Double Decimation and Sliding Vacua in the Nuclear Many-Body System

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Abstract

We propose that effective field theories for nuclei and nuclear matter comprise of “double decimation”: (1) the chiral symmetry decimation (CSD) and (2) Fermi liquid decimation (FLD). The Brown-Rho scaling identified as the parametric dependence intrinsic in the “vector manifestation” of hidden local symmetry theory of Harada and Yamawaki results from the first decimation, i.e., CSD. In a recent work we showed that in matter under conditions of high density or high temperature, dynamically generated hadron masses scaled with a common scale. Namely for low densities and temperatures the masses scaled as $m^*/m \simeq \left[\langle \bar{q}q\rangle^*/\langle \bar{q}q\rangle\right]^{1/2}$ whereas at higher densities and temperatures as $\left[\langle \bar{q}q\rangle^*/\langle \bar{q}q\rangle\right]$. In the present work we summarize new empirical evidence for Brown-Rho (BR) scaling and discuss in a general way its impact on the nuclear many-body problem. While the double decimation should be carried out, it has been a prevalent practice in nuclear physics community to proceed with the second decimation, assuming density independent masses, without implementing the first, chiral symmetry decimation. We show why most nuclear phenomena can be reproduced by theories using either density-independent, or density-dependent masses, a grand conspiracy of nature that is an aspect that could be tied to the Cheshire-Cat phenomenon in hadron physics. We go through some of the history of the chiral symmetry decimation (CSD) which involves BR scaling, in order to show that historically one had to look at very specific phenomena involving transition matrix elements to see that it is necessary. We identify what is left out in the Fermi liquid decimation (FLD) that does not incorporate the CSD. Experiments such as the dilepton production in relativistic heavy ion reactions, which are specifically designed to observe effects of dropping masses, could exhibit large effects from the reduced masses. However they are compounded with effects that are not directly tied to chiral symmetry. We discuss a recent STAR/RHIC observation where BR scaling can be singled out in a pristine environment.
1 Introduction

Brown and Rho adduced empirical evidence for BR scaling in terms of nuclear interactions and nuclear structure more than a decade ago [1, 2, 3]. One of the most convincing evidences was the decrease in the tensor interaction in nuclei: The $\pi$ and $\rho$ exchange enter into the tensor interaction with opposite signs. The $\rho$-coupling to the nucleon is about twice that of vector dominance, so its square is four times greater. Brown and Machleidt [4] showed this to follow unambiguously from the phase shifts from nucleon-nucleon scattering.

The large $\rho$ tensor coupling is important in free space, in that it cancels the otherwise extremely strong pionic coupling, which goes as $r^{-3}$ at short distances, at a distance of $\sim 0.6$ fm. Inside of this, the net tensor interaction is repulsive, but strongly cut down in net effect by the large repulsion from $\omega$-exchange which keeps the two nucleons apart. Because of the $\rho$ exchange, the nucleon-nucleon interaction never gets very strong, explaining why the effective nucleon-nucleon interaction obtained in the Sussex work [5], which neglected off-shell effects, worked well in reproducing the effective nucleon-nucleon interaction in nuclei. The very smooth behavior of the nuclear interactions in going off shell minimizes the difference between the on-shell interaction and the half off-shell one that is commonly calculated as effective interactions.

In medium, with the decrease in $m_\rho^*$ to $\sim 0.8 m_\rho$ by nuclear matter density $n_0$, the tensor interaction becomes even weaker, and this was the early evidence adduced by Brown and Rho for BR scaling.

Our work is prompted by two recent developments that are both profound and powerful for nuclear structure. One is the notion of “vector-manifestation fixed point” in effective field theory of hadrons discovered by Harada and Yamawaki [6, 7] and the other is the identification of nuclear matter as a Fermi liquid at its fixed point [8, 9, 10]. The first fixed point is dictated by the matching of effective field theories to QCD and the second accounts for the stability of strongly interacting many-body systems that possess a Fermi surface arising from a quantum critical phenomenon. We discuss how one could combine these two fixed point structures into an effective field theory of nuclei and nuclear matter.

The plan of this review is as follows.

In Section 2 we develop the concept of “double decimation” for describing nuclear phenomena based on the general strategy of effective field theories that can represent low-energy nonperturbative QCD. The first decimation deals with chiral symmetry scale and we suggest that the “parametric” density dependence encoded in BR scaling [3] is the consequence of the “chiral scale decimation” (CSD) tied to the “vector manifestation” à la Harada and Yamawaki [6, 7] in hidden local symmetry theory that is matched to QCD. The second decimation deals with the Fermi momentum scale and corresponds to (Landau) Fermi liquid description of nuclear matter. We shall refer to this as “Fermi liquid decimation (FLD).”

In Section 3, one way of doing effective field theory in nuclear physics is presented with a focus on how to incorporate the standard nuclear physics approach (SNPA) that has had a spectacular success, into the framework of modern effective field theory that we shall refer to as more effective effective field theory (MEEFT in short). Both the effective nucleon interaction $V_{\text{low}-k}$ and the electroweak matrix elements of few-nucleon systems are described in this formalism.

In Section 4 we discuss in a general way why the need for density dependent hadron masses was missed for so many years in nuclear phenomena; namely, the description of most phenomena with density independent masses could fit experiments. We show that with a common scaling
with density of hadron masses, the kinematics of the nuclear many-body problem are not changed very much in the region of densities up to nuclear matter density and we illustrate this with the highly successful Dirac phenomenology.

As cases where BR scaling can be singled out in a specific way in nuclear processes, we discuss in Section 5 how the scaling enters in what is known as Warburton’s $\epsilon_{MEC}$ factor in axial charge transitions in heavy nuclei as well as in deeply bound pionic atoms which have been recently studied experimentally.

In Section 6, we suggest that the vector manifestation fixed point $a = 1$ of Harada and Yamawaki is “precociously” reached when baryons are present and hence in nuclear medium, as a consequence of which the photon couples half and half to the vector meson and the nucleon core (or skyrmion). This picture may be relevant to the electromagnetic form factors of the proton recently measured at the J-Lab as well as to the longitudinal and transverse form factors of nuclei with the meson and nucleon masses dropping à la BR scaling.

While it is now generally accepted that the CERES dileptons can be explained by modified properties of hadrons in dense and hot matter, it is difficult to single out BR scaling as the principal mechanism for the shifted spectral distribution as temperature and density effects are compounded in the process. We sketch in Section 7 what we believe to be the correct way of interpreting the dilepton data.

It is discussed in Section 8 how the situation can be a lot clearer in the RHIC experiments. We discuss in particular how a direct measurement of medium-dependent vector-meson mass $m_V^*$ could be made in a pristine environment where temperature effects are small and where the density can be well reconstructed.

A discussion of the (absence of) need for “double decimation” in nuclear structure makes up Section 9 where we discuss effective interaction between nucleons in nuclei. The Kuo-Brown interactions [11] have undergone a renaissance with the finding from the renormalization-group (RG) formulation that all nucleon-nucleon interactions which fit the two-body phase shifts give the same effective interactions $V_{low-k}$ in nuclei. The RG cutoff $\Lambda$ is chosen to be at the upper scale to which phase shifts have been obtained from nucleon-nucleon scattering; i.e. $\Lambda \sim 350$ MeV.

That $V_{low-k}$ provides a unique effective interaction which works well in reproducing effective forces in nuclei is unquestionable, given the excellent fits to all sorts of data. Nonetheless, there are questions:

- (1) Where does the modification in the tensor interaction, which is lowered by the decreasing mass of the $\rho$-meson which contributes with opposite sign to the pion (whose mass is presumably unchanged in medium at low density), come in? It is well known that certain states, such as the ground state of $^{14}$N, depend sensitively on the tensor interaction. Also in the nuclear many body problem, the decrease in the tensor force will give less binding energy and less help with saturation, the second order tensor force giving a large contribution to the binding energy of nuclear matter. In fact, this second order term which has the form

$$V_{eff} \simeq \left(\frac{3 + \vec{\sigma}_1 \cdot \vec{\sigma}_2}{4}\right) V(r)$$

(1)

is large as we discuss. In an old paper, Feshbach and Schwinger [12] describe the difference between $^1S$ and $^3S$ potentials in terms of the above $V_{eff}$ which acts only in the latter state.

- (2) In addition to the modifications in both tensor and spin-orbit interactions, which have not been put into the Fermi liquid decimation, there is the question of off-shell energies of interme-
mediate particles \(^1\) which are taken to be plane waves. The conclusion following from the Bethe reference spectrum \([13]\) was that there was a small amount of binding in the intermediate states with energies just above the Fermi surface, referred to as “dispersion” correction. Generally taking the intermediate states to be free, as in the Schwinger interaction representation, seemed to be a good approximation. This is what is done in the RG calculations, in the Fermi liquid decimation.

Concluding remarks are given in Section 10 with a reference to the Cheshire Cat Principle revisited. We postulate that the missing “smoking gun” for BR scaling in nuclear structure physics is an aspect of the Cheshire Cat Principle discovered in the baryon structure as reflected in many-nucleon systems.

2 The Double Decimation

Our ultimate objective is to describe in a unified way finite nuclei, nuclear matter and dense matter up to chiral restoration. For this we introduce the approximation – double decimation – by which the phase structure in the hadronic sector can be drastically simplified. The effective field theory (EFT) approach to few-nucleon systems described below does not require both decimations but if one wants to correctly describe phase transitions of hadronic systems under extreme conditions of density and/or temperature, one needs at least two decimations which we now describe. We propose that the procedure consists of what we call “chiral symmetry decimation (CSD)” and “Fermi liquid decimation (FLD).”

2.1 The “intrinsic dependence (ID)” from chiral symmetry decimation

As we have explained in a series of recent papers \([14, 15, 16, 17]\), a candidate effective field theory relevant in the hadronic sector matched to QCD at a scale near the chiral scale \(\Lambda_\chi\) is hidden local symmetry (HLS) theory with the vector manifestation (VM) found by Harada and Yamawaki \([6, 7]\). This theory, denoted HLS/VM in short, with the vector mesons \(^2\rho, \omega\) etc. treated as light on the same footing as the (pseudo) Goldstone pions as originally suggested by Georgi \([18]\), turns out to yield results consistent with chiral perturbation theory \([19]\). Given this theory valid at low energy in matter-free space, one can then construct the chiral phase transition in both hot \([20]\) and dense \([21]\) media recovering BR scaling \([3]\) near chiral restoration. This theory makes unambiguous predictions, at least at one-loop order, on the vector and axial-vector susceptibilities and the pion velocity at the transition point that seem to be different from the standard scenario and that can ultimately be tested by QCD lattice measurements \([22]\).

In the rest of this section we will confine out consideration on finite density at zero temperature \(^3\).

An early, seminal attempt to arrive at nuclei and nuclear matter in chiral Lagrangian field theory that models QCD was made by Lynn \([23]\). What we are interested in here is to exploit the recent development of HLS/VM and arrive at a field theoretic description anchored on QCD. In studying nuclear systems, the EFT and QCD have to be matched via current correlators at a suitable scale \(\Lambda_M\) in the background of matter characterized by density \(n\). The matching defines the “bare” HLS Lagrangian, giving to the parameters of the Lagrangian the “intrinsic” density

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\(^1\)The holes are on-shell in the Brueckner-Bethe theory.

\(^2\)We are assuming \(U(2)\) flavor symmetry in the two-flavor case.

