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Non Local Effects Induced by the Particle-Vibration Coupling

O. Civitarese*
Departamento de Fisica, UNLP, La Plata, Argentina

R.P.J. Perazzo*, S.L. Reich, and M. Saraceno
Departamento de Física, Comisión Nacional de Energía Atómica
Buenos Aires, Argentina

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Dynamical corrections to the single particle potential are (perturbatively) evaluated. The corrections to m^*/m resulting from the coupling of independent (effective) fermions to collective vibrations (low energy modes) are analyzed within the NFT formalism. The results for the ^{208}Pb are discussed.

1. Introduction

The motion of nucleons within a nucleus can be described to a reasonable degree of approximation as independent particles moving in an average static potential. The average (Hartree-Fock) potential introduces through its exchange part a non local term in the single particle field. In limiting situations (nuclear matter) this non locality is approximately described by dressing the particles with an effective mass [1, 4].

For finite nuclei, on the other hand, the coupling of individual particles to the collective motion of the system, renormalizes the one body propagators through a highly non local, energy dependent interaction [5, 7]. The corrections to the single particle motion that are thus produced have a dynamical character.

In this work we concentrate on the study of these latter effects. In early calculations [8], it was found that the low energy (empirical) single particle spectrum is systematically compressed as compared to the eigenvalues of a static Wood Saxon well. As related quantities, both the "E-mass" and the "k-mass" have been calculated recently [2, 9, 11]. The relevance of the dynamical effects due to correlated motions, was again brought up recently in connection with the structure of the giant collective modes in finite nuclei [12, 14]. In Ref. 12 an estimate is made of the additional corrections to the m^*/m-ratio that should be introduced as a consequence of the dynamical effects in the RPA description of one phonon states. The "building-blocks" picture of fermionic and bosonic degrees of freedom [15] can be applied, within the framework of (effective) perturbation theories, in order to account for those dynamical effects.

In the present paper, we discuss to which extent the particle-phonon coupling can approximately be described, at low energy, by a redefinition of the static properties of the independent particle motion.

In Sect. 2 we review the formalism that leads to the definition of an effective mass, consistent with the particle-phonon coupling, which has to be regarded as the signature of the above mentioned dynamical effects (thus similar to the one introduced, e.g. by Migdal [16]). In Sect. 3 we illustrate the formalism for the case of a simplified model. The situation prevalent in the region of ^{208}Pb is analized in Sect. 4. Conclusions are drawn in Sect. 5.

2. Review of the Formalism

The Schrödinger equation for a single particle (s.p.) in the many body system involves, in general, a non local potential energy term. To fix ideas we may think that $H_0(\mathbf{r})$ is an average s.p. potential and that $F(r, r', e)$ describes the coupling to collective modes, thus

* Member of the CONICET, Argentina

$$H_0(r)\psi(r)+\int d^3r' F(r,r',\varepsilon)\psi(r')=\varepsilon\psi(r). \quad (1)$$

To account perturbatively for the non local character of F, the integral term of (1) is expanded in terms of the moment operators [17]

$$F^{(s)}_{\alpha,\alpha_2\ldots\alpha_s}(r,\varepsilon)=\frac{1}{S!}\int F(r,r',\varepsilon)\prod_{i=1}^{s}(r-r')_{\alpha_i}d^3r';$$
$$\alpha_i=1,2,3 \quad (2)$$

to get

$$\int F(r,r',\varepsilon)\psi(r')d^3r'=[F^{(0)}(r,\varepsilon)+F^{(1)}_{\alpha_1}(r,\varepsilon)\partial_{\alpha_1}$$
$$+F^{(2)}_{\alpha_1\alpha_2}(r,\varepsilon)\partial_{\alpha_1}\partial_{\alpha_2}+\ldots]\psi(r) \quad (3)$$

(∂_{α_i} denotes a partial derivative operator and a summation over repeated indices is assumed).
The isotropic term originated in the second moment operator reads

$$F^{(2)}(r,\varepsilon)=\int F(r,r',\varepsilon)(r'-r)\cdot(r'-r)d^3r' \quad (4)$$

