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## IN-MEDIUM NEUTRON-PROTON MASS DIFFERENCE WITH NUCLEAR CHIRAL MODELS

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# In-medium neutron-proton mass difference with nuclear chiral models


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The neutron-proton mass difference in nuclear matter is calculated within the context of chiral models involving nucleon and meson degrees of freedom. The neutron and proton self-energies are calculated in the Hartree-Fock approximation. Exchange terms are crucial to obtain different contributions for the neutron and proton self-energies. Density dependence of meson masses and coupling constants are taken into account. We find that the neutron-proton mass difference in nuclear matter increases as the density increases, contrary to the predictions of several quark models and of QCD sum rules at finite density.



21.65.+f, 11.30.Hv, 21.60.Jz, 21.10.Sf




## I. INTRODUCTION

The Nolen-Schiffer anomaly [1] (NSA) is a long-standing problem in nuclear physics. It is related to the failure of theory to explain the experimental mass differences of mirror nuclei. The mass difference of mirror nuclei (Z,N) and (N,Z), where a proton is replaced by a neutron, can be written as

$$\Delta M_{mirror} \equiv M(Z,N) - M(N,Z) = \Delta M_{``em"} - \Delta M_{np}, \qquad (1)$$

where $\Delta M_{np} = M_n - M_p$ is the free space neutron-proton mass difference and $\Delta M_{``em"}$ is the difference of eletromagnetic self-energies. $\Delta M_{``em"}$ is the Coulomb energy difference if charge independence of the nuclear forces is assumed. The NSA amounts to the fact that the experimental values of $\Delta M_{mirror}$ are systematically larger than the calcuted ones. The discrepancy between theory and experiment increases with the mass number $A$ and for $A \sim 209$ it can reach the value of 900 $KeV$.

Several nuclear structure effects have been invoked to solve the problem without definite success [2]. Recently, Henley and Krein [3] suggested a possible resolution to the NSA based on effects occurring at the level of the quark substructure of the nucleons, namely effects related to a partial restoration of chiral symmetry (PRCS) in nuclei. Using the Nambu–Jona-Lasinio model [4], supplemented with the Isgur and Karl model [5], they showed that as a consequence of PRCS, the value of $\Delta M_{np}$ in nuclei is smaller than its value in free space. Clearly, as $\Delta M_{np}$ decreases, the anomaly decreases.

The result of Henley and Krein was subsequently reobtained in an approach based on QCD sum rules [6]. In this approach, $\Delta M_{np}$ is connected directly with vacuum properties of QCD, namely with the scalar quark density $< \bar{q}q >$. This approach avoids the intermediate step of a phenomenological quark model for the nucleon, the only important input being the decrease of the quark scalar density in the medium.

In the past, charge-symmetry breaking effects were invoked to resolve the anomaly. In some sense, the suggestion of Ref. 3 is opposite to this, since it is based on the fact that the proton and neutron become isosymmetrical in nuclei as a consequence of PRCS. An important issue is to test whether the result $\Delta M_{np} \rightarrow 0$ in medium can also be obtained with chiral models involving nucleons and mesons only, without the recurrence to quark degrees of freedom. Such a study allows to gain insight on the role played by a PRCS realized at the level of the effective nuclear degrees of freedom on the subtle phenomenon of $\Delta M_{np} \rightarrow 0$. In this paper, we investigate the density dependence of the neutron-proton mass difference within the context of a simple nuclear chiral model, the linear $\sigma$-model [8], augmented with the $\omega$-meson to simulate the short-range nucleon-nucleon repulsion. Although the results obtained using either QCD or nuclear degrees of freedom are, so far, necessarily model dependent, there is the hope that common properties of the models lead to the same qualitative results, independent of the detailed dynamics. One example where such a hope is frustrated was shown by Griegel and Cohen [7]. They showed that the decrease of the effective $\sigma$-meson mass in medium, which is claimed to follow from the PRCS, is not a generic property of chiral models; it depends on the details of the model and on the way the mass is defined.

Although the initial motivation for the study of $\Delta M_{np}$ in nuclei was its possible relevance for the NSA, we think that such a study is interesting in itself. The consequences of $\Delta M_{np} \rightarrow 0$ should, in principle, be seen in other experimental situations, as for instance in (p,n) or (p,pn) scattering experiments. Of course, much remains to be done until the correct in-medium QCD condensates and their consequences are obtained; while this does not become available, one has to rely on the insights gained by using simple models.