\(^3\)Implementing temperature effects is straightforward and hence will be ignored in what follows in this section.
dependence (IDD). This means that the parameters of the Lagrangian such as hidden gauge coupling constant $g^*$, gauge boson mass $m^*$, etc will depend intricately not only on $\Lambda_M$ but also on density $n$. This intrinsic density dependence is generally missing in the calculations that do not make the matching to QCD. Next to account for quantum effects, we need to decimate the degrees of freedom and excitations from the matching scale $\Lambda_M$ to a scale commensurate with the presence of a Fermi sea, $\Lambda_{FS}$. The latter is the scale relevant to nuclear physics, typically $\Lambda_{FS} \lesssim 1.5$ fm$^{-1}$. What governs this first decimation is chiral symmetry.

2.1.1 Sliding vacua

The HLS/VM at the present stage of development can tell us what happens only at $n \approx 0$ and at $n \approx n_c$. It does not tell us how the intrinsic density dependence (IDD) interpolates from $n = 0$ to $n = n_c$. What we are interested in, however, involves a wide range of densities, from zero to nuclear matter and ultimately to chiral restoration. How to do this is at present poorly known. There are however two approximate ways to address this issue. One is based on extending NJL model so as to simulate the vector manifestation effect associated with vector-meson degrees of freedom. This is the approach we have proposed before [17] which we shall follow in the rest of the paper. The other which we shall not employ in this paper is however quite instructive on how the “vacuum change” induced by matter is manifested in many-nucleon systems. It is based on a skyrmion description of dense hadronic matter which we briefly describe.

The skyrmion approach to dense matter developed recently [24, 25, 26] makes transparent the crucial role of the intrinsic density dependence (IDD) in nuclear processes. The power of the skyrmion model is that a single effective Lagrangian provides a unified description of mesons and baryons and treats single-baryon and multi-baryon interactions on the same footing, thereby describing finite nuclei as well as infinite nuclear system – nuclear matter. Although the notion of skyrmions as modelling of QCD is fairly well established [27], one does not yet know at present how to write down a fully realistic skyrmion Lagrangian, and hence cannot do a quantitative calculation but one can gain a valuable insight into the structure of the “sliding vacua” that we are interested in.

Following [24, 25, 26], one considers a skyrmion-type Lagrangian with spontaneously broken chiral symmetry and scale symmetry, associated, respectively, with nearly massless quarks and trace anomaly of QCD. The effective fields involved are the chiral field $U = e^{i\pi/f}$ with $\pi$ the (pseudo)-Goldstone bosons for the former and the scalar “dilaton” field $\chi$ for the latter. Such a theory which may be considered to be an $N_c \rightarrow \infty$ approximation to QCD can, albeit approximately, describe not only the lowest-excitation, i.e., pionic, sector but also the baryonic sector and massive vector meson sector all lying below the chiral scale $\Lambda_\chi$ \(^4\). What is even more significant, it can treat on the same footing both single-baryon and multi-baryon systems including infinite nuclear matter [25].

We are interested in how low-energy degrees of freedom in many-body systems behave in dense matter, in particular as the density reaches a density at which QCD predicts a phase transition from the broken chiral symmetry to the unbroken chiral symmetry, i.e., chiral restoration. The simplest possible Lagrangian with the given symmetry requirements consistent with QCD

\(^4\)This naturally implies that the gluonium degrees of freedom lying higher than $\Lambda_\chi$ are to be integrated out leaving only the “soft” quarkonium degrees of freedom. As shown in [24], these are the degrees of freedom intricately tied to the spontaneous breaking of chiral symmetry.
The Lagrangian is of the form [24]

\[ \mathcal{L} = \frac{f_\pi^2}{4} \left( \frac{\chi}{f_\chi} \right)^2 \text{Tr}(\partial_\mu U^\dagger \partial^\mu U) + \frac{1}{32e^2} \text{Tr}([U^\dagger \partial_\mu U, U^\dagger \partial_\nu U])^2 + \frac{f_\pi^2 m_\pi^2}{4} \left( \frac{\chi}{f_\chi} \right)^3 \text{Tr}(U + U^\dagger - 2)
\]

\[ + \frac{1}{2} \partial_\mu \chi \partial^\mu \chi - \frac{1}{4} \frac{m_\chi^2}{f_\chi^2} \left[ \chi^4 \left( \ln(\chi/f_\chi) - \frac{1}{4} \right) + \frac{1}{4} \right]. \]  
\[(2)\]

We have denoted the vacuum expectation value of \( \chi \) as \( f_\chi \), a constant which describes the decay of the scalar \( \chi \) into pions in matter-free space. The QCD trace anomaly can be reproduced by the last term of (2), i.e., the potential energy \( V(\chi) \) for the scalar field, which is adjusted so that \( V = dV/d\chi = 0 \) and \( d^2V/d\chi^2 = m_\chi^2 \) at \( \chi = f_\chi \).

The vacuum state of the Lagrangian at zero baryon number density is defined by \( U = 1 \) and \( \chi = f_\chi \). The fluctuations of the pion and the scalar fields about this vacuum, defined through \( U = \exp(\vec{\tau} \cdot \vec{\phi}/f_\pi) \), and \( \chi = f_\chi + \tilde{\chi} \) give physical meaning to the model parameters: \( f_\pi \) as the pion decay constant, \( m_\pi \) as the pion mass, \( f_\chi \) as the scalar decay constant, and \( m_\chi \) as the scalar mass.

The coupled classical equations of motion for the \( U \) and \( \chi \) fields of (2) give rise to solitons for given winding numbers corresponding to given baryon numbers with their structure constrained by the classical scalar field. Nuclear matter is then described classically by an FCC crystal and possibly a Fermi liquid when quantized. When the system is strongly squeezed, there is a phase transition from the FCC state to a half-skyrmion state representing chiral restoration in dense system [24]. It is assumed that this transition remains intact when the initial state is in Fermi liquid. Fluctuations on top of the background crystal structure describe excitations of the pionic and scalar quantum numbers, the properties of which are then governed by the mean field values – denoted \( \chi^* \) – of \( \chi \) at a given density fixed by the unit cell size \( L \) of the crystal. What is found in [24] is that the mean field of the scalar \( \chi \) mostly – though not entirely – governs the scaling behavior of the parameters – such as the pion decay constant – of the Lagrangian indicative of the sliding vacuum structure. The results are found to be in qualitative agreement with what is found in BR scaling [3] as elaborated more precisely below in NJL. The VM structure is not seen explicitly in the skyrmion picture of [24] since there are no vector mesons in the Lagrangian but it does not appear difficult to obtain it once vector mesons are suitably implemented into the model Lagrangian.

### 2.1.2 Nuclear matter from chiral symmetry decimation

Given an HLS/VM theory decimated to \( \Lambda_{FS} \), how does one go about obtaining nuclei and nuclear matter? The first thing we need is a Fermi surface characterized by the Fermi momentum \( k_F \sim \Lambda_{FS} \) and the Fermi surface arises in effective field theories as a quantum critical phenomenon. As sketched in [27] and developed in [25, 26], in the skyrmion picture, a nucleus of mass number \( A \) arises as a topological soliton of winding number \( W = A \). Nuclear matter is then given by \( W = \infty \). When quantized, the extended soliton system will naturally possess a Fermi surface characterizing the filled Fermi system.

In confronting nuclei and nuclear matter in nature, however, it is more advantageous to work with explicit nucleon fields rather than with multi-winding number skyrmions. When nucleons are explicitly present in the theory, a nucleus will no longer emerge as a topological
object. Instead it must arise as a non-topological soliton in a way conjectured by Lynn [23]. Nuclear matter will then be more like a chiral liquid in Lynn’s language. We would like this soliton to emerge in a simple way from an effective Lagrangian endowed with the parameters of the Lagrangian intrinsically density-dependent. A systematic derivation of such a soliton structure that is realistic enough is lacking at the moment. However there is a short-cut approach to this and it relies on Walecka’s mean field theory of nuclear matter [28]. The principal point we put forward is that the Walecka mean-field solution of certain effective Lagrangian (specified below) can be identified as the soliton – topological or non-topological – solution described above.

Now, with the nucleon fields explicitly incorporated, there are two (equivalent) ways to write down such an effective Lagrangian of Walecka type that results from decimating down to \( \Lambda_{FS} [14] \). One is the type-I approach in which the heavy-meson degrees of freedom of the HLS Lagrangian are integrated out and the other is the type-II one in which relevant heavy-meson degrees of freedom are retained. The two versions give equivalent descriptions of the same physics for the ground state and low-frequency fluctuations.

The type-I Lagrangian has the form

\[
L_I = \bar{N} [i\gamma_\mu (\partial^\mu + iv^\mu + g^*_A \gamma_5 a^\mu) - M^*] N - \sum_i C^*_i (\bar{N} \Gamma_i N)^2 + \cdots
\]

where the ellipsis stands for higher dimension and/or higher derivative operators and the \( \Gamma_i \)'s Dirac and flavor matrices as well as derivatives consistent with chiral symmetry. The star affixed on the masses and coupling constants represents the intrinsic parametric dependence on density relevant at the scale \( \Lambda_{FS} \). The induced vector and axial vector “fields” are given by

\[
v_\mu = -i/2(\xi^\dagger \partial_\mu \xi + \xi \partial_\mu \xi^\dagger) \quad \text{and} \quad a_\mu = -i/2(\xi^\dagger \partial_\mu \xi - \xi \partial_\mu \xi^\dagger).
\]

In (4) only the pion (\( \pi \)) and nucleon (\( N \)) fields appear explicitly: all other fields have been integrated out. The effect of massive degrees of freedom that are integrated out and that of the decimated “shells” will be lodged in higher-dimension and/or higher-derivative interactions. The external electro-weak fields if needed are straightforwardly incorporated by suitable gauging.

To write the Lagrangian for the type-II approach, we need to pick the appropriate heavy degrees of freedom we want to consider explicitly. This Lagrangian will be essentially the HLS Lagrangian implemented with nucleon fields and chiral scalar heavy mesons that are not taken explicitly into account. The relevant heavy mesons for nuclear physics are a vector meson in the \( \omega \) channel which is a chiral scalar and a flavor scalar \( \sigma \) that plays an important role in Walecka-type model for nuclear matter. Assuming that \( U(N_f) \) (\( N_f = 2 \) for low-energy nuclear physics) symmetry holds in medium, we can put the \( \rho \) and \( \omega \) in the \( U(N_f) \) multiplet and write a Harada-Yamawaki HLS Lagrangian with the parametric dependence suitably taken into account. To make the discussion simple, let us consider symmetric nuclear matter in which case the type-II Lagrangian can be written in the form of Walecka linear theory [28] with the parametric density dependence represented by the star,

\[
L_{II} = \bar{N} [i\gamma_\mu (\partial^\mu + ig_\omega^\mu \omega^\mu) - M^* + h^* \sigma] N
\]

\[
- \frac{1}{4} F^2_{\mu\nu} + \frac{1}{2} (\partial_\mu \sigma)^2 + \frac{m_\omega^2}{2} \omega^2 - \frac{m_\sigma^2}{2} \sigma^2 + \cdots
\]

where the ellipsis denotes higher-dimension operators. We should stress that (5) is consistent with chiral symmetry since here both the \( \omega \) and \( \sigma \) fields are chiral singlets. In fact the \( \sigma \) here
has nothing to do with the chiral fourth-component scalar field of the linear sigma model except near the chiral phase transition density; it is a “dilaton” connected with the trace anomaly of QCD. The possibility that this dilaton turns in-medium into the fourth component of the chiral four-vector in the “mended symmetry” way à la Weinberg [29] near chiral restoration has been discussed in [30]. Since we are far from the critical density at which chiral restoration takes place, the vector mesons are massive and the would-be Goldstone scalars (the longitudinal components of the vector mesons) are absent. In (5), the pion and ρ-meson fields are dropped since they do not enter in the mean field approximation but they can be put back if needed (as for fluctuations in the pionic channel) in HLS symmetric way.