The r.h.s. of (3) can be written in hermitian form [17], provided that the second and higher order derivatives of all momenta of order $s\geqq 4$ can be neglected, and, under this assumption, (1) can be recast into the form

$$H_0^{\mathrm{eff}}(\mathbf{r})\psi(\mathbf{r})=\varepsilon\psi(\mathbf{r}) \quad (5)$$

where the s.p. (effective) potential $H_0^{\mathrm{eff}}(\mathbf{r})$ includes the corrections to $H_0(\mathbf{r})$ both in the kinetic energy term and in the potential term arising from the moment expansion of F [16, 17], namely
i) The effective kinetic energy term, with the mass-ratio

$$\mu(r,\varepsilon)\equiv\frac{m}{m^*}=\left(1-\frac{m}{\hbar^2}F^{(2)}(r,\varepsilon_l)\right)\cdot\left(1-\left(\frac{\partial F^{(0)}}{\partial\varepsilon}\right)_{\varepsilon=\varepsilon_l}\right)^{-1} \quad (6)$$

and,
ii) The effective one body potential

$$V^{\mathrm{eff}}(r)=\left\{V_0(r)+F^{(0)}(r,\varepsilon)-\varepsilon_l\left(\frac{\partial F^{(0)}}{\partial\varepsilon}\right)_{\varepsilon=\varepsilon_l}\right\}$$
$$\cdot\left\{1-\left(\frac{\partial F^{(0)}}{\partial\varepsilon^i}\right)_{\varepsilon_l}\right\}^{-1} \quad (7)$$

where ε_l denotes the unperturbed s.p. energy that is closer to the eigenvalue ε of (5).
The denominator appearing in (6) can be interpreted as the residue of the Green's function at the one quasiparticle pole (that is the spectroscopic amplitude to populate the s.p. state in a ony body transfer process); referred to as the E-mass [2]. The numerator is essentially related to the non local character of F and is usually referred to as the k-mass [2]. Both contributions (the E-mass and the k-mass) have been studied in detail for the case of infinite media [2, 4].
At this point, a decision has to be made in order to select a proper basis for the expansion of F. We thus recourse to the Nuclear Field Theory (NFT) [18], that is different to the usual framework used for infinite systems (i.e., Brueckner-Hartree-Fock or low density approximations [2]). Within the NFT it is assumed an expansion in powers of the degeneracy that allows to take into account the coupling of s.p. and low energy collective degree of freedom.
In what follows, we illustrate the coupling mechanism between particles and phonons, within the NFT, as a source of the dynamical corrections to the s.p. motion. (For a detailed description of the theory, we refer the reader to the review paper [18]). Within this context, we attempt to evaluate the validity of (6) for the F-kernel which proved to be suitable for the description of the s.p. spectra around ^{208}Pb [6].

3. A Schematic Model

In this section we discuss the dynamical corrections to the effective mass within a schematic model. We consider N fermions moving in a one-dimensional harmonic oscillator well of frequency w_0 that we take as the self consistent one body potential. Each harmonic single particle level can be occupied by up to Ω particles. If we allow for collective oscillations, the lowest order dynamical correction to the one body propagator (Fig. 1 A) corresponds to the process of (Fig. 1 B). The collective modes are taken to be of the particle-hole nature and are evaluated within the RPA. The residual two-body interaction that we assume to be active has the form*

$$V(x,x')=V_0\exp[-(x^2+x'^2)/a^2]. \quad (8)$$

The RPA frequencies W_s that correspond to the "breathing" modes, are the roots of the dispersion relation:

$$\frac{1}{V_0\Omega}=\sum_{n\bar{n}}\frac{|V_{n\bar{n}}|^2\,2\varepsilon_{n\bar{n}}}{\varepsilon^2_{n\bar{n}}-W_s^2}. \quad (9)$$

Throughout, we use the notation $n>n_F$, $\bar{n}\leqq n_F$; n_F labels the last filled shell. The energy $\varepsilon_{n\bar{n}}=\varepsilon_n-\varepsilon_{\bar{n}}=(n$

