In-medium modification of the isovector pion-nucleon amplitude

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We study the in-medium modification of the isovector $\pi N$ amplitude using a non-linear representation of the sigma model but keeping the scalar degree of freedom. We check that our result does not depend on the representation. We discuss the connection with other approaches based on chiral perturbation theory.

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I. INTRODUCTION

The experimental program on deeply bound pionic states has revived the interest on the interaction of pions with nuclei. Recently the attention has focused on the charge exchange scattering length $b_1$ [1, 2, 3, 4, 5, 6, 7]. The fit to these data [1] suggests a large enhancement of $b_1$ in nuclei, as anticipated in Ref. [3]. It is the purpose of this letter to investigate this question within a non-linear representation of the sigma model keeping the scalar degree of freedom. The advantage of this model is that it allows an explicit evaluation of the contribution of the nuclear pion gas in a way so much on a comparison with data but rather on the derivation itself of the in-medium value for $b_1$ and on some important questions of principle arising in comparison with other theoretical works.

II. THE S-WAVE PION SELF-ENERGY

Our starting point is based on the linear sigma model but reformulated as in Ref. [8]. The original sigma field $\sigma$ and pion fields $\pi$ are eliminated in favor of a chiral invariant scalar field, $S = f_\pi + s$, and a new pion field $\phi$ according to:

$$\sigma + i \bar{\pi} \cdot \pi = S \mu \equiv (f_\pi + s) \exp \left[ \frac{i \bar{\pi} \cdot \phi}{f_\pi} G \left( \frac{\phi^2}{f_\pi} \right) \right]. \tag{1}$$

As proposed in Ref. [8], the fluctuating scalar field $s$, which physically describes the fluctuations of the radius of the chiral circle around its vacuum value $f_\pi$, can be identified with the sigma meson of relativistic theories of the Walecka type. The function $G(X^2) = 1 + \alpha X^2 + \ldots$ selects the particular realization of the model. Changing $G$ amounts to a redefinition of the pion field, which should not affect physics. In the following we keep $\alpha$ arbitrary and check that the final results for physical observables do not depend on this parameter. The Lagrangian as given by eqs.(19-30) of Ref. [8] writes (notice that the notations $\Theta, \theta$ for the chiral invariant scalar field are now replaced by $S, s$):

$$\mathcal{L} = (f_\pi + s)^2 \text{Tr} \partial^\mu U \partial_\mu U^\dagger + \frac{1}{2} \partial^\mu s \partial_\mu s - \frac{m_\pi^2 - m_\eta^2}{8f_\pi^2} \left( s^2 + 2f_\pi s + \frac{2f_\pi^2 m_\pi^2}{m_\eta^2 - m_\pi^2} \right)^2$$

$$+ i \bar{N} \gamma_\mu \partial_\mu N - MN \left( 1 + \frac{s}{f_\pi} \right) \bar{N} N$$

$$+ \bar{N} \gamma_\mu \nu^\mu N + \left( 1 - (1 - g_A) \left( 1 + \frac{s}{f_\pi} \right)^2 \right) \bar{N} \gamma_\mu \gamma_5 A_\mu^\nu N$$

$$+ i \frac{1 - g_A}{2f_\pi} \left( 1 + \frac{s}{f_\pi} \right) \bar{N} \gamma_\mu \gamma_5 \partial_\mu s + \mathcal{L}_{\chi_{SB}}, \tag{2}$$

where we have defined (with $\xi^2 = U$):

$$\nu^\mu = \frac{i}{2} (\xi \partial_\mu \xi + \xi^\dagger \partial_\mu \xi) \quad A_\mu^\nu = \frac{i}{2} (\xi \partial_\mu \xi^\dagger - \xi^\dagger \partial_\mu \xi). \tag{3}$$
The scalar contribution is given by:

\[ \mathcal{L}_{\chi_{SB}} = f_{\pi} m_{\pi}^2 \left( f_{\pi} + s \right) \cos \left( \phi_{\pi} \frac{G}{f_{\pi}^2} \right) . \]  

