Chiral symmetry restoration and parity mixing

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Abstract

We derive the expressions of the vector and axial current from a chiral Lagrangian restricted to nucleons and pions. They display mixing terms between the axial and vector currents. We study the modifications in the nuclear medium of the coupling constants of the axial current, namely the pion decay constant and the nucleonic axial one due to the requirements of chiral symmetry. We express the renormalizations in terms of the local scalar pion density. The latter also governs the quark condensate evolution and we discuss the link between this evolution and the renormalizations. In the case of the nucleon axial coupling constant this renormalization corresponds to a new type of pion exchange currents, with two exchanged pions. We give an estimate for the resulting quenching. Although moderate it helps explaining the quenching experimentally observed.

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Introduction

The problem of the restoration in dense or hot matter of the chiral symmetry of the strong interactions, which is spontaneously violated in the QCD vacuum has been extensively addressed. The interest has largely focused on the quark condensate, considered as the order parameter. For independent particles the evolution of the quark condensate with density or temperature is governed by the sigma commutator of the particles present in the system with the simple following expression:

$$\frac{<\overline{q}q(\rho)>}{<\overline{q}q(0)>} = 1 - \sum_n \frac{\rho_n^s\, \Sigma_n}{f_\pi^2 m_\pi^2} \tag{1}$$

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where the sum extends over the species present in the medium, ρ^s is their scalar density and Σ their sigma commutator. Pions play a crucial role in the restoration process especially in the heat bath where they enter as the lightest particles created by the thermal fluctuations. In the nuclear medium the main ingredients are the nucleons, with some corrections from the exchanged pions. At normal nuclear density the magnitude of the condensate has dropped by about 1/3, a large amount of restoration. It is essentially the effect of the nucleons adding their effects independently, the corrections due to the interaction being small. Such a large amount of restoration raises the question about manifestations directly linked to the symmetry. If there is no spontaneous violation of the symmetry, *i.e.*, if it is realized in the Wigner mode, the hadron masses vanish or there exist parity doublets, each hadronic state being degenerate with its chiral partner. It is therefore legitimate to believe that the large amount of restoration at normal density manifests itself either by a decrease of the hadron masses, or by a mixing between opposite parities. A link between the evolution of the hadron masses and the amount of restoration has been suggested [2]. But it cannot be a straightforward one. Indeed the density or temperature evolution of the masses cannot have a direct relation to that of the condensate , as follows from the works of several authors [3, 4, 5]. On the other hand the significance of chiral symmetry restoration for the parity mixing was first established by Dey et al. [6] for the thermal case. They showed that in a pion gas a mixing occurs between the vector and axial correlators. It arises from the emission or absorption of s-wave thermal pions, which changes the parity of the system. The mixing goes along with a quenching effect of the correlators, which, to first order in the pion density, equals 4/3 of the quenching of the quark condensate. These points were also made by Steele et al. [7]. The extension of the formalism of Dey et al. to finite densities has been attempted by Krippa [8].

The aim of this work is the discussion of the implications of chiral symmetry restoration in the nuclear medium, in a world restricted here to nucleons and pions. The only transitions allowed in the nucleus are then nuclear transitions or pion production. We give the explicit expressions of the axial and vector current in a formalism based on chiral lagrangians. We will show that the nuclear pions renormalize the coupling constants of the axial current and that the renormalization can be expressed in terms of the pion scalar density. This last quantity also enters in the quark condensate evolution. However the complexity of the nuclear interactions bars a simple link between this evolution, which is an average concept, and the renormalizations. For instance, for the axial coupling constant g_A the detailed spatial structure of the pion scalar density is needed.

Our article is organized as follows. In section 1 we derive the expressions of the axial and vector currents from the chiral lagrangians. In section 2 we use these expressions to study the renormalization of the pion decay constant in the hot pion gas and in the nuclear medium. The thermal case is only introduced as an illustration of the method since the results are already known. In section 3 we apply the same technique to the axial coupling constant. To account for the nucleon-nucleon correlations we express the renormalization in the traditional treatment by the meson exchange currents. We give an estimate of the quenching of the axial coupling constant. We also discuss the renormalization of the Kroll-Ruderman matrix element of pion photoproduction.

