CBPF-NF-024/95

Remarks on Derivative and Nonlinear Scalar Couplings for Hadronic Models

by

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Abstract

The equivalence for a class of derivative and nonlinear coupling for hadronic models is established and connected with an effective Walecka model. We show that effective coupling constants can be implemented in this case. For a family of new models, which couple the mesonic scalar and vector fields, the results for infinite nuclear matter exhibit an improvement to the case where only scalar effective couplings are present.

Key-words: Nuclear Physics.

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1 Introduction

The linear $\sigma - \omega$ model (hereafter called Walecka model) [1] satisfactorily explains many properties of nuclear matter and finite nuclei. A shortcoming of this model is the prediction of a high value for the compression modulus $K = 550$ MeV. The introduction of nonlinear scalar self-coupling terms [2] has brought K to a reasonable value of 250 MeV in a theory with four free parameters. Recently, some attempts have been made in order to modify the Walecka model aiming to keep only two free parameters and expecting to have a softer equation of state [3]-[6]. These models are not renormalizable and have to be understood as effective models. They are constructed by modifying the usual covariant derivative term in such a way that, after an appropriate rescaling, the Lagrangian describes the motion of a baryon with an effective mass $M^* = m^* M$ instead of the bare mass M. This information goes to the meson-baryon coupling, modifying it to an effective scalar coupling constant while the vector coupling constant remaining the same. This kind of model (we refer as Model 1) can give a good result for the incompressibility K, a value for M^* compatible to the Skyrme force model (showing its modest relativistic content), and a poor spin-orbit splitting for finite nuclei calculations [5]. A generalization of this kind of model does not change much these features [6].

Qualitative different models can be obtained if, instead the modification of the covariant derivative term, one simply modifies in the Lagrangian the kinetic baryonic term. In the same way, as before, after an appropriate rescaling, the Lagrangian also describes a baryon of mass M^*. This information manifests modifying not only the scalar-baryon coupling but also the vector-baryon coupling. It generates a coupling mechanism between the vector and the scalar fields. This class of models (we refer as Model 2) is, from our point of view, more rich than those contained in Model 1. The idea of such a possible model was suggested in the Appendix of Ref.[3] and implemented for a particular case in Ref.[7]. Model 2 is not well known as Model 1 and up no now no generalization of it has been done.

In this work we shall be commited with an unified discussion of generalized Model 1 and Model 2 which are, in general, obtained from derivative coupling models. We show that these models can be constructed from the usual Walecka model by redefining in it the coupling constant to an effective one. Connected in this way, the effective coupling constants scale as a function of $m^* = M^*/M$. Zimmanyi and Moszkowski (ZM) [3] proposed a specific choice to modify the scalar coupling trough $m^*(\sigma) = m^*_{ZM}(\sigma) = (1 + g_\sigma \sigma/M)^{-1}$ that, if applied for Model 1, turns out that their model is completely equivalent to the Walecka model if one scales $(g^*_\sigma/g_\sigma) = m^*$. Any other choice for $m^*(\sigma)$ [4]-[6] contained in Model 1 changes this scaling. Model 2 can be interpreted as the Walecka model if we perform $g_\sigma \to g^*_\sigma$ and $g_\omega \to g^*_\omega$. Irrespective to the choice of $m^*(\sigma)$, we have $g^*_\omega/g_\omega = m^*$. In the particular case of $m^*(\sigma) = m_{ZM}(\sigma)$, effective vector and scalar constants scale identically as $g^*_\sigma/g_\sigma = g^*_\omega/g_\omega = m^*$.

Here the generalization of Model 1 and Model 2 means the introduction of a new free parameter α. We do it by imposing the following scaling in the effective Walecka model:

$$\text{Model 1}: \qquad \frac{g^*_\sigma}{g_\sigma} = {m^*}^\alpha \quad , \quad \frac{g^*_\omega}{g_\omega} = 1 \tag{1}$$

$$\text{Model 2}: \qquad \frac{g^*_\sigma}{g_\sigma} = {m^*}^\alpha \quad , \quad \frac{g^*_\omega}{g_\omega} = m^* \tag{2}$$

Notice that by doing so we do not need to explicit any especific function $m^*(\sigma)$, it is automaticly contained in the scaling procedure. The free constants g_σ and g_ω are chosen to fit the nuclear matter binding energy ($E_b = -15.75$ MeV) at the saturation density ($\rho_o = 0.15 fm^{-3}$). We have kept α as a continuum variable for both models. These models are then studied regarding the obtaining of M^*, K, the scalar potential (S), and the vector potential (V). The difference between the last two quantities is intrinsicly related to the spin-orbit splitting for finite nuclei. We show that Model 1, unlike Model 2, can not simultaneously achieves good performance calculating the above observables.

