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RHO-OMEGA MIXING AND NEUTRON-PROTON SELF-ENERGIES IN THE WALECKA MODEL

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Abstract

We use the Walecka model to investigate the influence of the charge symmetry breaking $\rho^0 - \omega$ mixing interaction on the neutron-proton self-energy difference in nuclear matter. Using $2m_\rho < \rho^0|H|\omega >= -4500\ MeV^2$, and employing the Dirac-Hartre-Fock approximation, we find that the neutron-proton self-energy difference is a decreasing function of the nuclear density, and that it has a value of the order of 700 KeV at the normal density. The results indicate that the Nolen-Schiffer anomaly might be explained by means of relativistic nuclear models in a similar way as it is explained by means of non-relativistic models.

The experimental values of the masses of mirror nuclei (or analog states) heavier than $A = 3$ are consistently larger than the theoretically calculated ones. The discrepancy, known as the Nolen-Schiffer anomaly [1] (NSA), increases with the mass number A and for $A \sim 209$ it can reach the value of 900 KeV. Several nuclear structure effects such as correlations, core polarization and isospin mixing have been invoked to solve the problem without definite success [2]. Because of this failure, explanations were searched outside conventional nuclear structure, and naturally quark models were invoked to solve the problem (for a recent review see Ref. [3]). However, it is largely known that charge symmetry breaking (CSB) forces of class III (pp-nn) and class IV (pn) [4] can affect the binding energy differences of mirror nuclei [5, 6]. In this context, recently Blunden and Iqbal [7] (BI) performed a systematic and detailed calculation of binding energy differences of mirror nuclei in the range of $A = 11$ to $A = 41$. BI used in their calculation nucleon-nucleon CSB potentials derived from $\rho^0 - \omega$ and $\pi^0 - \eta$ mixings and included the effects of the neutron-proton mass difference in OPEP and TPEP. Within the context of a Schrödinger equation calculation, these authors concluded that CSB effects can explain about 75% of the NSA. In addition, their calculation showed that the $\rho^0 - \omega$ mixing class III potential is by far the most important contributor to the CSB effect on the calculated binding energy differences.

In a more recent publication, Miller [8] recalculated the binding energy differences in the BI approach using a larger value for the $\rho^0 - \omega$ mixing matrix element than the one used by BI. The larger value of the matrix element is the result of a recent precise $\rho^0 - \omega$ mixing experiment [9]. Using this new value for the mixing, Miller showed that the net result of the calculations is that no significant anomaly remains.

The $\rho^0 - \omega$ mixing effective potential used in a Schrödinger equation framework is the Fourier transform of the non-relativistic approximation of the relevant one boson exchange graph of the NN scattering matrix [10, 4]. Owing to the importance of the possible resolution of the long-standing NSA by means of the CSB $\rho^0 - \omega$ mixing interaction, it would be worthwhile to investigate this matter in a relativistic framework, where the non-relativistic approximation to the interaction is avoided. In particular, relativistic nuclear models based on the original Walecka model [11] have been successful in describing several nuclear properties and it would be interesting to investigate the CSB effects of the $\rho^0 - \omega$ mixing mechanism in the context of these models.

In this letter we evaluate the influence of the $\rho^0-\omega$ mixing interaction on the density dependence of the neutron-proton self-energy difference in nuclear matter using the Walecka model with pions and rhos. This calculation should be considered as a first step towards a more complete calculation of the binding energy differences of mirror nuclei with CSB interactions in relativistic models. To our knowledge [3], the only relativistic calculation of the NSA is the one of Ref. [13]. Although the relativistic approach of Ref. [13] helps in reducing the anomaly as compared with simple non-relativistic calculations, they give comparable results to that obtained with non-relativistic density-dependent Hartree-Fock calculations. The mechanism responsible for the resolution of the NSA in the non-relativistic approach is the fact that the $\rho^0-\omega$ mixing interaction gives more binding to the neutrons than to the protons in medium and it is an increasing function of the nuclear density. We see in the following that in the absence of the $\rho^0-\omega$ mixing interaction in the Walecka model, the neutron-proton self-energy difference is almost constant as a function of the nuclear density. On the other hand, including the $\rho^0-\omega$ mixing and using the same value for the mixing matrix element as the one used by Miller [8], we obtain a neutron-proton self-energy difference which *decreases* with density, and at the normal nuclear density it is of the order of 700 KeV. This value is consistent with the NSA in the region of $A\sim 200$. Our results therefore indicate that a relativistic calculation on the lines of Ref. [13] including the $\rho^0-\omega$ mixing interaction might solve the NSA in a similar way as the non-relativistic calculation of Blunden and Iqbal [7] solves it.