In Sec. II we present the formalism used to derive the neutron-proton mass difference in nuclear matter. We use the relativistic Hartree-Fock approach to the nucleon propagator, in which tadpoles and exchange contributions are summed to all orders with Dyson's equation. The neutron and proton self-energies are specified by coupled nonlinear integral equations which we solve self-consistently in the so-called Dirac-Hartree-Fock approximation [9]. In Sec. III we present the numerical results for the momentum-dependent self-energies and the definition of effective nucleon masses is discussed. We will show that within the context of the nuclear models used in this paper, we obtain results contrary to the ones obtained with models dealing with quark degrees of freedom. Discussions and conclusions are presented in Sec. IV.

## II. NEUTRON AND PROTON SELF-ENERGIES

In this section we derive the equations for the effective masses of the neutron and of the proton in nuclear matter. Our starting point is the following Lagrangian density

$$\mathcal{L} = \mathcal{L}_S + \mathcal{L}_{SB}\,, \tag{2}$$

where $\mathcal{L}_S$ is the standard chirally symmetric $\sigma$-model Lagrangian, augmented with a neutral vector meson,

$$\mathcal{L}_S = \bar{\psi}(i\gamma_\mu\partial^\mu - g(\sigma + i\gamma_5\vec{\tau}\cdot\vec{\pi}) - g_v\gamma_\mu V^\mu)\psi + \frac{1}{2}(\partial_\mu\sigma\partial^\mu\sigma + \partial_\mu\vec{\pi}.\partial^\mu\vec{\pi})$$
$$-\frac{\lambda}{4}(\sigma^2 + \pi^2 - \sigma_0^2)^2 - \frac{1}{4}G_{\mu\nu}G^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu, \tag{3}$$

and $\mathcal{L}_{SB}$ is an explicit chiral symmetry breaking density

$$\mathcal{L}_{SB} = c\sigma - \bar{\psi}m_0\psi - \frac{1}{2}\mu_0^2\pi^{0^2} - \mu^2\pi^+\pi^-. \tag{4}$$

Here, $\psi$, $\sigma$, $\pi$ and $V^\mu$ stand respectively for the nucleon, the scalar-isoscalar meson, the pseudoscalar-isovector meson and the vector-isoscalar meson and $G^{\mu\nu} = \partial^\mu V^\nu - \partial^\nu V^\mu$. The vector field is added to simulate the short range nucleon-nucleon repulsion. In the explicit symmetry breaking Lagrangian we considered, in addition to the usual $c\sigma$, terms involving nucleon and pion masses. These account for the breaking of the isospin symmetry of the nucleon and of the pion. $m_0$ is a diagonal matrix in isospin space given by

$$m_0 = \begin{pmatrix} m_p & 0 \\ 0 & m_n \end{pmatrix}. \tag{5}$$

Chiral symmetry is realized in the Nambu–Goldstone mode when the $\sigma$ field acquires a nonzero vacuum expectation value, i.e., $<\sigma>_{vac} = u$. Shifting the $\sigma$ field as



$\sigma = u + s$, such that $<s>_{vac} = 0$, the Lagrangian density in Eq. (2) can be rewritten as

$$\mathcal{L} = \bar{\psi}(i\gamma_\mu\partial^\mu - M - g(s + i\gamma_5\vec{\tau}\cdot\vec{\pi}) - g_v\gamma_\mu V^\mu)\psi + \frac{1}{2}(\partial_\mu s\partial^\mu s - m_\sigma^2 s^2)$$
$$+\frac{1}{2}\partial_\mu\vec{\pi}.\partial^\mu\vec{\pi} - \frac{1}{2}m_{\pi^0}^2\pi^{0^2} - m_\pi^2\pi^+\pi^- - \frac{1}{4}G_{\mu\nu}G^{\mu\nu} + \frac{1}{2}m_v^2 V_\mu V^\mu$$
$$-g_1 s(s^2 + \pi^2) - g_2(s^2 + \pi^2)^2. \tag{6}$$

