IU-NTC-92-32
Fermilab Library
0 1160 0019654 7

IU/NTC 92-32
OSU-0901

# Finite Nuclei in Relativistic Models with a Light Chiral Scalar Meson

R. J. FURNSTAHL
*Department of Physics*
*The Ohio State University, Columbus, Ohio 43210*

BRIAN D. SEROT
*Department of Physics and Nuclear Theory Center*
*Indiana University, Bloomington, Indiana 47405*

November, 1992

## Abstract

Relativistic chiral models with a light scalar meson appear to provide an economical marriage of successful relativistic mean-field theories and chiral symmetry. The scalar meson serves as both the chiral partner of the pion and the mediator of the intermediate-range nucleon–nucleon (NN) attraction. However, while some of these models can reproduce the empirical nuclear matter saturation point, they fail to reproduce observed properties of finite nuclei, such as spin-orbit splittings, shell structure, charge densities, and surface energetics. These deficiencies imply that this realization of chiral symmetry is incorrect. An alternative scenario, which features a heavy chiral scalar and dynamical generation of the NN attraction, is discussed.

*Submitted to Physical Review C*

Relating the observed properties of hadrons and nuclei to the underlying theory of quantum chromodynamics (QCD) is a major challenge, particularly since most existing nuclear data reflect the properties of QCD at low energies and large distances. While present and future experiments that probe extreme conditions of density and temperature, or high four-momentum transfer, should help to elucidate the connections to QCD, theoretical models to interpret these experiments and make these connections must still be calibrated to low-energy data. Furthermore, extrapolations from observed nuclear properties to extreme conditions will be reliable only if general principles of physics, such as quantum mechanics, special relativity, and microscopic causality, are maintained. To meet these requirements, one naturally turns to field-theoretic models based on low-energy degrees of freedom—hadrons—that are constrained by the underlying symmetries of QCD.

In this letter, we consider hadronic models with chiral symmetry, which is a symmetry of the QCD lagrangian with massless $u$ and $d$ quarks, and ask whether they can quantitatively describe basic features of nuclear phenomenology. Both the soft-pion theorems and the partial conservation of the axial-vector current (PCAC) in weak interactions imply that (approximate) chiral symmetry is a genuine feature of low-energy hadronic interactions. Furthermore, to construct models with reasonable pion dynamics, which is necessary for a detailed understanding of nuclear structure, this symmetry must be enforced.

Relativistic chiral models with a light scalar meson would seem to be an economical marriage of successful relativistic mean-field theories and chiral symmetry. In these models, the scalar meson serves as both the chiral partner of the pion and the mediator of the intermediate-range nucleon–nucleon (NN) attraction. Several variations have been proposed, all based on the linear sigma model plus a neutral vector meson ($\omega$); they are solved in the mean-field (Hartree) approximation, with or without zero-point vacuum contributions included. It is possible (in some cases) to fix the parameters in these models so that the empirical nuclear matter binding energy and equilibrium density are reproduced. This would appear to calibrate the models to low-energy bulk nuclear properties, so that meaningful extrapolations can be made to extreme conditions of temperature or density.

We argue, however, that fitting to infinite nuclear matter is insufficient, and that these models fail *generically* to reproduce observed properties of finite nuclei. In fact, the chiral models are overconstrained and contain strong many-nucleon forces, which preclude a light enough scalar meson if the empirical pion coupling is adopted. This leads to fundamental deficiencies in the description of spin-orbit splittings, shell structure, charge densities, and surface energetics. These deficiencies imply that this implementation of chiral symmetry is incorrect. On the other hand, we contend that phenomenologically successful mean-field models are consistent with an alternative scenario for chiral symmetry in hadronic models, which features a heavy chiral scalar and dynamical generation of the NN attraction.

The appeal of the mean-field approximation is based largely on the phenomenological success of the Walecka model mean-field theory (MFT),[1] which contains nucleons, neutral scalar mesons, and neutral vector mesons. The parameters of the model can be fit to reproduce the empirical binding energy and saturation density of infinite nuclear matter; when subsequently applied to spherical finite nuclei, the Walecka model provides a realistic

description of many bulk properties, such as charge densities, rms radii, and energy spectra.[2]

The basic physics responsible for this success is that the NN interaction in this MFT contains strong attractive Lorentz scalar and repulsive Lorentz vector components, which tend to cancel in the central interaction that generates the binding, but which add to produce a strong spin-orbit force that is consistent with the observed single-particle spectra. The vector meson is identified with the $\omega$ and assigned a mass of roughly 780 MeV, while the scalar mass is determined by fitting nuclear properties and typically takes values around 500 MeV. Although such a light scalar meson is not observed as a true resonance in nature, scalar meson exchange is an economical way to simulate more complicated processes (such as the exchange of two correlated pions) that are responsible for the mid-range NN attraction. This procedure is also consistent with successful modern boson-exchange models of the NN interaction, all of which require large contributions in the channel with scalar-isoscalar quantum numbers.[3–6] Thus the Walecka model MFT provides a simple, appealing picture of nuclear saturation and nuclear properties that is consistent with more detailed studies of the underlying NN interaction. Nevertheless, chiral symmetry does not play an obvious role in the Walecka model.