To start with, we have used the following Lagrangian density[11],[12]:

$$\begin{aligned}\mathcal{L} \;=\;& \bar{\psi}(i\gamma_\mu\partial^\mu - M + g_s\phi - g_v\gamma_\mu V^\mu - \frac{g_\pi}{2M}\gamma^\mu\gamma_5\vec{\tau}\cdot\partial_\mu\pi - \frac{g_\rho}{2}\gamma_\mu\vec{\tau}\cdot\vec{\rho}^\mu)\psi \\ &+\frac{1}{2}(\partial_\mu\phi\partial^\mu\phi - m_s^2\phi^2) - \frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}(\partial_\mu\pi\partial^\mu\pi - m_\pi^2\vec{\pi}^2) \\ &-\frac{1}{4}\vec{L}_{\mu\nu}\vec{L}^{\mu\nu} + \frac{1}{2}m_v^2V_\mu V^\mu + \frac{1}{2}m_\rho^2\vec{\rho}_\mu\vec{\rho}^\mu + \lambda\rho_0^\mu V_\mu\,. \end{aligned} \tag{1}$$

Here, ψ, ϕ, π, V^μ and ρ^μ stand respectively for the nucleon, the scalar-isoscalar meson, the pseudoscalar-isovector meson, the vector-isoscalar meson and the vector-isovector meson, $F^{\mu\nu}=\partial^\mu V^\nu - \partial^\nu V^\mu$, $\vec{L}^{\mu\nu}=\partial^\mu\vec{\rho}^\nu-\partial^\nu\vec{\rho}^\mu$ and $\lambda=-2m_\rho<\omega|H|\rho^0>$ is the $\rho^0-\omega$ mixing parameter.

The relativistic Hartree-Fock equations are obtained by using Dyson's

equation to sum to all orders the self-consistent tadpole and exchange contributions to the baryon propagator

$$G(k) = G^0(k) + G^0(k)\Sigma(k)G(k), \tag{2}$$

where Σ is the proper self-energy. The self-energy is composed of a momentum independent tadpole term Σ^T and an exchange term $\Sigma^X(k)$:

$$\Sigma(k) = \Sigma^T + \Sigma^X(k). \tag{3}$$

The exchange term, when the neutron and the proton have different masses, gives different contributions to the neutron and proton self-energies, whereas the tadpole contributes equally to the proton and neutron self-energies.

In order not to repeat standard formulae which can be found e.g. in Refs. [14],[15],[16], we show here just some equations which are important for the understanding of the present work. Because of the translational and rotational invariances in the rest frame of the infinite nuclear matter and the assumed invariance under parity and time reversal, the self-energy may be written as [14],[12]

$$\Sigma(k) = \Sigma^s(k) - \gamma_0\Sigma^0(k) + \vec{\gamma}\cdot\vec{k}\Sigma^v(k)\,. \tag{4}$$

The $\rho^0-\omega$ mixing contribution for the proton, shown in Figure 1, is given by [17]

$$\Sigma_p^{\rho^0-\omega}(k) = ig_\rho g_v\int\frac{d^4q}{(2\pi)^4}\frac{\lambda\gamma^\mu G_p(q)\gamma^\mu}{[(k-q)^2-m_\rho^2+i\epsilon][(k-q)^2-m_v^2+i\epsilon]} \tag{5}$$

and a similar equation for the neutron can be read off from the above one by changing $p\rightarrow n$ and its sign.

We solve the coupled integral equations for the self-energies in the so-called Dirac-Hartree-Fock approximation [14]. This approximation amounts to keep in the baryon propagators the contributions from real nucleons in the Fermi sea only. The effects of the medium on virtual nucleons and anti-nucleons are neglected. This yields the familiar Hartree-Fock approximation of non-relativistic many body theory when the assumptions of non-relativistic kinematics and static meson exchange are made. The nucleon propagator in a Fermi sea with Fermi momentum k_F is then written as (the nuclear density is $\rho_0 = 2k_F^3/3\pi^2$)

$$G_b(k) = (\gamma_\mu k_b^{\mu *} + M_b^*(k))\frac{\pi i}{E_b^*(k)}\delta(k_b^0 - E_b(k))\theta(k_F - |\vec{k}|) \,, \tag{6}$$

where the subscript b stands either for proton or for neutron and

$$k_b^{\mu *} = k^\mu + \Sigma_b^\mu(k) = \left(k^0 + \Sigma_b^0(k),\ \vec{k}(1+\Sigma_b^v(k))\right) \,, \tag{7}$$

$$E_b^*(k) = \sqrt{(\vec{k}_b^*)^2 + M_b^*(k)^2}, \quad M_b^*(k) = M_b + \Sigma_b^s(k) \,, \tag{8}$$