1 The Lagrangian and the currents

Our starting point is the chiral Lagrangian in the form introduced by Weinberg. We use, as in our previous work of ref. [9], the version of Lynn [10], which allows one to obtain the nucleon sigma commutator in the tree approximation. The Lagrangian writes:

$$\begin{aligned}\mathcal{L} &= -\frac{1}{2}m_\pi^2\frac{\boldsymbol{\phi}^2}{1+\boldsymbol{\phi}^2/4f_\pi^2}+\frac{1}{2}\frac{\partial_\mu\boldsymbol{\phi}\cdot\partial^\mu\boldsymbol{\phi}}{(1+\boldsymbol{\phi}^2/4f_\pi^2)^2}\\ &\quad +2\sigma_N\overline{\psi}\psi\frac{\boldsymbol{\phi}^2/4f_\pi^2}{1+\boldsymbol{\phi}^2/4f_\pi^2}+\overline{\psi}(i\gamma_\mu\partial^\mu-M)\psi\\ &\quad -\frac{1}{4f_\pi^2}\frac{\overline{\psi}\gamma_\mu(\boldsymbol{\tau}\times\boldsymbol{\phi})\cdot\partial^\mu\boldsymbol{\phi}\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}+\frac{g_A}{2f_\pi}\frac{\overline{\psi}\gamma_\mu\gamma_5\boldsymbol{\tau}\cdot\partial^\mu\boldsymbol{\phi}\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}\ .\end{aligned}\tag{2}$$

We have to specify the quantity σ_N associated with the nucleon density in eq. (2). The free nucleon sigma commutator Σ_N cannot be entirely attributed to the pion cloud. We define σ_N to be the difference between the total and pionic contributions:

$$\Sigma_N=\sigma_N+\frac{1}{2}m_\pi^2\int d\boldsymbol{x}\langle N|\boldsymbol{\phi}^2(\boldsymbol{x})|N\rangle.\tag{3}$$

For instance in a description of the nucleon in terms of valence quarks and pions, the pionic contribution is approximatively 1/2 to 2/3 of the total value [11, 12].

From the Lagrangian of eq. (2) we derive the expressions of the axial and isovector vector currents:

$$\begin{aligned}\boldsymbol{\mathcal{A}}_\mu &= f_\pi\frac{\partial_\mu\boldsymbol{\phi}}{1+\boldsymbol{\phi}^2/4f_\pi^2}-\frac{1}{2f_\pi}\frac{[(\boldsymbol{\phi}\times\partial_\mu\boldsymbol{\phi})\times\boldsymbol{\phi}]}{(1+\boldsymbol{\phi}^2/4f_\pi^2)^2}\\ &\quad +\frac{g_A}{2}\overline{\psi}\gamma_\mu\gamma_5\boldsymbol{\tau}\psi+\frac{g_A}{4f_\pi^2}\frac{\overline{\psi}\gamma_\mu\gamma_5[(\boldsymbol{\tau}\times\boldsymbol{\phi})\times\boldsymbol{\phi}]\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}-\frac{1}{2f_\pi}\frac{\overline{\psi}\gamma_\mu(\boldsymbol{\tau}\times\boldsymbol{\phi})\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}\end{aligned}\tag{4}$$

$$\begin{aligned}\boldsymbol{\mathcal{V}}_\mu &= \frac{(\boldsymbol{\phi}\times\partial_\mu\boldsymbol{\phi})}{(1+\boldsymbol{\phi}^2/4f_\pi^2)^2}\\ &\quad +\frac{1}{2}\overline{\psi}\gamma_\mu\boldsymbol{\tau}\psi+\frac{1}{4f_\pi^2}\frac{\overline{\psi}\gamma_\mu[(\boldsymbol{\tau}\times\boldsymbol{\phi})\times\boldsymbol{\phi}]\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}-\frac{g_A}{2f_\pi}\frac{\overline{\psi}\gamma_\mu\gamma_5(\boldsymbol{\tau}\times\boldsymbol{\phi})\psi}{1+\boldsymbol{\phi}^2/4f_\pi^2}\ .\end{aligned}\tag{5}$$

