Instituto de Física
Universidade de São Paulo

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Braghin, F. L.

Nuclear Theory and Elementary Particle Phenomenology Group,
Instituto de Física, Universidade de São Paulo, São Paulo, Brazil

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Collective modes in hot asymmetric nuclear matter at variable densities

Fábio L. Braghin *

Nuclear Theory and Elementary Particle Phenomenology Group,

Instituto de Física, Universidade de São Paulo

C.P. 66.318; CEP 05315-970; São Paulo - SP, Brazil.

Abstract

A nearly exact expression for the response function of hot asymmetric nuclear matter is derived for Skyrme type effective interactions and the resulting strength distribution in the four channels of the particle-hole interaction are analysed (isovector, scalar, spin and spin-isospin). Several proton-neutron asymmetries are considered as well as different total densities. In the isovector channel the linear response shows a very collective behaviour (zero sound type) typical from correlated systems which becomes still more collective with increasing asymmetry. The other channels become collective only for a highly asymmetric nuclear matter. In the spin and spin-isospin channels zero sound modes are found for higher enough p-n asymmetries. The critical temperature at which the zero sound modes disappear is given for several cases whenever it occurs. The static limit of the polarizabilities are explicitly given yielding the symmetry energy coefficient for isovector, spin and spin-isospin channels. The dependence of the polarizabilities on p-n asymmetry and temperature is analysed, in particular in the isovector channel which is of interest for the supernovae mechanism.

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*e-mail: braghin@axpfep1.if.usp.br

1 Introduction

The study of symmetric and asymmetric nuclear matter provides relevant information concerning systems like large nuclei, dynamics of neutron stars and the supernovae mechanism, as well as effects present in high/medium energy heavy ion collision. Among many usually employed approaches, the linear response method for non relativistic symmetric and asymmetric nuclear matter at zero and finite temperatures has been investigated in references [1, 2, 3, 4, 5]. It corresponds to a summation of chain diagrams and it is well-suited for the study of collective modes taking into account the medium/long range correlations [6].

In references [1, 2] there were motivations from the possibility of getting information about the effective nucleon-nucleon interaction and about the theoretical description of the nuclear isovector dipole giant resonance (IVDGR) on the ground state and with increasing excitation energy. Indeed, an investigation of the dependence of the collective modes on the effective interaction (Skyrme type) was done. In these works the response function of the symmetric nuclear matter to small external isovector perturbation using Skyrme interactions was calculated and a zero sound collective mode was found at low temperatures for reasonable values for the effective mass when density dependent forces were considered. This phonon has the following dispersion relation: $\omega_{res} = c_0 q$, where c_0 is the zero sound velocity and q the momentum transfer. The energy of the mode is mainly determined by the effective mass together with a function V_0 which is written in terms of the Skyrme interaction parameters (and is directly proportional to the respective Landau parameter [3]): the higher the effective mass (and V_0) the lower (higher) is the mode frequency, being more collective in both cases. This phonon typical from correlated Fermi liquids indicates that many-body correlations are not at all negligeable in nuclear matter/heavy nuclei, since heavier the nucleus more collective the mode. The model worked out in these references was based on the Steinwedel-Jenssen one [7], considering that the transfered momentum between protons and neutrons is related to the mass number of the analised nucleus. In particular, for heavy nuclei one should consider small $\mathbf{q}$. With the increase of temperature the collective mode couples to the particle-hole spectrum making zero sound damped, until its disappearence by the time when the temperature is of the order of some MeV. This critical temperature above which there is no more zero sound depends strongly on the used effective interaction.

The charge longitudinal response of nuclei has been investigated for large momentum transfer e.g. [3, 6] and a Fermi gas coherent behaviour (typical from an independent particle approximation) has been found for light nuclei. An investigation of the other three channels of the response function was done by Hernández *et al* in [4].

The case of asymmetric hot nuclear matter was investigated in the isovector channel in reference [5]. where two prescriptions for the density fluctuations were analysed. The one which leads to a more collective behaviour was choosen corresponding to a sounder physical picture. This effect is observed in nuclei.

In the following some motivations for the present work are discussed. Firstly, in heavy/neutron rich nuclei there is a non neglible asymmetry in the neutron-proton number and it is interesting to evaluate its influence on the collective motion and consequently on the nuclear dynamics. Thus one can hope to extract information for the effective interaction dependence on the neutron proton asymmetry.

As far as collective modes are concerned, pieces of information about giant resonances are expected to be found. In particular the Isovector Dipole Giant Resonance (IVDGR) has been observed in light exotic nuclei [8] (containing the so-called 'soft mode'), and an halo region has been recently predicted for the exotic ^{122}Zr nucleus [9] which would contain up to 6 neutrons. In this calculation a pairing interaction with density dependent effective interaction of zero range was used. As pointed out in [5] one can hope the IVDGR as well as its soft mode to take place in exotic heavier nuclei which would possess a larger neutron halo. However the IVDGR (a good review may be found in [10] is not the only collective dipolar mode which takes place in nuclei: scalar [11], spin-isovector [12] and spin [13] excitations have also been observed as well as the double-IVDGR has [14]. Nuclear matter collective/coherent modes (in the four channels) are expected to be, at least, qualitatively related to these resonances in nuclear matter.

Besides that asymmetric nuclear matter and neutron matter have been extensively studied with several astrophysical motivations [15]. For instance, during the collapse and explosion of massive stars (supernovae), neutrino scattering from nuclei produces its trapping for dense nuclear matter. The inelastic contribution from particle-hole excitation (via Z_0) and consequently from its phonons (density fluctuations) are important phenomena which need a consistent quantum treatment. In neutron stars, a strong correlated system, neutrino scattering by density fluctuations influentiates the neutrino mean free path. In this system one may expect zero sound in the spin-isovector and scalar channels for densities $\rho_{NM} \geq 2.\rho_0$ [16].

Notwithstanding, in collapsing supernovae the nuclear matter density is assumed to be $\rho \simeq .915\rho_0$, where ρ_0 is the saturation density and the neutron-proton asymmetry is taken to be $\alpha = 2\rho_n - \rho_0 = 1/3$, i.e. in terms of the asymmetry coefficient used in this paper $b = 1$, with pressure $P = 0$. The density used for these systems are not the equilibrium density in large nuclei/ nuclear matter. Thus a study with variable density, smaller and bigger than $\rho_0 \simeq .17 fm^{-3}$, appears to be interesting. The low density regime is also interesting since it sheds light on the growth of instabilities at finite temperatures, which are responsible for fragmentation [17]. In the context of heavy ion collisions thermally excited nuclear matter is studied in the determination of the amount of energy which can be absorbed by nuclear matter with increasing temperature. In this work, however, nucleon-densities are considered to not depend on the temperature.

Another important aspect of the response function of nuclear matter is that its static limit is directly related to the energy symmetry coefficient in the corresponding channel. For instance in [2] this was shown for the isovector channel which gives place to the neutron-proton symmetry energy coefficient of (symmetric) nuclear matter (a_τ). In the present paper this is worked out in the four channels. If symmetry energy increases with temperature electron capture in pre-supernovae phase is hindered and this causes an increase in electron trapping reducing the size of the core mantle. This makes the shock wave need less energy to stop collapsing and it exploses earlier. With the usually assumed temperature dependence of the symmetry energy the calculated shock wave is nearly half of the observed [19, 18]. It is suggested below that the static polarizability of asymmetric nuclear matter may be more suitable than the standard coefficient a_τ at zero and finite temperatures.

In the present article a detailed derivation of a nearly exact expression for the response function of asymmetric hot nuclear matter using Skyrme effective interactions (showed in reference [5]) is done and it is employed for a further investigation in the four channels of the effective interaction: scalar, spin, isovector and spin-isovector, for different neutron-proton asymmetry coefficients. Zero sound modes are found in the isovector (as in symmetric nuclear matter [2]), spin and spin-isospin channels for high enough n-p asymmetry at different nuclear densities. An investigation of the response dependence on the density of nucleons is done for $\rho = .5\rho_0$ and $\rho = 2\rho_0$, where ρ_0 is the saturation density. The static polarizabilities are explicitly written for symmetric and asymmetric nuclear matter with increasing temperature. The static polarizabilities are obtained in the limit $\omega/q \to 0$ and its dependence with temperature and neutron-proton asymmetry are studied. Some results are shown and related to other works where the dependence of the symmetry energy coefficient with temperature was calculated using different techniques [18, 19] that take into account that the (total) effective mass of nucleons contains the frequency dependence which comes from the coupling to collective nuclear vibrations [20]. It is discussed the relevance of the static isospin polarizability for the supernovae mechanism. In the spin and spin-isospin channels the respective (spin and spin-isospin) symmetry energy coefficients also correspond to the static polarizabilities and they are calculated at finite temperatures and considering neutron excess. In the scalar channel the polarizability defines a "dipolar compressibility " which is directly relate to the usual volume compressibility K as it will be shown below. Effects of the proton-neutron asymmetry are discussed. The energy-weighted sum rule is verified and well satisfied in all cases which have been analysed.

2 Linear response calculation

The time-dependent Hartree-Fock equation, shown below, determines the temporal evolution of the one-body density matrix ρ. In the presence of an infinitesimal external perturbation, the equation reads, in natural units:

$$i\partial_t \rho = [W + V_e, \rho], \tag{1}$$

where W is the Hartree-Fock energy of protons or neutrons and V_e is the small amplitude (ϵ) external field of the form:

$$V_e = \epsilon \hat{O} e^{-i\mathbf{q}.\mathbf{r}} e^{-i(\omega+i\eta)t}. \tag{2}$$

In the above equation η is an infinitesimal number which corresponds to the adiabatic switching on of the external source. The operator $\hat{O}$ may be τ_3 (third isospin Pauli matrix), σ_3 (third spin Pauli matrix), $\tau_3\sigma_3$ or an unit matrix, respectively for the isovector, spin, spin-isovector and scalar channels. Each of these perturbations induces density fluctuations around the static solution so that equation (1) can be linearized. In the following, the calculation of the response function for the isovector channel with an asymmetry in the neutron-proton densities is showed, but the procedure is the same for the other channels.

