonic mass term) to motivate taking an MIT Bag description for the η wavefunction, and then coupling the light (up and down) quark and antiquark fields in the η to the scalar σ field in the nucleus working in mean-field approximation [30]. The strange-quark component of the wavefunction does not couple to the σ field and η – η' mixing is readily built into the model.

The mass for the η in nuclear matter is self-consistently calculated by solving for the MIT Bag in the nuclear medium [30]:

$$m_\eta^*(\vec{r}) = \frac{2[a_p^2 \Omega_q^*(\vec{r}) + b_p^2 \Omega_s(\vec{r})] - z_\eta}{R_\eta^*} + \frac{4}{3} \pi R_\eta^{*3} B, \quad (14)$$

$$\left. \frac{\partial m_j^*(\vec{r})}{\partial R_j} \right|_{R_j=R_j^*} = 0 \quad (j = \eta, \eta'). \quad (15)$$

Here Ω_q^* and Ω_s are light-quark and strange-quark Bag energy eigenvalues, R_η^* is the Bag radius in the medium and B is the Bag constant. The η – η' mixing angle θ is included in the terms

$$a_P = \frac{1}{\sqrt{3}} \cos \theta - \sqrt{\frac{2}{3}} \sin \theta, \quad b_P = \sqrt{\frac{2}{3}} \cos \theta + \frac{1}{\sqrt{3}} \sin \theta,$$

and can be varied in the model. One first solves the Bag for the free η with a given mixing angle, and then turns on QMC to obtain the mass-shift. Results for the η' are obtained by interchanging $a_P \leftrightarrow b_P$. In Eq. (14), z_η parameterizes the sum of the center-of-mass and gluon fluctuation effects, and is assumed to be independent of density [31]. The current quark masses are taken as $m_q = 5$ MeV and $m_s = 250$ MeV.²

² This is the strange-quark mass needed to reproduce the Lambda and Sigma masses in the model. While larger than the values for m_s quoted at momentum scales relevant perturbative QCD, the Bag model approximates QCD at a very low scale (well below 1 GeV)—a region where renormalization group evolution would make the running masses much larger than at 2 GeV².

Table 1

Physical masses fitted in free space, the bag masses in medium at normal nuclear-matter density, $\rho_0 = 0.15$ fm^{−3}, and corresponding meson–nucleon scattering lengths (see below)

	m (MeV)	m^* (MeV)	$\text{Re } a$ (fm)
η_8	547.75	500.0	0.43
η (-10°)	547.75	474.7	0.64
η (-20°)	547.75	449.3	0.85
η_0	958	878.6	0.99
η' (-10°)	958	899.2	0.74
η' (-20°)	958	921.3	0.47

The coupling constants in the model for the coupling of light-quarks to the σ (and ω and ρ) mean-fields in the nucleus are adjusted to fit the saturation energy and density of symmetric nuclear matter and the bulk symmetry energy. The Bag parameters used in these calculations are $\Omega_q = 2.05$ (for the light quarks) and $\Omega_s = 2.5$ (for the strange quark) for free hadrons with $B = (170 \text{ MeV})^4$. For nuclear matter density we find $\Omega_q^* = 1.81$ for the 1s state. This value depends on the coupling of light-quarks to the σ mean-field and is independent of the mixing angle θ . Likewise, Ω_q and Ω_s are determined by solving for light and strange quarks in the MIT Bag potential and are independent of θ .

For the η and η' mesons the ω vector mean-field couples with the same magnitude and opposite sign to the quarks and antiquarks in the meson, and therefore cancels. Increasing the mixing angle increases the amount of singlet relative to octet components in the η . This produces greater attraction through increasing the amount of light-quark compared to strange-quark components in the η and a reduced effective mass. Through Eq. (2) increasing the mixing angle also increases the η –nucleon scattering length $a_{\eta N}$. We quantify this in Table 1 which presents results for the pure octet ($\eta = \eta_8$, $\theta = 0$) and the values $\theta = -10^\circ$ and -20° (the physical mixing angle).

The values of $\text{Re } a_\eta$ quoted in Table 1 are obtained from substituting the in-medium and free masses into Eq. (2) with the Ericson–Ericson denominator turned-off, and using the free mass $m = m_\eta$ in the expression for b . The effect of exchanging m for m^* in b is a 5% increase in the quoted scattering length. The QMC model makes no claim about the imaginary part of the scattering length. The key observation is that η – η' mixing leads to a factor of two increase in the mass-shift and in the scattering length obtained in the model.

The QMC model is calibrated by fixing the coupling constants to the observed properties of nuclear matter or finite nuclei. So, even though it is mean-field (no correlations) it does fit observed binding energies. When one applies the same model with the same (quark level couplings) to the binding of etas the natural belief is that it should give the physical binding energies. From these one can extract an effective scattering length. Because the QMC model has been explored mainly at the mean-field level, it is not clear that one should include the Ericson–Ericson–Lorentz–Lorenz term in extracting the corresponding η nucleon scattering length. If one substitutes the scattering lengths given in Table 1 into Eq. (2) (and neglects

the imaginary part which is not predicted by the model) one obtains resummed values $a_{\text{eff}} = a/(1 + b(1/r))$ equal to 0.44 fm for the η and 0.28 fm for the η' for the physical mixing angle $\theta = -20$ degrees. (Here we take $\langle 1/r \rangle \simeq m_\pi$ for nuclear matter density [6].)