The conservation law of the vector current can be shown, using the equations of motion for the nucleon and the pion fields. The divergence of the axial current instead satisfies the following relation:

$$\partial^\mu \boldsymbol{\mathcal{A}}_\mu = -f_\pi m_\pi^2 \, \frac{\boldsymbol{\phi}}{1+\boldsymbol{\phi}^2/4f_\pi^2} \, (1 - \sigma_N \frac{\overline{\psi}\psi}{f_\pi^2 m_\pi^2}) \, . \tag{6}$$

Some comments on the expressions (4) and (5) are in order. Let us first discuss the free case. We recognize in some of the terms the usual expressions for the vector or axial current coupled to a free nucleon or pion. In addition the axial current can create one or more pions, either in free space (first terms of eq. (4)), or when it acts on the nucleon via a term (last one of eq. (4)) which is the equivalent for the axial current of the Weinberg-Tomozawa term of π-N scattering. Similarly the vector current acting on the nucleon can create one (or more) pion via the Kroll-Ruderman term, *i.e.* the contact piece of photoproduction (last term of eq. (5)).

Let us now turn to the case of a hadronic medium. The expressions (4) and (5) illustrate in a striking fashion the way in which the axial and vector current mixing occurs. Indeed, in the heat bath any of the pions can be a thermal one. As an example, consider the Kroll-Ruderman term of the vector current, ignoring at this stage the denominator. The creation or annihilation of a thermal pion of momentum q in this term takes care of the pion field, leaving a factor $e^{\pm iqx}$ and we are left with a current of opposite parity, to be taken at the momentum transfer $k \pm q$ where k is the photon momentum, as in the formalism of ref. [7]. Similarly the pion production or annihilation by the Weinberg term of the axial current introduces the vector current nuclear matrix element. To the extent that the Weinberg term is mediated by the rho meson and the Kroll-Ruderman one by the A_1 meson, these expressions include the effects, at low momenta, of the $\rho - A_1$ mixing.

It is interesting to observe on expressions (4) and (5) that the Kroll-Ruderman term itself can be obtained from the fourth term of the axial current by suppression of one of the pion fields (representing creation or annihilation of a thermal pion). Thus the three terms containing g_A in eqs (4) and (5) are linked together by suppression or addition of one pion field. The same is true for the three purely pionic terms and for the three terms in γ_μ as well. Thus a grouping three by three of the various terms naturally emerges from our expressions.

In the nuclear medium the virtual pions can be seen as a pion bath and similar considerations about the mixing might apply. However the pions do not come from an external reservoir but fully belong to the nucleus. Strictly speaking there is no mixing. However the mixing terms of the currents can pick a pion from the cloud of a nucleon introducing a similarity with the heat bath as displayed in fig. 1 in the case of the Kroll-Ruderman term. The corresponding process is the excitation of high lying nuclear states (2p-2h). In the case of the third isospin component of the current, it is part of the well known quasi-deuteron photoabsorption cross-section. Another example is the influence of the Weinberg-Tomozawa term on the time part of the axial current, which enters via the Pauli correlations and sizeably increases the time-like axial coupling constant [13]. The present approach puts these effects, where the mixing terms of the currents pick a pion from the cloud, in a perspective linked to chiral symmetry.

The mixing goes along with a renormalization of certain coupling constants, such as the axial one, that will now be discussed.

2 The pion decay constant

We start with the case of the hot pion gas. This is meant as an illustration of our method as no new result is reached. It serves to introduce quantities such as the residue γ (*i.e.*, the wave function renormalization) that will be used later. The production of a pion by the axial current is governed by the first two terms of the expression (4). Limiting the expansion to first order in the squared pion field we obtain:

$$\boldsymbol{\mathcal{A}}_\mu = f_\pi \partial_\mu \boldsymbol{\phi}\,(1 - \frac{7}{12}\frac{\boldsymbol{\phi}^2}{f_\pi^2})\;. \tag{7}$$