The mean energy for Skyrme- type effective interaction [21] is expressed in terms of the nucleon density (ρ_i), kinetic energy density (τ_i) and momentum density ($\mathbf{j}_i$), where i stands for protons or neutrons. The resulting approximated equation for the total fluctuation density in momentum space, $\delta\rho = \rho_n - \rho_p$, is:

$$\begin{aligned} i\partial_t \langle \mathbf{k}|\delta\rho|\mathbf{k}'\rangle = \;& (1+ac).(\epsilon'_{0p}(\mathbf{k}) - \epsilon'_{0p}(\mathbf{k}'))\langle \mathbf{k}|\delta\rho|\mathbf{k}'\rangle + 2(f_p(\mathbf{k}) - f_p(\mathbf{k}'))\langle \mathbf{k}|\delta W_p|\mathbf{k}'\rangle + \\ & +2(f_n(\mathbf{k}) - f_n(\mathbf{k}'))\langle \mathbf{k}|\delta W_n|\mathbf{k}'\rangle + 2\epsilon\,(f_n(\mathbf{k}') - f_n(\mathbf{k}) + f_p(\mathbf{k}') - f_p(\mathbf{k}))\,\delta(\mathbf{k}' - \mathbf{k} - \mathbf{q})e^{-i(\omega+i\eta)t}, \end{aligned} \tag{3}$$

where $\epsilon'_{0p}(\mathbf{k}) = \mathbf{k}^2(1+ac)/2m_p^*$, and m_p^* is the proton effective mass. In this expression $f_i(\mathbf{k})$ is the fermionic occupation number and three asymmetry coefficients have been used:

$$a = \frac{m_n^*}{m_p^*} - 1, \qquad b = \frac{\rho_{0n}}{\rho_{0p}} - 1, \qquad c = \frac{\delta\rho_n}{\delta\rho}. \tag{4}$$

The coefficient b is related to a frequently used asymmetry coefficient: $\alpha = 2\rho_{0n} - \rho_0$, by: $b = 2\alpha/(1-\alpha)$.

For the induced difference between proton and neutron density distribution we consider the following time-dependent prescriptions, analogously to the form of V_e:

$$\begin{aligned}
\langle \mathbf{r}|\delta\rho|\mathbf{r}\rangle &= \alpha e^{-i\mathbf{q}\cdot\mathbf{r}} e^{-i(\omega+i\eta)t}, \\
\langle \mathbf{r}|\delta\tau|\mathbf{r}\rangle &= \beta e^{-i\mathbf{q}\cdot\mathbf{r}} e^{-i(\omega+i\eta)t}, \\
\langle \mathbf{r}|\delta\mathbf{j}|\mathbf{r}\rangle &= \gamma \mathbf{q} e^{-i\mathbf{q}\cdot\mathbf{r}} e^{-i(\omega+i\eta)t}.
\end{aligned} \tag{5}$$

Due to the definition of the used densities [21] the above prescriptions are required to satisfy the following equalities:

$$(\alpha,\beta,\gamma) = \int \frac{\mathrm{d}^3k}{(2\pi)^3} \left(1, \mathbf{k}\cdot(\mathbf{k}+\mathbf{q}), \frac{1}{q^2}(2\mathbf{k}+\mathbf{q})\cdot\mathbf{q}\right) \langle \mathbf{k}|\delta\rho(t=0)|\mathbf{k}+\mathbf{q}\rangle. \tag{6}$$

From here on the index (s,t) for each (spin,isospin) channel is used. The deviation in the energy density of neutrons (and protons) can be written as:

$$\begin{aligned}
W_n - W_{0n} = \delta W_n = \; & 2\left(V_0^{0,1} + V_2^{0,1}\right)\delta\rho_n(\mathbf{r},t) + 2V_1^{0,1}\nabla\cdot\delta\rho_n(\mathbf{r},t)\nabla + \\
& +2V_1^{0,1}\delta\tau_n(\mathbf{r},t) + 2iV_1^{0,1}\left(\nabla\cdot\delta\mathbf{j}_n + \delta\mathbf{j}_n\cdot\nabla\right).
\end{aligned} \tag{7}$$

These functions $V_i^{0,1}$ are related to the parameters of the Skyrme interaction by the expressions of Appendix B.

In order to conclude this calculation another asymmetry coefficient is needed, a kinetic energy density dependent coefficient:

$$d = \frac{\beta_n}{\beta_p}. \tag{8}$$

In reference [5] two prescriptions were considered for coefficients c (from expression (4)) and d. They were taken either to be equal to 1/2 (prescription A) or calculated in terms of the saturation densities at zero temperature (prescription B). As discussed in that paper, prescription A leads to more reasonable results and will be considered. This means that density fluctuations should be proportional to the respective density of nucleons instead of being proportional to the perturbation amplitude. One considers:

$$c = \frac{1+b}{2+b}, \qquad d = \frac{1}{1+(1+b)^{2/3}}. \tag{9}$$

From equation (3) it is straightforward to show that the matrix elements at $t=0$ satisfy:

$$\begin{aligned}\langle \mathbf{k}|\delta\rho(t=0)|\mathbf{k}+\mathbf{q}\rangle &= \tfrac{1}{\epsilon_p'(\mathbf{k}+\mathbf{q})-\epsilon_p'(\mathbf{k})+\omega+i\eta}\Big\{4(V_0^{0,1}+V_2^{0,1}+V_1^{0,1}\mathbf{k}\cdot(\mathbf{k}+\mathbf{q}))\alpha(\delta f_n c-\delta f_p(c-1))+\\ &+2\epsilon(\delta f_n-\delta f_p)+V_1^{0,1}\beta(\delta f_n d-\delta f_p(d-1))+\\ &-V_1^{0,1}\gamma(2\mathbf{k}+\mathbf{q})\cdot\mathbf{q}(\delta f_n c''-\delta f_p(c''-1))+2\epsilon\delta f_n\Big\}.\end{aligned} \tag{10}$$

For this expression an ansatz for the total nucleon density fluctuation has been considered:

$$\delta\rho(t) = \delta\rho(t=0)e^{-i(\omega+i\eta)t}, \tag{11}$$

In expression (10) $\delta f_i = f_i(\mathbf{k}') - f_i(\mathbf{k})$, and another asymmetry coefficient is used:

$$c'' = \gamma_n/\gamma_p. \tag{12}$$

Its value was considered to be $c''=c$. The resulting response function does not depend sensitively on c''.

Multiplying equation (10) respectively by 1, $\mathbf{k}.(\mathbf{k}+\mathbf{q})$ and $(2\mathbf{k}+\mathbf{q}).\mathbf{q}$ and integrating them over $\mathbf{k}$ a set of linear equations for α, β and γ is obtained, namely:

$$\begin{aligned}\alpha &= (V_0^{0,1}+V_2^{0,1})P_{0c}\alpha+V_1^{0,1}P_{2c}\alpha+V_1^{0,1}P_{0d}\beta+2\gamma M_p^*\omega V_1^{0,1}\left(P_{0c''}+\frac{\epsilon}{2}(\Pi_{0n}+\Pi_{0p})\right),\\ \beta &= (V_0^{0,1}+V_2^{0,1})P_{2c}\alpha+V_1^{0,1}P_{4c}\alpha+V_1^{0,1}P_{2d}\beta+2\gamma M_p^*\omega V_1^{0,1}\left(P_{2c''}+\frac{\epsilon}{2}(\Pi_{2n}+\Pi_{2p})\right),\\ \gamma &= -\frac{2M_p^*\omega}{q^2\left\{1-4V_1^{0,1}M_p^*(\rho_{0n}c''+(1-c'')\rho_{0p}\right\}}\alpha.\end{aligned} \tag{13}$$

Solving the above linear system for α leads to the expression of the retarded response function $\Pi^R(\omega,q)$. This function is the polarizability, i.e., the ratio of the density fluctuation to the field strength: $\Pi^R(\omega,\mathbf{q}) = \alpha/\epsilon$. It results the following expression for the polarizability in the particular (s,t) channel:

$$\Pi_R^{s,t}(\omega,q) = \frac{\frac{1}{2}(\Pi_0^n+\Pi_0^p)(1-V_1^{s,t}P_{2d})+\frac{1}{2}V_1^{s,t}(\Pi_2^n+\Pi_2^p)P_{0d}}{1-\overline{V_0^{s,t}}P_{0,c}-V_1^{s,t}(P_{2,c}+P_{2,d})+V_1^{s,t}\overline{V_0^{s,t}}(P_{0,c}P_{2,d}-P_{2,c}P_{0,d})+(V_1^{s,t})^2(P_{2,c}P_{2,d}-P_{4,c}P_{0,d})}. \tag{14}$$

In the above expression we have used $M_p^* = m_p^*/(1+ac)$ and:

$$P_{2i,v} = v\Pi_{2i}^n(\omega,q)+(1-v)\Pi_{2i}^p(\omega,q). \tag{15}$$

In this expression $v = c, d, c''$ (they act as a weight for the proton and neutron Lindhard functions) and $i = 0, 1, 2$. The functions Π^i_{2N} are referred to as generalized Lindhard functions. They are defined as:

$$\Pi^i_{2N} = \Pi^i_{2N}(\omega, \mathbf{q}) = \frac{4}{(2\pi)^3} \int \mathrm{d}^3k \frac{f_q(\mathbf{k}+\mathbf{q}) - f_q(\mathbf{k})}{\omega + i\eta - \epsilon'_p(\mathbf{k}) + \epsilon'_p(\mathbf{k}+\mathbf{q})} \left(\mathbf{k}.(\mathbf{k}+\mathbf{q})\right)^N, \tag{16}$$

Expressions for $\Pi^i_{2N}(\omega, q)$ are shown in the Appendix.

In equation (14) we have used a modified coefficient $\overline{V_0^{0,1}}$ defined by:

$$\overline{V_0^{s,t}} = V_0^{s,t} + V_2^{s,t} - \left(\frac{2M_p^*\omega}{q}\right)^2 \frac{2V_1^{s,t}}{1 - 4V_1^{s,t} m_p^* \rho_{0p}}. \tag{17}$$

This modified coefficient arises because the change in the momentum density induced by the external field and also it is also due to the asymmetry between protons and neutrons. The complete expressions for all the four channels in terms of Skyrme parameters are shown in appendix B. Expression (14) generalizes that of references [2] to which it reduces when considering a symmetric nuclear matter.