2 The Models

We start by introducing the nonlinear Lagrangian density for the Model 1,

$$\mathcal{L}_{NL} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \bar{\psi}\left\{ \not{D}(g_\omega) - m^*(\sigma)\, M \right\}\psi + \frac{1}{2}\, m^2_\omega\, \omega_\mu\, \omega^\mu + \frac{1}{2}\left(\partial_\mu \sigma\, \partial^\mu \sigma - m^2_\sigma\, \sigma^2 \right) , \tag{3}$$

where $\not{D}(g_\omega) \doteq \gamma^\mu D_\mu(g_\omega) = \gamma^\mu(i\partial_\mu - g_\omega\omega_\mu)$, is the usual covariant derivative and the degrees of freedom are baryon fields (ψ), scalar meson fields (σ), and vector meson fields (ω^μ). The real function $m^*(\sigma)$ is to be defined within each model under consideration, with the condition that for zero density goes to 1 and vanish for higher densities, because the effective mass must approaches zero asymptotically. Indeed, the Dirac equation obtained from the Lagrangian density (3) gives

$$m^*(\sigma) = \frac{M^*}{M} \ , \tag{4}$$

where M and M^* are the bare and effective baryonic mass respectively.

Now we proceed to show that performing a spinor field transformation, $\mathcal{L}_{NL}$ can be obtained from a Lagrangian density with derivative scalar coupling (DSC). The proposed Lagrangian density [3] which generates $\mathcal{L}_{NL}$ is given by

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \bar{\psi}\Big\{[m^*(\sigma)]^{-1}\not{D}(g_\omega) - M\Big\}\psi + \frac{1}{2}m_\omega^2\,\omega_\mu\,\omega^\mu + \frac{1}{2}\Big(\partial_\mu\sigma\,\partial^\mu\sigma - m_\sigma^2\,\sigma^2\Big) \ . \tag{5}$$

Introducing the rescaled barionic field $\psi \rightarrow [m^*(\sigma)]^{\frac{1}{2}}\psi$, we obtain from Eq.(5) the rescaled Lagrangian density

$$\mathcal{L}_R = \mathcal{L}_{NL} + \Im \ , \tag{6}$$

where $\Im$ is an imaginary contribution given by

$$\Im = \frac{i}{2}(\bar{\psi}\gamma_\mu\psi)\,\partial^\mu\ln(m^*(\sigma)) \ . \tag{7}$$

This term does not carry any physical meaning and its appearance can be avoided by starting from the correct Hermitian Lagrangian density. To this end we replace the baryonic kinetic term in Eq.(3) by $\frac{i}{2}\{\bar{\psi}\gamma_\mu\partial^\mu\psi - (\partial^\mu\bar{\psi})\gamma_\mu\psi\}$, such that the imaginary $\Im$ contribution cancels after the field scaling. Taking this into account, the field rescaling is equivalent to the replacement

$$\{[m^*(\sigma)]^{-1}\not{D}(g_\omega) - M\} \rightarrow \{\not{D}(g_\omega) - m^*(\sigma)\,M\} \ . \tag{8}$$

Therefore, the Lagrangian densities given by Eq.(3) and Eq.(5) are completely equivalent; they provide the same physical content irrespective of the fact that we deal with infinite nuclear matter or finite nuclei. In other words, the equations of motion obtained from Eqs.(3) and (5) describe the same hadronic dynamics. The form of Eq.(5) is not arbitrary since it has the physical meaning (from Eq.(4)) that the modified kinetic energy describes the motion of a baryon of mass M^* instead of the bare mass M. Eq.(3) just carries now this information to the scalar-baryon coupling fields.

The Walecka model can be obtained as a particular case of the model defined by $\mathcal{L}$ or $\mathcal{L}_{NL}$ and its Lagrangian density is recovered making the choice $m^*(\sigma) = (1 - g_\sigma\sigma/M)$. In fact, all the recent nonlinear models [3]-[6] can be obtained from DSC model whose rescaled Lagrangian density can be interpreted as follows

$$\mathcal{L}_{NL} \equiv \mathcal{L}_{Walecka}(g_\sigma \rightarrow g_\sigma^*) \ , \tag{9}$$

where g_σ^* (hereafter we will interpret $*$ as refering to effective coupling constant in the medium) is now a function of σ, related to $m^*(\sigma)$ by

$$m^*(\sigma) = 1 - g_\sigma^*\sigma/M \ , \tag{10}$$

which establishes by itself a class of models, since $m^*(\sigma)$ is general. In the usual Zimanyi-Moszkowski (ZM) model [3],

$$m_{ZM}^*(\sigma) = (1 + g_\sigma\sigma/M)^{-1} \ , \tag{11}$$

and $g_\sigma^* = g_\sigma\,m_{ZM}^*(\sigma)$. In the same way, for the other models [4]-[6] the identification of g_σ^* and $m^*(\sigma)$ can easily be done.

