Quantum Hadrodynamics: Evolution and Revolution

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Abstract

The underlying philosophy and motivation for quantum hadrodynamics (QHD), namely, relativistic field theories of nuclear phenomena featuring manifest covariance, have evolved over the last quarter century in response to successes, failures, and sharp criticisms. A recent revolution in QHD, based on modern effective field theory and density functional theory perspectives, explains the successes, provides antidotes to the failures, rebuts the criticisms, and focuses the arguments in favor of a covariant representation.

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I. INTRODUCTION

Quantum hadrodynamics (QHD) refers to relativistic field theories for nuclei based on hadrons, in which the representation is manifestly covariant. The distinguishing empirical feature of covariant QHD is the presence of large (several hundred MeV), isoscalar, Lorentz scalar and vector mean fields (optical potentials) in nuclear matter at normal nuclear densities. During the last quarter century, QHD calculations have had numerous successes, but also apparent failures that arose when the dynamics of the quantum vacuum were computed with the same degrees of freedom used to describe the valence-nucleon physics. Moreover, there have been sharp criticisms over the years (some of which appeared in earlier issues of Comments) based on the relevance of a covariant approach to the nuclear many-body problem at observable densities, the use of relativistic, renormalizable quantum field theories with nucleon fields to describe composite hadrons and the quantum chromodynamics (QCD) vacuum, and the apparent lack of pion dynamics and chiral symmetry in the empirically successful calculations.

The underlying philosophy of QHD has evolved over the past twenty-five years in response to these successes, failures, and criticisms. Ultimately, the original motivation for renormalizable QHD lagrangians yielded to a more general approach based on the modern ideas of nonrenormalizable, effective relativistic field theories. This change in QHD philosophy is nothing short of a “revolution” that occurred during a two- or three-year span in the mid 1990s. The new approach, based on effective field theory (EFT) and density functional theory (DFT), allows us to understand the successful mean-field calculations of nuclear properties and how chiral symmetry works in QHD. It provides antidotes to earlier failures through a consistent, systematic, covariant treatment of the nuclear many-body problem, and it provides rebuttals to the previous criticisms.

In this Comment, we trace the QHD evolution and revolution, with special attention to how past criticisms and deficiencies have been nullified.

II. EVOLUTION

A. Original Motivation

Calculations of nuclear many-body systems based on relativistic hadronic field theories have existed for many years; one can trace their history back at least as far as the seminal work of Schiff in 1951 [1]. An important advance was made roughly 25 years ago, when Walecka enumerated the philosophy underlying a covariant description of nuclear matter, together with the rules for calculating within such a framework and a systematic program for investigating the observable consequences of the approach [2,3].

\[\text{Footnote: At the time of Walecka’s original paper, several other groups were studying relativistic, hadronic, field-theoretic approaches to the nuclear many-body problem, using similar but slightly different philosophies. See Ref. [3] for discussions of other approaches and results.}\]
To paraphrase the motivation presented so clearly by Walecka: to discuss neutron stars, it is necessary to have an equation of state that describes matter from observed terrestrial densities upward; a consistent theory should include mesonic degrees of freedom explicitly to allow for extrapolation to high densities; as the density of the matter is increased, relativistic propagation of the nucleons (and the retarded propagation of the virtual mesons) must be included; and causal restrictions on the propagation of excitation modes of the interacting system must be automatically contained in the theory. The basic conclusion following from these ideas is that, “The only consistent approach ... which meets these [objectives] is ... a local, relativistic, many-body quantum field theory.”

The systematic program to develop QHD as defined by Walecka focused on the nuclear many-body problem and relied on renormalizable lagrangian densities to define the models. This restriction was motivated by the desire to extrapolate away from the empirical calibration data in a manner that did not introduce any new, unknown parameters. Perhaps more importantly, the faithful pursuit of this framework would reveal whether renormalizable QHD was feasible and practical or not [3,4].

The original model contained neutrons, protons, and isoscalar, Lorentz scalar and vector mesons. In the nonrelativistic (Yukawa potential) limit (which was never actually used in the calculations), single-meson exchange generates basic observed features of the NN interaction: a strong, short-range repulsion and a medium-range attraction. It was assumed by fiat that the possible (renormalized) nonlinear couplings between the scalar fields were zero, for simplicity, although this restriction was relaxed quite early in the development by other practitioners. Moreover, no attempt was made to reconcile the model with the spontaneously broken, approximate $SU(2)_L \times SU(2)_R$ chiral symmetry of hadronic interactions.

The model contains large couplings, so a practical, nonperturbative approximation is needed as a starting point to describe the nuclear equation of state (EOS). Walecka argued that at high enough density, fluctuations in the meson fields could be ignored, and they could be replaced by their classical expectation values or mean fields. He also assumed that these conditions were sufficiently valid at ordinary nuclear density, so that the model could be calibrated in the mean-field approximation, and that the mean-field contributions would dominate the high density (e.g., neutron-star) EOS. Corrections to the mean-field approximation were calculated, and it was indeed found that in the context of renormalizable, Walecka-type models, the “stiff” EOS predicted by the mean-field theory (MFT), in which the pressure approaches the total energy density from below, becomes increasingly accurate as the nuclear density increases.