The pion field is expanded in terms of creation and annihilation operators B and $B^\dagger$ for a quasi-pion in the medium:

$$\boldsymbol{\phi}(x) = \gamma^{1/2}\int \frac{d\boldsymbol{k}}{(2\pi)^3}\frac{1}{(2\omega_k^*)^{1/2}}(\boldsymbol{B}_{\boldsymbol{k}} + \boldsymbol{B}_{-\boldsymbol{k}}{}^\dagger)e^{i(\boldsymbol{k}\cdot\boldsymbol{x} - \omega_k^* t)}\;, \tag{8}$$

where ω_k^* is the energy of a quasi-pion of momentum $\mathbf{k}$, $\omega_k^* = \sqrt{\boldsymbol{k}^2 + m_\pi^{*2}}$ with m_π^* the effective pion mass. The quantity γ is the residue of the pion pole. Since the derivative of the pion field gives no contribution when it acts on the pions of the bath, the matrix element for production of a quasi-pion by the axial current reduces to:

$$\langle 0|\boldsymbol{\mathcal{A}}_\mu(0)|\tilde{\pi}\rangle = \frac{\gamma^{1/2}}{(2\omega_k^*)^{1/2}} i f_\pi k_\mu (1 - \frac{7}{12}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle) = \frac{1}{(2\omega_k^*)^{1/2}} i f_\pi^* k_\mu \tag{9}$$

where the second equation defines the renormalized pion decay constant f_π^*. In a pion gas the residue γ has been derived by Chanfray et al. [14]. To first order in the quantity $\boldsymbol{\phi}^2$, equivalently the pion density, it writes, in the Weinberg representation:

$$\gamma = (1 - \frac{1}{2}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle)^{-1}\;. \tag{10}$$

The renormalized pion decay constant then reads:

$$f_\pi^* = f_\pi \gamma^{1/2}(1 - \frac{7}{12}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle) \approx f_\pi(1 - \frac{1}{3}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle)\;. \tag{11}$$

On the other hand, the temperature evolution of the condensate in a hot pion gas is, to first order in the quantity $\boldsymbol{\phi}^2$, as given in ref. [14]:

$$\frac{<\overline{q}q>_T}{<\overline{q}q>_0} = 1 - \frac{1}{2}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle_T \ . \tag{12}$$

Thus to first order in $\langle\boldsymbol{\phi}^2\rangle$ the renormalization of f_π follows the evolution of the condensate but with the coefficient 2/3, in agreement with chiral perturbation results and other works [6, 14, 15, 16]. Note that this renormalization applies to both space and time components of the axial current. This agrees with the findings of ref [17] where it is shown that to order T^2 Lorentz invariance is preserved.

We now turn to the dense medium. Formally we can follow the same procedure. The presence in the nuclear medium of a pion scalar density, in the form of an expectation value of the quantity $\boldsymbol{\phi}^2$, renormalizes the pion decay constant. Formally the expression is the same as previously, $f_\pi^* = f_\pi\gamma^{1/2}(1-\frac{7}{12}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle)$. If we treat the nuclear medium as a pion gas, the residue γ entirely arises from π-π interactions and is the same as given previously in eq. (10). In this simplified treatment f_π^* is given by eq. (11). It is linked to the pion scalar density, *i.e.* the expectation value $\boldsymbol{\phi}^2$. Even with the simple form (11), the renormalization of f_π does not follow 2/3 of the condensate one. The reason is that the condensate evolution in the nuclear medium is governed by the full nucleon sigma commutator Σ_N, which is not entirely due to $\boldsymbol{\phi}^2$. There exists also the non pionic contribution embodied in σ_N, as discussed previously. Thus the two renormalizations do not follow each other. This result is general and applies as well to the axial coupling constant g_A.

This is not the only restriction which prevents a simple link to the condensate in the nuclear medium. The residue γ itself is not entirely due to π-π scattering. There exist other sources for the energy dependence of the s-wave π-N interaction, such as the Δ excitation. The medium renormalization of f_π cannot be written in the simple form (11). This illustrates the complexity of the dense medium as compared to the hot pion gas. A more phenomenological approach has been followed by Chanfray et al. [18] who linked the in-medium pion decay constant through the nuclear Gell-Mann-Oakes-Renner relation to the evolution of the pion mass, itself obtained empirically from the s-wave pion-nucleus optical potential.