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

$$E_b(k) = [E_b^*(k) - \Sigma_b^0(k)]_{k_b^0 = E_b(k)} \,. \tag{9}$$

Performing the q^0 and angular integrals in the expressions for the various components Σ^s, Σ^0, $\cdots$ of the self-energy, we obtain the following coupled nonlinear integral equations

$$\begin{aligned}
\Sigma_p^s(k) \;=\; & \Sigma_s^T + \frac{1}{4\pi^2 k}\int_0^{k_F} dqq \left\{ \frac{g_s^2}{4}\frac{M_p^*(q)}{E_p^*(q)}\Theta_{\sigma p}(k,q) - g_v^2 \frac{M_p^*(q)}{E_p^*(q)}\Theta_{vp}(k,q) \right. \\
& -\frac{g_\pi^2}{4M^2}\left[\frac{M_p^*(q)}{E_p^*(q)}\left(\frac{m_\pi^2}{4}\Theta_{\pi p}(k,q) - kq\right) + 2(p\to n)\right] \\
& -\frac{g_\rho^2}{4}\left[\frac{M_p^*(q)}{E_p^*(q)}\Theta_{\rho p}(k,q) + 2(p\to n)\right] \\
& \left. + g_v g_\rho \frac{\lambda}{m_v^2 - m_\rho^2}\frac{M_p^*(q)}{E_p^*(q)}\left[\Theta_{\rho p}(k,q) - \Theta_{vp}(k,q)\right] \right\}
\end{aligned} \tag{10}$$

$$\begin{aligned}
\Sigma_p^0(k) \;=\; & -\frac{2}{3\pi^2}\frac{g_v^2 m_\rho^2}{(m_v^2 m_\rho^2 - \lambda^2)}k_F^3 - \frac{1}{8\pi^2 k}\int_0^{k_F} dqq \left\{ \frac{g_s^2}{2}\Theta_{\sigma p}(k,q) + g_v^2\Theta_{vp}(k,q) \right. \\
& +\frac{g_\pi^2}{4M^2}\left[\left((E_p(k)-E_p(q))^2 - \frac{m_\pi^2}{2} + \frac{qq_p^*(E_p(k)-E_p(q))}{E_p^*(q)}\right)\Theta_{\pi p}(k,q) \right. \\
& \left. - 2kq_p^*\frac{(E_p(k)-E_p(q))}{E_p^*(q)}\Phi_{\pi p}(k,q) + 2(p\to n)\right] \\
& +\frac{g_\rho^2}{4}\left[\Theta_{\rho p}(k,q) + 2(p\to n)\right] \\
& \left. - g_v g_\rho \frac{\lambda}{m_v^2 - m_\rho^2}\left[\Theta_{\rho p}(k,q) - \Theta_{vp}(k,q)\right] \right\}
\end{aligned} \tag{11}$$

$$\begin{aligned}
\Sigma_p^v(k) \;=\; & -\frac{1}{4\pi^2 k^2}\int_0^{k_F} dqq \left\{ \frac{q_p^*}{E_p^*(q)}\left(\frac{g_s^2}{2}\Phi_{\sigma p}(k,q) + g_v^2\Phi_{vp}(k,q)\right) \right. \\
& +\frac{g_\pi^2}{8M^2}\left[k\left(E_p(k) - E_p(q) + \frac{qq_p^*}{E_p^*(q)}\right)\Theta_{\pi p}(k,q) \right. \\
& - \left(\frac{q_p^*}{E_p^*(q)}[k^2 + q^2 + (E_p(k)-E_p(q))^2] \right. \\
& \left.\left. +2q(E_p(k)-E_p(q))\right)\Phi_{\pi p}(k,q) + 2(p\to n)\right] \\
& +\frac{g_\rho^2}{4}\left[\frac{q_p^*}{E_p^*(q)}\Phi_{\rho p}(k,q) + 2(p\to n)\right] \\
& \left. -g_v g_\rho \frac{\lambda}{m_v^2 - m_\rho^2}\frac{q_p^*}{E_p^*(q)}\left[\Phi_{\rho p}(k,q) - \Phi_{vp}(k,q)\right] \right\}
\end{aligned} \tag{12}$$