The hope was that nonrenormalizable and vacuum (short-range) effects, which had not been included in the corrections noted above, would be small enough to be described adequately by the long-range degrees of freedom of a renormalizable field theory, through the systematic evaluation of quantum loops. This hope was not fulfilled by explicit calculation of short-range loop effects. Furthermore, enlarging the nonlinear meson self-couplings to allow for effective, nonrenormalizable terms alters the high-density nuclear EOS qualitatively [3].

B. Successes

At the same time as the neutron-matter EOS was being studied, the MFT model was applied to the bulk and single-particle properties of doubly magic nuclei. The most impor-
tant conclusion of this early work was that calibration to the empirical equilibrium point of ordinary nuclear matter produces scalar and vector mean fields of roughly several hundred MeV at equilibrium density. When extended to finite nuclear systems, the resulting single-particle spin-orbit potential is roughly the same size as the observed spin-orbit potential, and thus the relativistic MFT predicts the existence of the nuclear shell model, without any ad hoc adjustments to the spin-orbit force. Moreover, MFT models that incorporated nonlinear (i.e., cubic and quartic) scalar field couplings reproduced bulk and single-particle nuclear observables as well as or better than any other concurrent models.

Relativistic MFT calculations have been performed for nuclei throughout the Periodic Table, with similarly realistic results and predictions; the reader is directed to the cited review articles for discussions of these numerous calculations [6–12]. When the MFT nuclear densities were folded with the free NN scattering matrix to compute proton–nucleus scattering observables (this is called the relativistic impulse approximation or RIA [13–17]), excellent descriptions of existing data and predictions for upcoming data were found—far superior to nonrelativistic calculations at the same level of approximation.

The vast majority of successful QHD predictions rely on the important observation that there are large, isoscalar, Lorentz scalar and vector mean fields in nuclei. Moreover, the successful calculations explicitly include only valence nucleons and long-range, many-body dynamics; the QHD degrees of freedom are designed precisely to describe this type of dynamics. It is also possible to study nuclear excited states in a random-phase approximation (RPA) that involves only long-range dynamics, but that still maintains the underlying symmetries of the lagrangian [18].

To explicitly include two-nucleon correlations, one uses the so-called Dirac–Brueckner–Hartree–Fock (DBHF) theory. With a covariant NN kernel that is fitted to two-body data and that contains large Lorentz scalar and vector (and pionic) components, one can simultaneously reproduce the nuclear matter equilibrium point at the two-hole-line level [19–21]. Moreover, although the correlation corrections produce changes in the MFT binding energy that are of the same order as the binding energy itself, the corrections to the large MFT scalar and vector self-energies (optical potentials) are small [22,23]. Nevertheless, to our knowledge, numerous approximations in the DBHF approach have never been quantitatively tested, and the systematic inclusion of contributions from the quantum vacuum was an unsolved problem [11].

C. Difficulties

Along with the numerous successes, there were also various difficulties that can be traced to a common source: the requirement that the QHD models be renormalizable. The difficulties fall into two basic classes: those arising from the computation of quantum loops (short-range physics) using long-range degrees of freedom, and those arising from attempting to maintain both renormalizability and chiral symmetry simultaneously.

Although the relativistic MFT results are encouraging, the QHD program sought a more complete description of the nuclear many-body system, which requires the development of reliable techniques to extend these calculations. Quantum loops are important for several reasons: loops ensure the unitarity of scattering amplitudes, baryon loops containing valence nucleons incorporate familiar many-body effects, loops introduce effects arising from the
modification of the quantum vacuum in the presence of valence nucleons, and meson loops in particular generate contributions to the extended structure of the nucleon.

Not all loops in QHD are problematic. For example, loops involving fluctuations of the pion field generate long-range effects (since the pion is light) and should be accurately described by explicit calculation within the QHD model (provided that approximate chiral symmetry is maintained). The strong, mid-range NN attraction arises from pion rescattering loops when two pions in the scalar, isoscalar channel are exchanged between nucleons. These long-range dynamical effects are much more efficiently described with hadrons than with QCD quarks and gluons; in fact, most QHD models go one step further and simulate the mid-range NN attraction using a Yukawa coupling to an explicit scalar, isoscalar field.

Problems arise with loops when one attempts to describe short-range dynamics using the heavier QHD degrees of freedom (nucleons and non-Goldstone bosons). In a renormalizable theory, a finite result can be obtained for the Casimir effect, and its addition to the MFT produces what is usually called the relativistic Hartree approximation (RHA). Although the new contributions are finite, they degrade the agreement of the nuclear predictions with experiment, particularly when one examines spin-orbit splittings for single-particle levels near the Fermi surface. These results imply that the QHD treatment of the quantum vacuum at the one-baryon-loop level is, at best, inadequate; although higher-order corrections might reduce the size of the one-loop terms, this can occur only through sensitive cancellations between relatively large contributions.

In fact, contributions from higher-loop terms within the renormalizable QHD framework do not improve the situation. Explicit calculations of the nuclear matter energy density at the two-loop level found enormous contributions that altered the description of the nuclear ground state qualitatively \[24\]. The conclusion was that the loop expansion does not provide a reliable approximation scheme in renormalizable QHD.

