Nuclear Matter Properties in Derivative Coupling Models Beyond Mean-Field Approximation

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The structure of infinite nuclear matter is studied with two of the Zimanyi - Moszkowski (ZM) models in the framework of a relativistic approximation which takes into account Hartree terms and beyond and is compared with the results which come out of the relativistic Hartree - Fock approach in the linear Walecka model. The simple treatment applied to these models can be used in substitution to the more complicated Dirac - Brueckner - Hartree - Fock method to perform future calculations in finite nuclei.

Conventional many- body calculations which do not consider mesonic degrees of freedom are not reliable for the study of nuclear matter where high densities are present. For these purposes, a relativistic calculation must be performed. Quantum chromodynamics (QCD) is the fundamental theory of the strong interaction and hence it should explain possible modifications of hadron properties in the nuclear medium. However, typical nuclear phenomena at intermediate and low energies cannot be analitically derived from QCD although one hopes that QCD will be solved numerically on the lattice in a near future. Meanwhile we are left with the construction of phenomenological models in order to try to describe nuclear phenomena and bulk properties. Walecka and collaborators used such a kind of theory for the first time around 1974 to describe the nucleon - nucleon interaction [1]. Since then the Walecka model [2] has been widely used to describe the properties of nuclear matter as well as some properties of finite nuclei. This model is based on a phenomenological treatment of the hadronic degrees of freedom and, because of this fact, it is also known as QHD-I (quantum hadrodynamics), which consists in a renormalizable relativistic quantum field theory. The early version of this model considers a scalar (σ) meson field and a vector (ω) meson field coupled to the baryonic (nucleon) field. Besides the relativistic mean field calculation, the Walecka model has also been used in a more complete treatment, the relativistic Hartree- Fock approximation [2], [3], [4].

Some of the drawbacks of the model are that the effective nucleon mass obtained at high densities is too small and its incompressibility at the energy density saturation is too large. To avoid this problem, Zimanyi and Moszkowski [5] introduced an alternative coupling between the scalar meson and the nucleon (which is a scalar derivative coupling) and another coupling between the scalar and vector mesons. The inclusion of these couplings renders the model non renormalizable. We should stress that a microscopic foundation for derivative scalar coupling models has been derived from the relativistic $SU(6)$ model [6]. It has already been in-

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vestigated, in a mean- field approximation [7],[8], how this derivative coupling Zimanyi - Moszkowski (ZM) model differs from the usual Walecka model when the effective nucleon mass, the energy density, the incompressibility and other important features are calculated.

To try to solve the above mentioned drawbacks of the Walecka model, more sophisticated treatments have also been developed. Boguta and Bodmer [9] introduced two non - linear terms (cubic and quartic terms in the scalar field) to the original Walecka Lagrangian and obtained more reasonable results for the compression modulus and the effective mass at the saturation density. In their work they used the semi - classical Thomas - Fermi approximation and showed that it is equivalent to the quantal mean field approximation. Another work [10] refers to a relativistic Hartree - Fock approximation in a non - linear Walecka model for nuclear matter and finite nuclei which considers σ, ω, π and ρ mesons. In this work, the non linear character of the Hamiltonian makes the calculation of the exchange contributions to the energy rather complicated, which required an approximation for the calculation of the Fock - like terms, where the σ - meson mass was considered to be momentum independent. Even though, the inclusion of the exchange terms and the isovector mesons (π, ρ) turned out to be important for the description of the energy and the incompressibility in nuclear matter and of the spin - orbit interaction. Another successful treatment utilizes the so called Relativistic Dirac-Brueckner-Hartree-Fock method (RDBHF) [11]. In this case the coupling constants for the sigma and omega mesons become density dependent and they actually decrease with increasing energy. The computational effort also increases considerably since self - consistency for the self energies is obtained by considering different coupling constants at each density. Nevertheless this sort of calculation has proved very good when applied to finite nuclei [12] and the authors concluded that Fock terms are small but important and should not be neglected. The difficulties of the original Walecka model and the way they are corrected in the nonlinear treatment is also discussed in [13], where the incompressibility and the nucleon effective mass are obtained consistently with the required description of spin- orbit splittings and nuclear deformation.