In order to reproduce the experimental value of the isoscalar scattering length it is necessary, as in chiral perturbation theory, to introduce additional chiral invariant pieces contributing to order \( \mathcal{O}(Q^2) \):

\[ \Delta \sigma_{\pi N}^{(2)} = c_3' \tilde{N}(u \cdot u) N + c_2' \tilde{N}(v \cdot u)^2 N , \]  

where \( u_{\mu} = i \xi^\dagger \partial_{\mu} U \xi^\dagger \) and \( v_{\mu} \) is the four-velocity of the nucleon. These \( c_2' \) and \( c_3' \) parameters have to be empirically determined. Finally the introduction of a form factor at the p-wave scalar piece of the sigma term. Fig. 1c is immediately obtained from the Lagrangian given in eq. (5) and fig. 1d is it involves the pionic piece of the sigma term. As for the scalar exchange contribution (fig. 1b) it is related to the

The pion loop contribution (fig. 1a), including vertex corrections not shown in Fig. 1a, is derived in Refs. [8, 12, 13].

In order to link it to the chiral perturbation result we expand it in orders of the pion mass:

\[ \Sigma_N^{(\pi)} = -4 c_1' m_{\pi}^2 + \langle \Sigma_N^{(\pi)\text{LNAC}} \rangle + ... . \]  

The coefficient \( c_1' \) depends on the model (i.e., on the form factor). It is related to the coefficient \( c_1 \) of chiral perturbation theory by \( c_1' = c_1 - M_{\pi}/4m_{\pi}^2 \). The leading non analytical contribution, \( \langle \Sigma_N^{(\pi)} \rangle \), is instead given by chiral symmetry alone:

\[ \langle \Sigma_N^{(\pi)\text{LNAC}} \rangle = - \frac{9 g_{\text{A}}^2}{64 \pi f_{\pi}^2} m_{\pi}^3 . \]  

Its numerical value is \(-22 \text{ MeV}\), to be compared with the value of the full sigma term which is \( \sigma_N \simeq +45 - 50 \text{ MeV}\). We stress that such a model pion cloud calculation has been successfully used to extrapolate lattice data into the low quark mass, i.e., to the low pion mass sector [13]. Introducing the scalar density of nuclear pions, \( \langle \phi_{\pi}^2 \rangle \), the relative amount of restoration from the pion cloud is given to leading order in density by:

\[ \frac{\Sigma_N^{(\pi)} \rho}{f_{\pi}^2 m_{\pi}^2} = \frac{\langle \phi_{\pi}^2 \rangle}{2 f_{\pi}^2} . \]  

We now turn to the s-wave isoscalar pion self-energy for pions of zero three-momentum. To linear order in density it is simply given by the vacuum \( \pi N \) \( T \)-matrix multiplied by the density. The various contributions are depicted in Fig. 2. The pion loop contribution (fig. 1a), including vertex corrections not shown in Fig. 1a, is derived in Refs. [8, 12, 13].

It involves the pionic piece of the sigma term. As for the scalar exchange contribution (fig. 1b) it is related to the scalar piece of the sigma term. Fig. 1c is immediately obtained from the Lagrangian given in eq. (4) and fig. 1d is the Born term contribution with pseudo-vector coupling. The final result to leading order in density reads:

\[ \Pi(\omega) = \rho \left\{ - \frac{\sigma_N}{f_{\pi}^2} - \frac{4 \Sigma_N^{(\pi)}}{3 f_{\pi}^2} \right\} + \omega^2 \left( \frac{g_{\text{A}}^2}{4 M_N f_{\pi}^2} - \frac{c_2' + c_3' - \Sigma_N^{(\pi)} / m_{\pi}^2}{f_{\pi}^2} \right) - \frac{2 \beta}{3} \left( \omega^2 - m_{\pi}^2 \right) \Sigma_N^{(\pi) / f_{\pi}^2 m_{\pi}^2} \} . \]  

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Note that the pion self-energy is not directly an observable and it depends on the representation through the $\beta$ factor:

$$\beta = 1 + 10 \left( \alpha - \frac{1}{6} \right). \quad (13)$$

![Diagram](image.png)