2.2 Fermi-liquid decimation

Given the type-I (4) or type-II (5) Lagrangians, the next step is to do the Fermi liquid decimation (FLD). In [31], Schwenk, Brown and Friman build on the effective interactions $V_{\text{low-k}}$ a decimation scheme to arrive at Fermi-liquid parameters, namely the Landau effective mass for the quasi-nucleon $m_{N}^{*}$ and the quasiparticle interactions $\mathcal{F}$. In doing this, the intrinsic density/temperature dependence coming from the chiral symmetry decimation was not included. It also makes the implicit assumption that $n$-body interactions for $n>2$ are negligible. In any event, it has been shown that both $m_{N}^{*}$ and $\mathcal{F}$ are fixed-point quantities. This approach purports to “calculate” the parameters of the “penultimate” effective Lagrangian from first principles.

An alternative method which we follow in this article is based on the mapping of the effective field theory Lagrangians (4) and (5) to the Fermi liquid fixed point theory as beautifully explained in a review by Shankar [8]. The mapping relies on the work done by Matsui [32] who showed that Walecka linear mean field theory is equivalent to Landau Fermi-liquid theory. Using this argument, it is then immediate to map the mean-field solution of (4) or (5) to Landau EFT: In the mean field approximation, the two theories, (4) and (5), yield the same results [33, 34]. This chain of reasoning, developed in [10, 35, 36], has shown that the fluctuations on top of the Fermi liquid fixed point – represented by the mean field of (4) or (5) – in response to EW external fields can be related to the BR scaling parameter given by the ratio $\Phi = m^{*}/m$.

2.3 The vector mass and gauge coupling scaling

In [14], we discussed on how the vector meson mass $m_{V}^{*}$ and the gauge coupling $g^{*}$ in HLS theory with the vector manifestation – both of which are the essential elements of our EFT – drop in medium as one goes from below to above nuclear matter density. At very low density, chiral perturbation theory provides relevant information and near the critical density, the VM tells us how they scale. In between, little could be said at the moment. Here we extrapolate the scaling behavior from zero density (or temperature modulo Dey-Eletsky-Ioffe low-temperature theorem [37]) to the critical density (or temperature). Our picture is summarized in Fig.1.

Up to near nuclear matter density relevant to the chiral symmetry decimation, the BR scaling of the mass is expected to go like $\sqrt{\langle \bar{q}q \rangle^{*}}$ and the vector coupling (more precisely the HLS gauge coupling) to stay more or less unchanged. The former is deduced from the GMOR relation for in-medium pion mass which should hold in HLS theory for small density, with the assumption that the pion mass does not change in density as indicated, e.g., by recent $SU(2)_{c}$ lattice results [38]. The latter generically follows from chiral models and is consistent with the weakening tensor forces in nuclei. We do not know precisely how the scaling changes as we
Figure 1: Schematic scaling behavior for $m_\rho^*$ and $g^*$ vs. density. The cross-over region which is not understood is marked by ??.

go above the matter density. However as we approach the critical density, we learn from HLS theory matched to QCD that the vector manifestation takes place, which means that both the hadronic mass and the gauge coupling constant should scale linearly in $\langle \bar{q}q \rangle^*$ [21]. Thus near the chiral transition point, $g^*/m_\rho^*$ should go to a constant of density as required for quark number susceptibilities [39].

3 MEEFT or More Effective Effective Field Theory

3.1 Doing effective field theory in nuclear physics

In this section, we consider the situation where one does not need to go through the double decimation procedure described above. Suppose one is interested only in describing low-energy nucleon-nucleon scattering. For this, there are a variety of equivalent ways to proceed [40]. First one writes down an effective Lagrangian using the fields describing the degrees of freedom relevant for the process in question. For instance, if one is looking at the S-wave two-nucleon scattering at low energy, say much less than the pion mass, one can simply take only the nucleon field as an explicit degree of freedom, integrating out all others, including the pion. The pions which are important due to the broken chiral symmetry may be introduced perturbatively. A systematic power counting can then be developed and used, in conjunction with an appropriate regularization (e.g., the power divergence substraction (PDS) scheme [40]), to compute the scattering amplitude by summing to all orders a particular set of diagrams. In principle, the parameters of the pionless Lagrangian could be determined for a given scale by lattice QCD. In this sense, such a systematic higher-order calculation in this EFT can be considered to be equivalent to doing QCD. In practice, however, the parameters are obtained from experiments. For two nucleon systems, two classes of parameters are to be determined from experiments.
One is single-particle vertex and this is given by on-shell information. The other is intrinsically two-particle in nature requiring data on two-nucleon processes. When the parameters are fully determined, this procedure with the pionic contributions taken into account perturbatively does lead to sensible results. For instance, it correctly reproduces such standard nuclear physics results as the effective-range formula.

The EFT which adheres strictly to order-by-order consistency in the power counting (that we shall refer to as “purist’s EFT”) however suffers from the lack of predictivity. The number of unknown parameters increases rapidly as the number of nucleons involved in the process increases. Thus even if two-nucleon systems are well described by the EFT in question, treating systems involving more than two or three nucleons becomes prohibitively difficult, if not impossible. Furthermore treating pions as perturbative misses the power of the “chiral filter mechanism” [41] that plays an important role in predictive calculations [42]. At present, going to nuclear matter is out of reach by this EFT approach.

Even if such a method were available so that we could write down an effective Lagrangian on top of a Fermi sea, the parameters of the penultimate effective theory for, say, nuclear matter, would be very far from the first principles, QCD: Calculating them would be somewhat like calculating the boiling point of the water starting from QED. We propose here how to circumvent this impasse. Our approach proposed here is admittedly indirect and drastically simplified. Broadly speaking, it involves a two-step decimation starting from the chiral Lagrangian that is matched to QCD at the chiral scale $\Lambda_\chi \sim 1$ GeV. The key ingredient that allows this feat is the highly refined “standard nuclear physics approach (SNPA)” and the objective is to marry the SNPA to an EFT.

As an illustration of our strategy, we briefly discuss here how the marriage can be effectuated. As summarized recently [43, 14], a thesis developed since some time posits that by combining the standard nuclear physics approach (SNPA) based on potentials fit to experiments with modern effective field theory, one can achieve a more predictive power than the purist’s EFT alone can. The idea was recently given a test in a variety of electroweak processes in nuclei, in particular, the solar $pp$ fusion and $hep$ processes where highly accurate predictions could be made free of parameters [44]. As alluded above, a closely related mechanism has been found at work in nuclear effective interactions that figure in nuclear structure calculations [45]. We briefly sketch these two recent developments which are closely connected to the issue at hand.

### 3.2 Effective interactions in nuclei

The calculation of the EW response functions that will be the subject of the next subsection relied on an implicit assumption on the effective forces that enter in the calculation of the wave functions. A recent development provides a support to this assumption.

Consider two nucleon scattering at very low energy, a process very well understood in SNPA. The relevant $T$ matrix for the scattering is a solution of the Lippmann-Schwinger equation with a “bare” potential $V_{\text{bare}}$ figuring as the driving term. The long-range part of the bare potential $V_{\text{bare}}$ is governed by chiral symmetry, namely, by a pion exchange and hence is unambiguous. But the short-range parts are not unique. Even if the potentials are determined accurately by fitting scattering data up to say $E_{\text{lab}} \sim 350$ MeV, those potentials that give the equivalent phase shifts can differ appreciably, in particular in the short-distance parts. In terms of effective field theory, what this means is that while the long-range parts given by “low-order” expansion are the same for all the realistic potentials, shorter-range terms that are given by
higher order terms can differ depending upon how they are computed. They will depend upon how the power counting is organized, what regularization is used etc. In practice, those “higher-order” terms are fixed by fitting to experimental data. The examples for such realistic potentials are the Paris potential, the Bonn potential, the chiral potential etc. Suppose in the integral equation satisfied by the scattering $T$-matrix, one integrates out the momentum scales above a given cutoff $\Lambda$ sufficiently high to accommodate the relevant degrees of freedom and probe momentum but low enough to exclude the massive scales that do not explicitly figure in the theory, with the requirement that the resulting effective potential reproduce the phase shifts while preserving the long-range wave function tails as given by the half-on-shell (HOS), $T(k', k; k^2)$. To the extent that the HOS $T$ matrix is a physical quantity, it should be independent of where the cutoff is set

$$\frac{d}{d\Lambda} T(k', k; k^2)_{\Lambda} = 0. \quad (6)$$

This condition leads to a renormalization group equation (RGE) for the effective potential, denoted $V_{\text{low}-k}$. It is important to note that the fit to experiments defines the complete theory for the probe momentum $k \ll \Lambda$. Now since the HOS $T$ is fit to experiments, integrating out the momentum component $p > \Lambda$ transfers the physics operative above the cutoff into the counter terms that are to be added to the bare potential to give the effective one, $V_{\text{low}-k}$. Bogner et al [45] have shown that the resulting $V_{\text{low}-k}$ is independent of the bare potential one starts with as long as it is consistent with the chiral structure, i.e., the long-range tails of the wave functions, and the $T$ is fit to experiments. It is important to recognize that this strategy is none other than that of the SNPA.

The results of Ref.[45] are reproduced in Fig.2. The figure on LHS shows the variety of bare potentials that are fit to experiments and the one on the RHS the effective potentials collapsing into a universal curve at $\Lambda = 2$ fm$^{-1}$ after the integrating-out procedure. How the collapse occurs from different $V_{\text{low}-k}$'s for different potentials is shown in Fig.3. For $\Lambda > 2$ fm$^{-1}$, $V_{\text{low}-k}$'s can be different for different potentials but they “collapse” to one universal curve...
Figure 3: The collapse of the diagonal momentum-space matrix elements of $V_{\text{low}-k}$ as the cutoff is lowered to $\Lambda = 2.1 \text{ fm}^{-1}$ in the $^3S_1$ partial wave.

for $\Lambda = 2.1 \text{ fm}^{-1}$. The reason for this is easy to understand: that the phase shift analyses from which the two-body interactions in Fig. 3 were fit were carried out for experiments up to laboratory energies of $\sim 350 \text{ MeV}$, which corresponds to a c.m. momentum of $2.1 \text{ fm}^{-1}$. Although the phase shifts were determined on-shell, as we noted in Section 1 the interaction is not far off-shell when used in nuclei. This is how one can understand that the half-off-shell $V_{\text{low}-k}$ reproduced the same diagonal matrix elements for the various potentials fit to experiments to the given cut-off scale. The point is that off-shell effects are unimportant in considerations of the smooth parts of the shell-model wave functions. Of course, these effects can be large if the wave functions are forced far off shell by short-range two-body interactions, but these will involve a scale higher than our $2.1 \text{ fm}^{-1}$ not constrained by the experimental data.

That while the bare potentials are widely different for all momenta, all realistic potentials give the identical effective potential illustrates the power of the strategy that combines the accuracy achieved by the SNPA and chiral effective field theory of QCD into a “more effective effective field theory (MEEFT).” The key lesson from this result is that the short-distance part of the potential which represents higher orders in EFT power counting schemes which differ for different counting schemes may be different from one “realistic potential” to another but when suitably regularized taking into account the constraints by experimental data, the resulting effective potential comes out unique. For instance, the “chiral potential” which is consistent with the chiral counting, and hence presumed to be more in line with the tenet of EFT and the $v_{18}$ potential which, apart from the long-range part, is not, can differ at “higher orders” but

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\[^5\text{At about this energy substantial inelasticity sets in, complicating analysis.}\]
gives the same $V_{\text{low} - k}$. This reflects how the MEEFT works.