Let us now discuss the possible modified versions of the Model 1 we will refer as Model 2. We keep the generalized $m^*(\sigma)$ but following the suggestion given in the appendix of Ref.[3], we shall restrict the $m^*(\sigma)$ dependence in $\mathcal{L}$ to the fermionic kinetic term. To this end we introduce the modified covariant derivative $\not{D}_{m^*}(g_\omega) = \left([m^*(\sigma)]^{-1}\, i\,\not\partial - g_\omega \not\omega\right)$ and write the Lagrangian in the form

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} + \bar\psi\left\{\not{D}_{m^*}(g_\omega) - M\right\}\psi + \frac{1}{2}m_\omega^2\,\omega_\mu\,\omega^\mu + \frac{1}{2}\left(\partial_\mu\sigma\,\partial^\mu\sigma - m_\sigma^2\,\sigma^2\right) . \tag{12}$$

Since under the previous scaling $\psi \to [m^*(\sigma)]^{\frac{1}{2}}\,\psi$, we have

$$\{\not{D}_{m^*}(g_\omega) - M\} \to \{\not{D}(g_\omega^*) - m^*(\sigma)M\} \ , \tag{13}$$

then a new class of nonlinear models can be generated, and again the connection with the Walecka model is performed through

$$\mathcal{L}_{NL} \equiv \mathcal{L}_{Walecka}(g_\sigma \to g_\sigma^* \ ; \ g_\omega \to g_\omega^*), \tag{14}$$

in which g_σ^* is connected with $m^*(\sigma)$ by Eq.(10) and

$$\frac{g_\omega^*}{g_\omega} = m^*(\sigma) \ . \tag{15}$$

Note that Eqs.(9) and (14) simplify the understanding of different kinds of nonlinear couplings since they are in fact effective Walecka models. An interesting point here is that once the Lagrangian is given by Eq.(12), Eq.(15) imposes the scaling of the effective vector coupling constant in the medium.

We can now investigate whether this same scaling can be extended to the effective scalar coupling constant in the medium too. Initially we look for a function $m^*(\sigma)$ such that the following constraint is satisfied:

$$\frac{g_\sigma^*}{g_\sigma} = \frac{g_\omega^*}{g_\omega} = m^* \ . \tag{16}$$

It turns out that the only function that fullfils the above requirement is $m^*(\sigma) = m^*_{ZM}(\sigma)$. This modified ZM model (ZMM) is the first hadronic model which exhibits this property. This kind of model couples the σ field to the ω field and some results we present here show how it changes the usual ZM model.

The coupling constants for the usual ZM model (ZM) now interpreted as given by Eq.(9) are presented in Ref.[3]. The model given by Eq.(14) (ZMM) saturates the infinite nuclear matter at the density $\rho_0 = 0.15\,fm^{-3}$ with the binding energy $E_b = -15.75$ MeV for $C_\sigma^2 = g_\sigma^2 M^2/m_\sigma^2 = 443.3$ and $C_\omega^2 = g_\omega^2 M^2/m_\omega^2 = 305.5$ [7]. We present, for both models, the incompressibility K, the scalar potential S, the vector potential V (in MeV) and the baryonic effective mass m^*:

Model	m^*	K	S	V
ZM	0.85	225	-141	82
ZMM	0.72	156	-267	204

Compared to ZM, ZMM presents the peculiarity of giving simultaneously a smaller m^* and a smaller K. This is not what occurs when we compare ZM model with the Walecka model. This peculiarity may be explained now by the nonlinear scalar-vector coupling contained in ZMM model. In reference [5], ZM is implemented for a finite nuclei calculation showing a poor result regarding the spin-orbit splitting which, as expected, is strongly dependent on the quantity $V - S$. In ZMM model, this quantity is larger than the double of that obtained for ZM model, suggesting that ZMM model may improve upon ZM model in this particular direction. It is also interesting to remark that for low energy the slope of the real optical potential (given by $1 - m^*$) provides information regarding the expected value of m^*. Experimental values [4, 9], in the limits of infinite mass number and zero radius gives m^* around 0.6. It does not support strongly ZM model, but tends to favour the Walecka model and ZMM model instead [7]. In ZMM model, K is smaller than the "empirical" prediction $K = 210 \pm 30$ MeV. Regarding this point, unlike the nonlinear $\sigma - \omega$ model, which presents unphysical behavior for $K < 200$ MeV [2], ZMM model 2 does not present any anomaly for the equation of state with such small value of K.