Chiral symmetry can be implemented at the hadronic level in two different ways. In the linear realization (*e.g.*, the linear sigma model[7,8]), an auxiliary scalar field is introduced as the chiral partner of the pion. The scalar and pion are combined into a chiral four-vector, and these fields mix under chiral transformations.[9] The symmetry prohibits a mass term for the nucleon in the lagrangian, and spontaneous symmetry breaking is used to generate a nucleon mass. In the nonlinear realization,[10] there is no scalar field, and the appropriate transformations are realized by the pions alone. For applications involving low-energy nuclear phenomena, the linear sigma model is particularly compelling, since there is an explicit scalar meson that can play the same role as the scalar contained in the Walecka model. But is chiral symmetry actually realized this way in nature? Should the chiral partner of the pion really be identified with the low-mass scalar field that provides the bulk of the mid-range attraction between nucleons? Or, to put it more bluntly, can we achieve a viable nuclear phenomenology with a single scalar meson serving in both roles?

Our answer is negative: chiral mean-field models using a light scalar meson *cannot* produce realistic results for finite nuclei. We survey chiral models that are capable of reproducing the empirical saturation point of nuclear matter and find that they fail generically to reproduce observed nuclear systematics. The most serious problems are that the spin-orbit interaction is much too small, which leads to unrealistic single-particle spectra and incorrect shell closures ("magic numbers"), the nuclear densities exhibit large oscillations that are unobserved, and the total binding energies (as well as the partitioning of the energy between volume and surface terms) are unphysical. Indeed, in most models, heavy spherical nuclei like $^{208}$Pb do not even exist.

To highlight these deficiencies, we compare the chiral results with those of successful calculations obtained in "nonchiral" mean-field models.[11–14] A more exhaustive comparison will be given elsewhere;[15] here we summarize the results. To accomodate a wide range of candidate chiral mean-field models, we start with a general lagrangian containing nucle-

ons, neutral scalar and vector mesons, and pions, which is invariant under (linear) chiral tranformations when $m_\pi = 0$.

After following conventional procedures to introduce a nonzero expectation value for the scalar field and then shifting the field, we obtain

$$\begin{aligned}\mathcal{L} = {} & \overline{\psi}[i\gamma_\mu\partial^\mu - g_{\rm v}\gamma_\mu V^\mu - (M - g_\pi\sigma) - ig_\pi\gamma_5\boldsymbol{\tau}\cdot\boldsymbol{\pi}]\psi \\ & + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma - m_\sigma^2\sigma^2) + \frac{1}{2}(\partial_\mu\boldsymbol{\pi}\cdot\partial^\mu\boldsymbol{\pi} - m_\pi^2\boldsymbol{\pi}^2) \\ & + \frac{g_\pi}{2M}(m_\sigma^2 - m_\pi^2)\sigma(\sigma^2 + \boldsymbol{\pi}^2) - \frac{g_\pi^2}{8M^2}(m_\sigma^2 - m_\pi^2)(\sigma^2 + \boldsymbol{\pi}^2)^2 \\ & - \frac{1}{4}(\partial_\mu V_\nu - \partial_\nu V_\mu)^2 + \frac{1}{2}m_{\rm v}^2 V_\mu V^\mu + \frac{1}{4!}\zeta g_{\rm v}^4 (V_\mu V^\mu)^2 \\ & - \eta^2\left(\frac{g_{\rm v}^2}{g_\pi}\right) M V_\mu V^\mu \sigma + \frac{1}{2}\eta^2 g_{\rm v}^2 V_\mu V^\mu(\sigma^2 + \boldsymbol{\pi}^2)\ . \end{aligned} \tag{1}$$

This lagrangian is similar to the one proposed by Boguta,[16] which uses a vector–scalar coupling to generate the vector meson mass. We generalize to allow for different vector–scalar and vector–nucleon couplings,[17] the possibility of a "bare" vector mass, and a $(V_\mu V^\mu)^2$ interaction. The result is the most general lagrangian with non-derivative couplings through dimension four that is consistent with linear chiral symmetry (for $m_\pi = 0$). (The justification for stopping at dimension four is given below.) The chiral symmetry breaking is determined entirely by the pion mass.