Similarly, vacuum contributions in the summation of ring diagrams produced unphysical poles at spacelike momenta in the meson propagators (sometimes called “ghosts”), which signal either an inconsistency in the QHD framework or an inadequate level of approximation. It was proposed that vertex corrections within the theory could solve these problems \[25,26\], but complete calculations involving vertex insertions proved to be impractical, and to our knowledge, no systematic, reliable approximation scheme for incorporating both long-range and short-range loop effects in renormalizable QHD theories has ever been found.

The second class of difficulties arises when one tries to embed the approximate, spontaneously broken, $SU(2)_L \times SU(2)_R$ chiral symmetry of QCD in a renormalizable QHD theory. The original Walecka model made no mention of chiral symmetry, but there were chiral models in use at that time based on the well-known Sigma model of Schwinger \[27\] and of Gell-Mann and Lévy \[28\]. In fact, if one takes the original Walecka model, adds massless pions that couple to nucleons with a pseudoscalar ($\gamma_5$) Yukawa coupling, and demands that the theory be Lorentz covariant, parity invariant, isospin and chiral invariant, and renormalizable, one is led to a lagrangian that is simply the Sigma model with an additional isoscalar vector meson. The nonzero nucleon and scalar masses are generated through the familiar spontaneous symmetry breaking, and given the similarity to the Walecka-model lagrangian, it is natural to identify the scalar field (which is the chiral partner of the pion) with the scalar field in the Walecka model.

The MFT for the chiral model can be motivated precisely as before, except that there
are now cubic and quartic scalar self-couplings that are not free parameters, but that are specified by the chiral symmetry. Unfortunately, the assumption that the chiral scalar field is the same as the Walecka scalar leads to dire consequences. First, it is impossible to reproduce the empirical nuclear matter equilibrium point in the MFT. Including quantum loops does not help, since one either generates extremely large contributions or arrives at uncertain results due to the appearance of unphysical poles in the meson propagators. An extensive mean-field analysis shows that it is impossible to generate realistic results for finite nuclei within this framework (see Refs. [29–31]). The conclusion is that the standard form of spontaneous chiral symmetry breaking, implemented in a model with a linear realization of the symmetry, cannot produce successful nuclear phenomenology at the mean-field level, if the chiral scalar field is identified with the scalar field in the Walecka model. This failure of the Sigma model is evidence that the simultaneous constraints of renormalizability and linear chiral symmetry are too restrictive.

Some progress was made on this problem by following the work of Weinberg and by making field transformations of the nucleon, pion, and scalar fields [32]. In addition to changing the form of the pion–nucleon interaction to include a pseudovector ($\gamma_\mu \gamma_5$) coupling, the transformation allows the introduction of a new scalar, isoscalar field that is not the chiral partner of the pion. This field plays the same role as in the Walecka model: it simulates important $\pi \pi$ and NN interactions that must be included from the outset to generate a realistic description of nuclear matter and nuclei.

A profound change has occurred, however, because in contrast to the original proposal of QHD as a renormalizable field theory, we are now forced to consider the new scalar field as an effective degree of freedom and the new chiral lagrangian as a nonrenormalizable, effective lagrangian. Thus the earlier philosophy of QHD must be generalized to include nonrenormalizable, effective field theories.

III. REVOLUTION

A. EFT/DFT Perspective

The revolution in QHD started with the reinterpretation of QHD lagrangians as nonrenormalizable EFT lagrangians. An effective lagrangian consists of known long-range interactions constrained by symmetries and a complete set of generic short-range interactions. The division between long and short is characterized by the breakdown scale $\Lambda$ of the EFT. While it is not possible at present to derive an effective hadronic theory directly from the underlying QCD, the EFT perspective implies that this is not necessary. If one constructs a general lagrangian that respects the symmetries of QCD: Lorentz covariance, parity conservation, time-reversal and charge-conjugation invariance, (approximate) isospin symmetry, and spontaneously broken chiral symmetry, then the EFT is a general parametrization of observables below the breakdown scale.

The EFT perspective, with the freedom to redefine and transform fields, implies that there are infinitely many representations of low-energy QCD physics. But they are not all equally efficient or physically transparent. One of the possible choices is between Lorentz covariant and nonrelativistic formulations. (In the context of EFT, these can be related by the heavy-baryon expansion [33].) Recent developments in baryon chiral perturbation
theory support the consistency (and utility) of a covariant EFT, with Dirac nucleon fields in a Lorentz invariant effective lagrangian density \[34,35\]. A similar framework underlies QHD approaches to nuclei.

For QHD, we identify \( \Lambda \) with the scale of non-Goldstone-boson physics (roughly 600 MeV). At momenta small compared to \( \Lambda \), short-distance physics (such as the substructure of nucleons) is only partially resolved and so may be incorporated into the coefficients of operators organized as a derivative expansion. The coefficients of these short-range terms may eventually be derived from QCD, but at present, they must be fitted by matching calculated and experimental observables. In principle, there are an infinite number of (non-renormalizable) terms, but in practice, the lagrangian or energy functional can be truncated to work to a given precision \[36\]. The EFT is useful if this truncation can be made at low enough order that the number of free parameters is not prohibitive.