In addition, there are the three equations for the neutron self-energy, obtained from the above equations by just exchanging the indices for protons with the indices for neutrons and changing the sign of the mixing term (the one which contains λ). The equation for Σ_s^T is

$$\Sigma_s^T = -\frac{g_s^2}{\pi^2 m_s^2}\int_0^{k_F} dqq^2 \left[\frac{M_p^*(q)}{E_p^*(q)} + \frac{M_n^*(q)}{E_n^*(q)}\right] \tag{13}$$

In the above equations $q = |\,\vec{q}\,|$, $k = |\,\vec{k}\,|$,

$$\Theta_{ib}(k,q) = ln\left|\frac{A_{ib}(k,q) + 2kq}{A_{ib}(k,q) - 2kq}\right| \quad \Phi_{ib}(k,q) = \frac{A_{ib}(k,q)\Theta_{ib}(k,q)}{4kq} - 1 \,, \tag{14}$$

and

$$A_{ib}(k,q) = \vec{k}^2 + \vec{q}^2 + m_i^2 - [E_b(q) - E_b(k)]^2 \,. \tag{15}$$

All self-energies are evaluated at the self-consistent single-particle energies, $q^0 = E(q)$.

Eqs.(10-12) are solved by a direct iteration procedure with mean-field self-energies as starting values. When the output values coincide within a difference of less than 10^{-8} with the input values at all points, we consider that the self-consistency is achieved.

For completeness, we also investigate the density dependence on the neutron-proton effective mass difference. This might be useful for comparison with other approaches which attempt to explain the NSA by means of medium modifications of the neutron-proton mass difference[3]. Horowitz and Serot [14] define the relativistic effective mass M^{eff} in analogy to the non-relativistic definition by the relation ($b = p, n$)

$$M_b^{eff}(q) = q\left[\left(\sqrt{\left(\frac{\partial E_b(k)}{\partial k}\right)^{-2}} - 1\right)_{k=q}\right] \tag{16}$$

which coincides with the definition of M^{eff} in the Hartree approximation. With this definition, we have for the in-medium neutron-proton mass difference (for $k = k_F$)

$$\Delta M_{eff}(k_F) = M_n^{eff}(k_F) - M_p^{eff}(k_F). \tag{17}$$

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. The pion and ρ meson masses, g_π and g_ρ are considered fixed at their experimental values, i.e., $m_\pi = 138$ MeV, $m_\rho = 770$ MeV, $g_\pi^2/4\pi = 14.4$ and $g_\rho^2/4\pi = 2.4$ respectively. The value of λ quoted in ref. [8] is $\lambda = 4500\ MeV^2$. To saturate the binding energy per nucleon at -15.75 MeV at the Fermi momentum of $1.38 fm^{-1}$ we use $g_s^2/4\pi = 7.2$ and $g_v^2/4\pi = 10.6$. The energy density is given by

$$\begin{aligned}
\mathcal{E} &= \frac{1}{\pi^2}\int_0^{k_F} k^2[E_p(k) + E_n(k)]dk - \frac{g_v^2 m_\rho^2}{2(m_v^2 m_\rho^2 - \lambda^2)}\rho_B^2 \\
&+ \frac{g_s^2}{m_s^2\pi^2}\int_0^{k_F} dq q^2 \frac{1}{2}\left[\frac{M_p^*(q)}{E_p^*(q)} + \frac{M_n^*(q)}{E_n^*(q)}\right] + g_s^2\left[I_\sigma(p,p) + I_\sigma(n,n)\right] \\
&+ 2g_v^2\left[I_v(p,p) + I_v(n,n)\right] + \frac{g_\pi^2}{4M^2}\left[I_\pi(p,p) + I_\pi(n,n) + 4I_\pi(p,n)\right] \\
&+ \frac{g_\rho^2}{2}\left[I_\rho(p,p) + I_\rho(n,n) + 4I_\rho(p,n)\right] \\
&- \frac{\lambda g_\rho g_v}{(m_\rho^2 - m_v^2)}\left[I'_\rho(p,p) - I'_v(p,p) - I'_\rho(n,n) + I'_v(n,n)\right].
\end{aligned} \tag{18}$$