From this Lagrangian, the free-space masses of the particles can easily be identified. The nucleon mass matrix is given by

$$M = m_0 + gu = \begin{pmatrix} M_p & 0 \\ 0 & M_n \end{pmatrix}. \tag{7}$$

The $\sigma$-meson mass is

$$m_\sigma^2 = \lambda(3u^2 - \sigma_0^2), \tag{8}$$

and the pion masses are

$$m_\pi^2 = \lambda(u^2 - \sigma_0^2) + \mu^2, \qquad m_{\pi^0}^2 = \lambda(u^2 - \sigma_0^2) + \mu_0^2\,. \tag{9}$$

The undetermined masses $m_p$, $m_n$, $\mu_0$ and $\mu$ are chosen in such a way that we obtain the standard definitions for the coupling constants $g_1$ and $g_2$

$$g_1 = \lambda\mu = \frac{g}{2M_p}(m_\sigma^2 - m_{\pi^0}^2), \qquad g_2 = \frac{gg_1}{4M_p}\,. \tag{10}$$

As usual, $u$ is chosen to minimize the meson effective potential, which gives the relation

$$c = \lambda u(u^2 - \sigma_0^2)\,. \tag{11}$$

Notice that we could have started with an isosymmetrical Lagrangian in Eq. (3-4) and broken the isospin after shifting the $\sigma$ field, obtaining Eq. (6).

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

$$G(k) = G^0(k) + G^0(k)\Sigma(k)G(k), \tag{12}$$

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

$$\Sigma(k) = \Sigma^T + \Sigma^X(k). \tag{13}$$

The exchange terms, where the neutron and the proton (and the pions) have different masses, give different contributions to the neutron and proton self-energies, whereas the tadpoles contribute equally to the proton and neutron self-energies.

In order not to repeat standard formulae which can be found e.g. in Ref. (9-11), we show here the contributions to the self-energies coming from the pions only. 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 [9]

$$\Sigma(k) = \Sigma^s(k) - \gamma_0\Sigma^0(k) + \vec{\gamma}\cdot\vec{k}\Sigma^v(k) \tag{14}$$

The pionic contributions, shown in Fig. 1, are given by

$$\Sigma_p^X(k) = -ig^2\int\frac{d^4q}{(2\pi)^4}\gamma_5\left[\frac{G_p(q)}{(k-q)^2 - m_{\pi^0}{}^2} + 2\frac{G_n(q)}{(k-q)^2 - m_{\pi}{}^2}\right]\gamma_5\,, \tag{15}$$

$$\Sigma_n^X(k) = -ig^2\int\frac{d^4q}{(2\pi)^4}\gamma_5\left[\frac{G_n(q)}{(k-q)^2 - m_{\pi^0}{}^2} + 2\frac{G_p(q)}{(k-q)^2 - m_{\pi}{}^2}\right]\gamma_5\,. \tag{16}$$

From these equations, the effect of different meson and nucleon masses on the difference of the neutron-proton self-energies becomes clear.



We solve the coupled integral equations for the self-energies in the so-called Dirac-Hartree-Fock approximation [9]. 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 $\rho_0 = 2k_F^3/3\pi^2$)

$$G_b(k) = (\gamma_\mu k_b^{\mu *} + M_b^*(k))\frac{\pi i}{E_b^*(k)}\delta(k_b^0 - E_b(k))\theta(k_F - |\vec{k}|)\,, \tag{17}$$

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

$$k_b^{\mu *} = k^\mu + \Sigma_b^\mu(k) = (k^0 + \Sigma_b^0(k),\ \vec{k}(1 + \Sigma_b^v(k)))\,, \tag{18}$$

$$E_b^*(k) = \sqrt{(\vec{k}_b^*)^2 + M_b^*(k)^2}, \quad M_b^*(k) = M_b + \Sigma_b^s(k)\,, \tag{19}$$