3 The axial coupling constant

We now turn to the axial coupling constant. Its renormalization is governed by the fourth term of eq. (4). After rearrangement with the Gamow-Teller current (third term), we get:

$$\frac{1}{2}g_A\overline{\psi}\gamma_\mu\gamma_5(\boldsymbol{\tau} + \frac{1}{2f_\pi^2}\frac{\boldsymbol{\phi\tau\cdot\phi - \tau\phi^2}}{1+\boldsymbol{\phi}^2/4f_\pi^2})\psi = \frac{1}{2}g_A\overline{\psi}\gamma_\mu\gamma_5\boldsymbol{\tau}\ \psi(1-\frac{1}{3}\langle\frac{\boldsymbol{\phi}^2/f_\pi^2}{1+\boldsymbol{\phi}^2/4f_\pi^2}\rangle_T) \ , \tag{13}$$

where on the right hand side the average is taken over the heat bath. On the other hand the condensate is obtained from the chiral symmetry breaking Lagrangian $\mathcal{L}_{sb} = -\frac{1}{2} m_\pi^2 \boldsymbol{\phi}^2/(1+\boldsymbol{\phi}^2/4f_\pi^2)$. Therefore the condensate evolution follows [14]:

$$\frac{<\overline{q}q>_{T,\rho}}{<\overline{q}q>_0} - 1 = -\frac{1}{2}\langle\frac{\boldsymbol{\phi}^2/f_\pi^2}{1+\boldsymbol{\phi}^2/4f_\pi^2}\rangle_{T,\rho} \,. \tag{14}$$

Hence the axial coupling constant renormalized by the pion loops (fig. 2a) can be written:

$$g_A^*/g_A = (1 - \frac{2}{3}\frac{<\overline{q}q>_T}{<\overline{q}q>_0}) \,. \tag{15}$$

Thus with this chiral Lagrangian, in a hot medium the axial coupling constant follows, to all orders in the pion density, 2/3 of the quark condensate evolution (as long as it is pion dominated). The factor 2/3 is easily understood here: only two charges out of three contribute to the renormalization while all three charge states participate in the condensate evolution. The quenching of g_A is in agreement with the universal behaviour of ref. [6] and with the former result of ref. [19]. We have checked the expected independence of our results on the particular representation of the non-linear Lagrangian.

We now turn to the case of finite density. The starting expression is the same as the left hand side of eq. (13). In the nuclear medium the pions originate from the other nucleons so that the nucleon-nucleon correlations cannot be ignored. Here it is useful to make the link between this renormalization and the traditional picture of meson exchange currents. We keep only the two-body terms which are the dominant ones and work to lowest order in the pion field. The corresponding graph is that of fig. 2b. This type of exchange graph with two pions is not usually considered in nuclear physics. It is dictated to us only by these chiral symmetry considerations. We have to express the triangular graph of the figure as an effective two-body operator to be evaluated between correlated two-nucleon wave functions. A simplification occurs in the static approximation where the pions do not transfer energy to the nucleon line. We are left with an integral over the squared pion propagator with leads to a simple form in x-space for the two-body operator:

$$O_{12} = -\frac{1}{6f_\pi^2} g_A(\gamma_\mu\gamma_5)_1(\boldsymbol{\tau_1} - \tfrac{\boldsymbol{i}}{\boldsymbol{2}}(\boldsymbol{\tau_1}\times\boldsymbol{\tau_2}))\varphi^2(\boldsymbol{x_1},\boldsymbol{x_2}) \,, \tag{16}$$

where $\varphi(\boldsymbol{x_1},\boldsymbol{x_2})$ is the Yukawa field, taken at the point $\boldsymbol{x_1}$, emitted by the nucleon located at the point $\boldsymbol{x_2}$ and we have made explicit the dependence in the isospin operator $\boldsymbol{\tau_2}$ of the emitting nucleon. The operator O_{12} has a direct and an exchange contribution. The latter contribution vanishes for the second piece of the two-body operator in the limit of zero momentum current. We will furthermore ignore the exchange term of the first piece (*i.e.* in $\boldsymbol{\tau_1}$) and consider only the short range correlations. We now focus on the direct terms and specialize to the charged currents. The isospin factors in eq. (16) reduce to the expression $3\tau_1^\pm - 2\tau_1^\pm\tau_2^\pm\tau_2^\mp = 2\tau_1^\pm(1\mp\tau_2^0/2)$.