Now we proceed to generalize the ZM and the ZMM models. Recall that both come from a particular case of Eq.(10) when $m^*(\sigma) = m^*_{ZM}(\sigma)$, this choice has the consequence that the scalar effective coupling constant scales as $(g_\sigma^*/g_\sigma) = m^*$. Different choices of $m^*(\sigma)$ can be done under the requirement of Eq.(10).

2.1 Model 1

From Eq.(9) we generate a family of models by choosing the scaling given by Eq.(1). The equation of state for such a model is given in a general expression together with Model 2. Model 1 have two particular cases: Walecka Model ($\alpha = 0$) and ZM model ($\alpha = 1$).

2.2 Model 2

The basis of this model is the Eq.(2 and Eq.(14) (the ZMM model case is achieved for $\alpha = 1$). Given α, $m^*(\sigma)$ is defined by Eq.(2) together with Eq.(10) as in the Model 1 case.

2.3 The Equations of State for the Models

When the meson fields in the Lagrangians are replaced by their mean values, we arrive at the mean field approximation (MFA). In Model 1 the scalar meson field equation implies σ to be a function of the scalar density (ρ_s) only. This is not which occurs in Model 2, once here σ becomes also a function of the baryonic density ρ_b. For rotationally and translationally invariant symmetric nuclear matter, the MFA equation for the scalar fields reads

$$\sigma = \frac{g_\sigma}{m_\sigma^2 M}\frac{m^{*\,\alpha+1}}{(1-\alpha)m^* + \alpha}\left[M\rho_s + \beta\left(\frac{g_\omega}{m_\omega}\right)^2 m^*\rho_b^2\right] , \tag{17}$$

where $\beta = 0$ and 1 for Model 1 and Model 2 respectively. This shows clearly how Model 2 and its generalization mixes the scalar and vector fields. The scalar and baryonic densities are related through

$$\frac{\rho_s}{\rho_b} = -\left(\frac{C_\omega^2}{C_\sigma^2}\right)\left(\frac{(1-\alpha)m^* + \alpha}{m^{*\,2\alpha+1-2\beta}}\right)\left(\frac{S}{V}\right) - \beta\left(\frac{V}{Mm^*}\right) , \tag{18}$$

where

$$S \;=\; -\,g_\sigma^*\,\sigma \;=\; -M(1-m^*)\,, \tag{19}$$

and

$$V \;=\; (C_\omega^2/M^2)\,m^{*\,2\beta}/\rho_b\,. \tag{20}$$

This ratio estimates the relativistic content of each model. Model 2 presents an additional term, favouring its relativistic contribution compared to Model 1.

The expressions for the energy density and pressure at a given temperature T can be found as usual by the MFA average of the energy-momentum tensor,

$$\mathcal{E} = \frac{C_\omega^2}{2M^2}{m^*}^{2\beta}\rho_b^2 + \frac{M^4}{2C_\sigma^2}\left(\frac{1-m^*}{m^{*\alpha}}\right)^2 + \frac{\gamma}{(2\pi)^3}\int d^3k\,E^*(k)(n_k + \bar{n}_k)\,, \tag{21}$$

and

$$p = \frac{C_\omega^2}{2M^2}{m^*}^{2\beta}\rho_b^2 - \frac{M^4}{2C_\sigma^2}\left(\frac{1-m^*}{m^{*\alpha}}\right)^2 + \frac{1}{3}\frac{\gamma}{(2\pi)^3}\int d^3k\,\frac{k^2}{E^*(k)}(n_k + \bar{n}_k)\,, \tag{22}$$

where

$$\rho_b = \frac{\gamma}{(2\pi)^3}\int d^3k\,(n_k - \bar{n}_k)\,. \tag{23}$$

Here γ is the degeneracy factor ($\gamma = 4$ for nuclear matter and $\gamma = 2$ for neutron matter), $\bar{n}_k$ and n_k stand for the Fermi-Dirac distribution for antibaryons and baryons with arguments $(E^* \pm \nu)/T$ respectively. $E^*(k)$ is given by

$$E^*(k) = (k^2 + M^{*2})^{\frac{1}{2}}\,, \tag{24}$$

while an effective chemical potential which preserves the number of baryons and antibaryons in the ensemble is defined by $\nu = \mu - V$, where μ is the thermodynamical chemical potential. According the Hugenholtz-van Hove theorem [8], the Fermi energy must be equal to the energy per baryon at the saturation density. Therefore, the following relation has to be satisfied,

$$\frac{\mathcal{E}}{\rho_o} = V + E^*(\rho_o)\,. \tag{25}$$

We finish this section by presenting an analytical general expression for the incompressibility valid for both models,

$$K = 9V + 3\frac{k_F^2}{(k_F^2+(M+S)^2)^{1/2}} + 9\left(\rho\frac{\partial S}{\partial \rho}\left(\frac{(M+S)}{(k_F^2+(M+S)^2)^{1/2}} + 2\beta\frac{V}{(M+S)}\right)\right)\,, \tag{26}$$

where all the quantities are calculated at the nuclear matter saturation density ρ_o and k_F is the fermi momentum.