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

$$E_b(k) = [E_b^*(k) - \Sigma_b^0(k)]_{k_b^0 = E_b(k)}\,. \tag{20}$$

Performing the angular and $k^0$ integrals in the expressions for the components of the self-energy, as those appearing in Eq. (15-16), we obtain the following coupled nonlinear integral equations

$$\begin{aligned}\Sigma_p^s(k) = \Sigma_s^T + \frac{1}{4\pi^2 k}\int_0^{k_F} dqq\Bigg[&\frac{M_p^*(q)}{E_p^*(q)}\frac{g^2}{4}(\Theta_{\sigma p}(k,q) - \Theta_{\pi^0 p}(k,q)) \\ &- \frac{M_p^*(q)}{E_p^*(q)}g_v^2\Theta_{vp}(k,q) - \frac{M_n^*(q)}{E_n^*(q)}\frac{g^2}{2}\Theta_{\pi n}(k,q)\Bigg]\end{aligned} \tag{21}$$

$$\Sigma_p^0(k) = -\frac{2}{3\pi^2}\frac{g_v^2}{m_v^2}k_F^3 - \frac{1}{4\pi^2 k}\int_0^{k_F} dqq\left[\frac{g^2}{4}(\Theta_{\sigma p}(k,q) + \Theta_{\pi^0 p}(k,q) + 2\Theta_{\pi n}(k,q)) + \frac{g_v^2}{2}\Theta_{vp}(k,q)\right] \tag{22}$$

$$\Sigma_p^v(k) = -\frac{1}{4\pi^2 k^2}\int_0^{k_F} dqq\left(\frac{q_p^*}{E_p^*(q)}\left[\frac{g^2}{2}(\Phi_{\sigma p}(k,q) + \Phi_{\pi^0 p}(k,q)) + g_v^2\Phi_{vp}(k,q)\right] + g^2\Phi_{\pi n}(k,q)\frac{q_n^*}{E_n^*(q)}\right). \tag{23}$$

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. The equation for $\Sigma_s^T$ is

$$\Sigma_s^T = -\frac{g^2}{\pi^2 m_\sigma^2}\int_0^{k_F} dqq^2\left(\frac{M_p^*(q)}{E_p^*(q)} + \frac{M_n^*(q)}{E_n^*(q)}\right) - \frac{3g_1}{gm_\sigma^2}(\Sigma_s^T)^2 - \frac{4g_2}{g^2 m_\sigma^2}(\Sigma_s^T)^3. \tag{24}$$

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

$$\Theta_{ib}(k,q) = ln\left|\frac{A_{ib}(k,q) + 2kq}{A_{ib}(k,q) - 2kq}\right| \quad \Phi_{ib}(k,q) = \frac{A_{ib}(k,q)\Theta_{ib}(k,q)}{4kq} - 1\,, \tag{25}$$

and

$$A_{ib}(k,q) = \vec{k}^2 + \vec{q}^2 + m_i^2 - (E_b(q) - E_b(k))^2\,. \tag{26}$$

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

Eq.(24) is the equation for the nucleon self-energy in a Hartree approximation to the linear $\sigma$-model [11], the last two terms coming from the potential energy in the Lagrangian of Eq.(6).



## III. NUMERICAL RESULTS

Eqs.(21-24) 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. It is known that such an iteration procedure does not converge [9,12], depending on the values of the parameters of the model, for nuclear densities larger than approximately two times the normal nuclear density. This is not a problem here, since we are interested in $\Delta M_{np}$ at low densities only, such as those occurring in medium to heavy nuclei.

In order to compare our calculation with the ones of Refs.(3,6), we need a definition for the in-medium nucleon mass. The natural guiding point is to look at the single-particle spectrum given by the pole position in the nucleon propagator in nuclear matter. In mean-field approximations (Hartree), the effective nucleon mass is defined as $M^{eff} = M + \Sigma^s_{MFT}$, where $\Sigma^s_{MFT}$ is the mean-field momentum-independent self-energy. This is a natural definition for the effective nucleon mass since the Dirac equation is then identical to that of a free particle with mass $M^{eff}$. In our case, the situation is not so simple because of the momentum dependence of the self-energies. Horowitz and Serot [9] define the relativistic effective mass $M^{eff}$ in analogy to the non-relativistic definition by the relation ($b = p, n$)

$$M_b^{eff}(q) = q\left[\left(\sqrt{\left(\frac{\partial E_b(q)}{\partial k}\right)^{-2} - 1}\right)_{k=q}\right]. \tag{27}$$