The resulting contributions depend on the relative number of protons and neutrons. In symmetric nuclear matter where they are equal, the factor, once summed over all the pion emitters (the nucleons with index 2), gives $2\tau_1^{\pm}$ multiplied by the nuclear density ρ. In the neutron gas instead, depending whether we consider neutron decay (the + component) or proton decay (the − one), we would get a factor $3\tau_1^+$ or $1\tau_1^-$ multiplied by the neutron density ρ_n. We can summarize these results by introducing an effective density, which depends on the charge of the current and on the neutron excess number:

$$\rho_{eff}^{+} = \frac{3N+Z}{2A}\rho \qquad \rho_{eff}^{-} = \frac{3Z+N}{2A}\rho \tag{17}$$

from which we recover the previous results. Sandwiching the whole operator O_{12} between two-nucleon wave functions, we thus obtain:

$$\delta g_A^{ex}/g_A = -\frac{1}{3f_\pi^2}\int d\boldsymbol{x_2}\rho_{eff}^{\pm}(x_2)[1+G(\boldsymbol{x_1},\boldsymbol{x_2}\ )]\varphi^2(\boldsymbol{x_1},\boldsymbol{x_2})\ , \tag{18}$$

where $G(\boldsymbol{x_1},\boldsymbol{x_2})$ is the short range nucleon-nucleon correlation function. In symmetric nuclear matter $\rho_{eff} = \rho$ whereas a similar formula would hold in the neutron gas, with the obvious replacement of ρ by the neutron density ρ_n and of the factor 1/3 in front by 1/2 and 1/6 for neutron and proton decay respectively. As is apparent on the expression (18) it is not the full pion field squared which acts in the renormalization of g_A, but only the part which extends beyond the range of the correlation hole. No such distinction occurred for the pion decay constant since the pion produced by the axial current can be anywhere in the nucleus. Thus the universality of the quenching which exists in the heat bath is lost.

In order to obtain an estimate for g_A^* in symmetric matter, we assume a total exclusion of other nucleons in a sphere of radius $r_0 = 0.6fm$. In order to facilitate the comparison of the quenching effect of g_A to that of the condensate which is governed by the nucleon sigma term, we introduce a quantity $(\Sigma_N)_{eff}$:

$$(\Sigma_N)_{eff} = \frac{1}{2}m_\pi^2\int d\boldsymbol{x}\theta(x-r_0)\boldsymbol{\varphi^2(x)}\ . \tag{19}$$

Numerically for point-like pion emitters, we find an effective value $(\Sigma_N)_{eff} \approx 21MeV$. It is interesting to compare this value with a model calculation in the quark picture. We have used the results of Wakamatsu [20] in a chiral soliton model. The quantity $\boldsymbol{\varphi}^2(\boldsymbol{x})$ is replaced by the sea quark density distribution according to: $\frac{1}{2}m_\pi^2\boldsymbol{\varphi^2(x)} \to 2m_q\bar{q}q\boldsymbol{(x)}$. This gives a very similar value $(\Sigma_N)_{eff} \approx 19MeV$.

However these numbers do not include the Pauli blocking effect which removes the occupied states in the process of pion emission. This effect has been calculated in refs. [22, 1] but for the whole space integral (*i.e.* without a cut-off) of the quantity $\boldsymbol{\phi}^2$. Expressed in terms of a modification of the sigma commutator it amounts to a reduction $(\Delta\Sigma_N)_{Pauli} = -2.6MeV$. The blocking effect, which is moderate, should be even less pronounced with the cut-off. We ignore it in the following.

Coming back to the renormalized axial coupling constant, we have:

$$g_A^*/g_A = 1 - \frac{2}{3}\frac{\rho(\Sigma_N)_{eff}}{f_\pi^2 m_\pi^2} \ . \tag{20}$$