The MFT energy density for nuclear matter in these models can be written as

$$\begin{aligned}\mathcal{E}(M^*, \rho_{\rm B}) = {} & g_{\rm v} V_0 \rho_{\rm B} - \frac{1}{2}m_{\rm v}^2 V_0^2 + \frac{1}{2}\eta^2\left(\frac{g_{\rm v}^2}{g_\pi^2}\right) V_0^2 (M^2 - M^{*2}) - \frac{\zeta}{24} g_{\rm v}^4 V_0^4 \\ & + \frac{m_\sigma^2}{2g_\pi^2}(M - M^*)^2 + \frac{\kappa}{6g_\pi^3}(M - M^*)^3 + \frac{\lambda}{24 g_\pi^4}(M - M^*)^4 \\ & + \gamma\int^{k_{\rm F}} \frac{{\rm d}^3 p}{(2\pi)^3}\,({\rm p}^2 + M^{*2})^{1/2}\ , \end{aligned} \tag{2}$$

where $V_0$ and $M^* \equiv M - g_\pi\sigma$ are determined by extremization, $\rho_{\rm B}$ is the baryon density, and the spin-isospin degeneracy is $\gamma = 4$. We introduce the parameters $\kappa$ and $\lambda$ as coefficients of the cubic and quartic $\sigma$ terms. In the chiral mean-field models studied here, they are constrained to be

$$\kappa = -\frac{3g_\pi}{M}(m_\sigma^2 - m_\pi^2)\ , \qquad \lambda = \frac{3g_\pi^2}{M^2}(m_\sigma^2 - m_\pi^2)\ . \tag{3}$$

These same parameters, without constraints, are included in the most successful mean-field models, which are extensions of the original Walecka model.[11–14]

The quartic vector term with parameter $\zeta$ was suggested by Gmuca.[18] In principle it provides additional flexibility for the chiral models, but in practice it has only a small effect,

and it does not modify our conclusions. (See Ref. 15 for further details.) We therefore set $\zeta = 0$ in all calculations considered here. Since the model in Eq. (1) is, in general, not renormalizable, we can also add arbitrary higher-order meson self-couplings. However, since the (dimensionless) scalar mean field is $g_\pi \sigma / M \approx 0.35$ in successful models, higher-order chiral interactions do not produce noticeable effects in nuclei unless their coupling constants are enormous. We feel that forcing the reproduction of nuclear properties through strong, unobserved, very-many-body forces is unphysical, and we therefore omit higher-order terms. Similar conclusions obtain for derivative couplings because the gradients are small over the nucleus.

We also allow the possibility of including chiral one-loop vacuum corrections, as in Refs. 19 and 20. By applying the renormalization scheme of Ref. 19, we obtain for the zero-point energy density

$$\Delta\mathcal{E}_c(M^*) = \frac{M}{3\pi^2}(M - M^*)^3 - \frac{1}{3\pi^2}(M - M^*)^4 + \Delta\mathcal{E}(M^*) \ . \tag{4}$$

The first two terms can be incorporated by modifying the tree-level values of $\kappa$ and $\lambda$ [Eq. (3)], while the third term, which starts at $\sigma^5$, is the usual contribution included with Walecka-type models solved in the Relativistic Hartree Approximation (RHA).[21,12,13] Note that different renormalization conditions simply change the effective values of $\kappa$ and $\lambda$ in the general model.

To provide a minimal calibration to nuclear properties, we require the models to reproduce nuclear matter saturation properties, which we define as a binding energy/nucleon of 15.75 MeV at a density corresponding to a Fermi momentum of $1.3\,\mathrm{fm}^{-1}$. (Our conclusions are not sensitive to the precise values.) We only consider models and approximations that can reproduce these conditions.

The linear sigma model plus neutral vector mesons ($\eta = \zeta = 0$) is a natural starting point. Kerman and Miller[22] showed, however, that it is impossible to reproduce nuclear matter saturation properties at the mean-field level. This failure occurs because the spontaneously broken chiral symmetry leads to large nonlinear scalar self-interactions that produce strong, attractive three-body forces between nucleons. This attractive force is *opposite* to that found in the successful MFT phenomenologies[11–14] and it precludes a successful description of saturation.

By adding the zero-point energy of the baryons in the Dirac sea [Eq. (4)], one can generate enough extra repulsion to stabilize the system and to reproduce the empirical nuclear matter saturation point with an appropriate choice of parameters.[19] Alternatively, the model proposed by Boguta[16] [which corresponds to $\zeta = 0$, $\eta^2 = 1$, and $g_\mathrm{v}^2 = g_\pi^2(m_\mathrm{v}^2/M^2)$] generates both the nucleon and $\omega$ masses through spontaneous symmetry breaking and is able to reproduce nuclear saturation at the mean-field level.

These saturating chiral models arise as special cases of the lagrangian in Eq. (1). We consider a wide range of such models and solve for finite nuclei. The mean-field (Dirac–Hartree) equations for finite nuclei are derived by following conventional procedures, and the equations are solved by standard iteration methods. (See Refs. 2, 12, and 21 for