3 Results and Discussions

We have implemented these generalized new models for $T = 0$, requiring $E_b = -15.75$ MeV at $\rho_o = 0.15\,\mathrm{fm}^{-3}$ for some values of α. We start this section by discussing the Model 1, asking whether by varying α simultaneous reasonable results for K, M^* and $V - S$ may be obtained. The answer for this question has to be found in the Fig.1, where these quantities are plotted together as a function of α. The answer is clearly no, since there is no region where we could pick up simultaneous reasonable results for the discussed quantities, supporting the conclusions extracted from a generalized " ansatz " implemented by Greiner and Reinhard [6]. We mean by reasonable results m^* around 0.6 as pointed out before, which, as we will see later, fixes the values of $V - S$ around 680 MeV. The experimental value of K , determined from the energy of the breathing mode of doubly magic nuclei [10], is 210 ± 30 MeV. G.E.Brown [11], using Fermi-liquid Landau theory, gives strong reasons for a lower value of K. Once this question sems to be under debate, we use the experimental value of Ref.[10] as an approximated upper limit.

The results for the Model 2 are presented in Fig. 2. The values of K are weakly sensitive in the region of $\alpha > 1$, and they reach the minimum value very close to $\alpha = 1$, the ZMM case. Particulary interesting are the results for the region $\alpha < 1$ in Fig. 2, where m^* and the quantity $V - S$ can improve the results of ZMM model given in the last section. As an example, for $\alpha = 0.9$ Model 2 furnishes $M^* = 558.3$ Mev, $V - S = 687.6$ Mev and the incompressibility $K = 166.2$ Mev. It is very interesting that for $\alpha = 0.88$ the values for $M^* = 507.0$ Mev, $V - S = 785.1$ Mev are approximately the same as the obtained in the usual linear Walecka Model but with a reasonable value for the incompressibility, $K = 181.3$ Mev.

Fig. 2 also shows, that for $\alpha < 0.79$ no nuclear matter saturation is achieved, including the limiting case of α=0 , which would correspond to the Walecka choice for $m^*(\sigma)$ in Eq.(10). Moreover, Fig.1 and Fig.2 show that M^* increase with α and, as a consequence of Eq.(25), the vector potential V decreases. We have allowed high values of α to obtain the curious situation where V vanishes, and the Model 1 and Model 2 degenerate. Nuclear matter saturation, only with scalar field, is then achieved. This occurs for $\alpha \approx 12.8$ which is the maximum value allowed for this parameter (beyond this, V becomes negative) with $C_\sigma^2 = 315.36$. The results for this especial case are: $M^* = 885.8$ Mev, $V - S = -S = 52.4$ Mev and $K = 65.8$ Mev.

After having gained some insights from each kind of model, we discuss here in which M^* could give model independent properties for the nuclear matter. M^* itself is a manifestation of the relativistic content of any particular model. On the other hands, if two different models present the same value of M^*, they have the same relativistic ratio given by Eq.(18). This quantity is presented in Fig. 3 for Model 1 and Model 2, for different values of α. This shows how Model 2 (unlike Model 1) can acquires extrem relativistic features. In particular, we see from Figs. 1 and 2 that Model 1 and Model 2 can give the same values for M^* (for different values of α) and, consequently, the same relativistic ratio too. If the situation is such that different models give the same M^*, the question is whether it can itself determine other observables. For instance, the values of S and V are fixed from Eq.(19) and Eq.(25) respectively if M^* is given. Therefore, the quantity $V - S$ and M^* are directly correlated and carry the same physical information for each particular model. Regarding the correlation between K and M^*, the situation differs and assumes a character of model dependence.

In order to have a better radiografy of K, we calculate separately the different terms which compose it, according Eq.(26). Notice that comparing to Model 1 ($\beta = 0$), Model 2 ($\beta = 1$) presents an additional fourth term (hereafter K_4). The first three terms of Eq.(25) we will refer as K_1, K_2 and K_3 respectively. We plot in Figs. 4 and 5 these quantities to see the isolated contribution of each one. K_1 and K_2 are completely determined by M^* since they directly depend on S and V. However, K_3 and K_4 depend not only on M^* but also on the slope of $M^*(\rho)$ which is negative and carries the information from the scalar field. For Model 1, we see from Fig.1 that when M^* increases the incompressibility K decreases. So, this explains why in Walecka model, where M^* is small, K is so high. In the usual ZM model occurs the opposite. In Model 2 we do not have this behavior, as we can see from Fig. 5, where we show how K_4 composes with K_3 to keep K almost constant when M^* increases.