The I_m and I'_m are integrals of the following form

$$\begin{aligned}
I_m(b,c) &= \frac{1}{(2\pi)^6}\int_0^{k_F}\frac{d^3k}{E_b^*(k)}\int_0^{k_F}\frac{d^3q}{E_c^*(q)} D_s^0(k-q)F_m(k,q)H_m(k,q) \\
I'_m(b,c) &= \frac{1}{(2\pi)^6}\int_0^{k_F}\frac{d^3k}{E_b^*(k)}\int_0^{k_F}\frac{d^3q}{E_c^*(q)} D_s^0(k-q)F'_m(k,q)H_m(k,q)
\end{aligned} \tag{19}$$

where the functions F_m and F'_m are given by

$$\begin{aligned}
F_m(k,q) &= \left[1/2 - (E_b(k) - E_c(q))^2 D_m^0(k-q)\right] \\
F'_m(k,q) &= \left[1 - (E_b(k) - E_c(q))^2 D_m^0(k-q)\right]
\end{aligned} \tag{20}$$

and the H_m's are given by

$$\begin{aligned}
H_\sigma(k,q) &= \left[k_{b\mu}^* q_c^{*\mu} + M_b^*(k)M_c^*(q)\right], \\
H_\pi(k,q) &= \left[2(k-q)_b^\mu q_{c\mu}^*(k-q)_c^\nu k_{b\nu}^* - (k-q)_\nu^2(M_b^*(k)M_c^*(q) + q_{b\mu}^* k_c^{*\mu})\right] \\
H_v(k,q) &= \left[k_{b\mu}^* q_c^{*\mu} - 2M_b^*(k)M_c^*(q)\right] \\
H_\rho(k,q) &= \left[k_{b\mu}^* q_c^{*\mu} - 2M_b^*(k)M_c^*(q)\right].
\end{aligned} \tag{21}$$

The D_m's are the free meson propagators

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

The energy per nucleon is shown in Fig. 2. The saturation point of the energy density is quite independent of λ. The curves for $\lambda = 0$ and $\lambda = 4500\ MeV$ are indistinguishable. This is because the $\rho^0 - \omega$ energy is very small and the mean-field energy is by far the dominant term in Eq. (18). Only for a (fictitious) value of $\lambda = 20000\ MeV$ (dotted line), the effect of the mixing would be visible on the energy density.

In Fig. 3, it is shown the neutron-proton self-energy difference $\Sigma_n^s - \Sigma_p^s$ as a function of k_F. For $\lambda = 0.0$ (no $\rho - \omega$ mixing considered), we observe that this difference remains almost unaltered with density. When the mixing term is included, $\Sigma_n^s - \Sigma_p^s$ decreases with increasing density. Contrary to the case of the energy density, the neutron-proton self-energy difference is very sensitive to the $\rho - \omega$ mixing.

In fig. 4, ΔM_{eff} is plotted as a function of k_F. For $\lambda = 0.0$, ΔM_{eff} increases with density and for $\lambda = 4500\ MeV^2$ the difference decreases in a way similar to the self-energy difference. Figs. 3 and 4 show very clearly the crucial role played by the $\rho^0 - \omega$ mixing interaction.

In summary, we underline the main features investigated in this work. We have studied the importance of the CSB $\rho^0 - \omega$ interaction on the neutron-proton self-energy difference in nuclear matter within the context of the Walecka model. By employing the self-consistent Dirac-Hartree-Fock approximation to the nucleon propagator, we have shown that the inclusion of the CSB $\rho^0 - \omega$ mixing interaction in the Walecka model produces a neutron-proton self-energy difference which is a decreasing function of the nuclear density. Using the currently accepted value for the $\rho^0 - \omega$ mixing matrix element, $2m_\rho < \rho^0|H|\omega >= -4500\ MeV^2$, the neutron-proton self-energy difference at the normal nuclear matter density has a value of the order of $700\ KeV$. The obtained value and density dependence of the neutron-proton self-energy difference are very encouraging to implement a more complete calculation of binding energy differences of mirror nuclei employing relativistic models.