| Model | $g_s^2$ | $g_v^2$ | $m_\sigma$ | $\kappa$ | $\lambda$ | $\eta$ | $M^*/M$ | $K$ |
|---|---|---|---|---|---|---|---|---|
| A | 109.71 | 190.61 | 520.0 | 0. | 0.0 | 0.00 | 0.54 | 550 |
| B | 95.11 | 148.93 | 500.8 | 5000. | −200.0 | 0.00 | 0.63 | 220 |
| C* | 54.21 | 102.58 | 458.0 | 0. | 0.0 | 0.00 | 0.73 | 450 |
| D* | 179.56 | 25.08 | 767.2 | −25195. | 359.5 | 0.00 | 0.89 | 110 |
| E | 68.18 | 47.41 | 650.5 | −11165. | 98.2 | 1.00 | 0.78 | 670 |
| F | 112.36 | 47.41 | 835.1 | −23618. | 266.6 | 1.28 | 0.78 | 670 |
| G* | 24.88 | 17.30 | 322.0 | −1652. | 8.8 | 1.00 | 0.90 | 150 |
| H* | 112.36 | 17.30 | 684.3 | −15856. | 179.0 | 2.13 | 0.90 | 150 |
| I | 43.58 | 61.85 | 500.0 | −5273. | 37.1 | 0.61 | 0.67 | 940 |

Table 1: Parameters and properties for a variety of mean-field models. The models are identified in the text; an asterisk means that zero-point contributions are included. Values for the compressibility $K$ (in MeV) and $M^*/M$ are given at equilibrium density. Both $\kappa$ and $m_\sigma$ are in MeV. For models E through H, $\eta^2 = g_s^2 m_v^2 / g_v^2 M^2$.

more details.) Some care is necessary to obtain convergence in these nonlinear models, for example, by heavily damping successive iterations. To make realistic comparisions to heavy nuclei, we also couple the nucleon to Coulomb and isovector $\rho$ fields as in Ref. 2. The strength of the $\rho$ coupling is set by requiring the nuclear matter symmetry energy at equilibrium to be 35 MeV in all models. When the zero-point terms are included, they are evaluated with a local-density approximation.[13,23] Finally, charge densities are obtained by folding the point proton density with the empirical proton form factor.

When solving chiral mean-field models, one must also be aware of the possibility of bound "anomalous" solutions, in which the scalar mean field interpolates between the minima of the effective potential. These solutions have been discussed by Boguta and collaborators.[16,24,25,20] In anomalous nuclei, the nuclear density is concentrated at the nuclear surface, in strong contradiction to experiment. Nevertheless, when $m_\pi = 0$, the anomalous solution is typically more deeply bound than the "normal" solution.[16,15] In contrast, we have verified by explicit calculation that including a finite pion mass pushes the energy of the anomalous solution above that of the normal solution. (The pion mass is not small on the scale of nuclear binding energies!) Thus, we do not consider these anomalous solutions further here and set $m_\pi = 0$. (The pion mass has negligible effect on "normal" solutions.)

In Table 1, parameters are given for a representative sample of mean-field models. A detailed survey will be presented elsewhere.[15] In these models, the nucleon mass is 939 MeV and the vector meson mass is 783 MeV. The first three models are not constrained by chiral symmetry [*i.e.*, Eq. (3) does not hold], while the others feature a chiral scalar meson.

We focus on predictions for the charge density (Figs. 1 and 2), single-particle spectra (Fig. 3), and energy systematics. The Walecka model (labelled A [Ref. 2] in Table 1) serves as a baseline for these predictions. In general, successful mean-field models (of which model B [Ref. 12] is typical) predict the nuclear shell structure with quantitatively