An identical mechanism is at work for the weak matrix elements relevant for the solar neutrino processes [44] as described below.

The $V_{\text{low} - k}$ is the basis of nuclear structure calculations replacing the role of G-matrix. For low-energy processes, it is insensitive to short-distance physics properly taking into account the standard short-range correlations. It is also the input for field theoretic calculation of the Landau parameters for nuclear matter [31] which are the Fermi liquid fixed point quantities.

3.3 Electroweak processes in nuclei

The next question we raise is: How can one do a unified effective field theory calculation which is truly predictive for the following processes?

\begin{equation}
    p + p \rightarrow d + e^+ + \nu_e, \tag{7}
\end{equation}

\begin{equation}
    ^3\text{H} \rightarrow ^3\text{He} + e^- + \bar{\nu}_e, \tag{8}
\end{equation}

\begin{equation}
    p + ^3\text{He} \rightarrow ^4\text{He} + e^+ + \nu_e. \tag{9}
\end{equation}

The processes (7) and (9) take place in the Sun, playing an essential role for the solar neutrino problem but have not been measured in the laboratories while the process (8) is accurately measured. Given the triton beta decay data, the objective then is to make an accurate prediction for the rates for (7) and (9). This highly non-trivial feat has been accomplished recently by Park et al [44] by an MEEFT. We note that such a predictive calculation is not yet feasible in other EFT strategies developed so far (e.g., the purist’s EFT).

The MEEFT strategy of [43] employed in the publications by Park et al. [44] goes as follows. First one picks a potential fit to on-shell NN scattering up to typically $E_{\text{lab}} \sim 300$ MeV plus many-body (typically three-body) potentials, the combination of which is to describe accurately many-nucleon scattering data. As stated, this potential is required to be consistent with chiral symmetry, the symmetry of QCD, which means that the long-range part of the potential is dictated by the pion exchange. An example for such a potential is the $v_{18}$ potential which has been used in the numerical calculation. Next one computes the weak current in a chiral chiral perturbation theory, e.g., heavy-baryon chiral perturbation theory, to an appropriate order in the power counting, say, $Q^n$ where $Q$ is the characteristic momentum scale. The current matrix element is then computed with this current and the “realistic” wave functions computed with the given potential. In doing this, one integrates out all momentum components above the cutoff $\Lambda$, the effect of which is a set of counter terms appearing in the effective current. Working to the next-to-next-to-next order (i.e., $Q^3$) relative to the leading matrix element given by the single-particle Gamow-Teller matrix element in (7)-(9), one encounters a number of counter terms but symmetry consideration reduces them to a single combination, denoted $d^R$ in [44], entering in all three processes. The counter term coefficient $d^R$ is cutoff-dependent and so is the matrix element of the current with the momentum component $p > \Lambda$ integrated out. The numerical values of these individual terms differ from one potential to another but the sum of the two does not. The strategy then is to exploit that the short-range component of the interactions must be universal, that is, independent of the mass number. This means that the coefficient $d^R$ can be determined for any given cutoff from the accurately measured process (8). Since there are no other unknowns to order $Q^3$, we have a parameter free theory to calculate all the matrix elements that figure in (7) and (9). Indeed this is what was done in [44]: The $S$
factor for the $pp$ fusion process was calculated within the accuracy of 0.4\% and that for the $hep$ process within the accuracy of 15\%. Both results, particularly the latter, are of unprecedented accuracy unmatched by other calculations \(^6\).

4 Saturation and Dirac Phenomenology

In detail the situation with dropping masses is complicated because the pion, both in lowest order and in two-pion exchange, contributes to the energy, but $m_π$ is most likely unscaled, whereas the scalar and vector mesons, two most important effective fields in nuclear medium, are. Rapp et al \cite{46} showed that given these scalings, saturation does come at the right density. The chief contributions to the binding energy come from the $σ$ and $ω$ exchange. Here we show schematically how these behave in the “swelled” world.

Consider a Hamiltonian

$$H = \sum_i \frac{1}{2m_N} \nabla_i^2 + \frac{1}{2} \sum_{ij} (g_ω \frac{e^{-m_ω r_{ij}}}{r_{ij}} - g_σ \frac{e^{-m_σ r_{ij}}}{r_{ij}})$$

with ground state energy

$$H\psi = E_0\psi.$$ \hspace{1cm} (11)

Now scale all $m$’s

$$m_N^* = \lambda m_N, \; m_σ^{*,ω} = \lambda m_{σ,ω}.$$ \hspace{1cm} (12)

Then

$$H^*(\lambda, \vec{x}) = \lambda H(1, \vec{r})$$ \hspace{1cm} (13)

where

$$\vec{x} = \lambda \vec{r}.$$ \hspace{1cm} (14)

Now

$$H^*(\lambda, r)\psi = E_0(\lambda)\psi$$ \hspace{1cm} (15)

or

$$\lambda H(1, x)\psi = E_0(\lambda)\psi.$$ \hspace{1cm} (16)

But $H(1, x)$ is just a relabelling of $H$ via $r \to x$, etc. so the solution of

$$H(1, x)\psi = E\psi$$ \hspace{1cm} (17)

for the lowest eigenvalue is $E_0$, or

$$E_0 = E_0(\lambda)/\lambda.$$ \hspace{1cm} (18)

\(^6\)We remind the readers that up to date, the calculated values for the $hep S$ factor varied by orders of magnitude, hence completely unknown.
Thus
\[ E_0(\lambda) = \lambda E_0, \]  
and the scaled system (with \( \lambda \)) is less bound than the original unscaled system since \( \lambda = 0.8 \) at nuclear matter density. The replacement of the correlated two \( \pi \) exchange, with the two \( \pi \)'s in a relative S-state by a \( \sigma \) with scaling mass is, however, not a good approximation as shown by Rapp et al [46]. Even with the loss in binding energy, because of the scaling, saturation does not occur in the correct region of densities, as can be seen from Fig.2 of the quoted paper. Although crossed channel exchange of a \( \rho \) with decreased mass increases the attraction, such a decreased \( \rho \) mass increases the (repulsive) part of the Lorentz-Lorenz correction of the pion coupling to \( NN^{-1} \) and \( \Delta N^{-1} \) bubbles. In addition, the form factor \( \Lambda_\rho^* \) of the t-channel \( \rho \) exchange must be scaled. Also repulsive contact interactions in the \( \pi\pi \) scattering needed to preserve chiral invariance are proportional to \( f_\pi^{-2} \); they balance the increase in attraction from t-channel \( \rho \)-exchange. When all of the constraints from chiral invariance and BR scaling are enforced, the \( 2\pi \) exchange potential – which at low densities behaves approximately like a scalar meson with BR-scaling mass – gives an effective scalar interaction corresponding to a \( \sigma \) in which the decrease in mass is slowed down. This decrease in the rate of dropping enables saturation at the correct density. From the references in [46] the interested reader will be able to construct the unphysical artifacts that arise from neglecting the chiral constraints in the two-pion exchange interactions and convince oneself that they are necessary on physical as well as on formal grounds.

Note that with BR scaling, and the above modifications of it, the mechanism of saturation is quite different from that of introduction of three-body forces, as in Pieper et al [47], who carry out an essentially exact calculation of binding energy of light nuclei in the conventional approach with bare hadron masses (that does not explicitly implement the intrinsic density dependence required by matching to QCD). We suggest that these may not be relevant to the physical situation.

From the foregoing, we see that the dynamics in the “sliding vacuum” substantially changed the binding energy, but in many other respects is only slightly different from that in the perturbative vacuum at low densities. Indeed, now it is clear why Brown and Rho found in their early work the change in tensor force as indicative of dropping masses; there the \( \rho \)-meson, which does drop in mass, beats against the \( \pi \)-meson, which does not.

It is quite remarkable that in much of nuclear structure, it does not seem to matter whether BR scaling is operative or not. For instance, consider one of the most successful predictions in nuclear physics, the Dirac phenomenology. The most complete theoretical paper on this was by Clark, Hama and Mercer [48]. Predictions of this theory as compared with experiment were incredibly successful, the rapid oscillation in the polarization and spin rotation as function of angle being reproduced in great detail. Brown, Sethi and Hintz [49] persuaded John Tjon and Steve Wallace to put a linear scaling of meson masses with density into their relativistic impulse approximation (RIA) calculation. We reproduce Figs. 12 and 13 from Brown, Sethi and Hintz as Fig.4 and Fig.5 here.

Looking at these figures with a magnifying glass, one can see that the fit to data is not worsened by putting in density-dependent masses \( 7 \). We now outline the calculations that Tjon

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\(^7\)As noted [17], the vector coupling and associated constants are presumed to be non-scaling up to nuclear matter density.
Figure 4: Results of Tjon and Wallace (see Brown, Sethi and Hinz [49]) for the differential cross section of 500 MeV protons scattered elastically from $^{40}$Ca. The dotted line gives the IA2 results. For the solid curve a linear scaling with density was assumed, with $m^*_\sigma(n_0)/m_\sigma = m^*_\omega(n_0)/m_\omega = m_\rho(n_0)/m_\rho = 0.85$. Note that the two are essentially on top of each other. For the dashed line $m^*_\sigma$ and $m^*_\omega$ scaled in this way but $m_\rho$ was held constant at its free-space value. The experimental points are given as solid round dots.
Figure 5: Results of Tjon and Wallace for the spin observable $s$ for 500 MeV protons scattered elastically from $^{40}$Ca. The lines represent the same as in Fig. 4. Note that the dashed line starts deviating strongly from the others beyond about 30 degrees.

and Wallace have carried out, scaling meson masses other than the pion as

$$m^*_M/m_M \approx 1 - 0.15n/n_0$$

and nucleon mass as

$$m^*_N/m_N \approx 1 - 0.2n/n_0.$$  \hspace{1cm} (20) \hspace{1cm} (21)

This scaling is somewhat milder than what we now advocate $m^*_M/m_M \approx 1 - 0.2n/n_0$ and $m^*_N/m_N \approx 1 - 0.3n/n_0$. The Tjon-Wallace calculation, putting in the local density, was incredibly complicated, with 128 components of Dirac wave functions. Tjon ran part of the program, Wallace the rest, and they had to be together to complete it. Such a daunting calculation is highly unlikely to be repeated in the foreseeable future. So we will draw what conclusion we can draw from their results.

Before proceeding, we must point out several caveats here in re-interpreting the Tjon-Wallace results in terms of our language.

Whereas the IA2 used by Tjon and Wallace to check the scaling masses in Dirac phenomenology, with results given in Brown, Sethi and Hintz did not include the effective mass
\( m_N^* \), it did keep positive and negative energy states in a plane wave decomposition, and handled pair theory correctly. As shown in \([50, 51]\), in perturbation theory this is equivalent to use of a nucleon effective mass in Walecka mean field theory.

The \( m_N^* \) obtained in the above way by Tjon and Wallace was \( \sim 0.8m_N \) at \( n = n_0 \), somewhat less than the \( 0.7m_N \) that we advocate. Also we now have the scaling meson masses as \( \sim 0.8m_M \) at \( n = n_0 \). Thus their \( m_N^* \) dropped \( \sim 4/3 \) times faster than the scaling meson masses. The linear scaling was, however, consistent with present theory, which as we return to later, has \( m_N^* \) scaling \( 3/2 \) times faster than \( m_M^* \).