In QHD, the only essential hadronic degrees of freedom are the nucleons and pions. The long-range pion–pion and pion–nucleon interactions are included in a nonlinear realization of chiral symmetry, which avoids dynamical assumptions inherent in linear representations. These interactions can be written down systematically (given a power-counting scheme) \[36\]. Low-mass vector mesons are typically included for phenomenological reasons, but are not required since their masses are of the order of \( \Lambda \); they are absent from point-coupling models, for example. In descriptions of NN scattering and of nuclear structure and reactions, the heavy bosons carry spacelike four-momenta and are “off the mass shell”; they therefore serve simply as a convenient way to parametrize the NN interaction in exchange channels with vector quantum numbers. This explains why it is useful to introduce collective degrees of freedom with other quantum numbers, such as a Delta baryon (with spin and isospin of \( 3/2 \)) to incorporate important pion–nucleon interactions. Because one must always truncate the lagrangian, these degrees of freedom can be efficient in the many-body problem whether or not they are actually observed as hadronic resonances.

A scalar, isoscalar mean-field in nuclei is an efficient way to include implicitly the effects of pion exchange that are the most important for describing bulk nuclear properties. Because chiral symmetry is realized nonlinearly, one can add a light scalar, isoscalar, chiral-singlet field to the theory and give it a Yukawa coupling to the nucleon, just as in the Walecka model. Nonlinear self-interactions of this new scalar must be included, with adjustable couplings that arise in part from the nucleon substructure. Since the expectation value of the pion field in nuclear matter vanishes at the mean-field level, one makes the remarkable observation that the MFT of Walecka-type QHD models is the same as the MFT of the chiral EFT model! Thus the MFT of Walecka-type models is consistent with chiral symmetry, provided we think in terms of a nonlinear realization of the symmetry. The light scalar, isoscalar field, which is not the chiral partner of the pion, plays the same role as in the Walecka model: it simulates important \( \pi\pi \) and NN interactions that must be included from the outset to generate a realistic description of nuclear matter and nuclei.

To make systematic calculations, the EFT approach exploits the separation of scales in physical systems, with the ratios of scales providing expansion parameters. A connection between appropriate QCD scales and nuclear phenomenology is made by applying Georgi and Manohar’s Naive Dimensional Analysis (NDA) and naturalness \[37,38\]. These principles prescribe how to count powers of the pion decay constant \( f_\pi \approx 94 \text{ MeV} \) and a larger mass scale \( \Lambda \) in effective lagrangians or energy functionals. The mass scale \( \Lambda \) is associated with the
new physics beyond the pions: the non-Goldstone boson masses or the nucleon mass. The signature of these low-energy QCD scales in the coefficients of a relativistic point-coupling model was first pointed out by Friar, Lynn, and Madland [39]. Subsequent analyses have extended and supplemented this idea, testing it in nonrelativistic mean-field models as well as in different types of relativistic models. Estimates of contributions to the energy functional from individual terms, based on NDA power counting, are quantitatively consistent with direct, high-quality fits to bulk nuclear observables [40,41]. Naturalness based on NDA scales has proved to be a very robust concept: nuclei know about these scales!

The successes of QHD mean-field phenomenology are, at first, rather mysterious from the EFT perspective alone, since the Hartree approximation is only the finite-density counterpart of the Born approximation at zero density. The density functional theory (DFT) perspective explains the successes of mean-field models and provides a new context for EFT power counting.

Conventional density functional theory is based on energy functionals of the ground-state density of a many-body system, whose extremization yields a variety of ground-state properties. In a covariant generalization of DFT applied to nuclei, these become functionals of the ground-state scalar density $\rho_s$ as well as the baryon current $B_\mu$. Relativistic mean-field models are analogs of the Kohn–Sham formalism of DFT [42], with local scalar and vector fields $\Phi(x)$ and $W(x)$ appearing in the role of relativistic Kohn–Sham potentials [12]. The mean-field models approximate the exact functional, which includes all higher-order correlations, using powers and gradients of auxiliary meson fields or nucleon densities.

The scalar and vector potentials are determined by extremizing the energy functional, which gives rise to a Dirac single-particle hamiltonian. The isoscalar part (for spherical nuclei) is

$$h_0 = -i \nabla \cdot \alpha + \beta (M - \Phi(r)) + W(r),$$

where $M$ is the nucleon mass and we define $M^* \equiv M - \Phi$. It is not necessary that $\Phi$ is simply proportional to a scalar meson field $\phi$. In fact, $\Phi$ could be proportional to $\phi$ (as in the original QHD models) or could be expressed as a sum of scalar and vector densities (as in relativistic point-couplings models) or could be a nonlinear function of $\phi$.

Density functional theory can provide a framework for the systematic incorporation of correlation effects, which are included exactly if the correct functional is identified. Mean-field models approximate this functional with powers and gradients of fields or densities, with the truncation determined by power counting. The inclusion of vacuum contributions, which was such a difficulty before, now becomes simple with the realization that all necessary counterterms are already present. The convergence of the EFT/DFT expansion is reasonable, but is slow enough that there are still too many terms to calibrate accurately by fitting to nuclear data. We rely on existing phenomenology as a guide to truncating the lagrangian most efficiently in light of ill-determined coefficients. In addition, while it is known that correlation corrections modify the scalar and vector self-energies by only a small amount (“Hartree dominance”), the mean-field energy functional omits possible nonanalytic terms; a combination of EFT and DFT may show us how to systematically include them [43,44].
B. Antidotes, Rebuttals, and Reinterpretations

In this section, we revisit past criticisms of QHD, many taken from earlier Comments \[45–47\] and others commonly expressed at conferences or elsewhere in the literature. We present without attribution a series of paraphrased statements in boldface that criticize different aspects of QHD and its predictions. Each statement is followed by a resolution based on the modern EFT/DFT perspective of QHD. We find that each criticism is either addressed and answered, revealed to be incorrect, or rendered moot.