It is important to stress that the ZM models also belong to the category of descriptions involving density-dependent coupling constants [8]. The description underlying the approach known as relativistic density-dependent Hartree-Fock [14] reproduces finite nuclei and nuclear matter saturation properties using coupling constants that are fitted, at each density value, to the RDBHF self- energy terms. The good agreement obtained for the ground state properties of spherical nuclei lends support to this sort of description involving density dependent coupling constants. Recently, a finite nuclei calculation has been performed with the ZM models and the energy levels and ground-state properties of the ^{16}O, ^{40}Ca, ^{48}Ca, ^{90}Zr and ^{208}Pb are in good agreement with the experimental results [15]. One of the main conclusions of this analysis is that a modified version of the model, refered in this paper as ZM3 model, produces better results than the original ZM model regarding the energy splitting of levels due to the spin-orbit interaction.

The ZM model has also been studied in a semi- classical approximation [16] and its results were compared with the non linear Walecka model suggested in [9]. In contrast to the non linear Walecka model, the non linearity in the ZM model does not introduce extra free parameters.

In the present work, we propose an alternative way of considering in a very simple manner direct and exchange terms in the infinite nuclear matter to obtain the properties of the ZM models mentioned above and compare them with the results which come out of the relativistic Hartree - Fock approximation for the Walecka model. We would like to point out that, to our knowledge, no complete quantum Hartree Fock calculation has been performed in a non linear relativistic model, although many works with quantum MFT (constant scalar and vector meson fields) or Thomas - Fermi approximation can be found yielding good results for nuclear matter bulk properties and/or finite nuclei (see some of the references already mentioned above).

In the recent literature three different possibilities have been considered for the coupling of the nucleon with the mesons in the Zimanyi - Moszkowski (ZM) model (for a review, check ref. [7]). In this letter we consider two of them. The first one is known as the original ZM model and the second one has been chosen because is the one which gives better results for the nuclear matter in the mean field approximation and we call it ZM3 (to be consistent with the definition in other published material). To start with, we write the Lagrangian density for the ZM and ZM3 models respectively as [5]:

$$\mathcal{L}_{zm} = -\bar{\psi}M\psi + (m^*_{zm})^{-1}\bar{\psi}(i\gamma_\mu\partial^\mu - g_v\gamma_\mu V^\mu)\psi + \frac{1}{2}(\partial_\mu\phi\partial^\mu\phi - m_s^2\phi^2) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu, \quad (1)$$

and

$$\mathcal{L}_{zm3} = -\bar{\psi}M\psi + (m^*_{zm})^{-1}\bar{\psi}i\gamma_\mu\partial^\mu\psi - g_v\bar{\psi}\gamma_\mu V^\mu\psi + \frac{1}{2}(\partial_\mu\phi\partial^\mu\phi - m_s^2\phi^2) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu, \quad (2)$$

where

$$m^*_{zm} = (1 + \frac{g_s\phi}{M})^{-1} \quad (3)$$

and ψ, ϕ and V^μ are field amplitudes for the nucleon, the scalar-isoscalar meson, the vector-isoscalar meson, $F^{\mu\nu} = \partial^\mu V^\nu - \partial^\nu V^\mu$ and M is the bare nucleon mass. Notice that these derivative coupled Lagrangians are Lorentz invariant, but they are not renormalizable. The physical meaning of this modified couplings is that the kinetic fermionic term describes the motion of a particle with an effective mass M^* instead of the bare mass M present in the conventional Walecka model.

In ref. [5], a rescaled Lagrangian is obtained from the above equations, with the rescaling of the fermion wave function as $\psi \to \sqrt{m^*_{zm}}\psi$ in eqs.(1) and (2) and the rescaling of the field V_μ as $V_\mu \to m^*_{zm}V_\mu$ in eq.(2) and it reads

$$\mathcal{L}_R = \bar{\psi}(i\gamma_\mu\partial^\mu - M + {m^*_{zm}}^\beta g_s\phi)\psi + {m^*_{zm}}^\alpha(-g_v\bar{\psi}\gamma_\mu V^\mu\psi - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu) + \frac{1}{2}(\partial_\mu\phi\partial^\mu\phi - m_s^2\phi^2), \quad (4)$$

where, for the Walecka model $\alpha = 0$, $\beta = 0$, for ZM $\alpha = 0$, $\beta = 1$ and for ZM3 $\alpha = 2$, $\beta = 1$.