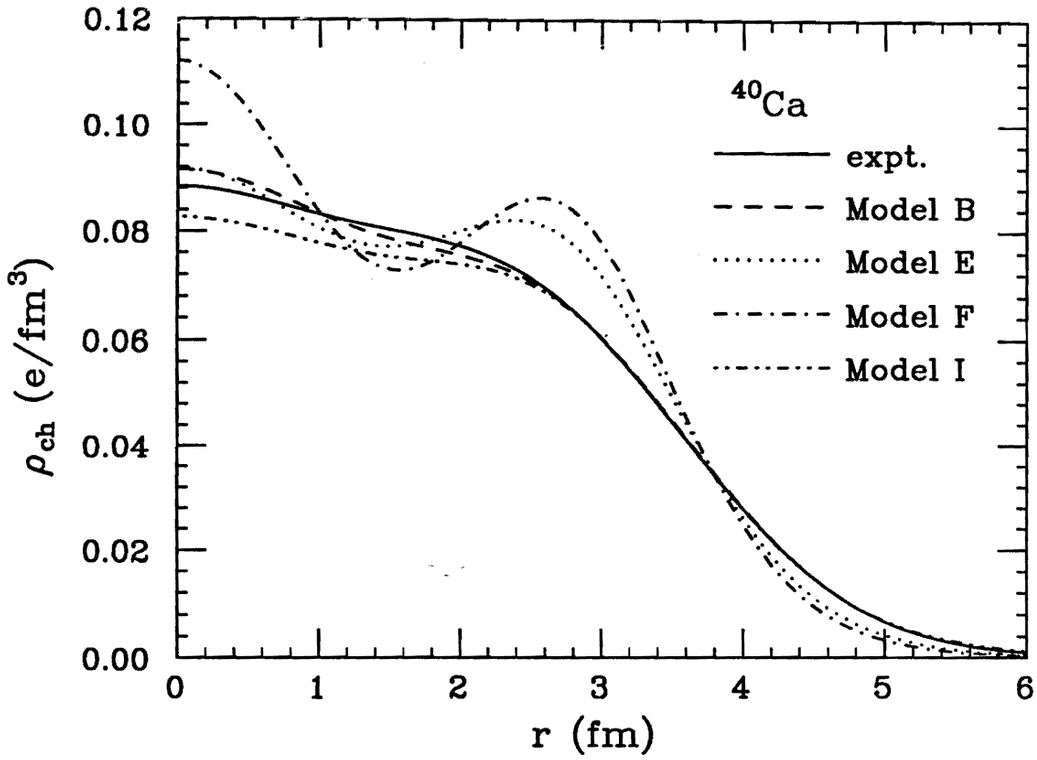


Figure 1: Charge density of $^{40}$Ca for several mean-field models specified in Table 1. The solid line is taken from experiment.[26] The dashed line is from model B, the dotted line is from model E, the dot-dashed line is from model F, and the double dot-dashed line is from model I.

accurate spin-orbit splittings (see Fig. 3), reproduce the systematic features of experimental charge densities (see Figs. 1 and 2), including the observed flatness for heavy nuclei, and generate binding energies and rms radii with realistic dependence on the mass number $A$ (see Ref. 13). Can a chiral mean-field model duplicate any or all of these successes?

Ideally, the scalar–nucleon coupling $g_s$ in a chiral mean-field model should be fixed at the empirical pion coupling $g_\pi \approx 13.4$ or at $g_\pi/g_A \approx 10.6$ through the Goldberger–Treiman relation. However, this condition effectively prevents a scalar meson mass light enough for reasonable finite nuclei. [Note that the mass by itself is irrelevant for static properties of uniform nuclear matter; see Eq. (2)]. A representative model of this type is model F [Refs. 16, 28], which is a variant of Boguta's model. The generic flaw is illustrated in Figs. 1 and 2: large scalar masses lead to large oscillations in the charge densities, which are not observed experimentally.

The scalar mass is reduced only marginally by adding zero-point contributions (models D [Ref. 19] and H [Ref. 17]), and thus there are still unacceptable oscillations in the charge densities. Furthermore, these contributions amplify another basic flaw of the chiral models: $M^*/M$ in nuclear matter is too close to unity. Vacuum contributions always drive the

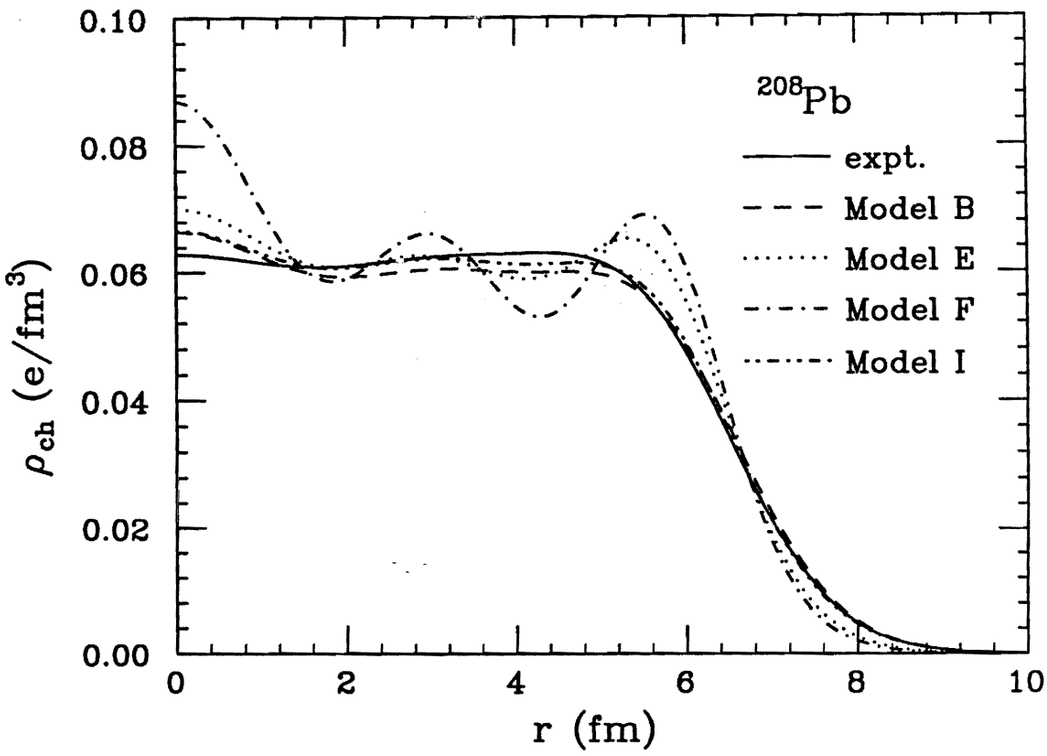


Figure 2: Charge density of $^{208}$Pb for several mean-field models. The solid line is taken from experiment.[26] The dashed line is from model B, the dotted line is from model E, the dot-dashed line is from model F, and the double dot-dashed line is from model I. Note that all lead nuclei are obtained by filling the "standard" set of single-particle orbitals, even though these may not be the lowest-energy levels in some chiral models (see Fig. 3).