3 Some properties of the dynamic - and static- polarizability

The strength distribution per unit volume, $S(\omega, q)$, corresponds to real transitions of the system, being then proportional to the imaginary part of the dynamic polarizability. It is directly related to the photoabsorption strength distribution S_{abs}, being a direct consequence from the fluctuation-dissipation theorem [22, 23]. Both are proportional to the imaginary part of the polarizability:

$$S(\omega, q) = (1 - e^{-\omega/T}) S_{abs}(\omega, q) = -\frac{1}{\pi} \Im m \Pi_R^{s,t}(\omega, q). \tag{18}$$

The denominator of the polarizability determines the existence of a collective mode. When there is a pole in $\Pi_R^{s,t}(\omega, q)$, the real part of the denominator yields the energy of the resonance while its imaginary part is directly related to the width.

As far as the momentum transfer is concerned with the resonant frequency it is worth to recall that in the model worked out in [1, 2] q is directly related to the radius of a nucleus as long as the nucleus is identified with a "box" inside nuclear matter. For a heavy nucleus the momentum transfer is small, in particular, for the nucleus of lead it was considered that: $q = \pi/(2R) = .23 fm^{-1}$. In this relation the usual nuclear radia were considered to be: $R = r_0 A^{\frac{1}{3}}$. This parametrization leads to the usual behaviour for medium/large nuclei: $\omega_{res} = BA^{-\frac{1}{3}}$.

It is important to check the results by considering the energy weighted sum rule (EWSR). For the dipolar modes considered here one is left with the following expression:

$$m_1^{s,t} = \frac{< 0|[\hat{O}^{s,t}, [H, \hat{O}^{s,t}]]|0 >}{2} \tag{19}$$

For the asymmetric nuclear matter it yields:

$$m_1^{s,t} = \int_0^\infty d\omega S^{s,t}(\omega, q)\omega = q^2 \left(\frac{\rho_{0p}}{m_p^*}(1 - 2V_1^{s,t} m_p^*) + \frac{\rho_{0n}}{m_n^*}(1 - 2V_1^{s,t} m_n^*) \right). \tag{20}$$

It is still interesting to remember that the following dispersion relation hold for the real (χ') and imaginary (χ'') parts of the retarded response function:

$$\begin{aligned} \chi' &= \int d\omega' S(q,\omega')\mathcal{P}\left(\frac{2\omega'}{\omega^2 - \omega'^2}\right), \\ \chi'' &= -\pi\left(S(q,\omega) - S(q,-\omega)\right). \end{aligned} \tag{21}$$

This dispersion relations were well studied in [4, 22].

As it will be shown in the results, at low frequencies ($\omega << qv_F$) the strength distribution of particle-hole spectrum have the general behaviour: $S(\omega) \propto \omega$.

It was pointed out in [2, 24] a simplified model can be used without loosing the main features of the complete response function given above. It corresponds to neglect the function $V_1^{s,t}$. This roughly corresponds to the idea that the main effect of the spacial non-locality (velocity-dependence) of the effective interaction is present in the effective mass calculation and almost absent in what concerns the collectivity. The expression of the response function becomes:

$$\Pi_R^{s,t}(\omega, q) = \frac{\frac{1}{2}(\Pi_0^n + \Pi_0^p)}{1 - \overline{V_0^{s,t} P_{0,c}}}. \tag{22}$$

This expression is tested in the section 4.1. A Thomas-Fermi type approximation also produces similiar results [2].

3.1 Static polarizabilities

In the static limit $\omega/q \to 0$ of the response function one obtains the static polarizabilities for the respective channel at finite temperature and for asymmetric matter. One obtains:

$$\begin{aligned} \Pi^{0,1}(\omega \to 0, q \to 0, T, b) &= \frac{\rho_0}{2A_\tau(T,b)}, \\ \Pi^{1,0}(\omega \to 0, q \to 0, T, b) &= \frac{\rho_0}{2A_\sigma(T,b)}, \\ \Pi^{1,1}(\omega \to 0, q \to 0, T, b) &= \frac{\rho_0}{2A_{\sigma\tau}(T,b)}, \\ \Pi^{0,0}(\omega \to 0, q \to 0, T, b) &= \frac{3\rho_0}{K_D(T,b)}. \end{aligned} \tag{23}$$

At zero temperature in the symmetric nuclear matter limit these static polarizabilities reduces to the known p-n symmetry energy coefficients: $a_\tau = A_\tau(T \to 0, b \to 0)$, the symmetry energy coefficient; $a_\sigma = A_\sigma(T \to 0, b \to 0)$ the spin symmetry energy coefficient, $a_{\sigma\tau} = A_{\sigma\tau}(T \to 0, b \to 0)$ the spin-isospin symmetry energy coefficient, $K_D = A_K(T \to 0, b \to 0)$ a dipolar compressibility of nuclear matter. This last coefficient is related to the usual nuclear matter volume compressibility (K_∞) by the following expression:

$$K_D = K_\infty + \frac{4}{5}T_F - 2V_1 k_F^2 \rho_0 + \frac{3}{4}t_3\rho 0^{\alpha+1}, \tag{24}$$

where k_F is the momentum at the Fermi surface. The values for these coefficients will be shown below.

The symmetry energy coefficients may then be written in the following form:

$$\begin{aligned}
a_\tau &= \frac{T_F}{3} + \frac{\rho_0}{2}V_0^{0,1} + k_F^2\rho_0 V_1^{0,1}, \\
a_\sigma &== \frac{T_F}{3} + \frac{\rho_0}{2}V_0^{1,0} + k_F^2\rho_0 V_1^{1,0}, \\
a_{\sigma\tau} &== \frac{T_F}{3} + \frac{\rho_0}{2}V_0^{1,1} + k_F^2\rho_0 V_1^{1,1}, \\
K_D &= 3. \left(\frac{2}{3}T_F + \rho_0 V_0^{0,0} + 2k_F^2\rho_0 V_1^{0,0}\right),
\end{aligned} \tag{25}$$

where T_F is the kinetic energy at the Fermi surface and V_0 and V_1 are given in appendix B.

4 Results

In this section the resulting strength distributions are plotted for several cases with different asymmetry coefficient and variable densities from $0.5\rho_0$ up to $2\rho_0$ in the four channels. In this work the SGII [25] is used in all but one figures. It was adjusted considering spin and spin-isospin observables as well as ground state observables like compressibility and, as other skyrme interactions, reproduces well several ground state properties. A non realistic skyrme force, SV, which has no density dependence is used in one exemple in order to stress the importance of this term.

4.1 Normal density asymmetric polarizabilities in the isovector channel

In figure 1 the strength distribution is shown for the case of an asymmetry coefficient corresponding to that of the nucleus of ^{208}Pb, i.e., b=.54 (equivalent to $\alpha = .23$) at several excitation energies: T from 0 to 6 MeV together with the zero temperature strength for the symmetric nuclear matter, see [2, 24]. The collectivity for such a small asymmetry coefficient does not change a lot but the energy of the resonance is nearly 1 MeV higher. The behaviour of the strength with the increase of the temperature is the same for the symmetric nuclear matter analysed in [1, 2]. Some cases of larger asymmetry are compared to the symmetric response function at zero temperature in figure 2. The dependence of $S(\omega, q)$ on the asymmetry is very strong and the resonance energy as well as the collectivity become much bigger for higher b. There is a small bump in the lower part of the spectrum (particle-hole excitations) for the case b=2 and one still smaller for b=8.

The strength for the simplified model discussed in section 3 is shown in figure 3 for b=0, .54 and b=8 at zero temperature. Comparing the collectivity by regarding the poles contribution to the total sum rule, shown in table 1, it is noted that collectivity diminishes without considering the function V_1 which is composed by non local terms. This is the opposite tendency of the results obtained for symmetric nuclear matter [24]. The energy of the mode also is a little smaller- to be compared with figure 2.

In figure 4 the case of a higher transfered momentum ($q = 2 \times .23 fm^{-1}$) is exhibited for the cases of b=0, .54 and 8 at zero temperature: it is twice the momentum associated with the nucleus of lead as discussed in section 3 (q=.23 fm^{-1}). This figure should be compared to figure 2. It is seen that the energy of the mode is twice the original one and (consequently) they become less collective. Also the collective mode contribution to the EWSR is nearly the same (greater the asymmetry closer are the percentual contribution of the modes due to q=.23 and q=.46 fm^{-1}). The energy of the resonance is twice the original making remember the double-IVDGR mode in nuclei [14].

The relevance of the effective mass and of the n-p coefficient b is explicited by fitting the resonant energies by the following expressions. In symmetric nuclear matter the energy of the resonance, which is mainly determined by the effective mass for a fixed V_0, can be parametrized nearly by $E^{res} \simeq Cq/m^*$, where C is a constant [2]. In asymmetric nuclear matter the energy shift due to the asymmetry may be written as: $\Delta E^{res}(b) \simeq ALog(4b) + 2b$, with $A \simeq 2.3$ for $b > 0$.

In figure 5 the response function calculated with the skyrme force SV, which contains no density dependent terms, is shown for b=.54 and exhibits no qualitative difference from that of the symmetric nuclear matter: there is no collectivity. Indeed, its effective mass is too low ($m^* \simeq .4m$).

For the very collective cases showed above there always a critical temperature at which the pole in the response function disappears. It is very interaction-sensitive and some values are given below. For b=0 this disappearence is found to happen at nearly $T \simeq 4 - 5MeV$. The absence of the interaction function V_1 makes it change a little to $T \simeq 5 - 6MeV$. For the Skyrme interaction SkM, not analysed in the rest of the paper, which differs from SGII mainly for the compressibility and symmetry energy coefficient (they are bigger than the values obtained with SGII) $T_c \simeq 7 - 8MeV$. Increasing the neutron-proton asymmetry makes the critical temperature greater, for instance, considering b= .54 and b=2 it is found $T_c \simeq 11MeV$ and $T_c \simeq 15MeV$ respectively. For the supernova case (b=1 and $\rho = .915\rho_0$) it is obtained $T_c \simeq 8-9MeV$ in the isovector channel and $T_c \simeq 6MeV$ in the spin-isovector one. Note that in these last channel for these values of b and ρ there is a collective zero sound which does not take place for a symmetric nuclear matter at normal density (shown in section 4.3).