The density dependence of the mass-shifts in the QMC model is discussed in Ref. [30]. Neglecting the Ericson–Ericson term, the mass-shift is approximately linear. For densities ρ between 0.5 and 1 times ρ_0 (nuclear matter density) we find

$$\frac{m_\eta^*}{m_\eta} \simeq 1 - 0.17 \frac{\rho}{\rho_0} \quad (16)$$

for the physical mixing angle -20° . The scattering lengths extracted from this analysis are density independent to within a few percent over the same range of densities.

Finally, we note that in the QMC treatment one assumes that the value of the mixing angle does not change in medium. As mentioned above this is not excluded and merits further investigation.

3.4. The $S_{11}(1535)$ resonance in nuclear matter

It is interesting to compare the different model predictions for the $S_{11}(1535)$ nucleon resonance which couples strongly to the η –nucleon system.³ In quark models the S_{11} is interpreted as a 3-quark state: $(1s)^2(1p)$. This interpretation has support from quenched lattice calculations [37] which also suggest that the $\Lambda(1405)$ resonance has a significant non 3-quark component. In the Cloudy Bag Model the $\Lambda(1405)$ is dynamically generated in the kaon–nucleon system [38]. Chiral coupled channels models with an octet $\eta = \eta_8$ agree with these predictions for the $\Lambda(1405)$ and differ for the $S_{11}(1535)$, which is interpreted as a $K\Sigma$ quasi-bound state [32].

Experiments in heavy-ion collisions [35] and η photoproduction from nuclei [33,34] suggest little modification of the $S_{11}(1535)$ excitation in-medium, though some evidence for the broadening of the S_{11} in nuclei was reported in [34]. Despite the different physics input, both QMC and the coupled channels models agree with this finding. In QMC the excitation energy is ~ 1544 MeV. This is obtained as follows. For a quark in the 1p state the Bag light-quark energy-eigenvalue in free space is $\Omega_q = 3.81$. In QMC at normal nuclear-matter density this is reduced to $\Omega_q^* = 3.77$. (Note the smaller mass shift compared to the s-wave eigenvalue.) The scalar mass term for the S_{11} is reduced to ~ 1424 MeV through coupling to the σ mean-field. The scalar attraction is compensated by repulsion from coupling to the omega mean-field, $\sim +120$ MeV, to give the excitation energy 1544 MeV. (For the η , the ω mean-field coupling to the quark and antiquark enters with equal magnitude and opposite sign and therefore cancels.) In chiral coupled channels calculations one finds a similar S_{11} excitation energy ~ 1560 MeV. Here the medium independence of the resonance excitation energy is interpreted as arising from the absence of

Pauli blocking of the $K\Sigma$ system in nuclear matter. We note that in QMC for all baryons the scalar attraction very nearly cancels the vector repulsion, leaving a small (few 10s of MeV) net attraction or repulsion.

4. Conclusions

η – η' mixing increases the flavour-singlet and light-quark components in the η . The greater the flavour-singlet component in the η , the greater the η binding energy in nuclei through increased attraction and the smaller the value of m_η^* . Through Eq. (2), this corresponds to an increased η –nucleon scattering length $a_{\eta N}$, greater than the value one would expect if the η were a pure octet state. Measurements of η bound-states in nuclei are therefore a probe of singlet axial U(1) dynamics in the η .

It will be very interesting to see the results from the new GSI experiment for η bound-states. Additional studies might be possible using η production in low-energy proton–nucleus collisions and in photoproduction. Here one might use an electromagnetic calorimeter, for example, WASA@COSY, to tag the two-photon decay of the η . However, unlike the GSI programme, one has to be careful in these experiments whether the η is produced inside the nucleus or on the surface. Possibilities to study η and η' -mesic nuclei in (γ, p) spectra are discussed in [39].