solution to small fields (see also model C [Ref. 21]), and correspondingly large $M^*/M$, resulting in spin-orbit splittings that are too small. The disruption of the single-particle structure, including the major shell closures, is illustrated in Fig. 3; for models D–H, the uppermost proton shell closes at uranium or thorium rather than lead! We emphasize that the predicted spin-orbit strength, as reflected in both the shell structure and in calculations of proton–nucleus spin observables, is one of the major phenomenological motivations for relativistic models.

Even if we allow $g_s$ to be a free parameter, which leads to a lower scalar mass, we cannot lower $M^*/M$ sufficiently, particularly with zero-point contributions included (for example, consider models E [Ref. 16] and G [Ref. 17] in Fig. 3). Thus, even if we eliminate this direct connection to real-world pion physics, the models have serious flaws. Other problems can also be documented,[15] for example, excessive surface energies, which are correlated with a large nuclear matter compressibility and which lead to poor energy and rms radii systematics for finite nuclei.

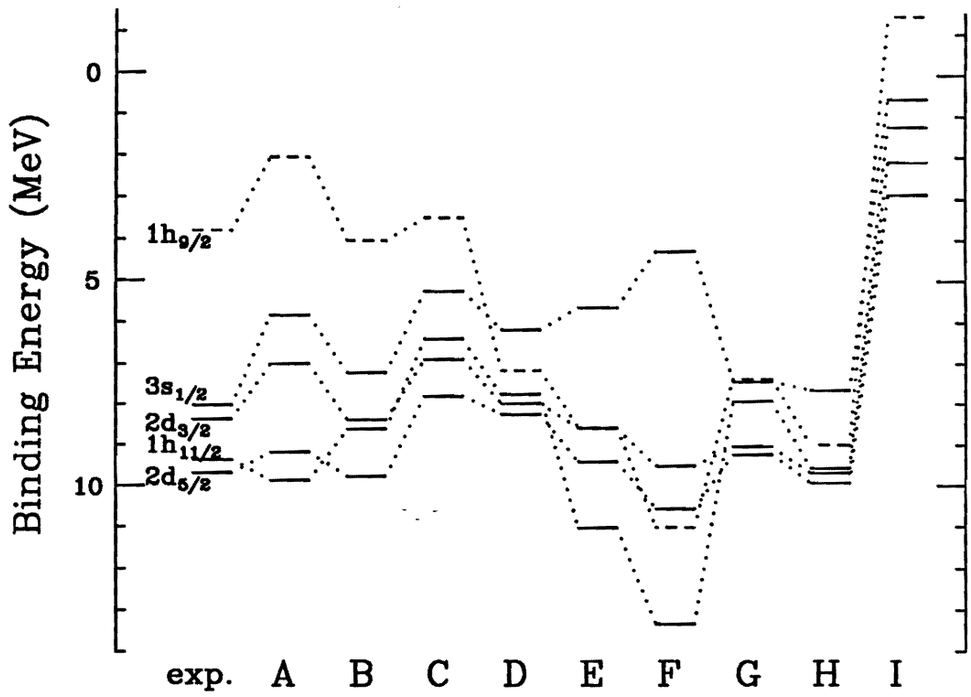


Figure 3: Predicted proton single-particle spectrum for $^{208}$Pb using models from Table 1. Only the least-bound major shell is shown. The left-most values are from experiment,[27] followed sequentially by the results for models A through I from left to right. Note that the $1h_{9/2}$ level, which is *unoccupied* in a real lead nucleus, is shown as a dashed line throughout.

Our calculations show that it is impossible to improve one or some aspects of a chiral model without degrading other features. Model I illustrates this point.[†] The full flexibility of Eq. (1) is used to set the scalar mass at 500 MeV, as preferred by successful models, and to reduce $M^*/M$ as much as possible. Even with an unrealistically small $g_s$, it is barely possible to reduce $M^*/M$ below 0.7 (successful models prefer 0.6 to 0.65), and this comes at the cost of an unacceptably large compressibility. The predicted binding energy/nucleon for $^{40}$Ca and $^{208}$Pb are 4.7 and 4.9 MeV, respectively (the experimental values are 8.5 MeV and 7.9 MeV), because the surface energies are roughly 70% larger than the empirical values. This is in contrast to an unconstrained mean-field model, such as B, which represents a "fine-tuning" of model A with simultaneous improvement in *all* characteristics of finite nuclei.