In a very simple description of the ZM and ZM3 models they can be viewed as models with non-linear effective scalar coupling constants g^*_s and g^*_v in such a way that the new Lagrangians and related quantities can be obtained simply by substituting the old coupling constants in the Walecka model by the new ones [8],i.e.:

$$\mathcal{L}_{zm} = \mathcal{L}_{Walecka}(g_s \to g^*_s)$$

and

$$\mathcal{L}_{zm3} = \mathcal{L}_{Walecka}(g_s \to g^*_s, g_v \to g^*_v),$$

where $g^*_s = m^*_{zm}\cdot g_s$ and $g^*_v = m^*_{zm}\cdot g_v$. Instead of the above substitutions, which can be cumbersome to be carried out in the self energy expressions, we utilize a new prescription which introduces modifications in the meson masses only [18]:

$$\mathcal{L}_{zm} = \mathcal{L}_{Walecka}(m_s \to m^*_s) \quad (5)$$

and

$$\mathcal{L}_{zm3} = \mathcal{L}_{Walecka}(m_s \to m^*_s, m_v \to m^*_v), \quad (6)$$

with $m^*_s = m_s/m^*_w$ and $m^*_v = m_v/m^*_w$, where, in this case,

$$m^*_w = 1 - \frac{g_s\phi}{M}. \quad (7)$$

The above prescription is exact only in the mean field approximation. In this work we apply this equivalence directly in the expressions obtained from the Walecka model, creating an *effective model* out of the ZM3 and the Walecka models.

In the Walecka model the relativistic Hartree-Fock equations [3],[4] are obtained by using Dyson's equation to sum to all orders the self-consistent tadpole and exchange contributions to the baryon propagator

$$G(k) = G^0(k) + G^0(k)\Sigma(k)G(k), \quad (8)$$

where Σ is the proper self-energy. Because of translational and rotational invariances in the rest frame of infinite nuclear matter and the assumed invariance under parity and time reversal, the self-energy may be written as [3],[17]

$$\Sigma(k) = \Sigma^s(k) - \gamma_0\Sigma^0(k) + \vec{\gamma}\cdot\vec{k}\Sigma^v(k). \quad (9)$$

Coupled integral equations are then solved for the self-energies in the so-called Dirac-Hartree-Fock approximation [3]. In this approximation just the contributions from real nucleons in the Fermi sea are kept in the baryon propagators. The effects of virtual nucleons and anti-nucleons on the medium are neglected [2], [3]. The nucleon propagator reads $G(k) = G_F(k) + G_D(k)$, where

$$G_F(k) = \frac{1}{\not{k}^* - M^*(k) + i\epsilon} \tag{10}$$

is the part which has to be taken into account if vacuum contributions are to be considered. In the Walecka model exchange contributions from the Dirac sea can be calculated at the two - loop level and they give better results when form factors (or vertex corrections) are considered [19], [20]. In this case renormalization is, of course, necessary. For renormalizable models the optimized δ - expansion is also a possibility for calculating vacuum contributions. Direct terms are easily computed but the inclusion of exchange terms requires a lengthy and cumbersome calculation [21]. In the ZM model, which is not renormalizable, a 'cut-off' would be required with the introduction of a renormalization prescription and this could raise controversial discussions. Hence, to avoid this technical difficulties and bearing in mind that we are looking for a method simpler than the DBHF, we have decided to leave vacuum contributions out of this calculation. Notice that vacuum effects would just add a small contribution to our results and they may be incorporated in the adjustment of the coupling constants.

The baryon propagator in a Fermi sea with Fermi momentum k_F is written as (the nuclear density is $\rho_B = 2k_F^3/3\pi^2$)

$$G_D(k) = (\gamma_\mu k^{\mu *} + M^*(k)) \frac{\pi i}{E^*(k)} \delta(k^0 - E(k)) \theta(k_F - |\vec{k}|), \tag{11}$$

where

$$k^{\mu *} = k^\mu + \Sigma^\mu(k) = \left(k^0 + \Sigma^0(k),\ \vec{k}(1 + \Sigma^v(k))\right) , \tag{12}$$

$$E^*(k) = \sqrt{(\vec{k}^*)^2 + M^*(k)^2} \tag{13}$$

$$M^*(k) = M + \Sigma^s(k), \tag{14}$$

and $E(k)$ is the single-particle energy, which is the solution of the transcendental equation