Nuclei are nonrelativistic systems because corrections to the kinetic energy are small. Relativistic phenomenology for nuclei has often been motivated by the need for relativistic kinematics when extrapolating to extreme conditions of density, temperature, or momentum transfer. Unfortunately, this motivation obscures the issue of Lorentz covariant vs. nonrelativistic approaches for nuclei under ordinary conditions. Relativistic kinematic corrections are indeed small for ordinary nuclear systems. The important aspect of relativity in these systems is not that a nucleon’s momentum is comparable to its rest mass, but that maintaining covariance allows scalars to be distinguished from the time components of four-vectors. This distinction is easy to see by expanding the self-consistent nuclear matter energy density \( \mathcal{E} \) in powers of the Fermi momentum \( k_F \) (see Ref. \[12\], p. 554):

\[
\frac{\mathcal{E}}{\rho_B} = M + \left[ \frac{3k_F^2}{10M} - \frac{3k_F^4}{56M^3} + \frac{k_F^6}{48M^5} - \frac{15k_F^8}{1408M^7} + \frac{21k_F^{10}}{3328M^9} + \cdots \right] \\
+ \frac{g_s^2}{2m_s^2} \rho_B - \frac{g_s^2}{2m_s^2} \rho_B + \frac{g_s^2}{m_s^2} \rho_B \left[ \frac{3k_F^2}{10M} - \frac{36k_F^4}{175M^3} + \frac{16k_F^6}{105M^5} - \frac{64k_F^8}{539M^7} + \cdots \right] \\
+ \left( \frac{g_v^2}{m_v^2M} \right)^2 \left[ \frac{3k_F^2}{10M} - \frac{351k_F^4}{700M^3} + \cdots \right] + \left( \frac{g_s^2}{m_s^2M} \right)^3 \left[ \frac{3k_F^2}{10M} - \cdots \right] + \cdots. \tag{2}
\]

The corrections to the nonrelativistic kinetic energy (contained in the first term in brackets) are indeed small at equilibrium density, but the velocity dependence inherent in the Lorentz scalar interaction introduces significant corrections of higher than linear order in \( \rho_B \). The leading correction in each order is repulsive, so these corrections are important in establishing the equilibrium point.

In the nuclear medium, a covariant treatment implies distinct scalar and four-vector nucleon self-energies or optical potentials. The relevant question is: What are their natural mean values? QHD phenomenology implies several hundred MeV in the center of a heavy nucleus.

The success of nonrelativistic approaches shows that covariant approaches are wrong or unnecessary. Historically, the successes of nonrelativistic nuclear phenomenology have been cited to cast doubt on the relevance of large scalar and vector potentials. But in a nonrelativistic treatment of nuclei, the distinction between a potential that transforms like a scalar and one that transforms like the time component of a four-vector is lost. Because the leading-order contributions of these two types are opposite in sign, an underlying large

\[\text{Here } \rho_B = \frac{2k_F^3/3\pi^2}{3} \text{ is the baryon density, and } g_i \text{ and } m_i \text{ denote the scalar and vector couplings and masses, respectively. Nonlinear meson interactions are omitted for brevity.}\]
scale characterizing individual covariant potentials would be hidden in the nonrelativistic central potential. (The scalar and vector terms add constructively in the nonrelativistic spin-orbit potential, producing an uncharacteristically large result.) Furthermore, the EFT expansion implies that even potentials as large as 300 to 400 MeV are sufficiently smaller than the nucleon mass that a nonrelativistic expansion should converge, if not necessarily optimally. Thus the success of nonrelativistic nuclear phenomenology provides little direct evidence about covariant potentials.

The primary focus of nuclear theory is to fit the two-nucleon data and then to solve the many-body problem. The EFT coefficients are fixed by any sufficient data set; NN data has no special significance. Indeed, a density functional is best determined by finite-density data. Furthermore, NN data by itself cannot be sufficient. Many-body forces are inevitable \[48,44\], and their size can be estimated and shown to be non-negligible at ordinary densities (verified phenomenologically \[50\]). In a mean-field density functional to be used for medium to heavy nuclei, at least one three-body and two four-body parameters are necessary \[41\].

Why use a field theory? Potential models are easier. Relativistic quantum field theory based on a local lagrangian density provides a general parametrization of experimental observables consistent with the essential physics of the strong interaction: quantum mechanics, special relativity, unitarity, causality, cluster decomposition, and the intrinsic symmetries of QCD \[50\]. Consequently, we can describe the physics of low-energy QCD (e.g., ordinary nuclei) using an effective field theory with hadronic degrees of freedom. Field theory offers advantages over conventional potential models by systematically accommodating relativistic corrections and by allowing the construction of complete, consistent operators to describe interactions with external probes.