Based on our survey of chiral models and the characteristics of phenomenologically successful mean-field models, we can list the ingredients necessary for a mean-field model

[†]Here we use an equilibrium binding energy of 16.5 MeV at a density corresponding to $k_F = 1.38\,\mathrm{fm}^{-1}$ to produce more favorable results for nuclei.

to reproduce nuclear systematics at what we consider to be a minimal level of accuracy:

1. The correct bulk density and binding energy systematics of finite nuclei are largely ensured by requiring that the model reproduce "empirical" nuclear matter saturation properties. We stress that this is necessary but not sufficient! Observed properties of finite nuclei provide further constraints.

2. The scalar mass must be light enough to avoid strong fluctuations in the charge density, which signal impending instabilities.[29] The preferred value from "best fit" models is around $m_\sigma = 500\,\mathrm{MeV}$; $700\,\mathrm{MeV}$ (for example) is simply too large.

3. The effective nucleon mass $M^*$ must be small enough to imply a large spin-orbit force. Values of $0.6 \lesssim M^*/M \lesssim 0.65$ are found in "best fit" models. This implies that the mean scalar and vector meson fields are large, which in turn constrains the way that $\sigma^3$ and $\sigma^4$ terms can enter in order to avoid unrealistic density dependence. We note that large nucleon self-energies are supported by recent QCD sum rule calculations for nucleons in nuclear matter.[30]

4. The nuclear matter compressibility $K$ must be low enough to reproduce the nuclear surface realistically. This is reflected in the energy and charge density systematics with $A$. A desirable range is difficult to determine precisely, but $K \approx 200$–$250\,\mathrm{MeV}$ is found in models with good systematics, and values greater than $400\,\mathrm{MeV}$ are problematic. Models with low compressibilities ($K \approx 100\,\mathrm{MeV}$) cannot reproduce the systematics of density with $A$.[15]

Our study shows that chiral models with a light scalar meson can meet only the first criterion in general, and in some cases the last one. We note that some of these ingredients are only marginally satisfied by the original Walecka model. In particular, the spin-orbit force is somewhat too *large* and the compressibility is definitely too high. These deficiencies, however, are not fatal. By enlarging the model to include scalar meson self-couplings,[31,11,12] which allow for small adjustments in the density dependence of the interactions, the imperfections of the Walecka model can be eliminated *simultaneously* with relatively minor "fine-tuning." In the end, one can reproduce the bulk and single-particle properties throughout the periodic table as well as any existing microscopic model.[13,14] Although there are serious issues surrounding the extension of these calculations beyond the mean-field theory,[23] this relativistic MFT provides a successful nuclear phenomenology.

An important difference between the successful mean-field models and the chiral models is that the former are not constrained by chiral symmetry to fix $\kappa$, $\lambda$, and $g_s = g_\pi$. The clear implication is that chiral models with a light scalar meson have too many constraints. There is not enough freedom to adjust the dynamics to simultaneously achieve successful results for both nuclear matter and finite nuclei. Although nuclear matter saturation properties may seem reasonable, there are problems with finite nuclei that cannot be "tuned away." This is not simply a matter of adding more parameters to play with; additional nonlinear chiral couplings (like $\zeta$) do not help.

Thus we conclude that a linear realization of chiral symmetry with a scalar meson playing a dual role does not work. The generic failures of the mean-field chiral models are caused by systematic problems in the many-nucleon forces that must arise from the constraints of chiral symmetry and its spontaneous breaking. Although solutions found for finite nuclei are not pathological, they are simply too poor a representation of empirical nuclei to provide useful input to other calculations, such as those describing electron–nucleus or proton–nucleus scattering. Furthermore, the deficiencies imply that the physics is incorrect. In short, one should not use the chiral partner of the pion to also generate the mid-range, Lorentz scalar attraction between nucleons.

A different realization of chiral symmetry seems more compatible with the observed properties of nuclei and with successful relativistic mean-field models. Here one takes the chiral scalar mass to be *large* to eliminate the unphysical nonlinearities and then generates the mid-range attractive force between nucleons *dynamically* through correlated two-pion exchange.[32–36] This picture can be implemented with a nonlinear realization of the chiral symmetry, with the heavy chiral scalar merely playing the role of a regulator to retain the renormalizability of the linear $\sigma$ model.[23] The strong scalar-isoscalar two-pion exchange can be simulated by adding a low-mass "effective" scalar field coupled directly to the nucleons. Since the pion mean field vanishes, the resulting chiral mean-field theory looks just like the mean-field Walecka model. Moreover, although the coupling strength $g_s$ of the light scalar is comparable to $g_\pi$, we no longer require $g_s = g_\pi$ and are free to adjust $g_s$ within a reasonable range. In short, the Walecka model is consistent with the underlying chiral dynamics of QCD, although its realization in nuclear physics is subtle. We also observe that similar dynamics would be obtained directly in the linear model by going beyond mean-field theory, since inclusion of two-nucleon correlations would naturally generate sufficient mid-range NN attraction (through correlated two-pion exchange), and a light chiral scalar would no longer be needed (or desired).