$$E(k) = [E^*(k) - \Sigma^0(k)]_{k^0 = E(k)} \ . \tag{15}$$

Performing the q^0 and angular integrals in the expressions for the various components Σ^s, Σ^0, Σ^v of the self-energy, three coupled nonlinear integral equations are then obtained. For the Walecka model, the equations for the self - energies are given in ref. [2], (pg. 131) and we do not reproduce them here. In the ZM model the very same steps are performed in order to obtain these three components of the self energy, with the exception that we have included an approximation. This approximation amounts to considering m^*, defined in eq. (7), as a function only of the momentum, i.e.,

$$m_w^*(k) = \frac{M^*(k)}{M}. \tag{16}$$

In MFT, the above expression is exact, since $\Sigma_s = -g_s\phi$. Because of this approximation the Hartree terms are exactly calculated, but the Fock ones are somewhat approximated as in ref. [10]. The exact calculation would imply in considering m^* as a function of the field operator ϕ and this is complicated. However, we still have an improvement on the Hartree calculation with the introduction of the exchange terms, albeit it is done in an approximate way, which amounts in considering the most important part of the exchange graphs, as in the chain approximation. The nonlinear integral equations for the ZM3 model are the same expressions as obtained in the Walecka model with the following modifications:

1.) All m_s are substituted by m_s^* and all m_v are substituted by m_v^*;

2.) As a consequence of 1.), the functions $\Theta_i(k,q)$ and $\Phi_i(k,q)$ are modified because of their dependence on $A_i(k,q)$, which becomes

$$A_i(k,q) = \vec{k}^2 + \vec{q}^2 + m_i^{*\,2} - [E(q) - E(k)]^2 \ , \tag{17}$$

where $i = \sigma$ or v.

For the ZM model just the modifications related to the mass of the sigma meson are carried out. All self-energies are evaluated at the self-consistent single-particle energies, $q^0 = E(q)$ and the equations for Σ^s, Σ^0 and Σ^v are solved by a direct iteration procedure.

We choose to normalize the model parameters using the bulk binding energy and saturation density of nuclear matter as usual. As normally done in calculations with the Walecka model, we identify the vector meson with the ω whose mass is $m_v = 783$ MeV and set $m_s = 550$ MeV for the scalar meson mass. For the nucleon mass we take $M = 939$ MeV. The energy density for the ZM3 model once the substitution given in eq. (6) is performed yields:

$$\mathcal{E} = \frac{2}{\pi^2}\int_0^{k_F} k^2 E(k)dk - \frac{g_v^2}{2m_v^*(k_F)^2}\rho_B^2 + \frac{2g_s^2}{m_s^*(k_F)^2\pi^4}\left(\int_0^{k_F} k^2\frac{M^*(q)}{E^*(q)}dk\right)^2 + 2g_s^2 I_\sigma + 4g_v^2 I_v\,, \tag{18}$$

where the I_i, $i = \sigma, v$ are integrals of the following form

$$I_i = \frac{1}{(2\pi)^6}\int_0^{k_F}\frac{d^3k}{E^*(k)}\int_0^{k_F}\frac{d^3q}{E^*(q)}D_i^0(k-q)F_i(k,q)H_i(k,q) \tag{19}$$

with the functions F_i and H_i given by

$$F_i(k,q) = 1/2 - (E(k)-E(q))^2 D_i^0(k-q), \tag{20}$$

$$H_\sigma(k,q) = k_\mu^* q^{*\mu} + M^*(k)M^*(q),$$
$$H_v(k,q) = k_\mu^* q^{*\mu} - 2M^*(k)M^*(q);\,, \tag{21}$$