**FIG. 1:** The various contributions to the s-wave pion self-energy. More explanations are given in the text.

### III. IN-MEDIUM MODIFICATION OF THE ISOVECTOR PION-NUCLEON AMPLITUDE

The energy dependence of the pion self-energy implies a wave-function renormalization of the in-medium pion, as noticed in several works [13, 14, 15, 16, 17]. The particular importance of the energy dependence for the interpretation of the deeply bound pionic states has recently been emphasized in Ref. [6]. The pion propagator has the form

$$D = \left( \omega^2 - m^2 - \Pi(\omega) \right)^{-1} = Z \left( \omega^2 - m^2 \right)^{-1}. \quad \text{From the expression for the self-energy, eq. (12), we obtain for the residue:}$$

$$Z = 1 + \left( \frac{g_A^2}{4M_N f^2_\pi} - 2 \frac{c'_2 + c'_3 - \Sigma^{(e)}_N}{f^2_\pi} - 2 \frac{3}{3} \frac{2 \beta}{f^2_\pi} \frac{\Sigma^{(\pi)}_N}{m^2_\pi} \right) \rho. \quad (14)$$

The value of the isoscalar scattering length, $b_0$, provides a relation between the different parameters entering this expression. For the nucleon it is given by:

$$4\pi \left( 1 + \frac{m_\pi}{M_N} \right) b_0 = -\frac{\sigma_N}{f^2_\pi} + \frac{4 \Sigma^{(\pi)}_N}{f^2_\pi} + \left( \frac{g_A^2}{4M_N f^2_\pi} - 2 \frac{c'_2 + c'_3 - \Sigma^{(e)}_N}{f^2_\pi} \right) m^2_\pi \simeq 0. \quad (15)$$

According to recent accurate data [18], this value is compatible with zero. This near-vanishing of $b_0$ (believed to be fortuitous) translates into the following expression for the residue:

$$Z \simeq 1 + \frac{\sigma_N \rho}{f^2_\pi m^2_\pi} - \frac{4 \Sigma^{(\pi)}_N \rho}{f^2_\pi m^2_\pi} - \frac{2 \beta}{f^2_\pi} \frac{\Sigma^{(\pi)}_N \rho}{m^2_\pi}. \quad (16)$$

The residue provides a first source for the renormalization of the isovector scattering length $b_1$. Notice that if the pion loop correction (the last two terms in eq. (16)) is ignored one recovers the original proposal of Ref. [3]. However this description would be incomplete since, as previously discussed, the pion cloud piece of the sigma term is certainly not negligible. In addition a description of the in-medium modification with the only influence of the residue is clearly unsatisfactory since the result then depends on the representation. Indeed there is another source of renormalization, which to our knowledge has previously been ignored. It arises from the pion loop correction of the isovector pion-nucleon amplitude, inherent to the non-linear realization. The isovector piece (Weinberg-Tomozawa) of the chiral Lagrangian is given by:

$$\mathcal{L}_{WT} = \frac{i}{2} \bar{N} \gamma^\mu \left( \xi \partial_\mu \xi^\dagger + \xi^\dagger \partial_\mu \xi \right) N \simeq -\frac{1}{4f^2_\pi} \bar{N} \gamma^\mu \left( 1 + 2 \left( \alpha - \frac{1}{24} \right) \frac{\sigma^2}{f^2_\pi} \right) \vec{\tau} \cdot (\vec{\phi} \times \partial_\mu \vec{\phi}) N, \quad (17)$$

where the second expression arises from an expansion to fourth order in the pion field. After the appropriate isospin averaging, one obtains an effective Weinberg-Tomozawa Lagrangian which is modified in the medium by one-pion...
loop correction (see fig. 2):

\[
\mathcal{L}_{WT}^{\text{eff}} = -\frac{1}{4 f_\pi^2} \left( 1 + \frac{10}{3} \left( \alpha - \frac{1}{24} \right) \frac{\langle \phi^2 \rangle}{f_\pi^2} \right) \bar{N} \gamma^\mu \vec{\tau} \cdot (\vec{\phi} \times \partial_\mu \vec{\phi}) N
\]

\[
= -\frac{1}{4 f_\pi^2} \left( 1 + \frac{20}{3} \left( \alpha - \frac{1}{24} \right) \frac{\Sigma_N^{(s)}(N)}{f_\pi^2 m_\pi^2} \right) \bar{N} \gamma^\mu \vec{\tau} \cdot (\vec{\phi} \times \partial_\mu \vec{\phi}) N. \quad (18)
\]

FIG. 2: Pion loop correction to the isovector pion-nucleon amplitude.