4 Conclusions

In summary, we have shown the equivalence between the derivative scalar coupling and some nonlinear hadronic models which now are understood as an effective Walecka models. In this case the effective coupling constants depend on the density and are completely determined by the effective nucleon mass m^*. Two class of models were studied. The first has only an effective scalar coupling g^*_σ related with m^* by $m^* = 1 - g^*_\sigma \sigma / M$. The second also includes a new effective vector coupling g^*_ω which always is given by $g^*_\omega = m^* g_\omega$. We have also shown that only for a particular choice of $m^* = m^*_{ZM}(\sigma)$ both effective meson coupling constants scale as m^*. We have generalized the scaling, to conclude that in a theory with three free parameters, Model 1 can not succeed to furnish good results for the finite nuclei splitting simultaneously with the values of K and m^*. From Fig. 2, however, it is possible to pick up values of α giving values of m^*, K and $V - S$ around those obtained from the nonlinear $\sigma - \omega$ model, that we assume as good. The aim of this work is, however, not the proposal of any particular alternative model, but a systematic study to see in which direction the Walecka model could be changed with some possibilities of success. In this way becomes clear that, to change only the scalar coupling (Model 1) one definitely can improve K but (supporting the Ref.[6]) with no hope to have good answers to the spin-orbit splitting for finite nuclei. This is a direct consequence from the high value of M^* that Model 1 furnishes when K goes in the right direction by variation of α. Model 2, instead, suggests that a modification of the Walecka model which includes a mixed coupling of the scalar and the vector fields may be the way to expect improvement in the calculation of the discussed observables.

Acknowledgements

The authors express their thanks to Conselho Nacional de Desenvolvimento Científico e Tecnológico (CNPq), Brazil, for partial financial support. One of the authors (M. C.) also thanks to Centro Latinoamericano de Física (CLAF) and CNPq for financial support and the kindly hospitality of CBPF.

Figure Caption

Figure 1: The nucleon effective mass (M^*), the difference between the vector and scalar potentials ($V - S$) and the incompressibility K as function of α for Model 1.

Figure 2: The nucleon effective mass (M^*), the difference between the vector and scalar potentials ($V - S$) and the incompressibility K as function of α for Model 2.

Figure 3: The relativistic ratio between scalar and baryonic densities (ρ_s/ρ_b) as a function of α for both models.

Figure 4: The components K_1, K_2 and K_3 of the incompressibility K as a fuction of α for Model 1.

Figure 5: The components K_1, K_2, K_3 and K_4 of the incompressibility K as a fuction of α for Model 2.

Figure 6: The effective nucleon mass as a function of the scalar sigma field ($u = g_\sigma \sigma / M$) for different values of α.