4.2 Varying nuclear density

The variation of nuclear matter density is achieved by adjusting the chemical potential present in the Fermi-Dirac occupation number which, by the way, is an important piece of the Lindhard functions. In figure 6 the response function for a half density of the equilibrium one ($\rho = \rho_0/2$) is shown for different asymmetry coefficients (b=0,.54,8). The mode become less collective as it can be seen from table 1, and its energy is smaller if compared to figure 2 (normal density case).

A higher density case ($\rho = 2\rho_0$) is shown in figure 7. The modes become more collective for big proton-neutron asymmetry and they are pushed to higher energies.

A specific exemple of astrophysical interest is shown in the next figure 8: as expected to happen in supernovae the density is taken to be $\rho = .915\rho_0$ and the asymmetry ($\alpha = 1/3$) b=1. The resonance follows the behaviour previously described as well as its temperature dependence.

As far as the temperature dependence is concerned some remarks on the damping mechanism are useful. For the zero sound mode there are two mechanisms of damping: firstly the collisions between nucleons which may be present even at low T, and also the Landau damping which occurs at low frequencies smaller than that of the Fermi surface $\omega \leq qv_F$. In spite of that, zero sound does not propagate in the full collison regime which takes place at non zero temperatures but in a (at least nearly) non equilibrium system [22]. In the results shown in this paper there is clearly only Landau damping.

4.3 Normal and double density asymmetric polarizabilities in the spin, spin-isospin and scalar channels

The case of the scalar channel is exemplified in figures 9 and 10. In the first one, the strength function for b=.54 is compared to that of symmetric nuclear matter at T=0 MeV and is a little less collective.

The opositte tendency is seen at a density $\rho = 2\rho_0$ in the figure 10.

The spin channel is analysed in figures 11 and 12. In figure 11 the normal density cases for b=0, .54 and b=8 are exhibited and, as expected, the mode becomes more collective with increasing asymmetry [16]. The double density case is shown in figure 12 and the energy of the coherent (which becomes more collective) mode is pushed away from the particle-hole spectrum. There is spin zero sound at higher n-p asymmetry b.

Finally the spin-isospin channel is shown in figures 13, 14 and 15. At normal density (figure 13) the general behaviour of the isovector channel is observed but it is much less collective. The case with a higher density is exhibited in figure 14 and the case suited for the supernovae medium (b=1, $\rho = .915\rho_0$) in figure 15. The collective mode is also of the zero-sound type.

In short there are coherent/collective motion in these three channels only for higher proton-neutron asymmetries or/and higher densities. For (very) small b, these channels exhibit rather a behaviour mainly identified with a non correlated Fermi gas which seems to be in qualitative agreement with experimental results from the (spin and scalar) giant dipole resonances. The coherent modes do not have the zero sound relation dispersion but the collective ones do have.

4.4 Static polarizabilities

In table 2 the denominators of the static polarizabilities are shown as well as their values for different asymmetric nuclear matter with interaction SGII. The temperature dependence is negligeable mainly because, in this work, the density is kept constant independently of the temperature for a fixed asymmetry coefficient b. It is also shown, in this table, the values for these coefficients after the analysis of [26] with the parametrization of Landau theory for Fermi liquids for the available experimental fits.

It is seen that the static polarizabilities associated with the energy symmetry coefficients in each channel depend strongly on the asymmetry coefficient b. One can expect these polarizabilites ($A_{s,t}$) to be more appropriate that the symmetry energy coefficients ($a_{s,t} = A_{s,t}(T = 0, b = 0)$) if a finite temperature asymmetric nuclear matter calculation is considered, as, for exemple, in the electron capture in hot asymmetric supernovae matter [16]. For this system, larger the symmetry energy contribution (a_τ or $A_{0,1}$) more quenched is the neutronization and, consequently, higher is the lepton fraction at trapping. This helps the supernovae explosion causing a stronger shockwave.

In reference [18] an approach from the Fermi gas model was considered (i.e. considering that a_τ has well defined potential and kinetic parts). An increase of nearly 2.5 MeV for the asymmetry coefficient was found with the augmentation of temperature from 0 to 1 MeV. This would happen because the total effective mass was considered, i.e. the coupling to the vibrational states of the nuclear medium yields a temporal non locality to the mass, increasing it at low temperatures. In contrast to this result a Monte Carlo calculation of [19] shows almost no variation for that coefficient. In the present work the temperature dependence is also absent, as seen in the last column of table 2, because matter density does not depend on T. Precise experimental fits to test these results as well as the spin and spin-isospin energy symmetry coefficients (a_σ and $a_{\sigma\tau}$) are needed.

5 Conclusions

A nearly exact expression for the response function of hot asymmetric nuclear matter with different densities of nucleons was calculated, extending the work of [5] and generalizing the case previously studied of symmetric nuclear matter with fixed density [2]. In the isovector channel the collectivity of the response was analysed as well as its evolution with increasing excitation energy until the disappearence of the zero sound whose critical temperature depends very strongly on the interaction and on the proton-neutron asymmetry coefficient b. At normal density the spin, and scalar channel show no collective behaviour that can happen for very asymmetric nuclear matter. The spin-isospin channel exhibit some coherence which increases a lot with the parameter b. In these last cases the increase of b (and ρ) causes the appearence of a zero sound in the respective channel at the respective channel.

The zero sound phonon is a coherent superposition of states that can be obtained by giving extra momentum $\hbar q$ to one of the quasiparticles. This disturbance propagates only at very low temperatures or very high frequency. It propagates in a collisionless regime, and collisions mainly start to take place with increasing energy excitation, destroying collectivity. The presence of such collective modes indicates that nuclear matter/large nuclei are quite correlated systems and these many-body correlations present in the linear response approach should be taken into account in the description of many important observables mainly at zero temperature. It is a behaviour typical from quantum liquids where zero point motion prevents from forming a solid. The critical temperature beyond which zero sound does not take place anymore as well as the momentum transfer between nucleons are expected to be measured.

It is important to keep in mind that the used Skyrme effective interactions have been fitted at normal density and, in principle, are not suitable neither for much higher densities, when relativistic effects are very important, nor for very low densities. This, however, does not seem to invalidate the main conclusions of this article. It would be thus of great interest to study the dependence of the effective interaction on the proton-neutron interaction, density of protons and neutrons in nuclear matter and temperature. These cases would provide new fits for the effective (not only Skyrme) interaction in order to have reproduced the respective phenomenological aspects for each channel. A detailed investigation of the dependence of the results of this paper on the used (existing) Skyrme interaction is left for a forthcoming publication [27].

Static polarizabilities were also explicitly calculated exhibiting almost no variation with the temperature, what happens mainly because the densities were considered to not depend on T. These static polarizabilities (with T and b dependence) may be more suitable than the usual zero temperature symmetry energy coefficient for phenomena where asymmetric nuclear matter is considered.

Acknowledgements

The author wishes to thank FAPESP, Brazil, for the financial support and D. Vautherin who introduced the astrophysical subject to him.

Appendix A: Expressions of the asymmetric generalized Lindhard functions

In this appendix we give the explicit expressions of the real and imaginary parts of the generalized asymmetric Lindhard functions defined in Eqs. (16). For this purpose we need to define the following integrals:

$$I^i_{2N} = \frac{2}{(2\pi)^3} \int \mathrm{d}^3k k^{(2N)} \frac{(f^i(\mathbf{k}+\mathbf{q}) - f^i(\mathbf{k}))}{\omega + i\eta + \epsilon'_p(\mathbf{k}+\mathbf{q}) - \epsilon'_p(\mathbf{k})}. \tag{A.1}$$

The imaginary parts of the generalized Lindhard functions are given by

$$\begin{aligned}
\Im m\Pi^i_0(\omega,\mathbf{q}) &= -\frac{m^*_i M^*_p}{\pi q \beta} \log \frac{1+e^{\beta(\mu_i - E_- \frac{M_p}{m^*_i})}}{1+e^{\beta(\mu_i - E_+ \frac{M^*_p}{m^*_i})}}, \\
\Im m\Pi^i_2(\omega,\mathbf{q}) &= -\frac{2m^{*2}_i M^*_p}{\pi \beta^2 q} \left(\beta \frac{M^*_p}{m^*_i} \sqrt{E_+ E_-} \log \frac{1+e^{\beta(\mu_i - E_- \frac{M^*_p}{m^*_i})}}{1+e^{\beta(\mu_i - E_+ \frac{M^*_p}{m^*_i})}} + \mathrm{Li}_2\left(1+e^{\beta(\mu_i - E_+ \frac{M^*_p}{m^*_i})}\right) \right. \\
&\quad \left. -\mathrm{Li}_2\left(1+e^{\beta(\mu_i - E_- \frac{M^*_p}{m^*_i})}\right)\right), \\
\Im m\Pi^i_4(\omega,\mathbf{q}) &= -q\sqrt{2m^*_i E_+}\Im m\Pi^i_2 + 2\Im m I^i_4 - 2\sqrt{2m^*_i E_+} q \Im m I^i_2,
\end{aligned} \tag{A.2}$$

where Li_2 is the Euler dilogarithmic function [28]

$$\mathrm{Li}_2(x) = \int_1^x \frac{\log(t)}{t-1}\mathrm{d}t, \tag{A.3}$$

and

$$E_\pm = \frac{M_p^*}{2q^2}(\omega \pm \frac{q^2}{2M_p^*})^2. \tag{A.4}$$

The expressions of the functions $\Im m I_2^i$ and $\Im m I_4^i$ are :

$$\Im m I_2^i(\omega,q) = -\frac{m_i^{*2}M_p^*}{\pi q \beta^2}\left(\beta E_+ \frac{M_p^*}{m_i^*}\log\frac{1+e^{\beta(\mu_i - E_- \frac{M_p^*}{m_i^*})}}{1+e^{\beta(\mu_i - E_+ \frac{M_p^*}{m_i^*})}} + \mathrm{Li}_2(1+e^{\beta(\mu_i - E_+ \frac{M_p^*}{m_i^*})}) - \mathrm{Li}_2(1+e^{\beta(\mu_i - E_- \frac{M_p^*}{m_i^*})})\right),$$

$$\Im m I_4^i(\omega,q) = -\frac{2M_p^{*3}m_i^{*2}}{\pi q \beta}E_+^2\left(\log\frac{1+e^{\beta(\mu_i - E_- \frac{M_p^*}{m_i^*})}}{1+e^{\beta(\mu_i - E_+ \frac{M_p^*}{m_i^*})}} + 2\int_1^\infty \mathrm{d}z z \log\frac{1+e^{-\beta(zE_+\frac{M_p^*}{m_i^*} - \omega\frac{M_p^*}{m_i^*} - \mu_i)}}{1+e^{-\beta(zE_+\frac{M^*}{m_i^*} - \mu)}}\right). \tag{A.5}$$