This definition coincides with the definition of $M^{eff}$ in the Hartree approximation and it is the one we adopt here. We will come back to this point later when we discuss the results. With this definition, we have for the in-medium neutron-proton mass difference (for $k = k_F$)

$$\Delta M_{np}(k_F) = M_n^{eff}(k_F) - M_p^{eff}(k_F). \tag{28}$$

The parameters of the model are $g$, $g_v$, $m_\omega$ and $m_\sigma$. Explicit one-loop calculations in the chiral limit of the pion self-energy [13] show that $u = f_\pi$, where $f_\pi$ is the physical pion decay constant. Therefore, if we use the experimental values for the proton mass and for the pion decay constant, from Eq. (11), the parameter $g$ has the value $g \sim 10$. It is then possible to reproduce the experimental value $g_{\pi NN} = 13.5$ in the linear $\sigma$-model using the Goldberger-Treiman relation

$$g = \frac{M_p}{f_\pi} = \frac{g_{\pi NN}}{g_A} \tag{29}$$

with $g \sim 10$. We assume that the Golberger-Treiman relation is valid in medium and so g is fixed. Thus, the only free parameters in Eq.(21-24) are $g_v$ and $m_\sigma$, once we take for the vector meson mass its free-space value $m_v = 783\ MeV$. In principle, $m_\sigma$ and $g_v$ can be fitted to known ground-state nuclear matter properties. However, in order to calculate, for example, the binding energy, it is necessary to include the meson loop terms in the meson self-energies [13]. Since we are not interested here in studying saturation properties of nuclear matter, we instead vary $m_\sigma$ and $g_v$ over a reasonably wide range of values; the saturation values must certainly be in the chosen interval.

In Fig. 2(a) we plot $\Delta M_{np}$ as a function of $k_F$ for $m_\sigma = 1000\ MeV$ and for several different values of $g_v$. In Fig. 2(b) we plot the same function for $g_v = 3.0$ and different values of $m_\sigma$. We note that the individual functions $M_n^*(k_F)$ and $M_p^*(k_F)$, which we do not plot, decrease with increasing density, similarly to the behavior obtained with NJL-based models, but their difference increases (at some density the difference of course will start to decrease). This shows that the self-energy corrections to the bare masses $M_n$ and $M_p$ are such that the neutron remains



heavier than the proton in medium. The definition of $M^{eff}$ through Eq.(27) does not spoil this fact. $M_n^{eff}$ and $M_p^{eff}$ decrease until the density reaches some value (which depends on the parameters used) and then start to increase. However, from Fig. 2 we note that $\Delta M_{np}$ *increases* as the density increases, in contrast to what is obtained in Refs. (3,7).

We have also studied in an approximate way the influence of possible density dependence of the meson masses and of the coupling constants on $\Delta M_{np}$. To calculate the density dependence of these quantities, we should solve coupled nonlinear integral equations for the meson self-energies and for the vertex functions, similar to the ones for the nucleon self-energies. This is clearly a formidable task. Here we simply parametrize this dependence as follows

$$m^i(\rho) = m_i(0)(1 - \alpha_i\ \rho/\rho_0), \tag{30}$$

$$g_v(\rho) = g_v(0)(1 - \alpha_v\ \rho/\rho_0), \qquad g(\rho) = g(0)(1 - \alpha\ \rho/\rho_0) \tag{31}$$

where $i = \sigma, \omega, \pi$, and $(\rho_0)$ $\rho$ is the (normal) nuclear density. The $\alpha$'s are varied arbitrarily. In Fig. 3 we plot $\Delta M_{np}$ as a function of $k_F$ for density dependent masses, with different combinations of $\alpha$'s. We observe that the tendency of an increasing neutron-proton mass difference is still obtained. Varying $g$ and $g_v$ according to Eq.(31), the same qualitative tendency of an increasing neutron-proton mass difference is obtained. Different combinations of $\alpha$'s do not change the qualitative results.

There has been some dispute in the literature[7,14] concerning the correct density dependence of meson masses and coupling constants. In Fig. 3 we have used only positive $\alpha^i$'s, i.e. decreasing masses as the density increases. We also studied the case of negative $\alpha^i$'s. The result is that we still get the same tendency, and the same order of magnitude, of an increasing neutron-proton mass difference. This is also true when negative $\alpha$'s are used for the coupling constants. Admitedly, the way we take into account the density dependence of the masses and the coupling constants is not entirely satisfactory, as momentum dependence of the vertices and self-energies have been neglected. However, we believe that the qualitative behavior we obtained for the neutron-proton mass difference in nuclear matter will not be modified in a more complete treatment.