Fig. 1

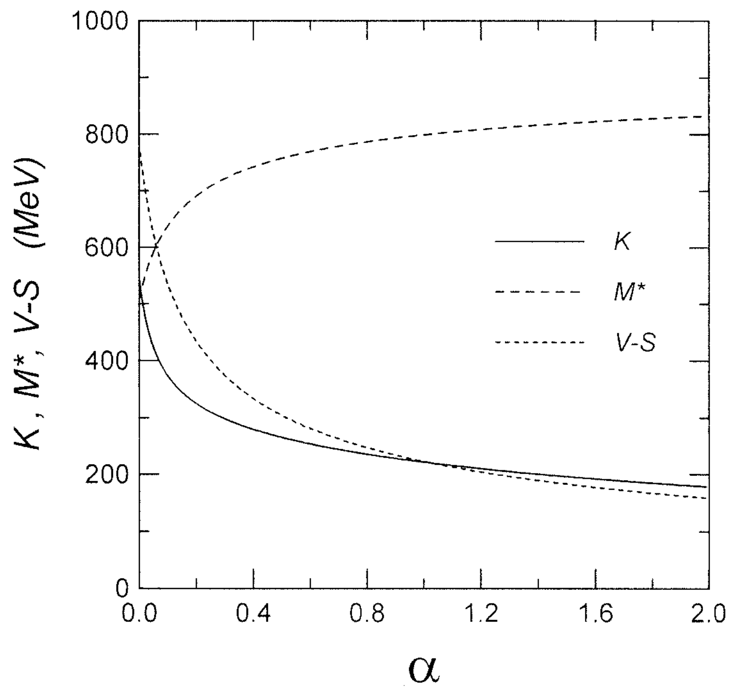


Fig. 2

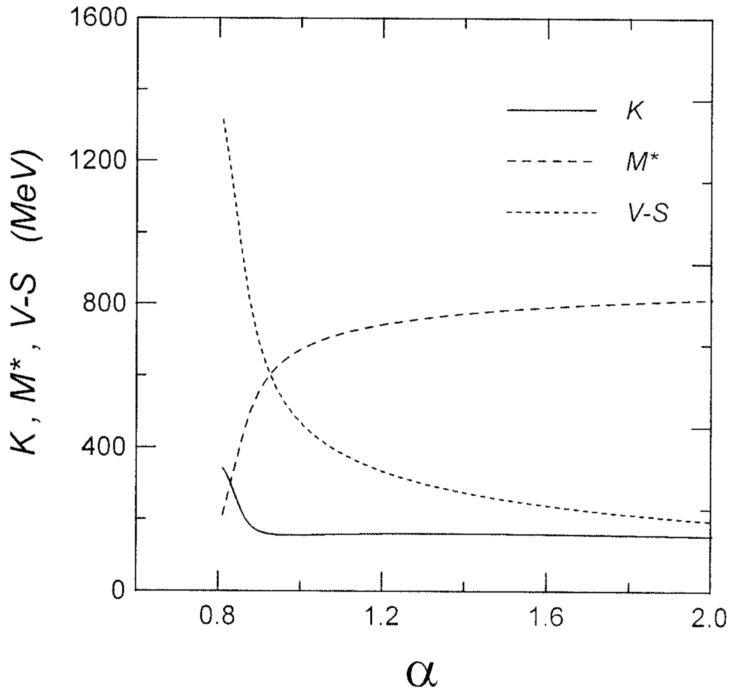


Fig. 3

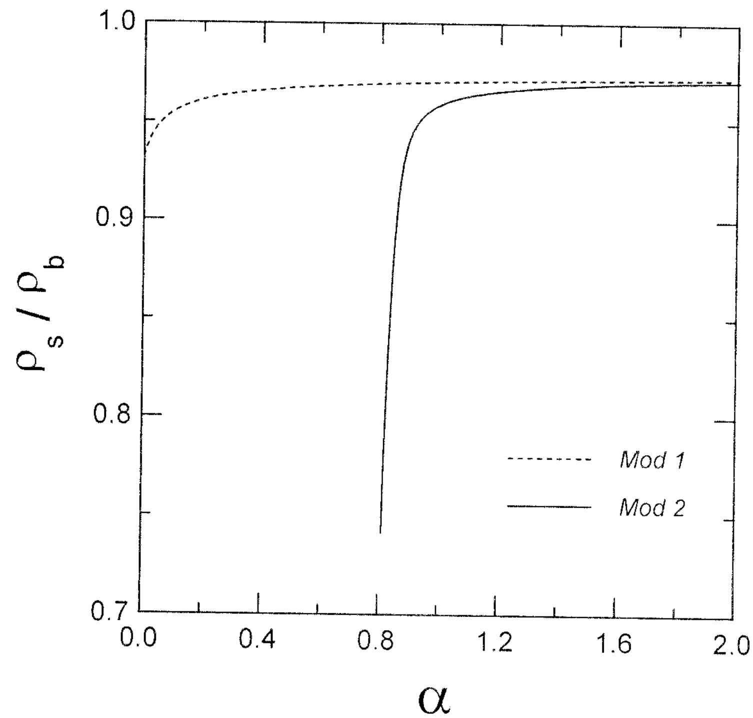
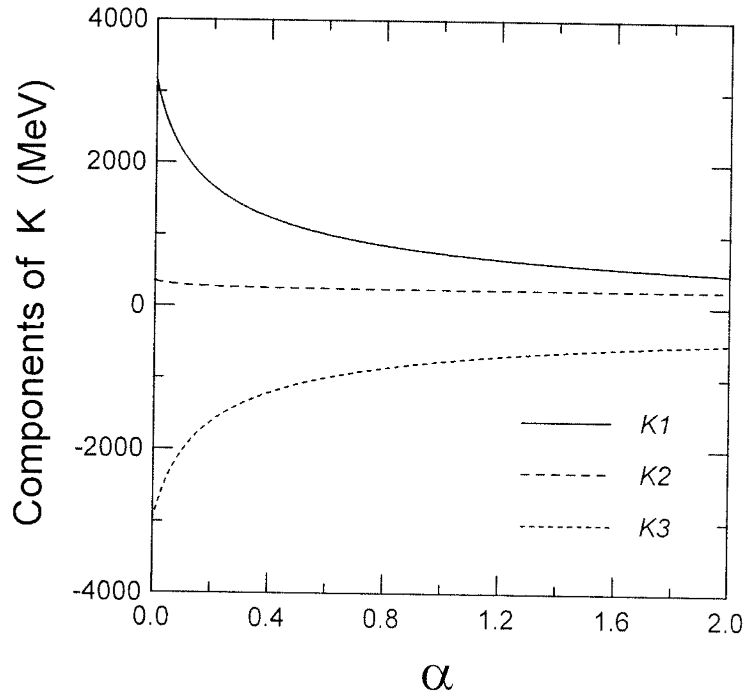