It is important to find specific observables that distinguish between relativistic and nonrelativistic theory. This pursuit will not be fruitful. There are field transformations that connect relativistic (covariant) and nonrelativistic theories, with the ratios of the fields to the nucleon mass acting as the parameters controlling truncation (see Refs. \[51\] and \[49\]). These parameters are small enough that a nonrelativistic approach should reproduce relativistic results, although not necessarily at the same level of approximation. The more appropriate question is: What is the most efficient representation?

Large potentials are an artifact of a relativistic formulation. We argue that the large potentials used in a covariant description of nuclear phenomenology are manifestations of the underlying mass scales of low-energy QCD, which are hidden in nonrelativistic treatments \[52\]. These QCD mass scales are inescapable if one considers the \(^1S_0\) NN phase shift (which becomes repulsive at about 250 MeV laboratory kinetic energy) together with the singlet scattering length of roughly \((8 \text{ MeV})^{-1}\) that signals an almost-bound state near zero energy \[53\].

It is more efficient to work with a nonrelativistic theory because large cancellations are built in. If there were an approximate symmetry that enforced the cancellation between scalar and vector contributions, then it would be desirable to build the cancellation into any EFT lagrangian or energy functional. (Chiral symmetry alone does not lead to scalar-vector fine tuning.) However, if the cancellation is accidental or of unknown origin, hiding the underlying scales may be counterproductive. We argue that nuclei naturally fall into the second category, with the relevant scales set not by the nonrelativistic binding en-
ergy and central potential (tens of MeV), but by the large covariant potentials (hundreds of MeV). The signals of large underlying scales are patterns in the data that are simply and efficiently explained by large covariant potentials, but which require more complicated explanations in a nonrelativistic treatment [52]. Examples of these are:

- The spin-orbit force, which appears automatically with the observed strength in a covariant formulation, but is not fully reproduced in even the most sophisticated non-relativistic calculations [54].

- Medium-energy proton–nucleus spin observables, which are reproduced by the relativistic impulse approximation with intuitive real optical potentials [13,16], while nonrelativistic treatments require full-folding and medium effects [55–58], and have nonintuitive potentials that change *qualitatively* with projectile energy.

- The energy dependence of the optical potential for nucleon–nucleus scattering up to 100 MeV, which is predicted at the relativistic mean-field level from the Lorentz structure of the interaction [3] (and higher-order corrections are small). In conventional nonrelativistic treatments, the energy dependence comes from the nonlocality of exchange corrections in a Hartree–Fock or Brueckner–Hartree–Fock approximation.

- The scalar, isoscalar part of the NN kernel below 1 GeV, which can be studied in an essentially model-independent way. Chiral symmetry, unitarity, and the natural strength of the $\pi\pi$ interaction imply an integrated strength that results in a large scalar single-particle potential [59]. We are not aware of a loophole here.

- The equilibrium of nuclear matter, which is not an ordinary, nonrelativistic Fermi liquid, since it is too dilute and too weakly bound. These characteristics arise in the mean-field energy/particle from an empirically small coefficient of the $k^3_F$ term in a density expansion [see Eq. (2)]. The cancellation between scalar and vector contributions to the nuclear matter binding energy account for this fine tuning.

**There is no experimental evidence of large scalar and vector fields.** There can be no *direct* experimental verification (or refutation) of any nuclear potentials. The evidence that a natural representation contains large fields, which is achieved only with a covariant formulation, comes from both empirical and theoretical analyses of NN scattering and nuclear properties [20,41,52]. As noted above, this manifestation of QCD scales translates in many instances into simpler, more efficient, more compelling explanations of nuclear phenomena than in nonrelativistic formulations.

Empirical support from nuclear properties comes from the study of covariant density functionals fit to nuclei [36]. A good fit to nuclear properties requires the local scalar and vector potentials to be roughly 300 MeV, and the hierarchy of energy contributions follow NDA predictions [41] (see Fig. 1). A more subtle argument is that the spread of 15 MeV or more among “realistic” nonrelativistic predictions of the nuclear matter equilibrium binding energy (the “Coester line”) [20] would be difficult to understand as calibration errors (“off-shell effects”), if the underlying scale of the two-body interaction were only 50 MeV. In contrast, large covariant two-body potentials in a relativistic formulation imply sizable three-body contributions in the corresponding nonrelativistic calculation that are consistent with
FIG. 1. Contributions to the energy/particle in $^{16}$O and $^{208}$Pb for two covariant, mean-field, point-coupling models. Absolute values are shown. The filled symbols are net values. The small symbols indicate estimates based on NDA, with the error bars corresponding to natural coefficients from 1/2 to 2. The equilibrium binding energy of nuclear matter is $\epsilon_0$. This spread. Finally, independent empirical support comes from fits of a covariant kernel for the NN interaction, which is used to calculate the NN scattering matrix. Every accurate fit has led to large, isoscalar, scalar and vector contributions of comparable magnitude, but of opposite sign, which translate in the medium into single-particle potentials of several hundred MeV at equilibrium density.

The pieces of evidence supporting a representation with large nucleon scalar and vector potentials, while not definitive when considered individually, collectively comprise a compelling positive argument.

We know that mean-field theory cannot be a correct description of nuclei because important long- and short-range correlations are omitted. Mean-field models are approximate implementations of Kohn–Sham density functional theory, which means that correlation effects are included in simple Hartree calculations. Moreover, the “Hartree dominance” of the single-particle potentials has been demonstrated, implying that short-range correlation corrections are no more than tens of MeV. The bulk properties of interest in mean-field phenomenology are primarily isoscalar observables that involve low resolution, so long-range pionic correlations are of minor importance; for other observables, EFT provides a systematic framework for explicitly including pionic contributions.