To finalize, we remark that although the situation regarding the resolution of the NSA by means of the $\rho^0 - \omega$ mixing interaction is very satisfactory, there remains a question of principle that might be worthwhile to investigate[18]. The question is related to the off-shell behavior of the $< \rho^0|H|\omega >$ matrix element. In constructing the CSB potential, or using the $\rho^0 - \omega$ "propagator" in nucleon self-energy diagrams, one uses for $< \rho^0|H|\omega >$ the value extracted at the ω pole. However, the four-momentum transfer carried by the exchanged mesons is spacelike. This matter was recently investigated[18] in a simple quark model for the $< \rho^0|H|\omega >$ amplitude. The conclusion was that off-shell effects are in fact very important and are such that the CSB potential is reduced in one order of magnitude compared to potential usually used. If this can be confirmed in more realistic calculations, the explanation of the NSA based on the $\rho^0 - \omega$ mixing mechanism, as well as the explanations of other[3] CSB phenomena, have to be re-examined and probably the quark substructure of nucleons and mesons have to be invoked.

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References

[1] J.A. Nolen and J.P. Schiffer, Annu. Rev. Nucl. Sci. 19 (1969) 414.

[2] N. Auerbach, Phys. Rep. 98 (1983) 273.

[3] G.A. Miller, B.M.K. Nefkens and I. Slaus, Phys. Rep. 194 (1990) 1.

[4] E.M. Henley and G.A. Miller, in: Mesons in nuclei, Vol. 1, eds. M. Rho and D.H. Wilkinson (North-Holland, Amsterdam, 1977) p. 405.

[5] P. Langacker and D.A. Sparrow, Phys. Rev. Lett. 43 (1979) 1559; Phys. Rev. C 25 (1982) 1194.

[6] S.A. Coon and M.D. Scadron, Phys. Rev. C 26 (1982) 562.

[7] P.G. Blunden and M.J. Iqbal, Phys. Lett. B 198 (1987) 14.

[8] G.A. Miller, Nucl. Phys. A 518 (1989) 345.

[9] L.M. Barkov et al. Nucl. Phys. B 256 (1985) 365.

[10] P.C. McNamee, M.D. Scadron and S.A. Coon, Nucl. Phys. A 249 (1975) 483.

[11] B.D. Serot and J.D. Walecka, Adv. Nucl. Phys. 16 (1985) 1.

[12] L.S. Celenza and C.M. Shakin, Relativistic nuclear physics: Theories of structure and scattering (World Scientific, 1986).

[13] Y. Nedjadi and J.R. Rook, Nucl. Phys. A 484 (1988) 525.

[14] C.J. Horowitz and B.D. Serot, Nucl. Phys. A399 (1983) 529.

[15] A. Bouyssy, J.-F. Mathiot, N.V. Giai and S. Marcos, Phys. Rev. C 36(1987)380.

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

[17] H.R. Christiansen, L.N. Epele, H. Franchiotti and C.A. Garcia-Canal, Phys. Lett. B 267 (1991) 164; L.N. Epele, H. Franchiotti, C.A. Garcia-Canal, G.A. Gonzalez Sprinberg and R. Mendez Galain, Phys. Lett. B 277 (1992) 33.

[18] T. Goldman, J.A. Henderson and A.W. Thomas, to be published in Few-Body Physics.

Figure 1

$\rho^0 - \omega$ mixing contribution to the neutron and proton self-energies.

Figure 2

Binding energy per nucleon as a function of k_F for $m_s = 550\ MeV$, $m_\pi = 138\ MeV$, $m_v = 738\ MeV$, $m_\rho = 770\ MeV$, $g_s^2/4\pi = 7.2$, $g_\pi^2/4\pi = 14.4$, $g_v^2/4\pi = 10.6$, $g_\rho^2/4\pi = 2.4$ and $\lambda = 0.0$ or $\lambda = 4500\ MeV^2$ (solid line) and $\lambda = 20000 MeV^2$ (dotted line).

Figure 3

Neutron-proton self-energy difference as a function of k_F for the same parameters as in Fig. 2 with $\lambda = 4500\ MeV^2$ (dashed line) and $\lambda = 0.0$ (solid line).

Figure 4

In-medium neutron-proton mass difference as a function of k_F for the same parameters as in Fig. 2 with $\lambda = 4500\ MeV^2$ (dashed line) and $\lambda = 0.0$ (solid line).