At zero temperature the real parts of the Π_{2N} are given by

$$\begin{aligned}
\Re e\Pi_0(T=0) = &\ \frac{M_p^* k_F}{\pi^2}\left(-1 + \frac{k_F}{2q}\left[\phi(x_+) + \phi(x_-)\right]\right),\\
\Re e\Pi_2(T=0) = &\ \frac{M_p^* k_F^3}{2\pi^2}\Big(-3 + x_+x_- + x_+^2 + x_-^2 + \\
&\ \frac{k_F}{2q}\left[(1 - x_+^2 - 2x_+x_-)\phi(x_+) + (1 - x_-^2 - 2x_+x_-)\phi(x_-)\right]\Big),\\
\Re e\Pi_4(T=0) = &\ 2\Big(\Re e I_4(T=0) - 2q\sqrt{2M_p^*E_+}\Re e I_2(T=0) + \\
&\ q^2 M_p^* E_+ \Re e\Pi_0(T=0) + \frac{1}{3\pi^2}M_p^* q^2 k_F^3\Big),
\end{aligned} \tag{A.6}$$

where

$$\begin{aligned}
\Re e I_2(T=0) = &\ \frac{M_p^* k_F^3}{4\pi^2}\left(-3 - x_+x_- - x_+^2 + x_-^2 + \frac{k_F}{2q}\left[(1+x_+^2)\phi(x_+) + (1 + x_-^2 + \frac{4m^*\omega}{k_F^2})\phi(x_-)\right]\right),\\
\Re e I_4(T=0) = &\ -\frac{M_p^* k_F^6}{2\pi^2 q}\Big(\big(\frac{1}{6}(1 + x_-^2 + x_-^4) + \frac{M_p^*\omega}{k_F^2}(1+x_-^2) + \frac{2M_p^{*2}\omega^2}{k_F^4}\big)\phi(x_-) + \\
&\ \frac{1}{6}(1 + x_+^2 + x_+^4)\phi(x_+) + \frac{5q}{3k_F} + \frac{1}{3}(\frac{1}{3}x_-^3 + \frac{1}{3}x_+^3 + \frac{2M_p^*\omega x_-}{k_F^2} + \frac{8q^2x_+}{k_F^2}) + \\
&\ \frac{x_-^5}{3} + \frac{x_+^5}{3} + \frac{2M_p^*\omega x_-^3}{k_F^2} + \frac{4M_p^{*2}\omega^2 x_-}{k_F^4}\Big),
\end{aligned} \tag{A.7}$$

with $x_\pm = \frac{q}{2k_F} \pm \frac{M_p^*\omega}{qk_F}$, and $\phi(x) = (1-x^2)\log|\frac{x-1}{x+1}|$.

For non zero temperature the expression of $\Re e\Pi_{2N}$ (N=0,1,2) is an average of the zero temperature functions calculated for the same values of ω and q, but with various values of the Fermi momentum

k_F distributed with a weight factor which is just the derivative of the Fermi occupation number. Explicitly one has the following formula:

$$\begin{aligned}
\Re e\Pi_0^i(\omega,q,T) &= -\int \Re e\Pi_0(\omega,q,T=0,k_F=k)\mathrm{d}f^i(k,T), \\
\Re e\Pi_2^i(\omega,q,T) &= -\int \Re e\Pi_2(\omega,q,T=0,k_F=k)\mathrm{d}f^i(k,T), \\
\Re e\Pi_4^i(\omega,q,T) &= -\int \Re e\Pi_4(\omega,q,T=0,k_F=k)\ \mathrm{d}f^i(k,T),
\end{aligned} \tag{A.8}$$

where $f^i(k,T)$ is the occupation number. For the case of zero temperature we have

$$\mathrm{d}f^i(k) = -\delta(k-k_F^i)\mathrm{d}k,$$

yielding the above expressions for these functions.

Appendix B: Functions $V_i^{s,t}$

In the isovector channel one obtains:

$$\begin{aligned}
V_0^{0,1} &= \left(-\frac{t_0}{2}\left(x_0+\frac{1}{2}\right)-\frac{t_3}{12}\left(x_3+\frac{1}{2}\right)\rho_0^\alpha-\frac{q^2}{16}(3t_1(1+2x_1)+t_2(1+2x_2))\right)(1+bc), \\
V_1^{0,1} &= \frac{1}{16}(t_2(1+2x_2)-t_1(1+2x_1)), \\
V_2^{0,1} &= t_3\left(\frac{1}{2}+x_3\right)\rho_0^{\alpha-1}(c\rho_{n0}-(c-1)\rho_{p0})/12,
\end{aligned} \tag{B.1}$$

where ρ_{0n}, ρ_{0p} and ρ_0 are the proton, neutron and total saturation densities of asymmetric nuclear matter. For the other channels:

$$\begin{aligned}
V_0^{0,0} &= \left(3\frac{t_0}{4}+(\alpha+1)(\alpha+2)\frac{t_3}{16}\rho_0^\alpha+q^2(9\frac{t_1}{32}-(5+4x_2)\frac{t_2}{32})\right)(1+bc), \\
V_1^{0,0} &= 3\frac{t_1}{16}+(5+4x_2)\frac{t_2}{16}, \\
V_2^{0,0} &= \frac{t_3}{12}(x_3+.5)(c\rho_n+(c-1)\rho_p\rho_0^{(\alpha-1.)}), \\
V_0^{1,0} &= \left(-(1-2x_0)\frac{t_0}{4}-(1-2x_3)\frac{t3}{24}\rho_0^\alpha-q^2(3(1-2x_1)\frac{t_1}{32}+(1+2x_2)\frac{t_2}{32})\right)(1+bc), \\
V_1^{1,0} &= -(1-2x_1)\frac{t_1}{16}+(1+2x_2)\frac{t_2}{16}, \\
V_2^{1,0} &= -\frac{t_3}{48}\rho_0^{\alpha-1}(2+\alpha)(\rho_nc+\rho_p(c-1)), \\
V_0^{1,1} &= \left(-\frac{t_0}{4}-\frac{t_3}{24}\rho_0^\alpha+q^2(-3\frac{t_1}{32}-\frac{t_2}{32})\right)(1+bc), \\
V_1^{1,1} &= \frac{-t_1+t_2}{16}, \\
V_2^{1,1} &= -\alpha\frac{t_3}{24}\rho_0^{(\alpha-1)}(\rho_nc-\rho_p(c-1)).
\end{aligned} \tag{B.2}$$

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Figure captions

Figure 1 Distribution of strength per unit volume for the operator $\tau \exp(i\mathbf{q}.\mathbf{r})$ (in fm^{-2}) - isovector channel - as a function of the energy ω (in MeV). Asymmetry $b = 0$ (solid line) at $T = 0$ MeV, and $b = .54$ for different values of the temperature $T = 0, 2, 4$ and 6 MeV, in the case of interaction SGII with $q = .23 fm^{-1}$.

Figure 2 Distribution of strength per unit volume - isovector channel. Parameters: b= 0, 2, 8 and 32 at T=0 MeV and with $q = .23 fm^{-1}$.

Figure 3 Distribution of strength per unit volume - isovector channel neglecting function of interaction V_1. Parameters: b=0, .54 and 8 at T=0 MeV.

Figure 4 The same of figure 1 for $q = 2 \times .23 fm^{-1}$.

Figure 5 The same of figure 1 for interaction SV, without the case b=0.

Figure 6 The same of figure 1 for different density $\rho = .5\rho_0$.

Figure 7 The same of figure 1 for different density $\rho = 2\rho_0$.

Figure 8 Distribution of strength per unit volume - isovector channel for b=1, density $\rho = .915\rho_0$ for $T = 0, 2, 4, 6$ MeV.

Figure 9 Distribution of strength per unit volume - scalar channel. Parameters b=0 and .54, transfered momentum $q = .23 fm^{-1}$.

Figure 10 The same of figure 9 but for density $\rho = 2\rho_0$.

Figure 11 Distribution of strength per unit volume - spin channel. Parameters b=0, .54 and 8, transfered momentum $q = .23 fm^{-1}$.

Figure 12 The same of figure 11 but for b=0 and .54 at density $\rho = 2\rho_0$.

Figure 13 Distribution of strength per unit volume - spin-isovector channel. Parameters b=0, .54 and 8, transfered momentum $q = .23 fm^{-1}$.

Figure 14 The same of figure 13 but for density $\rho = 2\rho_0$.

Figure 15 The same of figure 14 but: $\rho = .915\rho_0$ and b=1 for different temperatures T=0, 2, 4 and 6 MeV.