The isospin-antisymmetric amplitude is proportional to the pion energy \( q_0 \). In the vacuum \( q_0 = m_\pi \), and \( b_1 = -m_\pi/(4\pi f_\pi^2(1 + m_\pi/M_N)) \). For the quasi-particle in the medium the energy is the pion effective mass \( m_\pi^* \). Since we take \( b_0 \) to be zero, to first order in the density the effective and bare pion masses are equal. According to the previous discussion the in-medium renormalization of \( b_1 \) is then:

\[
\frac{b_1^*}{b_1} = Z \left( 1 + \frac{20}{3} \left( \alpha - \frac{1}{24} \right) \frac{\Sigma_N^{(s)}(N)}{f_\pi^2 m_\pi^2} \right),
\]

which gives:

\[
\frac{b_1^*}{b_1} = 1 + \frac{\sigma_N \rho}{f_\pi^2 m_\pi^2} - \frac{7}{6} \frac{\Sigma_N^{(s)}(N)}{f_\pi^2 m_\pi^2}.
\]

Notice that once the two influences are taken into account the representation dependence has disappeared, an important consistency check of our result. Coming to a numerical evaluation, the parameters of the model, in particular the sigma mass, can be fixed by a simultaneous fit to the saturation curve and to pion-pion phase shifts [19, 20]. A typical value is \( m_\sigma = 800 \text{ MeV} \), yielding \( \Sigma_N^{(s)} \simeq 28 \text{ MeV} \). Taking also \( \Sigma_N^{(\pi)} \simeq 25 \text{ MeV} \), we have:

\[
\frac{\Sigma_N^{(s)}}{f_\pi^2 m_\pi^2} \simeq 0.21 \frac{\rho}{\rho_0}, \quad \frac{\Sigma_N^{(\pi)}}{f_\pi^2 m_\pi^2} \simeq 0.18 \frac{\rho}{\rho_0}.
\]

This leads to the estimate:

\[
\frac{b_1^*}{b_1} \simeq 1 + 0.18 \frac{\rho}{\rho_0}.
\]

Even if the coefficient 0.18 may be slightly changed if the relative weight between the pion cloud and the scalar contribution to the sigma term is changed, the enhancement is more moderate than the original proposal [3] which ignores the pion loop correction (in that case the coefficient is of the order of 0.35 – 0.40 depending on the precise value of the sigma term). In our case this pion loop correction partly cancels the effect of the pion-nucleon sigma term.

**IV. COMPARISON WITH A CHIRAL PERTURBATION APPROACH**

The pion-self-energy has been calculated using a chiral perturbation approach. To leading order in density, it reads [6, 21]:

\[
\Pi(\omega) = \rho \left\{ -\left( \frac{\sigma_N}{f_\pi^2} - \frac{4}{3} \frac{\Sigma_N^{(s)}}{f_\pi^2} \right) + \omega^2 \left( \frac{g_A^2}{4 M_N f_\pi^2} - 2 \frac{c_2 + c_3}{f_\pi^2} \right) + \frac{\zeta}{3} \left( \omega^2 - m_\pi^2 \right) \frac{\Sigma_N^{(s)}}{f_\pi^2 m_\pi^2} \right\}. \quad (23)
\]
One can make a one-to-one correspondence between this CHIPT result and the result obtained in our model (eq. [22]). The first term ($\sigma_N$) and the third term (Born term) are identical in both expressions of the self-energy. The fourth terms in both expressions are also equivalent since the comparison between the two just defines the quantity $c_2 + c_3$ in the particular model that we use. As for the representation dependent coefficient $\zeta$ which appears in the last term of eq. [23] it has to be be identified with $-2\beta$. However there is an essential difference both at the conceptual and quantitative levels. The second term and the fifth one (representation dependent term) involve the LNAC piece of the pion loop, at variance with the model calculation where the full pion loop appears. We will return later to this difference and we study for the moment the consequences for the $b_1$ parameter. Following the same steps as in section III, with the constraint that the isoscalar scattering length vanishes, we obtain for the residue:

$$Z = 1 + \frac{\sigma_N \rho}{f_\pi^2 m_\pi^2} - \frac{4}{3} \frac{(\Sigma_N^{(\pi)})^{(LNAC)}}{f^2 m_\pi^2} - \frac{2}{3} \frac{(\Sigma_N^{(\pi)})^{(LNAC)}}{f_\pi^2 m_\pi^2}. \quad (24)$$

Taking, as in Ref. [1], $\zeta = 0$ (i.e. $\beta = 0$), the influence of the residue on the $b_1$ parameter becomes:

$$\left(\frac{b_1^*}{b_1}\right)^{(CHIPT1)} = Z = 1 + \frac{\sigma_N \rho}{f_\pi^2 m_\pi^2} - \frac{4}{3} \frac{(\Sigma_N^{(\pi)})^{(LNAC)}}{f^2 m_\pi^2} \rho \simeq 1 + 0.53 \frac{\rho}{\rho_0}, \quad (25)$$

which yields a larger enhancement than our model calculation. However this result is not representation independent. The pion-loop correction in the Weinberg-Tomozawa amplitude which is not taken into account in the CHIPT calculation should be added. If this is done in a consistent way one should retain only the LNAC piece of the pion loop in order to reach the representation independent following result:

$$\left(\frac{b_1^*}{b_1}\right)^{(CHIPT2)} = 1 + \frac{\sigma_N \rho}{f_\pi^2 m_\pi^2} - \frac{7}{6} \frac{(\Sigma_N^{(\pi)})^{(LNAC)}}{f^2 m_\pi^2} \rho \simeq 1 + 0.51 \frac{\rho}{\rho_0}. \quad (26)$$

This is numerically very close to what is obtained from the prescription used in Ref. [1] but still difficult to reconcile with our model result. The crucial difference is the fact that, in our model, the pion loop always appears as a whole and the LNAC can never appear as an isolated piece. In other words the model involves the combination $-4\zeta^1 m_\pi^2 + (\Sigma_N^{(\pi)})^{(LNAC)}$ rather than $(\Sigma_N^{(\pi)})^{(LNAC)}$ alone. The two quantities are of comparable magnitude but of opposite signs. The difference between the two approaches is associated with different off-mass shell behavior of the $\pi N$ amplitude. As a further test of our method, in the PCAC representation in which the relation $\partial_\mu \vec{A}^\mu = - f_\pi m_\pi^2 \vec{\phi}$ holds, the soft pion amplitude (obtained by taking $\omega = 0$) should be proportional to $\sigma_N$. In order to obtain the PCAC representation it is sufficient to send the sigma mass to infinity (the sigma term then has only the pionic contribution) and to take $\beta = 1$ (i.e., $\alpha = 1/6$). In that case one recovers at the soft pion point the venerable low-energy theorem since:

$$\left(\mathcal{M}(0, 0, 0, 0)\right)^{Model}_{PCAC} = - \frac{\sigma_N}{f_\pi} + 2 \frac{\Sigma_N^{(\pi)}}{f_\pi^2} = \frac{\sigma_N}{f_\pi^2}. \quad (27)$$

Instead with the prescription of eq. [23] for the CHIPT off-shell amplitude, one gets:

$$\left(\mathcal{M}(0, 0, 0, 0)\right)^{CHIPT}_{PCAC} = - \frac{\sigma_N}{f_\pi} + 2 \frac{\Sigma_N^{(\pi)}}{f_\pi^2} \quad (28)$$

The two terms on the r.h.s. being negative cannot cooperate to reproduce the positive sigma term, in conflict with the soft pion theorem.