Fig. 4


4000
2000
0
-2000
-4000
Components of K (MeV)
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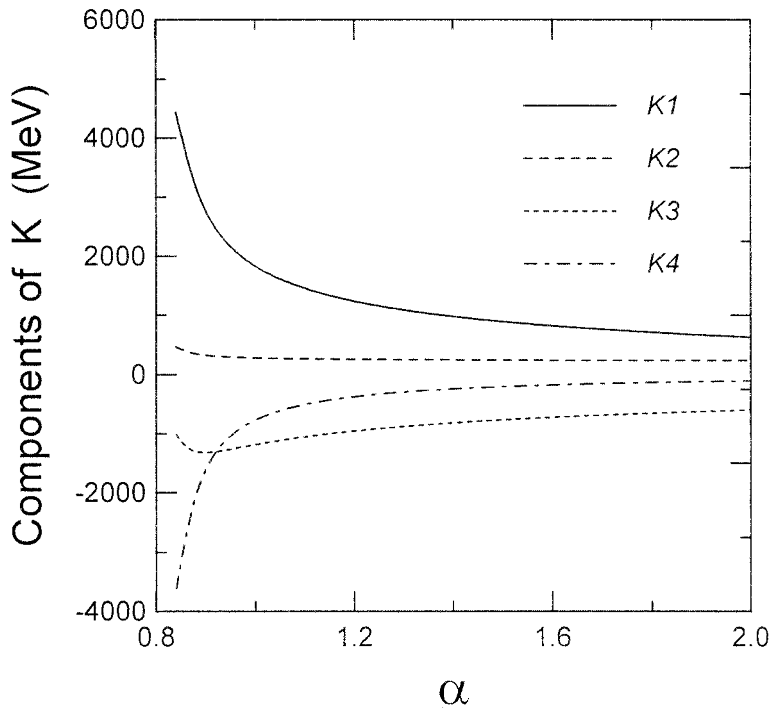
0.8
1.2
1.6
2.0
α
K1
K2
K3

Fig. 5


6000
4000
2000
0
-2000
-4000
Components of K (MeV)
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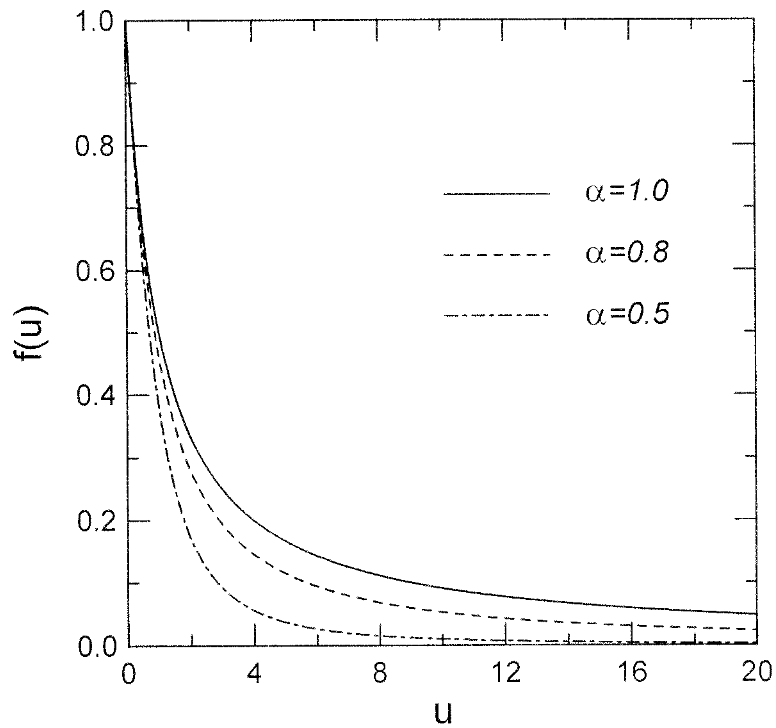
K2
K3
K4
0.8
1.2
1.6
2.0
α

Fig.6



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