Table 1: Right hand side (RHS) of the energy weighted sum rule (MeV $\times$ fm^{-3}) compared to the contribution of the particle-hole spectrum of the strength $S(\omega)$ obtained numericaly for $T = 0$, 3 and 6 MeV for some cases analysed in the Text (isovector channel). The missing strength is due to the pole.

b, $\frac{\rho}{\rho_0}$	RHS	$m_1(T=0)$	$m_1(T=3)$	$m_1(T=6)$	RHS($q = 2 * q_0$)	$m_1(T=0)$
0, 1	53	28	53	53	212	212
.54, 1	53	16	53	53	212	212
.54, 1, no V_1	45	14	45	45	–	–
8, 1	53	.44	.44	.47	212	2
.54, .5	22	6	22	22	–	–
.54, 2	141	94	140	141	–	–

Table 2: Static polarizabilities in the four channels: isovector channel ($A_\tau = \rho_0/2\Pi^{0,1}$), spin-isovector ($A_{\sigma\tau} = \rho_0/2\Pi^{1,1}$), spin ($A_\sigma = \rho_0/2\Pi^{1,0}$) and scalar channel (dipolar compressibility: $A_K = \rho_0/2\Pi^{0,0}$), for different proton-neutron asymmetry coefficients b in the case of interaction SGII. In the "exp" column it is shown most probable experimental values after the analyses of [26] based on Landau Fermi liquid theory. In the last column the variation of the coefficients for $\Delta T = 6MeV$ when b= .54.

Channel, ρ_0	b=0	b=.54	b=2	exp.	$\Delta A_{s,t}$
A_τ (MeV)	26.9	34.7	59.5	$28 \rightarrow 38$	.05
$A_{\tau,\sigma}$ (MeV)	23.4	29.5	48.4	$39 \rightarrow 41$	.06
A_σ (MeV)	15.7	15.9	55	$10 \rightarrow 15$	.01
A_K (MeV)	14.7	11.2	-2.6	–	.2