We have also studied the model dependence of the results. We have neglected the potential energy of the $\sigma$-model in Eq. (7) and obtained in this way the traditional Walecka model [10] (with pions). This implies that the last two terms in the tadpole equation, Eq. (21), are now absent. Since the model is not chirally symmetric, the constraint of equal pion-nucleon and sigma-nucleon couplings is not required. Using the adequate parameters for the saturation of nuclear matter [9], we have recalculated $\Delta M_{np}$ and obtained the same qualitative result for $\Delta M_{np}$ as before, i.e., an increasing function of the density. The potential energy affects mainly the $\sigma$-meson mass [7], and this we have already shown not to affect our qualitative results.

To end this section, we remark that the eletromagnetic self-energies are extremely small compared to the nuclear self energies, and do not affect our results.



## IV. CONCLUSIONS

We have investigated the neutron-proton mass difference $\Delta M_{np}$ in nuclear matter within the context of a chiral model involving nucleon and meson degrees of freedom only. Our calculation employed the self-consistent Dirac-Hartree-Fock approximation to the nucleon propagator. We found that $\Delta M_{np}$ increases as the density increases. Previous calculations, using models where the chiral symmetry is realized at the quark level, found that $\Delta M_{np}$ decreases with the density.

Within the context of the chiral $\sigma + \omega$ model we have used in this paper, the density dependence of meson masses and coupling constants was addressed in an approximate way. By neglecting the momentum dependence of the vertices and meson masses, the density dependence of these quantities was taken into account by postulating an arbitrary, but presumably plausible, function of the density. The general trend of an increasing function $\Delta M_{np}$ with the density was obtained, independently of increasing or decreasing masses. Varying the coupling constants does not change the qualitative results. In addition to the correct density dependence of the meson self-energies and vertices, the problem of the Dirac sea has to be addressed as well. This is a difficult and yet unsolved problem. All nuclear structure calculations using relativistic models in the Hartree-Fock approximation suffer from this deficiency. The renormalization of the exchange diagrams introduces complications related to the appearance of "ghost poles" [15,16] which do not have a clear physical interpretation. How dependent our results are on the polarization of the Dirac sea is an open problem.

To conclude, we observe that it might well be that in order to implement chiral symmetry at the nuclear level, which has the correct restoration properties of the QCD condensates, more elaborate chiral models have to be considered. However, if the correct implementation of chiral symmetry requires the introduction of, for example, the $\rho$ and delta ressonances, much of the beauty of the chiral symmetry arguments would probably be lost in the complications of the nuclear many-body problems which appear when such ressonances are considered. We would then have a situation where the use of the QCD degrees of freedom would be more economical and transparent for treating a nuclear phenomenon.

### V. ACKNOWLEDGEMENTS

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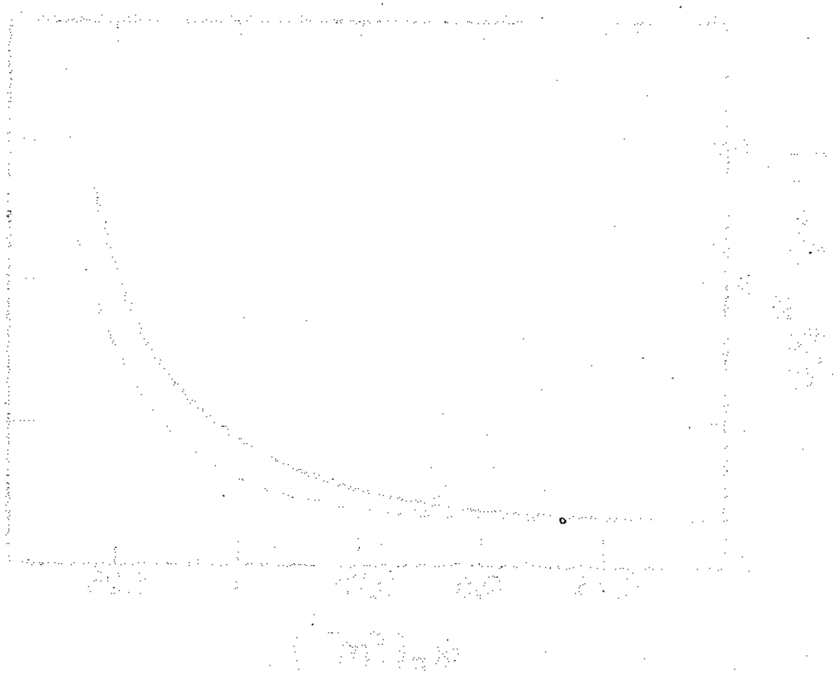