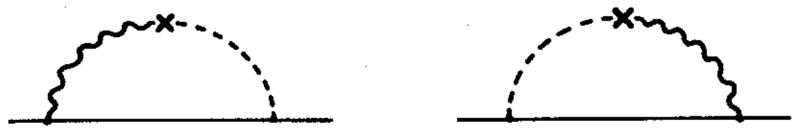

FIGURE 1

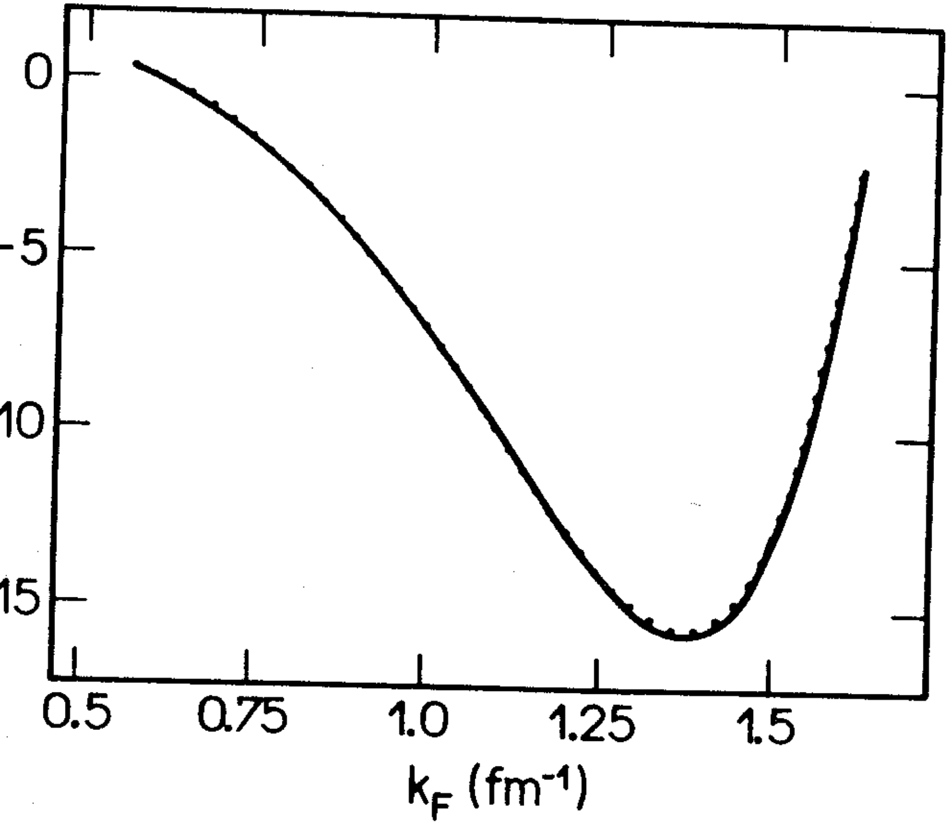

FIGURE 2