From the Tjon-Wallace results, we conclude that Dirac phenomenology is preserved well in all detail with the approximately linear scaling. In other words their fine structure of the polarization variables survives not only BR scaling, but even scaling with effective nucleon mass scaling somewhat faster than the meson masses. There is no discernable effect of the “breaking” of this scaling by the pion effective mass being kept constant. We thus find unfounded those claims that since conventional nuclear theory provides a successful description of the nuclear many-body problem, changes such as BR scaling should not be made in that it would upset the successful predictions. We have shown that much of nuclear structure – with a few notable exceptions which can be understood – are inert to the change of mass parameters.

5 In-Medium Pion Decay Constant

The fact that most of nuclear structure physics do not “see” the BR scaling does not mean that all nuclear observables are inert to it. We have mentioned that the tensor force in nuclei has the scaling \( \rho \) meson beating against the unscaling pion, which should leave a distinct imprint. In this section we discuss several observables that probe the scaling \( f^*_\pi/f_\pi \), the bona-fide order parameter of chiral symmetry in the hadronic language.

In section 2.1, the notion of parametric density dependence in HLS theory was described. Now when such HLS Lagrangian is applied to nuclear matter, the intrinsic density dependence factor \( \Phi \) can in turn be related to a Landau parameter which addresses many-body interactions as described in section 2.2, an indication that the change of QCD vacuum in medium is intricately tied to many-body properties. This demonstrates that it is most likely to be futile to try to separate what one would attribute to QCD from what amounts to many-nucleon dynamics.

Consider long wavelength fluctuations on top of the “vacuum” defined by such parameters. As effective degrees of freedom, we may pick the pions and the nucleons and integrate out vector mesons and other heavy hadrons including scalars from the HLS Lagrangian. The resulting Lagrangian will take the same chiral symmetric form as in the free space except that the parameters of the Lagrangian intricately depend on density. With this Lagrangian, the power counting will be formally the same as in ChPT based on the Lagrangian whose parameters are defined in matter-free space. According to the discussion given in Section 2.3, we have, up to near nuclear matter density,

\[
\Phi(n) \equiv m^*_V/m_V \approx F^*_\pi(n)/F_\pi \approx (\langle \bar{q}q(n) \rangle^*/\langle \bar{q}q \rangle)^{1/2}
\]

where \( F_\pi \) is the bare (parametric) pion decay constant. The (last) relation between the pion decay constant and the quark condensate comes from the Gell-Mann-Oakes-Renner mass formula for the pion in medium. The in-medium quark condensate can be estimated for low density in
terms of the pion-nucleon sigma term $\Sigma_{\pi N}$ using the low-energy theorem [52, 53]

$$\langle \bar{q}q(n) \rangle^*/\langle \bar{q}q \rangle = 1 - \frac{\Sigma_{\pi N}}{m_{\pi}^2 f_{\pi}^2} n + \cdots$$

(23)

Using the presently accepted value for the sigma term $\Sigma_{\pi N} \approx 45$ MeV, we get

$$\langle \bar{q}q(n) \rangle^*/\langle \bar{q}q \rangle \approx (1 - 0.36 n/n_0)^{1/2}.$$  

(24)

which gives $\langle \bar{q}q(n_0) \rangle^*/\langle \bar{q}q \rangle \approx 0.8$ at nuclear matter density. On the other hand, the quantity $\Phi$ of (22) can be obtained by relating the vector-meson mass scaling to the Landau parameter $F_1$ and then using the relation between the Landau parameter and the anomalous gyromagnetic ratio of heavy nuclei as well as the properties of nuclear matter [10, 36]. One gets

$$\Phi(n) \approx (1 + 0.28 n/n_0)^{-1}$$

(25)

where $n_0$ is normal nuclear matter density. At nuclear matter density, this gives $\Phi(n_0) \approx 0.78$. Thus the two relations (24) and (25) agree up to nuclear matter density. In what follows we shall adopt (25) since it directly reflects BR scaling. Furthermore the model-independent relation (24) depends sensitively on the value of the $\pi N$ sigma term, the precise value of which is still highly controversial and uncertain. For a recent summary, see [54].

Given (25), one can make a few independent predictions. The first quantity we can look at is Warburton’s $\epsilon_{MEC}$ [55] for axial-charge transitions in heavy nuclei. Using the chiral Lagrangian with the parametric dependence defined above, one can easily compute this in the leading chiral order $^8$, that is, in tree order [56]. The prediction is that [14]

$$\epsilon_{MEC}(n = n_0/2) \approx 1.63,$$

$$\epsilon_{MEC}(n = n_0) \approx 2.02.$$  

(26)

These should be compared with the observed values $\epsilon_{MEC}^{exp} = 1.60 \pm 0.05 (2.01 \pm 0.10)$ for the mass numbers $A = 50(208)$.

The next thing we consider is the recent experiment on deeply bound pionic atoms [57, 58]. There is an on-going discussion as to whether this experiment is signalling “partial restoration” of chiral symmetry. There can be a variety of ways to approach this problem. In the framework we are adopting here with the chiral Lagrangian based on HLS/VM, what the experiment provides is an information on the only scale dependent parameter in the theory, namely, the ratio $F_\pi^*(n_0)/F_\pi$ at $n \lesssim n_0$. In the tree order, this ratio is given by (25) which at nuclear matter density is

$$(F_\pi^*(n_0)/F_\pi)^2 = 0.61$$

(27)

with a theoretical uncertainty of $\sim 10\%$ inherent in nuclear $\delta g_l$. This agrees with the value extracted – in the tree order – from the pionic atom data [58]

$$(F_\pi^*(n_0)/F_\pi)^2 = 0.65 \pm 0.05.$$  

(28)

One can understand this result obtained in the leading order ChPT with HLS/VM in a way more familiar to nuclear physicists, as follows based on the works by Friedman [59] and

$^8$According to the chiral filter, the soft-pion term should dominate with next-order corrections suppressed for the axial-charge transitions.
Weise and coworkers [60, 61]. For instance, Friedman invokes density dependence in two places. One is to incorporate the ratio [61] \( f_\pi^*/f_\pi \approx \sqrt{\langle qq \rangle^*/\langle q\bar{q} \rangle} \approx 0.78 \) at \( n = n_0 \) which enters into the Weinberg-Tomozawa isovector term in the pion-nucleon scattering amplitude. This scaling can be identified [17] as the “intrinsic” density dependence required by matching to QCD. The other medium modification Friedman needs goes back to the relativistic impulse approximation used by Birbrair and collaborators [63] which gives small components of the Dirac wave functions for the nucleon enhanced by the factor

\[
F = \frac{m_N}{M(r)}
\]

where \( M(r) = m_N + \frac{1}{2}[S(r) - V(r)] \) with \( S(r) \) and \( V(r) \) the vector mean fields. In fact, if we write

\[
M(r) = m_N + S - \frac{1}{2}(S + V)
\]

and note that up to nuclear matter density \( (S + V)/2 \) turns out to be only about 10% of the \( S \) in magnitude, we see that \( M(r) \approx m^*_\pi(r) \), the nucleon effective (Landau) mass.

In Brown-Rho (BR) scaling [3], the nucleon effective mass scales more rapidly than \( m^*_\rho \) or \( f_\pi^* \) because of pionic effects in the solitonic background going as \( \sqrt{g^*_\rho/g_A} \). In the original formulation in terms of skyrmions [3], this factor entered when the Skyrme quartic term was taken into account for \( g_A \). Now the Skyrme quartic term is known to contain a lot more than what naively appears to result when heavy degrees of freedom, e.g., all heavy excitations of the \( \rho \) vector meson quantum numbers, are integrated out. It seems to represent physics ranging from part of pionic effects to extreme short-distance effects (e.g., proton decay [65]). (For discussions on some of these matters, see [27].) On the other hand, one can also calculate its effect from matching of chiral Lagrangian theory with sliding vacuum to Landau Fermi liquid theory as discussed in [10, 35, 36]. There this factor arose as a pionic contribution correlated with heavy modes in the vector-meson channel, given by the formula \( g_A^*/g_A = (1 - \frac{1}{3}\tilde{F}_1(\pi)\frac{m_\pi^*}{m_\pi})^{-2} \) where \( \tilde{F}_1(\pi) \) is the pionic contribution to the \( F_1 \) Landau-Migdal parameter which has the value \( \frac{1}{3}\tilde{F}_1(\pi) = -0.153 \) at nuclear matter density. With \( \Phi(n_0) = \frac{m_\pi^*(n_0)}{m_\pi} \approx 0.78 \), one obtains \( g_A^*/g_A \approx 0.80 \) or \( g_A^* \approx 1.0 \). It must be mentioned here that how these descriptions are related is not understood yet and remains an open theoretical issue. Now in finite nuclei the \( g_A^* \) is found almost universally to be unity [66]. Thus,

\[
m_N^*/m_N \approx \left( \frac{1}{g_A} \right)^{1/2} \frac{m^*_\rho}{m_\rho} \approx 0.7 \quad \text{at} \quad n = n_0.
\]

\[
\text{This actually goes back to the work of Lutz, Klimpt and Weise [62] for the value of } f_\pi^* \text{ obtained from the Gell-Mann-Oakes-Renner relation assumed to hold in matter with the pion mass unscaled by density. At low density, we expect this assumption to be valid. At present, the theoretical evidence for an unscaled pion mass comes from two-color (SU(2)\_c) QCD on lattice [38] although the error bars are a bit too big to confirm the constancy and of course the two-color QCD may not reflect the real QCD and from a skyrmion description [24].}
\]

\[
\text{Friedman invokes the in-medium pion decay constant expressed in terms of the pion-nucleon sigma term. As noted above, the numerical value turns out to be equivalent to } \Phi(n_0) \text{ used in this paper. The point we want to stress is that this quantity is the one that figures in BR scaling.}
\]

\[
\text{This observation is the basis of the suggestion made in [64] that fluctuations in nuclear matter be computed around the “shifted vacuum” at which } (S + V)/2 \text{ is exactly equal to zero. This “vacuum” at this point is close to the Fermi-liquid fixed point. Note that this is analogous to the notion of fluctuating around BR’s “sliding vacuum” [10, 24, 35, 36].}
\]
We suggest that this is equivalent to the nucleon mass scaling needed by Friedman. In our formulation, the phenomenological approach of Friedman corresponds to the leading-order treatment of our effective theory with the vector manifestation [7, 14] for which the only scaling parameter is the pion decay constant with the pion mass unscaled.\(^\text{11}\)

Note that we arrived at (27) without making direct reference to the quark condensate. It was based on the “many-body” relation presumably valid up to and near nuclear matter density and not beyond the Fermi-liquid fixed point. Relying on (22), one might naively conclude that the result signals a reduction of the quark condensate with density and hence might be taken as the signal that chiral symmetry is being restored. This inference is valid however only in the tree approximation used here and does not apply when one works at higher (loop) orders. The bare pion decay constant fixed at the matching scale does not necessarily follow the quark condensate near chiral restoration as density is increased unless one takes into account the parametric dependence that follows the renormalization group flow. This is so even though the physical pion decay constant vanishes at the chiral transition in the chiral limit. With the effective Lagrangian that most of the people in the field use, namely that in which the parametric dependence is absent, the behavior of the ratio as a function of density has little to do with chiral restoration, a point which seems to be often overlooked.

6 The Vector Manifestation and the Photon Coupling to the Nucleon

It turns out that in HLS/VM, the photon coupling to hadrons in medium can be quite different from that in free-space. For instance it has been shown by Harada and Sasaki [67] that the vector dominance in the photon coupling is strongly violated near the VM fixed point and hence near the chiral transition point. While experiments at Jefferson Lab probe densities near that of nuclear matter and hence rather far from the VM fixed point, the presence of nucleon in the medium is expected to drive the parameter \(a\) in HLS Lagrangian toward 1 from its value \(a = 2\) at the vector-dominance regime.\(^\text{12}\) This would have a strong ramification on properties of in-medium nucleon form factors. In this section, we briefly discuss what can be expected.