Figure 1

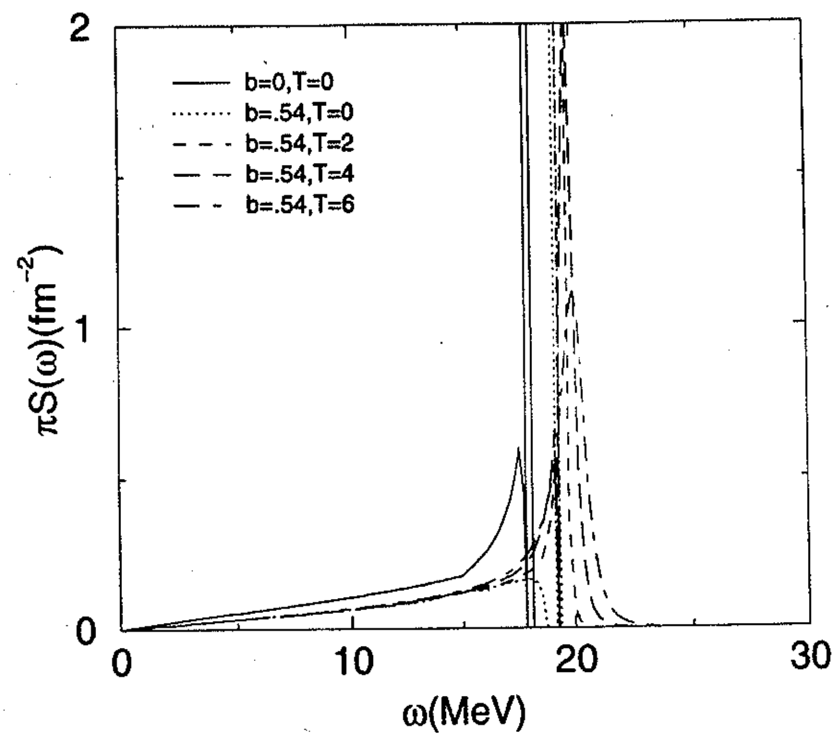

Figure 2

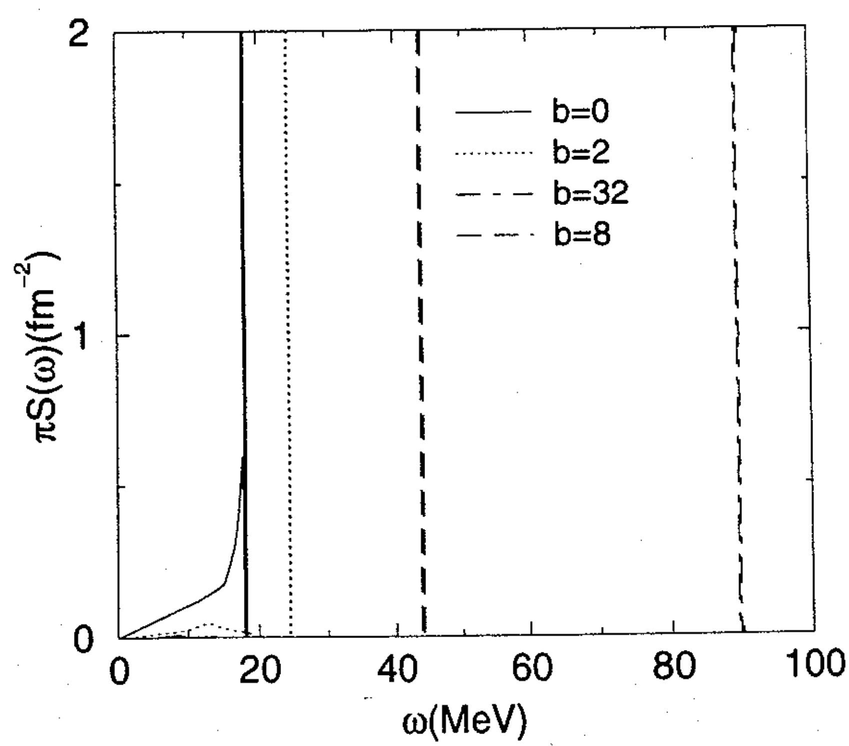

Figure 3

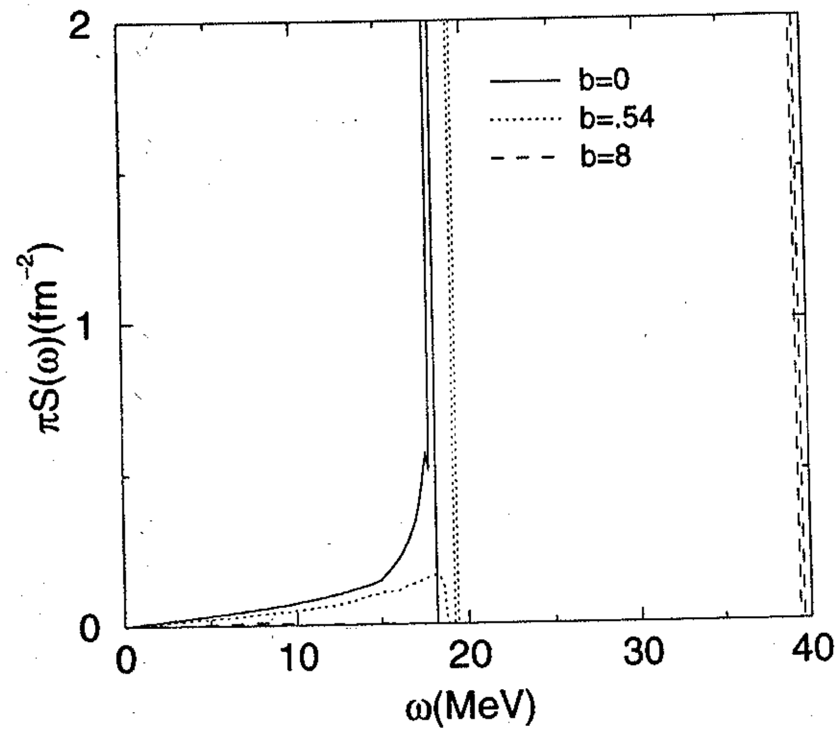

Figure 4

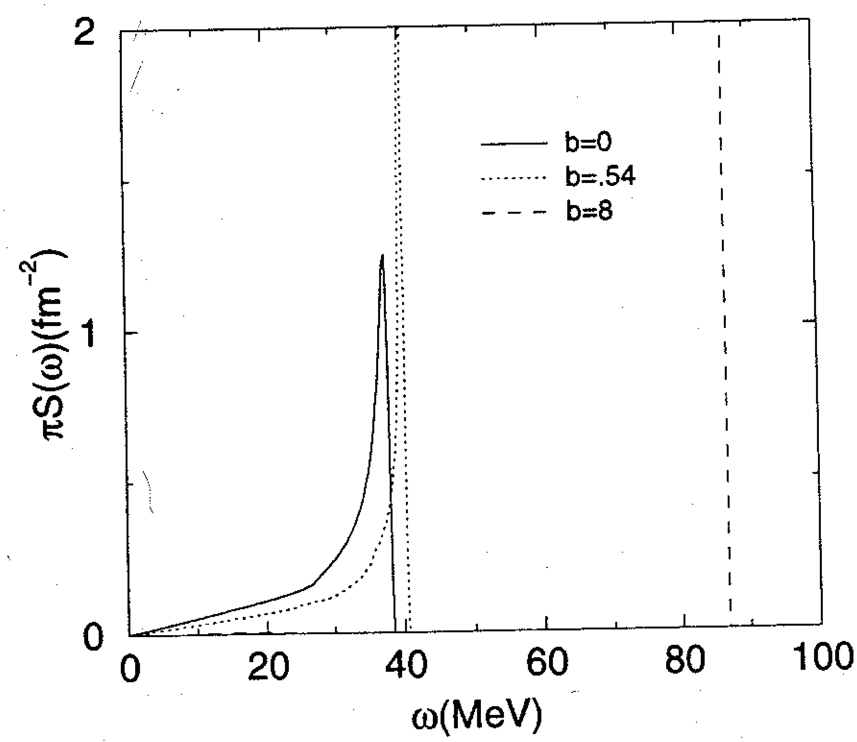

Figure 5

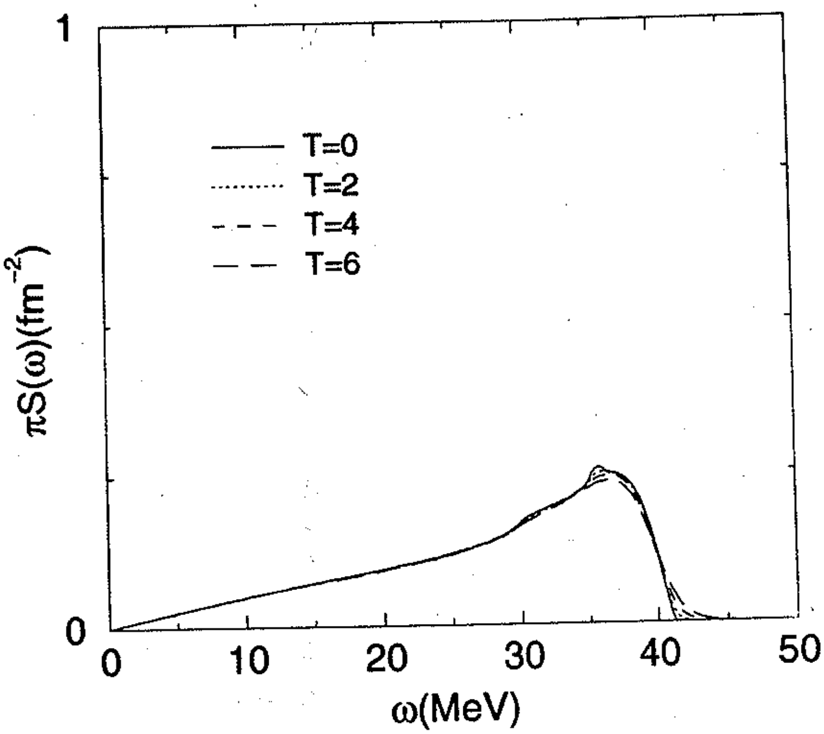

Figure 6

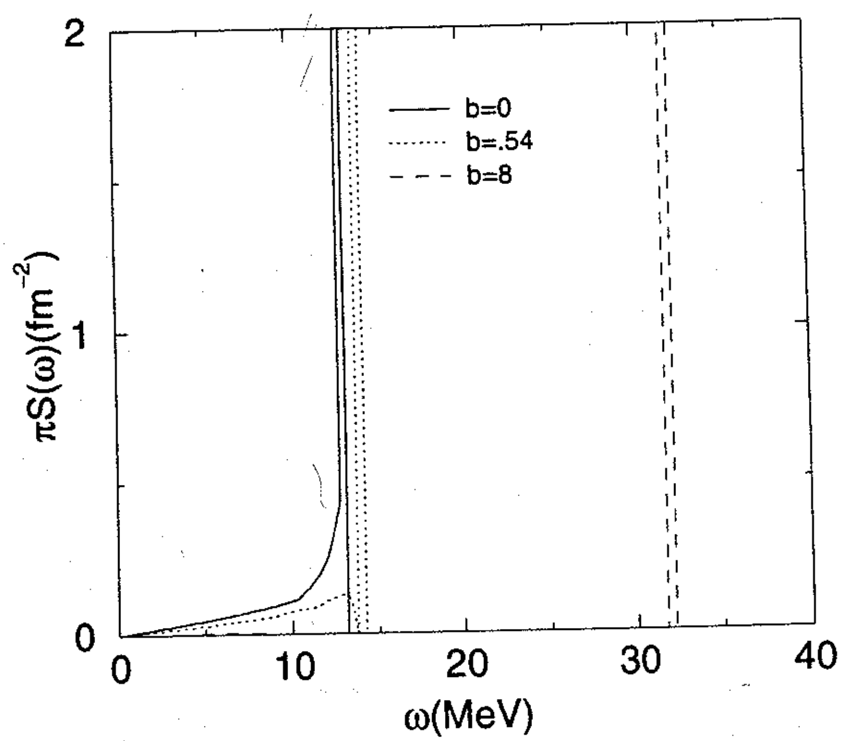

Figure 7

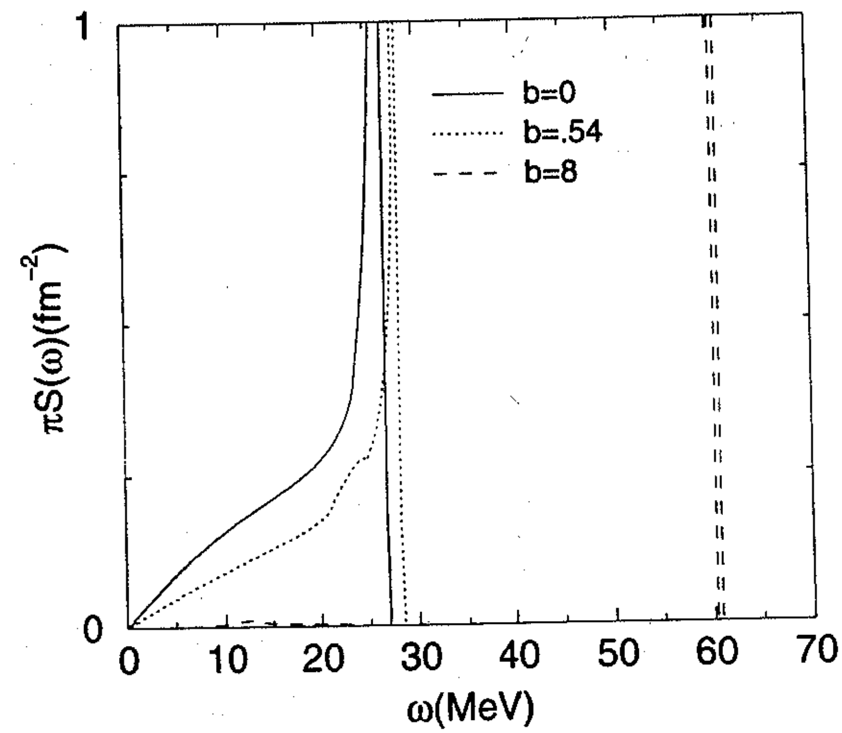

Figure 8