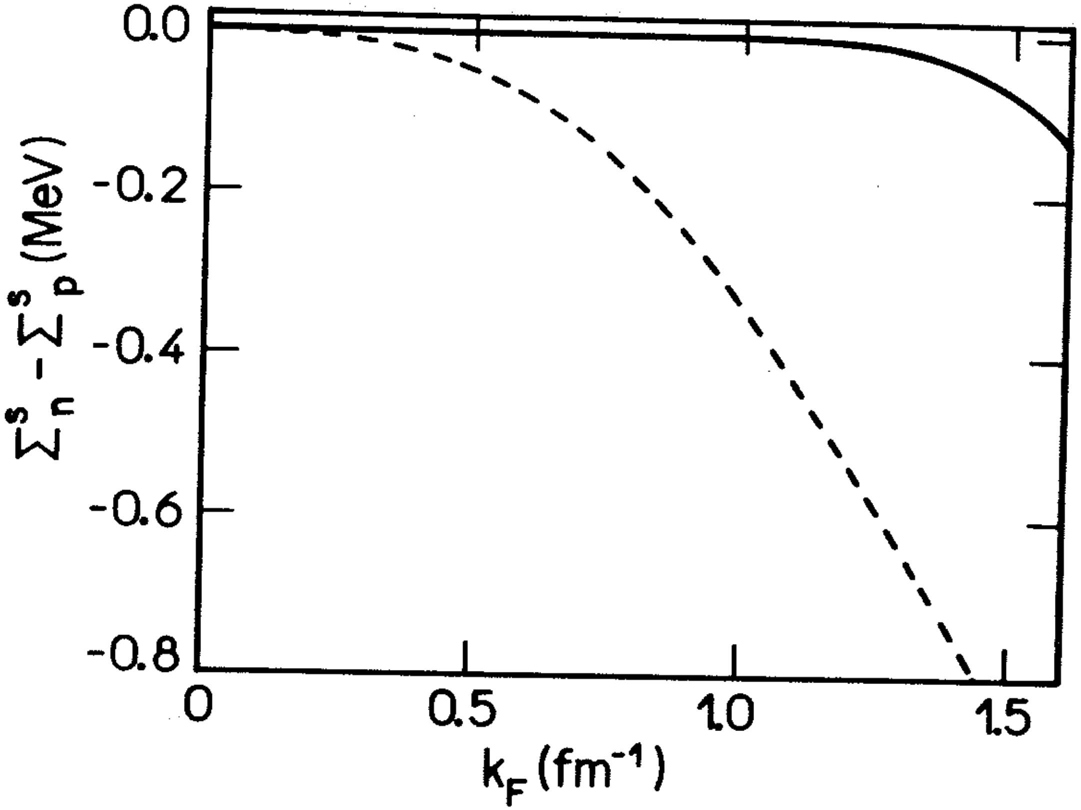

FIGURE 3

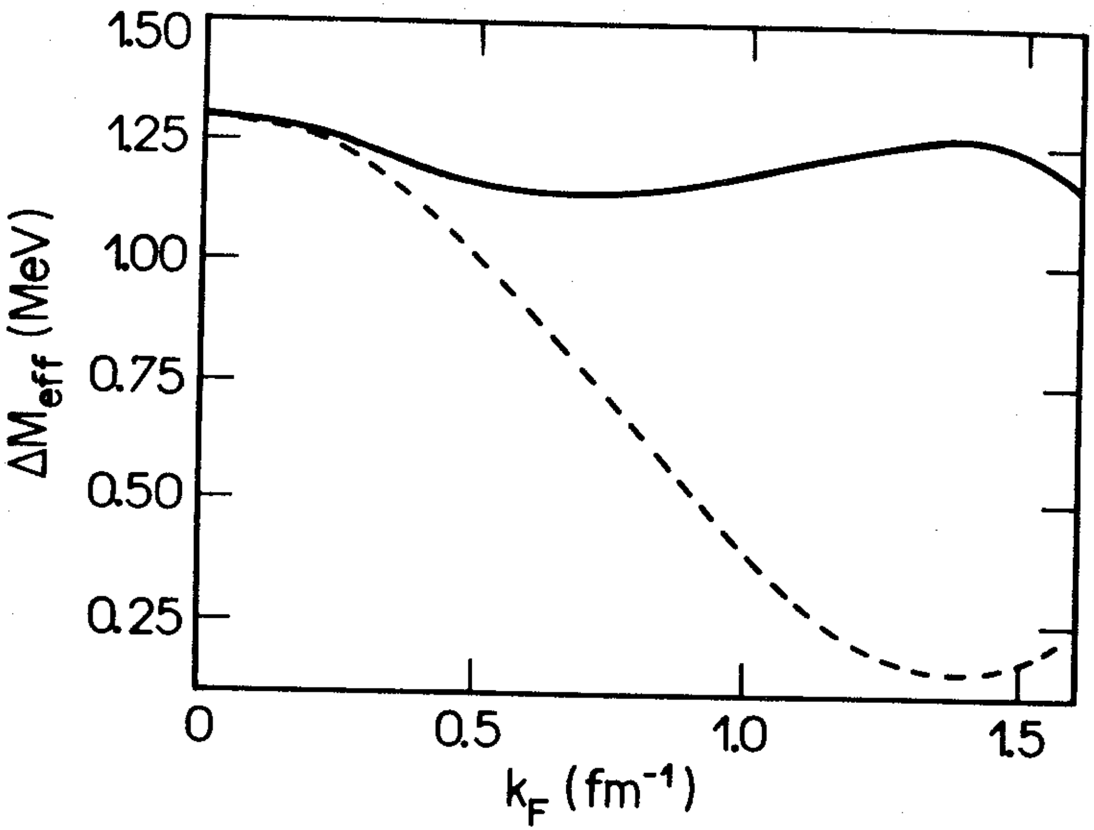

FIGURE 4