In HLS, the photon couples to the degrees of freedom involved in the theory as

\[
\delta\mathcal{L} = -2eagF_\pi^2 A^\mu \text{tr}[\rho_\mu Q] + 2ie \left(1 - \frac{a}{2}\right)A^\mu \text{tr}[J_\mu Q] \tag{32}
\]

where \(A^\mu\) is the photon field, \(Q\) is the electric charge matrix \(Q = \frac{1}{3} \text{diag}(2-1-1)\) and \(J_\mu\) is the vector current made up of the chiral field \(\xi\). In HLS theory, baryons could be thought of arising as skyrmions in which case \(J_\mu\) can be identified as the skyrmion current. Alternatively one could introduce “bare” baryon or quasiquark field into the HLS scheme in which case, one can think of \(J_\mu\) as a “bare” baryon current. The usual vector dominance (VD) picture corresponds to taking \(a = 2\) in which case there will be no direct coupling of the photon to the baryon just as there is none in the case of the pion (for which the second term of (32) would be of the form

\(^{11}\)In HLS/VM with the explicit vector degrees of freedom, the scaling parameters are the gauge coupling \(g^*\), the vector-meson mass \(m_\pi\), the pion decay constant \(F_\pi^*\) and the coefficient \(a\). However in medium \(a^* \approx 1\), \(g^* \approx g\), so \(m_\pi \) goes as \(F_\pi^*\) through HLS/VM relations. We are left with only one scaling factor as in the case where the vectors are integrated out.

\(^{12}\)The deviation from vector dominance in the photon coupling to the baryons has been known since some time. In fact, the chiral bag model [27, 68] provided a natural mechanism for such a deviation.
It turns out [69] that the vector dominance (VD) at $a = 2$ is on an unstable trajectory of renormalization-group flow of HLS theory with no connection to the trajectory that leads to the VM and that the fact that in nature the VDM seems to work in matter-free space and at $N_F = 3$ is merely an “accident.” In fact, in the presence of matter (temperature or density), the flow consistent with QCD is on the trajectory that leads to the VM fixed point $a = 1$. Even in the absence of matter, $a \approx 1$ seems close to nature although the VDM that works for the EM pionic form factor corresponds closer to $a = 2$. See [70].

The vector dominance picture is not a good one for the photon coupling to the nucleon, so $a = 2$ must hold poorly when nucleons are involved. Indeed, there is an indication that the photon coupling to a single nucleon is already near this fixed point. Historically a picture closely resembling this one was adopted by Iachello, Jackson and Lande in 1973 [71]. There the authors assume that the $\gamma$-ray couples to nucleon more or less equally (at low momentum transfer) through the vector meson and directly to a compact core. The presence of the small core in the proton, of $0.2 \sim 0.3$ fm is indicated also in the proton structure function in (deep) inelastic scattering [72]. In fact this is the 50/50 picture that arises at the “magic angle” in the chiral bag model [68, 27] used by Soyeur, Brown and Rho [73, 74] in analyzing nuclear form factors. In medium, we predict that the electromagnetic form factor will have the form for $a = 1$,

$$
\frac{e}{2} \frac{1}{1 - q^2/m_V^2} + \frac{e}{2} H(q^2)
$$

where $H$ is a slowly varying function of $q^2$ (with $q$ being the four-momentum transfer) with $H(0) = 1$ and $m_V$ is the parametric mass that enters at finite density. The photon point $q^2 = 0$ gives the correct charge. The $\gamma$-ray will couple to the dileptons in this half-way manner. Note that the dileptons discussed below will experience the same propagator suppression, namely, the correction factor

$$
F \approx \frac{1 + 1/(1 + Q^2/m_V^2)}{1 + 1/(1 + Q^2/m_V^2)}
$$

with $Q \equiv |q|$ but $Q$ is generally small, the dileptons being nearly back to back, so this is probably unimportant. Thus though vector dominance is violated (expected in the nucleon sector from other considerations than that of the VM), it is an adequate approximation in the dilepton calculation as we will discuss below.

7 “Sobar” Configurations and CERES Dileptons

There have been much discussions on the possible evidence for changes in hadron properties in the CERES dilepton data [75]. While the processes discussed above involve transition matrix elements with specific kinematics, the dilepton experiments measure spectral functions or more generally correlation functions averaged over density and temperature. For this purpose, we need to look at the spectral distribution in the hot and/or dense environment. In terms of the framework based on HLS/VM theory we are adopting in this paper, this means that we need to incorporate consistently into vector-vector correlation functions both quantum fluctuations with the parametric masses and coupling constants and thermal loop and/or dense loop effects generated in the renormalization-group flow. Now as pointed out in [76, 77], this means, for the dilepton processes, considering both BR scaling figuring in density-dependent parameters
that results from the CSD and the mixing to the “sobar” configuration $N^*N^{-1}$ computed as “fluctuations” on top of the soliton configuration. This requires that the double-decimation be consistently implemented. We shall refer to this procedure as “BR/RW fusion.” In the Rapp-Wambach approach (abbreviated as R/W) [75], the fusion has not been implemented: There, the second decimation is replaced by configuration mixing in lowest-order perturbation theory while the first decimation CSD is left out.

In several low-energy phenomena the R/W $\rho$-sobar provides most of the low-lying “$\rho$” strength as we shall outline, B/R coming in as the intrinsic effect to somewhat increase the effect. This is because the $\rho$-sobar has $\sim 20\%$ of the $\rho$-strength, at a low energy of 580 MeV. BR scaling can only move the remaining 80\% at the parametric $\rho$-mass to this energy at densities $>n_0$ (nuclear matter density), and such high densities have not been investigated experimentally. The $\rho$-sobar pushes the elementary $\rho$ up in energy, the states repelling each other; thus, in the region of the elementary $\rho$ the $\rho$-sobar and the $\rho$ “defuse.” We shall see in Section 7.1 and Section 8 that there is good empirical evidence for this, establishing that both R/W and B/R are present as required.

7.1 The role of the “$\rho$-sobar” in the $^3$He($\gamma, \rho^0$)ppn reaction

Before we treat the CERES dilepton production which will require the fusion of the “$\rho$-sobar” (or R/W) and BR scaling (i.e., B/R), we consider the process where the $\rho$-sobar plays the primary role. Two recent papers by Lolos et al [78, 80] have discovered the fact that the $[N^*(1520),N^{-1}]$ excitation decays $\sim 20\%$ of the time into a $\rho$-meson, in agreement with the 15 - 30 \% listed in the Particle Data Book and the $\sim 20\%$ found by Langgärtner et al [79]. We call the $[N^*(1520),N^{-1}]^{1-},J=1$ the “$\rho$-sobar” because when measured in finite nuclei, rather than when produced on a single proton by, e.g., $\pi + p$ interactions, it takes on a collective character with increased $\rho$-meson content due to the admixture of the elementary $\rho$ with density [76].

A clear evidence has been found that the $\rho$ production in the deuteron is dominated by the $N^*(1520)$, an element of the $\rho$-sobar. Of course the nucleon density of the deuteron is so low that one could hardly expect any appreciable medium dependence from this nucleus. Greater $\rho$ production should be seen in $^3$He, although the present experimental accuracy does not seem to be sufficient to show this. Here we give theoretical estimates for how much greater it should be with Rapp/Wambach alone and with Brown/Rho fused with Rapp/Wambach. These estimates can be easily extended to heavier nuclei.

From the $^3$He wave functions of Papandreou et al [81] we easily see that the average density in $^3$He is half that in nuclear matter. Although these authors could explain the experimental results with BR scaling alone, with our analysis we find that they had to decrease the $\rho$ mass twice too much by their mean field. As we shall develop, R/W which they also considered gives most of the experimental effects.

From [76] – that we shall refer to as BLRRW – we see that the mixture of the elementary $\rho$ into the $\rho$-sobar, which goes as the square of the matrix element divided by the energy difference increases linearly with the density $n$. This is because the $\rho$-sobar is a collective state, a linear combination of the excitations of all nucleons in the nucleus up into the $N^*(1520)$. The amount of $\rho$ admixed into the $\rho$-sobar is the same as the amount of $\rho$-sobar admixed into the $\rho$. This is displayed, as function of density, in Fig.1 of Kim, Rapp, Brown and Rho [77] in which it is shown that the sobar configurations can conveniently be incorporated into a massive Yang-Mills theory. For $n = n_0/2$, the $Z$ factor comes out to be $Z_\rho = 0.15$ for R/W; fusing with B/R
increases it to 0.23. This gives the increased $\rho$ content of the *in-medium* $\rho$-sobar, which then increases from 20% found by Langgärtner et al (15 - 25% in the Particle Data Booklet) to 40% with the fused R/W and B/R. If the same were to be obtained by R/W alone, the density would have to be increased by $\sim 50\%$.

The recent unpublished calculation on the Regina $^3$He($\gamma,\rho^0$)ppn reaction by Rapp show that compared with experiment [82], the R/W results are spread widely about the unperturbed zero-density $\rho$-sobar energy of 580 MeV, somewhat more upwards than downwards, by the large imaginary part which we estimate to be $\gtrsim 200$ MeV, 150 MeV zero-density width of the $\rho$ plus some in-medium width. The $\rho$-sobar width increases with energy spreading the strength more upwards. On the whole the R/W fit is good, except that too much $\rho$ strength is predicted in the region up to the elementary $\rho$ at 770 MeV. This will be improved by the fusing with B/R, the latter lowering the parametric $\rho$ mass in the Lagrangian by 77 MeV. The definite need for the *fusing* of R/W and B/R is seen, therefore, most clearly in the region of the elementary $\rho$ where the push upwards by R/W must be compensated for by the downward shift from B/R. Here R/W and B/R “defuse” for the elementary $\rho$. We return to this point in the next section on RHIC physics.

Asked by the experimentalists, Rapp extended his calculations to higher densities than present in $^3$He and $^{12}$C in order to increase the medium effect. As noted, an $\sim 50\%$ increase in R/W in the $\rho$-sobar region is effected by the fusion with B/R, which we believe to explain the greater collectivity apparently seen.

While the fusion of B/R with R/W does give a good description of the observed phenomena, the accuracy in measurements of the Langgärtner et al and Regina-Tokyo groups is not great enough yet to show the necessary increase in $\rho^0$ production cross section per nucleon in going from proton to $^3$He targets from medium effect. In Huber et al [80], the cross section for photo-production of the $\rho$ of $10.4 \pm 2.5$ $\mu$b on the three $^3$He nucleons is found for the 620 - 700 MeV range of photon energy ($m_\rho^* \sim 500$ MeV) whereas Langgärtner et find $\sim 3.5$ $\mu$b/nucleon for $\rho^0$ production. Possible uncertainties in the latter cross section are not given, but if they are comparable in magnitude to the Regina uncertainties, even with the upper limit of Regina and the lower limit of Langgärtner et al we could not achieve our estimated medium effect.

In the $^3$He ($\gamma,\rho^0$)ppn the low-mass $\rho$ strength is mostly explained by R/W, the B/R coming in to slightly enhance the low-mass strength, but chiefly to lower the high-mass $\rho$ strength, which in R/W alone fails to explain experiment. Thus the higher energy region in which R/W and B/R defuse shows definitively that both effects are present. (See the next section.)

### 7.2 The CERES: Fusion of B/R and R/W

We now go on to the fusion of the two effects described above in the dilepton production at CERN by the CERES collaboration in which R/W and B/R play about equal roles, the present error bar and binning in the experiment being too large to pin down their separate roles.