and the $D_i's$ are the meson propagators

$$D_i^0(k) = \frac{1}{k_\mu^2 - {m_i^*}^2 + i\epsilon}\,. \tag{22}$$

To saturate the binding energy per nucleon at -16.15 MeV at the Fermi momentum of $1.14 fm^{-1}$ we use $g_s^2 = 100$ and $g_v^2 = 60$ for the ZM model. In the ZM3 model, the energy saturates at -17.0 MeV at $1.40 fm^{-1}$ for $g_s^2 = 114$ and $g_v^2 = 119$. In the Walecka model, the energy saturates at -15.75 MeV at $1.42 fm^{-1}$ for $g_s^2 = 83.11$ and $g_v^2 = 108.05$. One of the reasons for the introduction of the derivative coupling used in the ZM and ZM3 models is the small value of the effective nucleon mass at saturation density obtained with the Walecka model ($M^*/M = 0.5$). We have observed that this problem is corrected in the new models. At the saturation densities found for each model, M^*/M is 0.8 in the ZM model and 0.69 in the ZM3 model. Σ^0 and Σ^v in terms of the momentum also decrease much less in the ZM and ZM3 models than in the Walecka model. We can produce results compatible with the ones suggested in ref. [13], i.e. an incompressibility between 180 and 360 MeV and M^*/M between 0.58 and 0.64 by choosing a different set of parameters and a slightly deeper energy density at the saturation point, but this does not improve our results for the scalar and vector potentials. (see below and also the table).

Table 1

Scalar, vector potentials and incompressibility at nuclear matter saturation density are shown for the Walecka and the ZM models. MF stands for mean-field and HF for Hartree-Fock approximation.

models	S (MeV)	V (MeV)	V + S (MeV)	V - S (MeV)	K (MeV)
Walecka - MF	-431.02	354.12	-76.87	785.18	550.82
Walecka - HF	-458.31	379.01	-79.30	837.32	585.00
ZM - MF	-140.64	82.50	-58.13	223.13	224.71
ZM - HF	-177.5	109.24	-68.26	286.74	298.47
ZM3 - MF	-267.00	203.71	-63.28	470.71	155.74
ZM3 - HF	-281.00	190.29	-90.71	471.71	174.38

Other important quantities to be analyzed are the scalar and vector potentials and the incompressibility. To derive expressions for the potentials, we start from the Dirac equation:

$$\left((1+\Sigma^v)\vec{\alpha}\cdot\vec{k}+\gamma_0(M+\Sigma^s)-\Sigma^0\right)\psi=E\psi, \tag{23}$$

and after some simple algebraic manipulations, it can be rewritten as

$$\left(\vec{\alpha}\cdot\vec{k}+\gamma_0(M+\frac{\Sigma^s-M\Sigma^v}{1+\Sigma^v})+\frac{-\Sigma^0+E\Sigma^v}{1+\Sigma^v}\right)\psi=E\psi, \tag{24}$$

from where we may define in a natural way

$$S=\frac{\Sigma^s-M\Sigma^v}{1+\Sigma^v} \tag{25}$$

as the scalar potential and

$$V=\frac{-\Sigma^0+E\Sigma^v}{1+\Sigma^v}. \tag{26}$$

as the vector potential, where E given by eq.(15) can be written as $E=\mathcal{E}/\rho$ at the saturation density. In table 1 these potentials are displayed for the Walecka and the ZM models in the mean field and Hartree - Fock calculations. The $V-S$ quantity is related to the spin-orbit splitting in finite nuclei and the $V+S$ quantity corresponds to the real part of the optical potential for the zero three-momentum. There is some evidence for the strong potentials of the Walecka model, and as we can see from table 1 the ZM3 potentials are stronger than the ZM ones explaining why the ZM3 model gives a better description of the nuclear spectra [15].

Concerning the incompressibility K, it has also been calculated for the three models, where

$$K=9\rho_0^2\frac{\partial^2}{\partial\rho^2}\frac{\mathcal{E}}{\rho}|_{\rho=\rho_0} \tag{27}$$

with $\rho_0(k)=\rho_B(k/k_F)^3$. The results are also shown in table 1. According to ref [22], the expected incompressibility value is $K=210\pm30$ MeV, which means that the results obtained for the ZM and ZM3 models are of great improvement in comparison with the results coming from the Walecka model.

To conclude, we would like to comment that we have calculated some relativistic features which are important for the understanding of nuclear matter within the context of the ZM and ZM3 models. For the present calculation we have taken into account the Hartree terms (related to the direct diagrams for the baryon propagator) and Fock - like terms (related to the exchange diagrams) given by eq.(16). We have calculated the effective mass, the incompressibility and the scalar and vector potentials and obtained good results out of the ZM3 model. The same treatment applied to these models can be extented to finite nuclei calculation in substitution to the more complicated relativistic Dirac - Brueckner - Hartree - Fock approach. This investigation is under way.

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