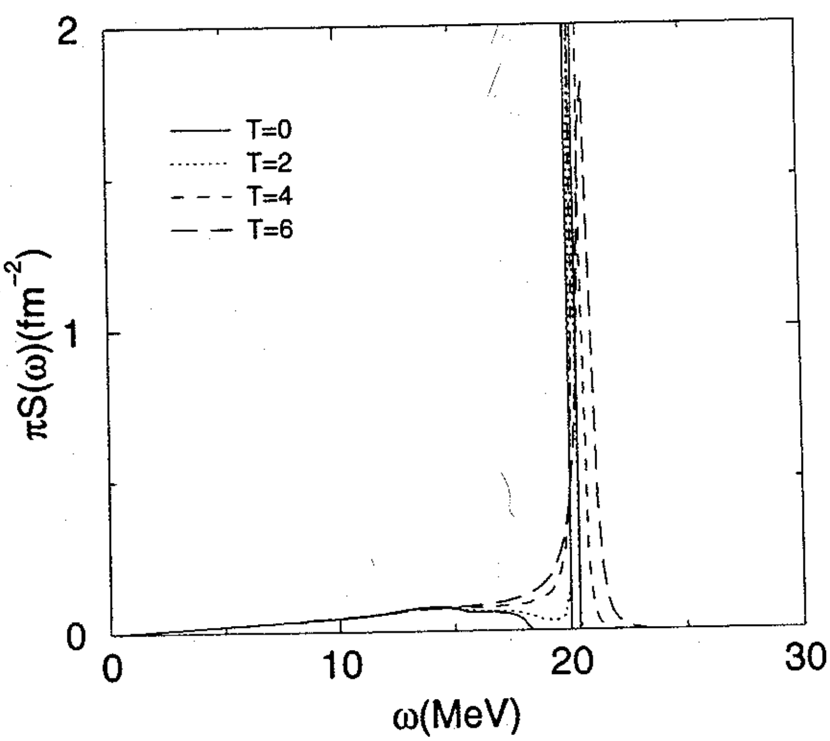

Figure 9

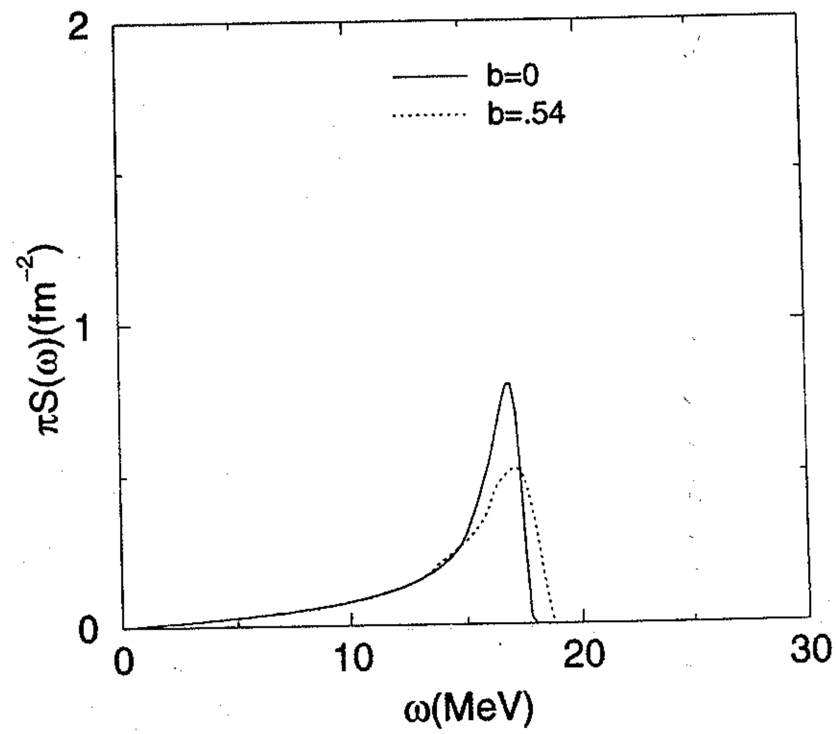

Figure 10

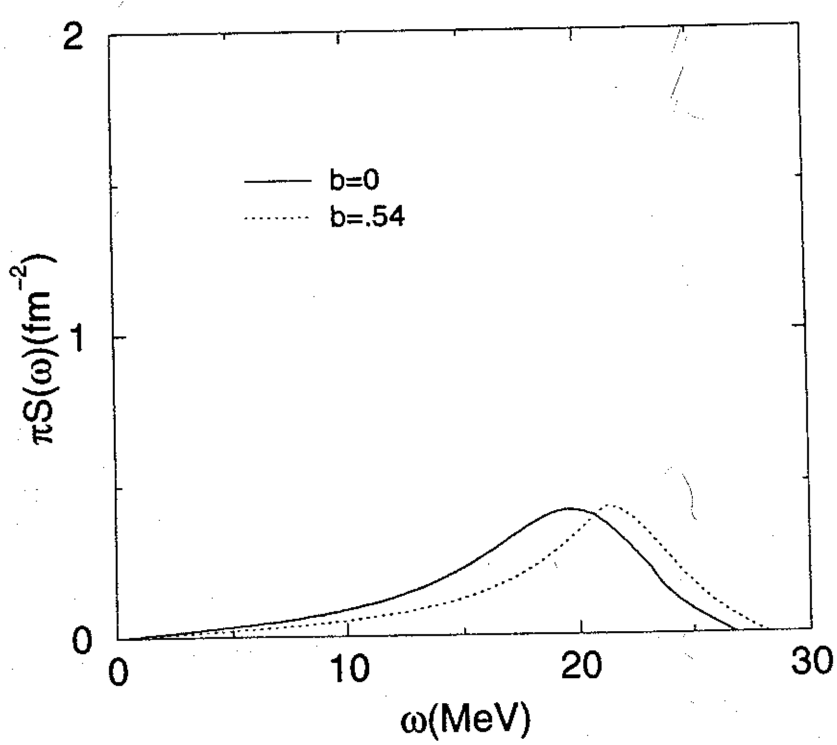

Figure 11

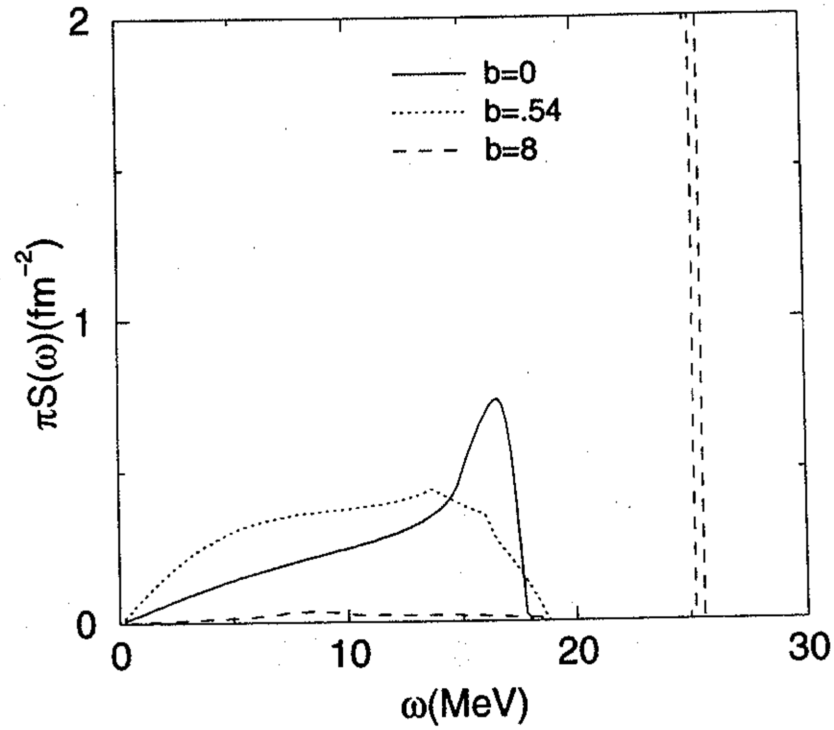

Figure 12

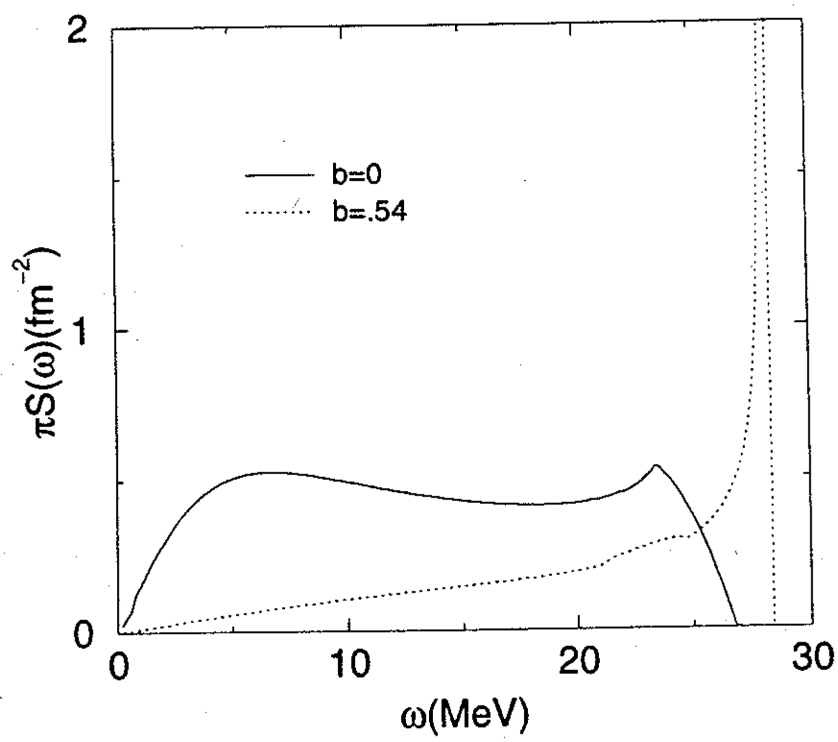

Figure 13

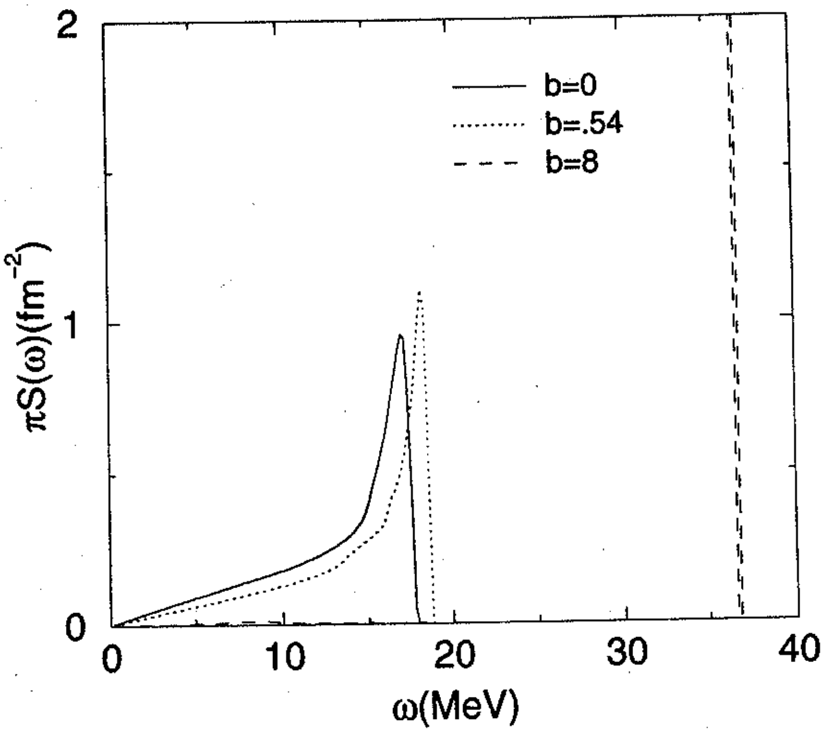

Figure 14

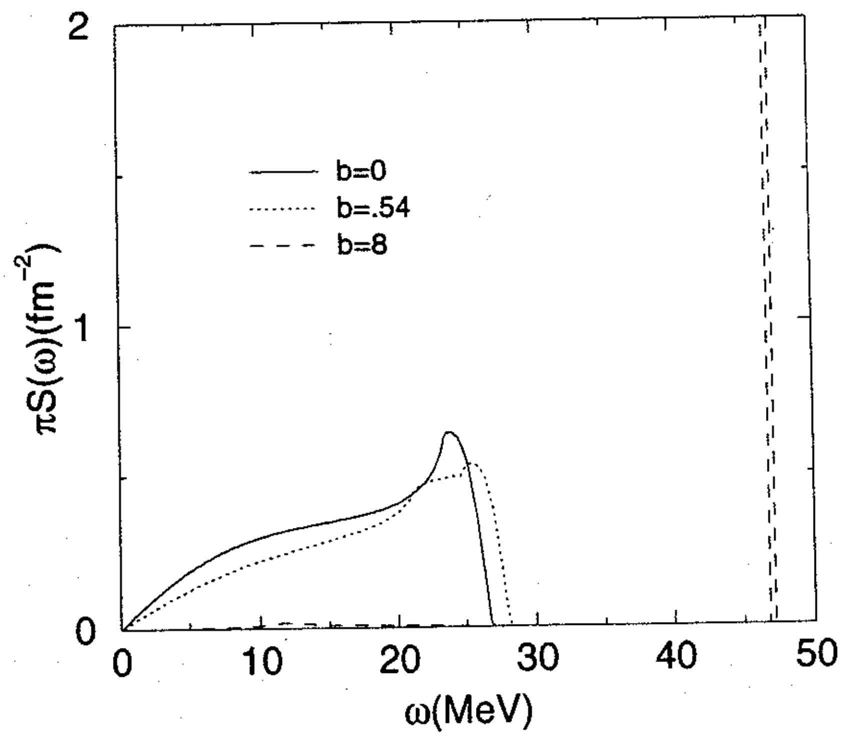

Figure 15