V. REMARKS AND CONCLUSION

We have studied the in-medium renormalization of the isovector amplitude in the linear sigma model. We use it in a non-linearized representation in which the radius of the chiral circle is not frozen, which preserves the sigma degree of freedom. In this approach the nucleon sigma commutator is built of a scalar meson exchange with the condensate and a two-pion exchange one, the pion cloud contribution. We find, as in our previous works on other in-medium quantities, that the independence on the representation is achieved through a combination of two effects: i) the residue of the pion propagator arising from the energy dependence of the pion self-energy, ii) the influence of the pion loops on the Weinberg-Tomozawa amplitude (which must also apply to the chiral perturbation approach of Refs. [6, 21] so as
to make it independent on the representation). Our description results in a more moderate renormalization of $b_1$: at normal nuclear density, an enhancement by 18%, which is less than the prediction of Ref. [3] or the one which can be derived from the results of Ref. [1,2]. The difference between our model and the CHIPT approaches lies in the role of the pion loops. Our treatment for these is identical to the one employed in Ref. [1] for the pion gas in a heat bath. The results of Ref. [2], when taken in the chiral limit, reproduce the chiral perturbation expansion up to second order. The only difference between the present calculation and the description of the pion gas in a heat bath amounts to the (physically natural) replacement of the scalar density of thermally excited pions by that of the nuclear pions, a positive quantity. On the contrary, in the chiral approach of Refs. [6,21] only the leading non-analytical term appears in the self-energy (beyond what is implicitly contained in the nucleon sigma commutator). Since $(\Sigma_N^{(\pi)})^{(LNAC)}$ is negative and comparatively large, the numerical consequences which follow explain the difference in the renormalization factor. It will be interesting to exactly clarify why, in the chiral approach, $(\Sigma_N^{(\pi)})^{(LNAC)}$, which is in fact a piece of the pion density, dissociates from the rest, while the whole pion density (with its simple physical interpretation) naturally enters in our approach. The nucleon size, which is crucial in any model evaluation of the nuclear pion loop, introduces a scale which does not seem to be fully incorporated in this CHIPT approach. The model calculations of the pion loop which explicitly takes into account the nucleon size (i.e., in practice the form factor) have proven their relevance and usefulness. For instance they give an excellent fit to the quark mass dependence of the lattice results allowing a convincing extrapolation to the physical mass region [11].

The actual fits of recent data on deeply bound pionic states suggest a larger enhancement of $b_1$ than our prediction [1,2,3]. In this context we want to point out that the comparison between our result and the the fit to pionic atom data is not direct. First of all we consider only strong interactions and ignore effects due to Coulomb forces, while these are crucial to bind the pion within the pionic atoms on which the fit is based. They are discussed in Ref. [1]. Moreover not all renormalization effects are included in our approach. For instance, we did not consider pion rescattering on two correlated nucleons. It is known to provide a sizeable part of the repulsive isospin-symmetric potential. Since it is not explicitly introduced in the fits to data for the charge exchange potential, it has to be included as a medium renormalization of $b_1$. This effect can be evaluated from the formulae given in Ref. [22]:

$$\delta b_1^{LPS} = -(1 + \frac{m_\pi}{M_N})b_1(b_0 + b_1)\frac{3}{2\pi}p_F,$$

(29)

where $p_F$ is the Fermi momentum of the nucleons. At normal density, it amounts to a decrease of $b_1$ by $\simeq 5\%$, i.e., it does not help to explain a larger enhancement. It would be interesting to attempt to fit the pionic atom data with a more moderate enhancement of the $b_1$ parameter. If the larger enhancement is confirmed it would be an incentive to push the investigations further. One might consider the influence of higher order effects in the density, which could upset the nearly total cancellations present in the isospin symmetric scattering length. For instance, the pionic piece of the sigma commutator should be replaced by the fully dressed one in which the pion line is dressed by nucleon-hole and Δ-hole excitations screened by short range correlations. At normal nuclear density the effect leads to a moderate enhancement (of the order of 25\%) of $\langle\phi^2\rangle$ [23]. The many-body effects in the other contributions should be similarly investigated, including those concerning the phenomenological parameters $c_2'$ and $c_4'$, which would require a model for these quantities. The question of the in-medium modification of the charge exchange $\pi N$ amplitude first raised by Weise is indeed quite challenging on the theoretical side as well as on the experimental one. It is likely that more physics remains to be understood in this problem.

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