The relativistic mean-field approximation may be valid at high densities, but nuclei are low-density systems. The modern view is that the successes of relativistic mean-field theory do not depend on the justification of the mean-field approximation at high density, but on the flexibility of the mean-field density functional near equilibrium density. The combination of EFT and DFT concepts and methods applied to mean-field models of nuclei reveals that:

- NDA provides an organizational principle for the EFT. Power counting and the limited number of bulk nuclear observables explain the success of conventional mean-field
models, which contain fewer parameters than the most general EFT models.

- Vacuum effects, chiral symmetry, and nucleon substructure are all included in general QHD models.
- Ground-state nuclear properties provide information at low resolution. Models with different degrees of freedom (e.g., four- vs. two-component nucleons or point-coupling vs. meson models) are simply different organizations of the EFT. All are consistent with NDA.

Relativistic many-body calculations have unquantifiable errors. The EFT framework based on NDA and naturalness provides an organizational scheme for truncating a lagrangian or energy functional and for making well-defined error estimates (see Fig. 1). Vacuum corrections, which disrupted early attempts at QHD expansion schemes, are innocuous in the EFT approach. (They are automatically absorbed into the coefficients.) Moreover, there are (in principle) no off-shell ambiguities.

QHD calculations apply perturbation theory, which is not sensible with large coupling constants. In fact, QHD does have a sensible expansion, which is not in powers of the couplings. We work instead with density functional theory, with NDA power counting identifying reasonable expansion parameters. It is true that the short-distance (ultraviolet) behavior may be incorrect, but the EFT can correct the behavior systematically for low-energy observables using a small number of parameters (verified phenomenologically [41]).

Calculations of magnetic moments in relativistic models are inconsistent with the data due to enhancements from a small effective nucleon mass. Naively, the baryon current of a nucleus with a single valence nucleon with momentum $p$ outside a closed shell is $p/M^*$, compared to the Schmidt current $p/M$. However, if the calculation is forced to respect Lorentz covariance and the first law of thermodynamics, the nuclear current is constrained to be $p/\mu$, where $\mu \approx M$ is the chemical potential [60]. Thus there is no enhancement in a consistent relativistic framework.

Relativistic theories have ghosts. No they don’t. Ghosts arise from improper treatment of the ultraviolet behavior in renormalizable theories; in EFT, this short-distance behavior is included systematically by fitting a small number of parameters to nuclear observables. Moreover, the long-distance instability known as “Brown–Ravenhall disease” [61] does not arise in EFT, because the EFT framework is not quantum mechanics with a fixed number of particles.

The successes of the relativistic impulse approximation are meaningless because correlation corrections are large. On the contrary, relativistic correlation corrections to the optical potential are small, in contrast to the nonrelativistic framework (“Hartree dominance”). Thus the success of RIA calculations is expected. We emphasize that covariant and nonrelativistic expansions can have very different rates of convergence. The covariant representation appears superior in most instances [55,58].

G-parity implies that a successful description of NN scattering leads to unphysical predictions in the NN sector. Explicit calculations show that absorptive processes dominate the $\bar{N}N$ optical potential [62]. Thus the consequences of G-parity are not directly observable, since the transformed NN scattering amplitude is known only at unphysical kinematics for the $\bar{N}N$ system. Moreover, in the EFT framework, there are no
on-shell antinucleons; there are only valence nucleons. A simple extension of QHD to the NN sector pushes the EFT expansion beyond its breakdown scale Λ, so this extrapolation is suspect.

**QHD does not have pions and chiral symmetry.** Modern QHD effective lagrangians include pions in a nonlinear realization of chiral symmetry [36]. Confusion about the apparent absence of pions arises because explicit pionic contributions do not contribute to mean-field energy functionals, although correlated pionic contributions are implicitly contained in the effective scalar field. Long-range pionic contributions can be included systematically, but do not qualitatively change mean-field phenomenology [33, 34, 43].

The factorization of nuclear amplitudes into a product of on-shell, single-nucleon form factors and many-body amplitudes is incorrect. The modern chiral EFT lagrangians of QHD do not require such a factorization. The single-nucleon structure is included explicitly in the lagrangian through a derivative expansion [36]. This produces results that are very similar to the standard “folding” procedure [3].

The anomalous moment of the nucleon is clearly a property of its internal quantum structure; by itself, this precludes the representation of the nucleon as a local field. This is directly refuted by the EFT lagrangian [36], which not only accommodates an anomalous moment, it requires it!

Local meson fields and “point” nucleons provide no possibility for quark substructure. This is simply incorrect. At energy and momentum scales small compared to the underlying QCD scale Λ, details of the quark substructure are not resolved. It follows that the substructure can be incorporated through a systematic expansion of nonlinear and gradient interactions in the effective lagrangian, with the dynamics encoded in the local hadronic couplings. This is the essence of the EFT approach. A clear example of this expansion is the single-nucleon structure included in modern QHD chiral lagrangians [36].