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References

- [1] P. Kienle, T. Yamazaki, Prog. Part. Nucl. Phys. 52 (2004) 85.
- [2] R.S. Hayano, S. Hirezaki, A. Gillitzer, Eur. Phys. J. A 6 (1999) 99.
- [3] K. Suzuki, et al., Phys. Rev. Lett. 92 (2004) 072302.
- [4] S.D. Bass, Phys. Scr. T 99 (2002) 96.
- [5] G.M. Shore, hep-ph/9812354.
- [6] T.E.O. Ericson, W. Weise, Pions and Nuclei, Oxford Univ. Press, Oxford, 1988.
- [7] G. Veneziano, Nucl. Phys. B 159 (1979) 213; E. Witten, Ann. Phys. 128 (1980) 363.
- [8] P. Di Vecchia, G. Veneziano, Nucl. Phys. B 171 (1980) 253.
- [9] J. Kogut, L. Susskind, Phys. Rev. D 11 (1975) 3594; E. Witten, Nucl. Phys. B 149 (1979) 285; I. Horvath, N. Isgur, J. McCune, H.B. Thacker, Phys. Rev. D 65 (2001) 014502.
- [10] G. 't Hooft, Phys. Rev. Lett. 37 (1976) 8; G. 't Hooft, Phys. Rev. D 14 (1976) 3432.
- [11] F.J. Gilman, R. Kauffman, Phys. Rev. D 36 (1987) 2761; F.J. Gilman, R. Kauffman, Phys. Rev. D 37 (1988) 3348, Erratum.
- [12] P. Ball, J.M. Frere, M. Tytgat, Phys. Lett. B 365 (1996) 367.
- [13] H. Leutwyler, Nucl. Phys. B (Proc. Suppl.) 64 (1998) 223; R. Kaiser, H. Leutwyler, hep-ph/9806336.
- [14] T. Feldmann, P. Kroll, B. Stech, Phys. Rev. D 58 (1998) 114006; T. Feldmann, P. Kroll, B. Stech, Phys. Lett. B 449 (1999) 339;

³ We refer to [36] for a recent discussion of the role of the $S_{11}(1535)$ in the η –nucleon optical potential.

- T. Feldmann, *Int. J. Mod. Phys. A* 15 (2000) 159.
- [15] P. Di Vecchia, F. Nicodemi, R. Pettorino, G. Veneziano, *Nucl. Phys. B* 181 (1981) 318.
- [16] S.D. Bass, E. Marco, *Phys. Rev. D* 65 (2002) 057503;
A.P. Szczepaniak, A.R. Dzierba, S. Tiege, *Phys. Rev. Lett.* 91 (2003) 092002.
- [17] F. Bissey, et al., *Nucl. Phys. B (Proc. Suppl.)* 141 (2005) 22.
- [18] T. Schäfer, *Phys. Rev. D* 67 (2003) 074502.
- [19] S.D. Bass, *Phys. Lett. B* 463 (1999) 286.
- [20] G.M. Shore, G. Veneziano, *Phys. Lett. B* 244 (1990) 75;
T. Hatsuda, *Nucl. Phys. B* 329 (1990) 376.
- [21] S.D. Bass, *Rev. Mod. Phys.* 77 (2005) 1257, hep-ph/0411005.
- [22] S.D. Bass, S. Wetzel, W. Weise, *Nucl. Phys. A* 686 (2001) 429.
- [23] S.J. Brodsky, I. Schmidt, G.F. de Teramond, *Phys. Rev. Lett.* 64 (1990) 1011.
- [24] A.M. Green, S. Wycech, *Phys. Rev. C* 60 (1999) 035208;
A.M. Green, S. Wycech, *Phys. Rev. C* 71 (2005) 014001.
- [25] COSY-11 Collaboration, P. Moskal, et al., *Phys. Rev. C* 69 (2004) 025203.
- [26] COSY-11 Collaboration, P. Moskal, et al., *Phys. Lett. B* 474 (2000) 416.
- [27] COSY-11 Collaboration, P. Moskal, et al., *Phys. Lett. B* 482 (2000) 356.
- [28] T. Waas, W. Weise, *Nucl. Phys. A* 625 (1997) 287.
- [29] T. Inoue, E. Oset, *Nucl. Phys. A* 710 (2002) 354;
C. Garcia-Recio, T. Inoue, J. Nieves, E. Oset, *Phys. Lett. B* 550 (2002) 47.
- [30] K. Tsushima, D.H. Lu, A.W. Thomas, K. Saito, *Phys. Lett. B* 443 (1998) 26;
K. Tsushima, *Nucl. Phys. A* 670 (2000) 198c.
- [31] P.A.M. Guichon, K. Saito, E. Rodionov, A.W. Thomas, *Nucl. Phys. A* 601 (1996) 349.
- [32] N. Kaiser, T. Waas, W. Weise, *Nucl. Phys. A* 612 (1997) 297.
- [33] M. Rößig-Landau, et al., *Phys. Lett. B* 373 (1996) 45.
- [34] T. Yorita, et al., *Phys. Lett. B* 476 (2000) 226.
- [35] R. Averbeck, et al., *Z. Phys. A* 359 (1997) 65.
- [36] D. Jido, H. Nagahiro, S. Hirenzaki, *Phys. Rev. C* 66 (2002) 045202;
H. Nagahiro, D. Jido, S. Hirenzaki, *Phys. Rev. C* 68 (2003) 035205.
- [37] W. Melnitchouk, et al., *Phys. Rev. D* 67 (2003) 114506;
D. Brömmel, et al., *Phys. Rev. D* 69 (2004) 094513;
N. Mathur, et al., *Phys. Lett. B* 605 (2005) 137.
- [38] E.A. Veit, B.K. Jennings, A.W. Thomas, R.C. Barrett, *Phys. Rev. D* 31 (1985) 1033.
- [39] H. Nagahiro, S. Hirenzaki, *Phys. Rev. Lett.* 94 (2005) 232503;
H. Nagahiro, D. Jido, S. Hirenzaki, *Nucl. Phys. A* 761 (2005) 92.