FIG. 1. The $\pi^0$ and $\pi^{\pm}$ exchange contribution to the neutron and proton self-energies

FIG. 2. In-medium neutron-proton mass difference as a function of $k_F$ for (a): $m_\sigma = 1000\ MeV$ and $g_v = 2.81$ (solid), $g_v = 0.5$ (dotted), $g_v = 15.0$ (dashed) and (b) for $g_v = 3.0$ and $m_\sigma = 1000\ MeV$(solid), $m_\sigma = 700\ MeV$ (dashed), $m_\sigma = 550\ MeV$ (dotted).

FIG. 3. In-medium neutron-proton mass difference as a function of $k_F$ for $g_v = 2.81$ and $m_\sigma$, $m_v$ and $m_\pi$ varying with density according to Eq.(31). The solid line is for $\alpha^i = 0$, the dashed line is for $\alpha_\sigma = \alpha_v = \alpha_\pi = 0.05$, and the dotted line is for $\alpha_\sigma = \alpha_v = \alpha_\pi = 0.1$.

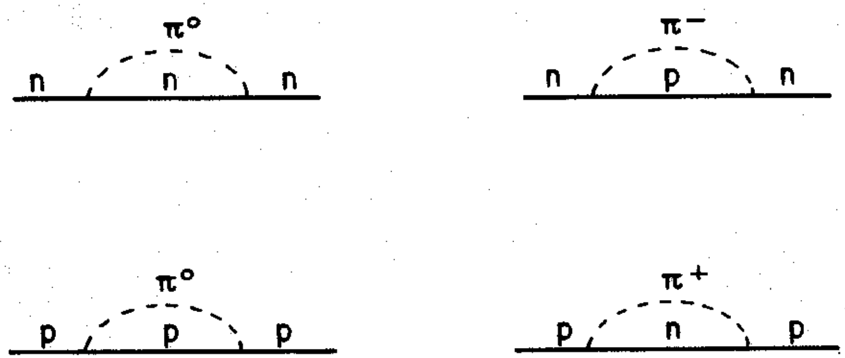



Fig. 1

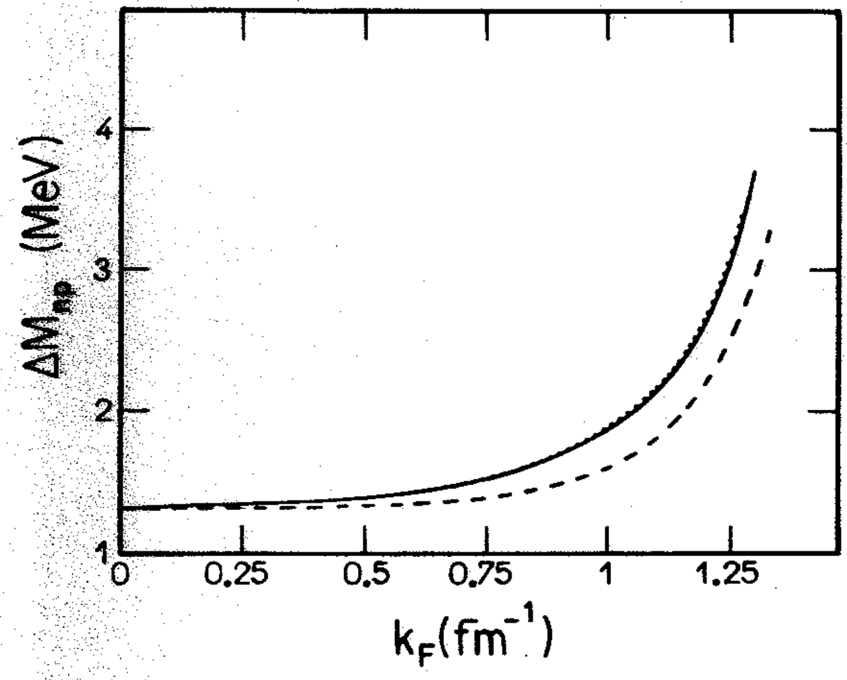
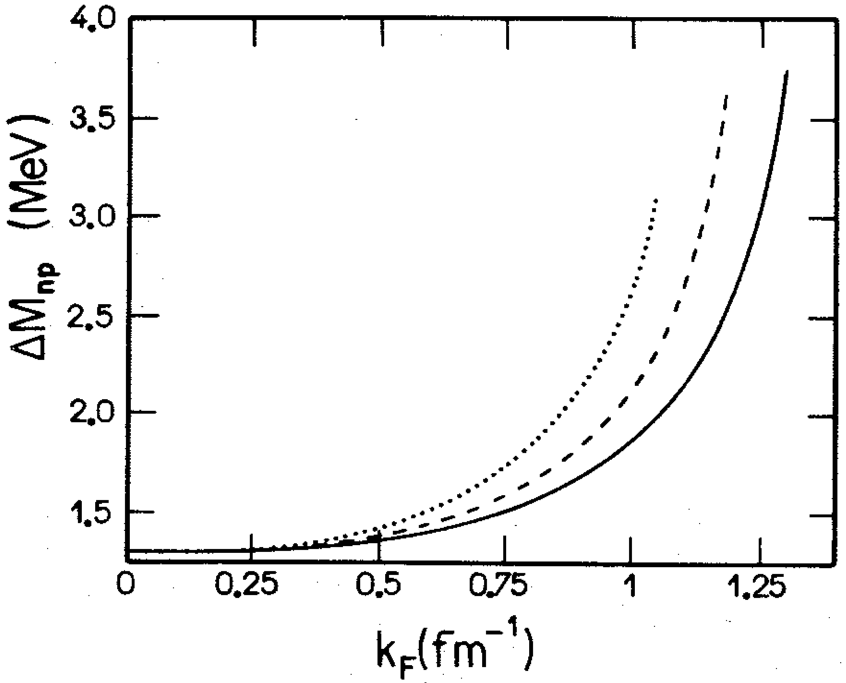
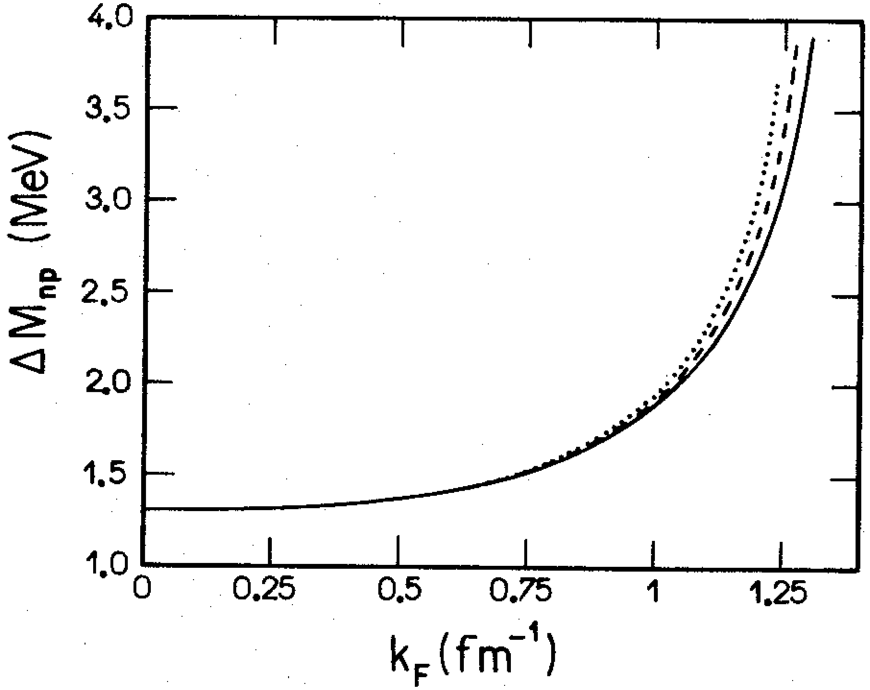



Fig. 2a



4.0
3.5
3.0
2.5
2.0
1.5
ΔM_np (MeV)
0
0.25
0.5
0.75
1
1.25
k_F(fm^-1)


Fig. 2b



4.0
3.5
3.0
2.5
2.0
1.5
1.0
ΔM_np (MeV)
0
0.25
0.5
0.75
1
1.25
k_F(fm^-1)


Fig. 3