Virtual nucleon–antinucleon Z graphs, which are essential to relativistic phenomenology, should be suppressed because the nucleon has substructure. A local Dirac field for the nucleon does not imply a physical point nucleon [3, 36]. Moreover, the virtual NN pair is far off shell, and the off-shell intermediate states in a particular representation cannot be interpreted in terms of on-shell physics. The EFT framework ensures that any incorrect short-distance dynamics can be corrected systematically with counterterms. Recent formulations of covariant chiral perturbation theory also verify that implicit Z graphs are not a problem, and that consistent power counting is possible [34, 35].

The QHD treatment of the vacuum neglects nucleon substructure, violates $N_c$ counting rules, and relies on unphysical NN contributions (Z graphs). The modern QHD treatment of vacuum dynamics has changed this discussion completely. Vacuum contributions in the EFT framework are not calculated explicitly but are implicitly and systematically contained in a small number of fitted parameters. Any physical consequences of hadronic substructure, for example, are automatically included. Furthermore, these implicit contributions are consistent with $N_c$ counting using NDA [64].

Hidden QCD color is relevant for low-energy nuclear physics and is not contained in QHD [45]. The EFT perspective says this objection must be irrelevant, because all observable amplitudes are color singlets (since color is confined). Regardless of the underlying QCD picture, the EFT must be valid at sufficiently low energies, without explicitly invoking colored degrees of freedom.
Quantum chromodynamics of quarks and gluons is the fundamental theory of the strong interaction, so we should describe nuclei in terms of quarks. For most ordinary nuclear phenomena, a description based on hadronic degrees of freedom is most appropriate: Hadrons are the particles actually observed in experiments and thus are more efficient. Hadronic calculations can be calibrated using empirical nuclear properties and scattering observables. Hadronic models have historically provided accurate descriptions of NN scattering and the bulk and single-particle properties of nuclei. It is better to match the effective hadronic theory to QCD to determine its coefficients and then use the EFT to calculate nuclear structure and reactions.

IV. DISCUSSION

A. Summary

The hadronic theory of QHD is truly a manifestation of QCD in the strong-coupling regime. It is currently impossible to construct this theory directly from QCD, but the effective field theory perspective shows that we can make progress regardless, without recourse to ad hoc models.

The history of QHD from 1974 shows an evolution driven by the successes and difficulties of the original approach. One of the cornerstones of this approach was the necessity for a consistent, microscopic treatment of nuclear systems using hadrons \(^\text{[2]}\). Although the mean-field theory was phenomenologically successful, the ultimate goal was to improve upon this approximation to incorporate both many-body and short-distance (quantum vacuum) effects. Unfortunately, this goal was sometimes overlooked, because numerous “improvements” degraded the quality of the mean-field-theory results; this often led to the imposition of arbitrary constraints or restrictions on QHD calculations. Fortunately, however, perseverance within the original framework ultimately led to the conclusion that the constraint of renormalizability is too restrictive, and alternatives to this requirement were sought.

The result is the modern viewpoint of QHD based on effective field theory and density functional theory (the “revolution”), which solves the most serious problems while preserving intact the successful predictions for bulk and single-particle nuclear observables. The EFT framework identifies the systematic “organizing principle” behind the successful QHD calculations: energy scales arising from the underlying QCD define the dimensional analysis for terms in the effective lagrangian. Naturalness and the size of the nuclear mean fields allow for a practical expansion and truncation, and also clarify the scope and limitations of QHD. Density functional theory and the Kohn–Sham formalism then explain why the truncated mean-field energy functional can be flexible enough to yield accurate results for certain nuclear observables.

Analyses based on QHD, as defined here, provide a correct description of baryonic systems at sufficiently large distances and low energies. But we argue further that a covariant formulation of the dynamics manifests the true energy scales of QCD in nuclei and provides an efficient and comprehensive explanation of observed bulk and single-particle systematics.
B. Outlook

The development of QHD is far from over, and there are many issues to be addressed. Whereas the majority of existing QHD calculations focus on isoscalar physics, the incorporation of pions using a nonlinear realization of chiral symmetry and the inclusion of the Delta baryon as a collective $\pi N$ degree of freedom should produce accurate results in the isovector sector; this merits further study. Calculations of excited states using a consistent, conserving approach to the random-phase approximation [65], coupled with new, more accurate data on nuclear breathing modes [66], could lead to a more precise determination of QHD parameters [41]. The situation for nuclear currents and magnetic nuclear form factors at low momentum transfer is still an open problem that must be re-examined in the context of modern EFT. It is also important to pursue the connection between QHD results for many-body systems and the recent calculations of few-nucleon systems within the EFT framework, including covariant chiral perturbation theory. A major challenge is to develop and apply systematic and consistent “power counting” schemes that lead to more general conserving approximations and to study renormalization-group methods that could determine the analytic structure of the ground-state energy functional [44]. In addition, we must learn how to extrapolate beyond the breakdown scale of the EFT description.

Ultimately we must answer the question: What is the best way to connect nuclear phenomenology to QCD? Quark models have been advocated, but they continue to lack a systematic framework and a direct connection to QCD. The EFT perspective makes explicit quark degrees of freedom nonessential. Nevertheless, finding an efficient, tractable, nonperturbative way to match the QCD lagrangian to the long-range, strong-coupling, effective field theory of QHD is a major goal for the future.

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