Finally, we emphasize an additional conclusion: Requiring a mean-field model to reproduce nuclear matter saturation properties is *not* sufficient to ensure a quantitative description of well-established nuclear phenomenology. Any model claiming phenomenological success must realistically reproduce spin-orbit splittings, the nuclear shell structure, basic features and systematics of the charge densities, and surface energetics. Unless these constraints are incorporated, extrapolation from nuclear matter properties cannot be considered meaningful.

We thank B. C. Clark and R. J. Perry for useful conversations and critical comments. This work was supported in part by the U.S. Department of Energy under contract No. DE–FG02–87ER–40365, the National Science Foundation under Grants No. PHY–9203145 and PHY–9258270, and the Sloan Foundation.

## References


[1]J. D. Walecka, Ann. of Phys. **83**, 491 (1974).

[2]C. J. Horowitz and B. D. Serot, Nucl. Phys. **A368**, 503 (1981).

[3]G. E. Brown and A. D. Jackson, *The Nucleon–Nucleon Interaction*, North-Holland, Amsterdam, 1976.

[4]R. Vinh Mau, in *Mesons in Nuclei*, vol. I, M. Rho and D. Wilkinson, eds., North-Holland, Amsterdam, 1979, p. 151.

[5]M. J. Zuilhof and J. A. Tjon, Phys. Rev. **C24**, 736 (1981).

[6]R. Machleidt, Adv. Nucl. Phys. **19**, 189 (1989).

[7]J. Schwinger, Ann. of Phys. **2**, 407 (1957).

[8]M. Gell-Mann and M. Lévy, Nuovo Cim. **16**, 705 (1960).

[9]B. W. Lee, *Chiral Dynamics*, Gordon and Breach, New York, 1972.

[10]S. Weinberg, Phys. Rev. Lett. **18**, 188 (1967).

[11]P. G. Reinhard, M. Rufa, J. Maruhn, W. Greiner, and J. Friedrich, Z. Phys. **A323**, 13 (1986).

[12]R. J. Furnstahl, C. E. Price, and G. E. Walker, Phys. Rev. C **36**, 2590 (1987).

[13]R. J. Furnstahl and C. E. Price, Phys. Rev. **C40**, 1398 (1989).

[14]Y. K. Gambhir, P. Ring, and A. Thimet, Ann. of Phys. **198**, 132 (1990).

[15]R. J. Furnstahl and B. D. Serot, in preparation.

[16]J. Boguta, Phys. Lett. **120B**, 34 (1983); **128B**, 19 (1983).

[17]S. Sarkar and S. K. Chowdhury, Phys. Lett. **153B**, 358 (1985).

[18]S. Gmuca, Nucl. Phys. **A547**, 447 (1992).

[19]T. Matsui and B. D. Serot, Ann. of Phys. **144**, 107 (1982).

[20]J. Boguta, Nucl. Phys. **A501**, 637 (1989).

[21]B. D. Serot and J. D. Walecka, Adv. Nucl. Phys. **16**, 1 (1986).

[22]A. K. Kerman and L. D. Miller, *Second High-Energy Heavy Ion Summer Study*, Lawrence Berkeley Laboratory report LBL–3675, (1974).

[23]B. D. Serot, Rep. Prog. Phys. **55**, (1992), in press.

[24]J. Kunz, D. Masak, U. Post, and J. Boguta, Phys. Lett. **169B**, 133 (1986).

[25]J. Kunz, D. Masak, and U. Post, Phys. Lett. **186B**, 124 (1987).

[26]H. de Vries, C. W. de Jager, and C. de Vries, Atomic Data and Nuclear Data Tables **36**, 495 (1987).

[27]A. Bohr and B. R. Mottelson, *Nuclear Structure*, Vol. I, Benjamin, New York, 1969.

[28]V. N. Fomenko, S. Marcos, and L. N. Savushkin, preprint, 1992.

[29]C. E. Price, J. R. Shepard, and J. A. McNeil, Phys. Rev. C **41**, 1234 (1990); Phys. Rev. C **42**, 247 (1990).

[30]T. D. Cohen, R. J. Furnstahl, and D. K. Griegel, Phys. Rev. Lett. **67**, 961 (1991); R. J. Furnstahl, D. K. Griegel, and T. D. Cohen, Phys. Rev. C **46** , 1507 (1992).

[31]J. Boguta and A. R. Bodmer, Nucl. Phys. **A292**, 413 (1977).

[32]A. D. Jackson, D. O. Riska, and B. Verwest, Nucl. Phys. **A249**, 397 (1975).

[33]J. W. Durso, M. Saarela, G. E. Brown, and A. D. Jackson, Nucl. Phys. **A278**, 445 (1977).

[34]J. W. Durso, A. D. Jackson, and B. J. Verwest, Nucl. Phys. **A345**, 471 (1980).

[35]W. Lin and B. D. Serot, Phys. Lett. **233B**, 23 (1989).

[36]W. Lin and B. D. Serot, Nucl. Phys. **A512**, 637 (1990).