This represents a 10% quenching at normal nuclear density in symmetric matter (15% for neutron decay in a neutron gas of the same density), while the condensate has dropped by 35%. Notice that the evolution of g_A is sizeably slower. This quenching applies to all the components, space or time, of the axial current. Other renormalization effects have to be added. They are known to act differently on the different components. For instance the Weinberg-Tomozawa term acts on the time component alone, producing a sizeable enhancement [13]. In the case of the space component the nucleon polarization under the influence of the pion field $N \to \Delta$ leads to the Lorentz-Lorenz quenching [23]. In the latter case the two renormalizations go in the same direction of a quenching. The extra reduction that we have introduced in this work could help to explain the large amount of quenching observed in Gamow-Teller transitions. To get an idea, we fictitiously translate the reduction by chiral symmetry into an equivalent Lorentz-Lorenz effect. We introduce an effective Landau-Migdal parameter $\delta g'_{N\Delta}$, to be added to the genuine one, so as to reproduce the 10% quenching. This corresponds to an increase $\delta g'_{N\Delta} \approx 0.16$, a significant increase. Indeed the quenching of the Gamow-Teller sum rule requires, if all attributed to the Lorentz-Lorenz effect, $g'_{N\Delta}$ to be as big as 0.6 - 0.7 while the favoured theoretical value is around 0.4 [21]. Hence the chiral induced quenching would help to fill the gap.

Closely related to the Gamow-Teller transition is the pion photoproduction at threshold through the Kroll-Ruderman term. To lowest order the nuclear transition is governed by the axial current. We want now to discuss how it is renormalized in the medium following chiral symmetry requirements. Expanding to first order in $\boldsymbol{\phi}^2/f_\pi^2$ and applying Wick theorem, the relevant current writes:

$$(\boldsymbol{\mathcal{V}}_\mu)_{KR} = -\frac{g_A}{2f_\pi}\overline{\psi}\gamma_\mu\gamma_5(\boldsymbol{\tau}\times\boldsymbol{\phi})\psi(1 - \frac{5}{12}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle) \ . \tag{21}$$

For the production of a quasi-pion in the medium, the renormalization r_{KR} of the amplitude involves again the residue γ:

$$r_{KR} = \gamma^{1/2}(1 - \frac{5}{12}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle) \ . \tag{22}$$

For illustrating the complexity of the situation we first assume that the residue is entirely given by $\pi - \pi$ scattering and take the value of eq. (10). Moreover we ignore the correlation complications. We obtain then:

$$r_{KR} = (1 - \frac{1}{6}\langle\frac{\boldsymbol{\phi}^2}{f_\pi^2}\rangle) \ , \tag{23}$$

which is 1/3 of the variation of the condensate, in contradistinction to the axial transitions where the factor is 2/3. This result does not contradict the general expressions of Dey et al. [6] as the Kroll-Ruderman term represents already a mixing of the axial current into the vector one. This reduction factor could apply to other mixing amplitudes, but we have not established it. The evolution as 1/3 of the condensate one would apply in the hot pion gas situation. In the nuclear medium all the complications mentioned previously occur: the role of the correlations, the link between the condensate evolution and the expectation value of $\boldsymbol{\phi}^2$ and the problem with the residue γ. This case cumulates all of the difficulties of the dense medium. In all instances the overall renormalization of the Kroll-Ruderman matrix element in the nuclear medium should be small.

4 Conclusion

In conclusion we have investigated the behaviour of the nuclear medium in relation with chiral symmetry restoration. We have focused on the extension of the parity mixing concept between the axial and vector correlators, which exists in the hot pion gas. In the nuclear medium there is no mixing *stricto sensu*. Indeed the pions, which induce the mixing, are not part of an external system, as in the thermal case, but they belong to the virtual pion cloud which is an integral part of the nucleus. We have shown that nevertheless certain consequences of the mixing survive. The nucleus behaves in certain respects as a pion reservoir. The virtual pion emitted by a nucleon acts, as illustrated in fig. 3, on the remainder of the nucleus, *i.e.* on the system of (A-1) nucleons, as the pion of the heat bath. The mixing which does not exist at the level of the whole nucleus is present only at the sublevel of the (A-1) nucleon system. This translates by the fact that, in the "mixing" cross-sections (such as the quasi-deuteron photoabsorption one), at least one nucleon has to be ejected: the emitter or absorber of the pion. As for the heat bath, this pion reservoir produces a quenching of the axial coupling constants. Since the pion originates from a neighbouring nucleon this renormalization is nothing else than a meson exchange contribution. It involves the exchange of two pions and has not been so far considered, to our knowledge. We have expressed the renormalizations in terms of the pion scalar density. The same quantity also enters in the quark condensate evolution. One can therefore think of a link between the two quantities, as occurs in the heat bath where the link is simple. There is however an important difference of the nuclear medium with respect to the heat bath: the renormalizations are not described by a universal quenching factor expressed in terms of the average squared pion field. The nucleonic observables such as g_A are renormalized differently due the sensitivity to nucleon-nucleon short range correlations. In this case only that part of the pionic field which is beyond the correlation hole enters in the renormalization. This prevents the link to the condensate evolution which instead involves the average scalar density. This is an illustration of the point made by T. Ericson [24] about the possible importance of the spatial fluctuations of the condensate. We have given an estimate for the quenching of the axial coupling constant arising from the requirements of chiral symmetry. Although it is not very large (about 10%), this additional quenching is significant and may help explain the large observed quenching of the Gamow-Teller sum rule. We have also discussed the photoproduction amplitude arising from the Kroll-Ruderman term. It represents a mixing term of the axial current into the vector one. We have shown that its evolution is slower than the axial coupling constant one.