Ralf Rapp (private communication) has carried out calculations fusing B/R and R/W, using vector dominance. Although it has been shown in HLS/VM [83] that vector dominance is violated at high temperature, as noted above, this should not affect the numerical results. From the Rapp results, one can say with some confidence that the fit is improved with the fusion of B/R and R/W. However due to the error bars, one cannot say firmly that the fusion is *required*. As it stands, both R/W and B/R fit the data just as well within the error bars. It is clearly difficult to differentiate their separate roles by experiment in the low-energy regime. However we
shall see from the next section that this separation is straightforward in the high-energy region where they defuse.

8 RHIC Resonances

It should be clear in our discussion of the dileptons given above, also $\rho$-mesons produced by $\gamma$-rays on $^3$He and other nuclei, that it is intricate to separate the role of BR scaling from that of Rapp-Wambach. It is therefore significant and exciting that a direct measurement of the $\rho$-meson mass $^{13}m_\rho^*$ could be made in a pristine atmosphere where temperature effects are small and where the density can be well reconstructed.

Figure 6: The invariant mass distributions in pp and mid-central AuAu of the $\pi^+\pi^-$ system, with a transverse momentum cut $0.2 < p_t < 0.9$ GeV. From Fachini [84]. Contributions from specific resonances are indicated by different lines.

We show in Fig.6 the STAR/RHIC results [84]. The (preliminary) fits gave

$$m_\rho = 0.698 \pm 0.013 \text{ GeV for } Au - Au,$$
$$= 0.729 \pm 0.006 \text{ GeV for } pp. \quad (35)$$

The collisions were measured at $\sqrt{s} = 200$ GeV. There is a long history of the detailed $\rho$ spectral shape. The difference between its appearance in elementary reactions, $e^+e^- \rightarrow \pi^+\pi^-$ or $\tau \rightarrow \nu_\tau\pi^+\pi^-$ and hadroproduction reaction is well known. The latest Review of Particle Physics [85] averages the $\rho$ mass for these two sets of experiments separately, with a clear systematic difference of the order of 10 MeV between them:

$$m_\rho^{\text{leptoproduction}} = 775.9 \pm 0.5 \text{ MeV},$$
$$m_\rho^{\text{hadroproduction}} = 766.9 \pm 0.5 \text{ MeV}. \quad (36)$$

It was noticed back in the 1960's by Hagedorn and others that the particle composition in $pp$ can be well reproduced by statistical models, see, e.g., ref.[86]. The fitted temperature

13The “mass” involved here is the pole mass. However the difference between the parametric and pole is of higher order in the power counting and can be taken to be negligible for the discussion.
is about the same as the chemical freezeout temperature of \( \sim 165 \) MeV found at RHIC. In some sense the \( \rho \)-meson must be born in a heat bath, with Boltzmann factors which cut off the high energy end of the rather broad \( \rho \). Some discussion of this can be found in Shuryak and Brown [87] and in Kolb and Prakash [88]. We expect that many papers will be written on this subject, which is not within the scope of this paper. We can only assume that perhaps half or less [88] of the drop in the \( \rho \)-mass found in the Au-Au experiments comes from the heat-bath-related effects, and that these are about the same as in the pp experiment.

Next we discuss shifts in the \( \rho \)-mass from forward scattering such as \( \rho + \pi \rightarrow a_1 \rightarrow \rho + \pi \). There are many of these which enter as principal values as discussed in [89, 90] as well as in [87]. It seems difficult to get more than a few MeV out of these, the sign probably attractive. We note that a substantial part of the upward push in these resonance calculations comes from the \( \rho \)-sobar \( (N^*(1520)N^{-1}) \) which gives the Rapp-Wambach effects. Thus the RW and BR “defuse” for the elementary \( \rho \). In other words, the BR scaling must substantially overcome the Rapp-Wambach configuration mixing in order to produce a substantial downward shift. In fact the BR scaling shift will be, from section 2.3 (see [17]),

\[
m_\rho^* \simeq m_\rho (1 + 0.28n/n_0)^{-1}
\]

where \( n \) is the total baryon density in nonstrange particles.\(^{14}\) This comes about because for low densities \( m_\rho^* \) scales as \( f_\pi^* \) while the quark mass and the pion mass do not scale.

As a final approximation in determining the density at which the \( \rho \)-meson decays we might consider thermal (kinetic) freezeout, since the two \( \pi \)'s must come unscattered to the detector. The kinetic freezeout at RHIC can be obtained from a hydro-based fit to the data\(^{15}\)

\[
T_k \approx 100 \text{ MeV, } \mu_\pi \approx 81 \text{ MeV, } \mu_N \approx 380 \text{ MeV, } \mu_K \approx 167 \text{ MeV}
\]

which translate into the pion density \( n_\pi \approx 0.06 \text{ fm}^{-3} \) and

\[
n_{N+\bar{N}} \approx 0.0075 \text{ fm}^{-3},
\]

or only 1/20 of nuclear matter density.

To obtain a better approximation we must optimize the product of the \( \rho \) resonance formation and its decay, as shown in Fig.7. Ignoring the lifetime of the \( \rho \) in comparison with the other times – which underestimates the density at which the \( \rho \) is formed since all the \( \rho \)'s will follow the fireball expansion – the parameters for resonance formation and decay are found by Shuryak and Brown [87] to be

\[
T_k \approx 120 \text{ MeV, } \mu_\pi \approx 62 \text{ MeV, } \mu_N \approx 270 \text{ MeV, } \mu_K \approx 115 \text{ MeV}
\]

\(^{14}\)In Shuryak and Brown [87], the mass shift was associated, in the language of linear sigma model, with the “amplitude-field” fluctuation of the chiral field – or radius field \( R \) – not with the “phase-field” or pionic fluctuation. It is perhaps worthwhile to point out that this way of looking at the mass shift is totally equivalent to the nonlinear sigma model approach of Brown and Rho [3] where the “dilaton” field of trace anomaly \( \chi \) plays the role of the amplitude field in [87]. This point has been further clarified in [24]. In fact, it is shown that the approach of [3] is closer to the modern treatment based on HLS/VM theory of Harada and Yamawaki. In the language of HLS/VM, it is the intrinsic parametric dependence on density resulting from matching to QCD that plays an essential role in the mass scaling. The pionic fluctuation in the theory is subject, of course, to low-energy theorems of chiral symmetry.

\(^{15}\)P. Kolb, private communication.
with densities about 1.4 times higher than at kinetic freezeout. Taking into account that the density of excited baryons is double that of $N + \bar{N}$, and the above conditions, we find the total density in nonstrange baryons to be $0.21n_0$.

We note that in BR scaling, basically the total density in nonstrange quarks and antiquarks should be used, whereas in RW configuration mixing the assumption of these authors is that the resonances with $\rho$ quantum numbers built upon excited baryons will be only $\sim 50\%$ as efficient in configuration mixing as those built upon the nucleon. Thus, at a total baryon density with 2/3 made up of excited baryons, they would add the 1/3 in $N + \bar{N}$. In this way BR becomes somewhat larger than (and of opposite sign to) RW under RHIC condition.

At the resulting density $n \approx 0.21n_0$, and with the $\rho$-mass dropping $\sim 22\%$ at $n_0$, we find a decrease of $\sim 43$ MeV in $m_\rho^*$. However, the $\rho$-mesons themselves – as well as the $\omega$’s, to a lesser extent – are also a source of scalar density, i.e., they are composed of constituent quarks

Considering the $\rho$’s, we note that equilibrium through $\rho \leftrightarrow 2\pi$ continues basically down to thermal freezeout, and with $\mu_\pi \approx 62$ MeV, the pion fugacity will increase the number of $\pi$’s by $e^{\mu_\pi/T}$ and the number of $\rho$’s by the square of this factor. With the greater multiplicity, but smaller Boltzmann factor than the $\pi$’s, we find $n_\rho/n_\pi \sim 0.15$, with $n_\pi \sim 0.08/\text{fm}^3$, or an $n_\rho \sim 0.013/\text{fm}^3$ which considering that the $\rho$ presents 2/3 of the scalar source as a nucleon, gives the additional $\sim 10$ MeV to the downward shift of the $\rho$-meson. Thus our total shift from BR scaling is $\sim 53$ MeV, roughly 80% of the total observed downward shift. This seems to account for most of the shift left over when the temperature-driven shift is subtracted away.

The 150 MeV width of the $\rho$ is unchanged. Since the decay is p-wave, the width should decrease with the cube of the momenta of the pions into which the $\rho$ decays. This would cause an $\sim 30\%$ decrease in width. The fact that the width does not change means that compensating increases must be furnished by the resonances, both those above and below the $\rho$ in mass contributing with equal sign to the width. From the upper part of Fig.1 of Rapp and Gale [90], one can see that their resonances contribute $\sim 50$ MeV to the total width for mass $M = 700$ MeV at $T = 150$ MeV whereas in [89], the increase is somewhat larger. In any case, the STAR results

\footnote{Since pion is not a source of scalar density, although composed of constituent quarks, cancellations must insure that its mass is not changed; these in turn insure that the pions are removed as source of scalar density.}

\footnote{We should note that the two scalings (24) and (25) which coincide for $n = 0$ and $n = n_0$ give results that differ by $\sim 13$ MeV in the shift at $n \approx 0.21n_0$ relevant to the process in question. The shift given here is somewhat larger than that quoted by Shuryak and Brown [87]. This difference cannot be taken seriously as we really do not know precisely how the scaling interpolates between $n = 0$ and $n = n_0$.}
give a nice check that the increase in width of $\sim 50$ MeV (collision broadening) is cancelled by the decrease in width from lower penetrability.

The RHIC work has the great advantage over the dileptons that the energy resolution separates off the upper region in which BR and RW defuse, showing that BR must first cancel RW and then add some net downward shift to the $\rho$-mass, the general size predicted by BR scaling fitting in nicely with the experimental results.

9 Effective Forces in Nuclei

The $V_{\text{low}-k}(r)$ has supplanted the Kuo-Brown G-matrix [11] as effective interaction to be used in nuclear structure calculations. Integrating-out of the high frequency parts of the two-body interaction supplants all of the paraphernalia of off-shell particle energies, etc [91]. Here we discuss how the CSD (chiral symmetry decimation) may change some of the results of the FLD (Fermi liquid decimation).

The most extensive, but still only partial, investigation of the problem we discuss here was carried out by Hosaka and Toki [92] for the $2s, 1d$-shell matrix elements (with $^{16}\text{O}$ as closed shell). In this region of nuclei, up to $^{40}\text{Ca}$, empirical two-body matrix elements which fit well the experimental spectra have been determined. The most definite and important result of Hosaka and Toki is that “the central G-matrix elements are already well reproduced by using the free-space parameters and that they are extremely sensitive to the masses of $\sigma$ and $\omega$ mesons. This indeed provides a strong constraint from the phenomenological side on the way the meson masses, particularly $m_{\sigma}^{*}$ and $m_{\omega}^{*}$ should scale in medium. In fact, it turned out that they have to be correlated such that $m_{\sigma}^{*}/m_{\sigma} \sim m_{\omega}^{*}/m_{\omega}$ as long as the coupling constants are kept unchanged.” Note that the latter conclusion is consistent with what we have found in [17] and with the type-II analysis of nuclear matter by Song [35]. As for the former, it can be explained by the argument presented in the introduction, namely the effect of the strong $\rho$-meson tensor coupling which more than cancels the strong pionic tensor coupling at short distances and keeps the nucleon-nucleon potential small in magnitude (except for the strong short-range repulsion which is counterbalanced by the short-range attraction and integrated out to give $V_{\text{low}-k}(r)$). Thus off-shell effects are small.