This work can be extended to enlarge the space. The first step is to include the Delta excitation. Another extension concerns the explicit introduction of the rho and the A_1 mesons, which the mixing of the axial and vector correlators allows to be excited either by the vector or by the axial current.

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Figure captions:

Fig.1: Illustration of a mixing effect in the vector correlator
(a) in the heat bath (denoted by a cross), (b) equivalent diagram in the nucleus with its translation (c) in many-body diagrams.

Fig.2: Renormalization of the nucleonic axial coupling constant
(a) - by a pion loop in the hot pion gas (the cross denotes the heat bath),
(b) - by the virtual pion cloud in the nucleus.

Fig.3: Illustration of a mixing effect in the nucleus by the Kroll-Ruderman term.

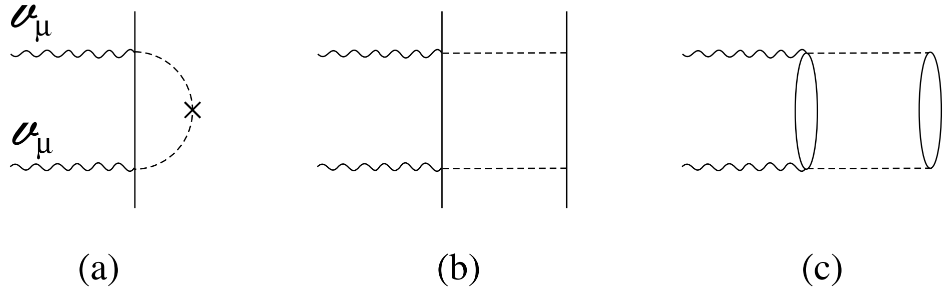

Figure 1: Illustration of a mixing effect in the vector correlator (a) in the heat bath (denoted by a cross), (b) equivalent diagram in the nucleus with its translation (c) in many-body diagrams.

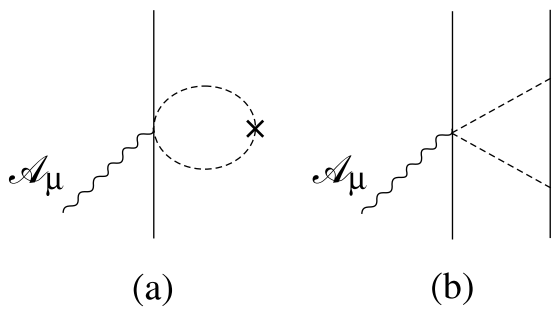

Figure 2: Renormalization of the nucleonic axial coupling constant (a) - by a pion loop in the hot pion gas (the cross denotes the heat bath), (b) - by the virtual pion cloud in the nucleus.

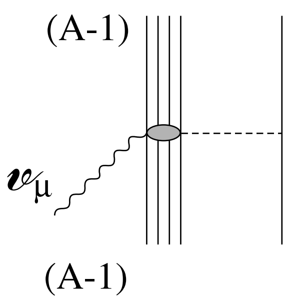

Figure 3: Illustration of a mixing effect in the nucleus by the Kroll-Ruderman term.