“Encouraged by this fact” (i.e., that with BR scaling the central G-matrix elements were well reproduced), Hosaka and Toki “have calculated the G-matrix elements using the various meson-nucleon masses scaled nearly the same way while keeping the pion strength unchanged. This time our interest is concerned with the LS and tensor matrix elements, since they are generally in poor agreement with nature if the free-space masses are employed.” By using the scaled masses, however, they “could not see significant improvements in the comparison of the calculated and empirical matrix elements.” They find that generally the LS matrix elements are enhanced by the increase of the $\sigma$, $\omega$ and $\rho$ contributions and the tensor matrix elements are suppressed by the increase in the repulsive components due to $\rho$ exchange.

The little improvement from the dropping masses in the spin-orbit channel is disappointing, given the excellent agreement found by Tjon and Wallace with or without dropping masses in Fig.4 and 5. However, the rapid oscillations characteristic of the Dirac phenomenology, as compared with the nonrelativistic equations, come about from the interference in contributions from different densities in the nucleus. Such an effect would be missing in taking the scaled masses to be constant, independent of local density, as Hosaka and Toki did.
It was actually the decrease in the in-medium tensor interaction found in the 1989 and 1990 Brown-Rho papers which provided an empirical guide for the dropping $m^*_\rho$.

The strong spin-orbit interaction, which is the basis for the shell model, was obtained in the Walecka-type relativistic mean field calculations, a factor $\sim 2$ at $n = n_0$ greater than found in nonrelativistic Brueckner calculations. This spin-orbit interaction is very important in the single-particle energies, giving the splitting between the spin-orbit partners in the shell model. The Walecka theory gets sufficient spin-orbit splitting by using nucleon effective mass $m^*_N \sim 0.6m_N$, somewhat smaller than what we find, $m^*_N \approx 0.7m_N$. We actually get $m^*_N$ this low by assuming $g^*_A \approx 1$ in medium. Since one does not expect that the Gamow-Teller coupling constant (globally) goes below 1, this is the lowest we can expect.

The spin-orbit interactions due to scalar and vector exchanges that Hosaka and Toki took are of the form

$$V_{LS}^S = -\frac{g^2_S}{4\pi} \frac{m_S}{2m_N} (\frac{m_S}{m_N})^2 Z_1(m_Sr),$$

$$V_{LS}^V = -\frac{g^2_S}{4\pi} m_V (\frac{3}{2} + \frac{2f_V}{g_V})(\frac{m_V}{m_N})^2 Z_1(m_Vr),$$

where

$$Z_1(x) = (1 + 1/x)e^{-x}.$$  \hspace{1cm} (42)

Note that the volume integrals of $V_{LS}^S, V_{LS}^V$ are independent of $m_S$ and $m_V$ respectively, the masses going into rescaling the $r^2dr$ factor in the integral. In mean field, energies etc. depend chiefly on the volume integrals. However, an additional $m^{-2}_N$ is present in the spin-orbit term (also in the tensor interaction). Taking $m_N \rightarrow m^*_N$ would strongly enhance both of these interactions, say, by a factor of $\sim 2$ at $n = n_0$. Of course, the average density in the $s,d$-shell lies below $n_0$.

As noted, the need for nucleon effective masses was clear in Walecka theory and these effective masses fixed up the spin-orbit interaction. If $m^*_N$ scaled like $m^*_S$ or $m^*_V$, the scale invariance noted in Section 4 would still hold, being violated only by the constancy of the pion mass. But because of the "loop correction"

$$\frac{m^*_N}{m_N} \approx \sqrt{\frac{g^*_A f^*_\pi}{g_A f_\pi}}$$

and since $g^*_A$ tends to unity precociously in nuclei, the $m^*_N$ scales more rapidly than $f^*_\pi$ just up to nuclear matter density. So the $\sqrt{g^*_A/g_A}$ factor breaks the scale invariance, but judging from the Tjon-Wallace results, not in a serious way.

We therefore conclude that the tensor force is the only component in nuclear force to be substantially changed by BR scaling, bearing out the 1990 work [2]. Along the same lines and even earlier [95] evidence had been found for an in-medium $\rho$ with a factor of 1.5 times that in the free G-matrix [96]. This came from a study of 447 sd-shell binding energy data. Assuming the strength of the $\rho$-exchange potential be proportional to strength times (range)$^2$, this would correspond to $m^*_\rho \sim 0.82m_\rho$ if the strength is kept constant and the range changed. In 2s,
1d-shell nuclei the average density is well below \( n_0 \), so one would not expect the \( m_\rho^* \) to drop quite so much. In any case Brown et al [95] found this change to improve agreement with data.

It appears difficult to find “smoking guns” in the structure of nuclei for 20% changes in hadron masses in going from \( n = 0 \) to \( n = n_0 \) implied by BR scaling. In the case of \( \omega \) - and \( \sigma \)-exchange we saw from the Hosaka-Toki work that the effects cancel each other. However, the \( \rho \) meson has no low-mass isoscalar partner, and one might look elsewhere than in the tensor interaction for effects, especially in the symmetry energy, which is so important in astrophysical applications.

In mean field theory the effect of the vector-coupled \( \rho \) on the symmetry energy at nuclear matter density would be an increase by \( (m_\rho/m_\rho^*)^2 \approx (0.8)^{-2} = 1.56 \). However, modifications in the \( \rho \)-mass also affect the second-order tensor interaction, which, as we remarked in the Introduction, gives the main difference between \( ^3\!S_0 \) and \( ^1\!S_0 \) states.

It was shown by Kuo and Brown [97] that the closure works well to approximate the second-order tensor interaction; i.e.,

\[
V_T \frac{Q}{E} V_T \approx \frac{V_T^2}{E_{\text{eff}}},
\]

with \( E_{\text{eff}} \approx 250 \text{ MeV}^{19} \). Now

\[
V_T^2 \approx -\frac{1}{250 \text{MeV}}(3 - \tau_1 \cdot \tau_2)(6 + 2\sigma_1 \cdot \sigma_2 - 2S_{12})V_T(\pi + \rho)(r) \tag{45}
\]

where the final factor is the square of the radial part of the tensor term. In getting (45), we have used the identity \( \sigma \cdot A\sigma \cdot B = A \cdot B + i\sigma \cdot [A \times B] \) and omitted the term linear in \( \sigma \) which will vanish when averaged over angle. We see that in the triplet state the second-order tensor gives the contribution

\[
\frac{24}{250 \text{MeV}}V_T^2(\pi + \rho)(r) \tag{46}
\]

which explains why the triplet state is bound, and the singlet not. However this is also the contribution to the symmetry energy

\[
\delta V_{\text{symm}} = \frac{6\tau_1 \cdot \tau_2}{250 \text{MeV}}V_T^2(\pi + \rho)(r), \tag{47}
\]

which comes to \( \sim 1/4 \) of the central contribution. Relativistic Brueckner-Hartree-Fock calculations with the strong \( \rho \)-coupling of the Bonn potential [98] show the \( \rho \)-contribution to \( \delta V_{\text{symm}} \) just cancels that from the vector coupling of the \( \rho \), leaving the pion exchange (mostly in second order) on the source of symmetry energy. Thus we seem to be foiled, again, in nuclear structure physics a strong “smoking gun” for the BR scaling (or parametric dependence on density) of the \( \rho \) mass.

The conclusion is that in all nuclear structure observables so far probed, the parametric density dependence symptomatic of BR scaling is masked in an intricate way. This strongly suggests a “hidden symmetry” which seems to induce almost exact cancellation of possible smoking-gun signals, a phenomenon that is very much reminiscent of the Cheshire-Cat phenomenon expounded in [27].

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\textsuperscript{19}Note that \( E_{\text{eff}} > 175 \text{ MeV} \), the c.m energy corresponding \( p = 2.1 \text{ fm} \) which is the cutoff that figures for \( V_{\text{low-k}} \). In EFT, this means that the second order tensor effect should reside predominantly in the local counter term.
10 Conclusion: Return of the Cheshire Cat

The main conclusion we have arrived at in this paper is that up to date, no “smoking gun” signals have been observed in nuclear structure physics for BR scaling indicative of the spontaneous breaking or restoration of chiral symmetry, a basic element of QCD. This outcome, although somewhat disappointing and perhaps unappealing, is however very much in line with what we have been arguing since some years, namely, that low-energy strong interactions are strongly governed by the (approximate) Cheshire-Cat Principle (CCP). Here we would like to enumerate a few cases where the CCP can be applicable.

On the most fundamental level is the structure of the nucleon. As noted as early as in 1984 [100], the bag radius in the bag model of the nucleon, naively interpreted as the confinement size, is a gauge artifact. Here the bag boundary provides the space-time point at which QCD and hadronic variables are matched. The CCP statement is that the physics should not depend on where the matching is made. How this comes about through the combination of quantum anomalies and topology of skyrmions is detailed in [27]. What the Cheshire-Cat phenomenon implies is that there is a continuous map between QCD degrees of freedom and hadronic degrees of freedom for low-energy hadronic properties.

One can think of the HLS/VM à la Harada and Yamawaki highlighted in this paper as a CCP in momentum space with an additional remarkable feature that is new, namely the presence and importance of the vector-manifestation fixed point. The vector manifestation (VM) may very well be present also in the chiral bag/skyrmion picture as conjectured in [24]. The important point regarding the VM fixed point is that while physical processes away from the phase transition critical point – a few of which were considered above – may be given by “fusing” or “defusing” or by some combination of the two of the soabar and elementary-particle modes, thereby obstructing the clear evidencing of the “smoking gun,” the VM governs the scaling behavior of the physical parameters that enter the processes, such as e.g. sending the vector meson mass to zero in the chiral limit. How the chiral symmetry restoration and nuclear interactions manifest themselves in physical processes is therefore irrelevant.

A case that is relevant at the next level of fundamental nature is the connection between the BR scaling Φ reflecting the property of QCD vacuum and the Landau parameter $F_1$ as discussed in [9, 35, 36]. This illustrates a possible mapping between the complex vacuum structure of QCD characterized by the quark condensate in medium and many-body nuclear interactions embodied in the Landau parameters. If this identification is correct, then it will suggest an inherent ambiguity in delineating QCD effects from many-body hadronic effects conventionally treated in the standard nuclear physics approach (SNPA) discussed in Section 3.

Although not worked out in detail, we expect the issue of Gamow-Teller strengths in nuclei to be in a way quite similar to the relation between the BR scaling Φ and the Landau parameter $F_1$. It has been debated [99] since many years as to whether the “quenching” of Gamow-Teller strengths in nuclei as observed in giant Gamow-Teller resonances is due to “conventional” multi-particle-multi-hole effects (e.g., core polarization [101]) or to “exotic” effects (e.g., Δ-hole excitations [102]). The resolution to this debate [103] is that both are relevant in a way analogous to what happens in the Dirac phenomenology discussed in Section 4. As noted in Section 9, the tensor force predominantly excites (particle-hole) states at an energy $\sim 300$ MeV which is comparable to the Δ-hole excitation energy. Because of subtle cancellations between various terms of the same scale involving multi-particle-multi-hole configurations and Δ-hole configurations, the effective Gamow-Teller coupling constant $g^*_A$ for the transition to the lowest
Gamow-Teller state can be calculated equally well by saturating – in the standard EFT language – the “counter term” by the $\Delta$-hole configurations or by the multi-particle-multi-hole states or by both. Decimated down to the vicinity of the Fermi surface, both should reside in the “counter” term. If one were to look at the excitation functions going over a range of energies, one of course would have to be careful with the multitude of scales that figure for the specific excitations involved so as not to encounter the breakdown of the particular EFT one is using.

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