* The interaction (8) mocks up a residual multipole-multipole interaction with a multipolarity that goes roughly as the inverse of the range a. This can be checked investigating the distribution of the sum rule,

$$S_T=\sum_S B_S=\sum_S W_S|\langle S|e^{-x^2/a^2}|b\rangle|^2.$$

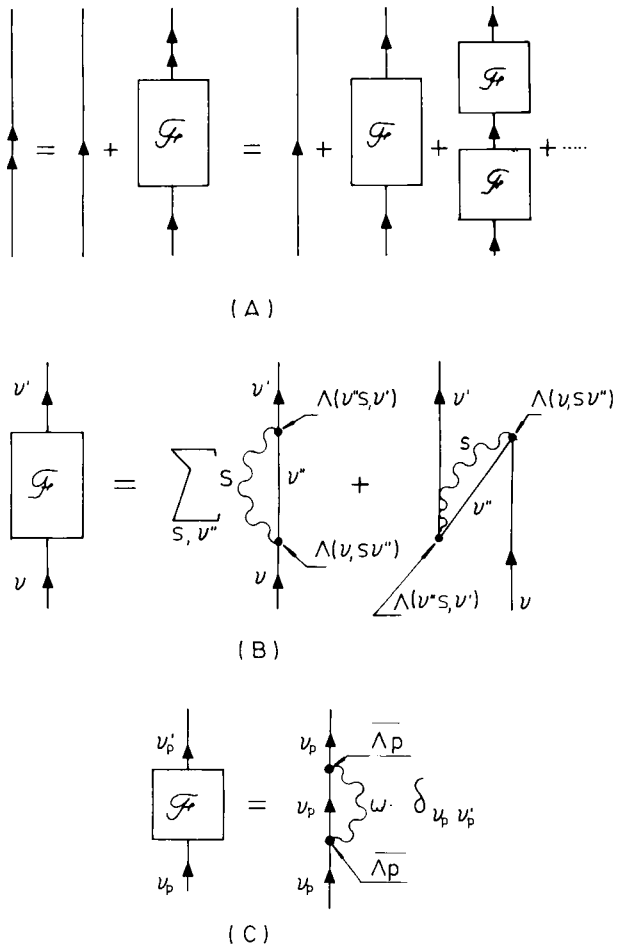

Fig. 1. **A** Diagrammatic representation of the Dyson equation. The box represents any irreducible process involving fermionic and bosonic lines. **B** Representation of the lowest order process renormalizing the free propagator. Particle-hole bosons are labelled by s and single particle states by ν. The vertex functions Λ couple fermionic and collective modes. **C** Particular case of **B** in which only one intermediate state is assumed to be active (see Sect. 3)

$-\bar{n})\hbar w_0$ is the uncorrelated particle-hole energy. The vertex functions Λ_s that couple the single particle and collective degrees of freedom are [6, 15],

$$\frac{1}{\Lambda_s^2}=\sum_{n\bar{n}}\frac{|V_{n\bar{n}}|^2 4\varepsilon_{n\bar{n}} W_s}{(\varepsilon_{n\bar{n}}^2-W_s^2)^2}. \tag{10}$$

Once the vertex functions Λ_s are defined we have all the ingredients to calculate the kernel F in the harmonic oscillator representation:

$$F_{nn'}(\varepsilon)=C_{nn'}\sum_{n'',s}\Lambda_s^2 V_{nn''}V_{n''n'}\cdot\left\{\frac{\theta_{n''}}{\varepsilon-\varepsilon_{n''}-W_s}+\frac{1-\theta_{n''}}{\varepsilon+\varepsilon_{n''}+W_s}\right\} \tag{11}$$

$$F_{\bar{n}\bar{n}'}(\varepsilon)=C_{\bar{n}\bar{n}'}\sum_{n'',s}\Lambda_s^2 V_{\bar{n}n''}V_{n''\bar{n}'}\cdot\left\{\frac{1-\theta_{n''}}{\varepsilon-\varepsilon_{n''}-W_s}-\frac{\theta_{n''}}{\varepsilon-\varepsilon_{n''}+W_s}\right\}$$

with

$$C_{nn'}=\tfrac{1}{2}(1-(-)^{n+n'});\qquad \theta_n=\begin{cases}0 & n>n_F\\ 1 & n\leqq n_F\end{cases}.$$

After expressing F in the coordinate representation, the effective mass ratio $\mu(x,\varepsilon)$ can readily be calculated. In Figs. 2A and 2B we plot $\mu(x,\varepsilon)$ for different values of V_0 and a. The energy dependence of $\mu(x,\varepsilon)$ is given by the energy denominators of $F_{nn'}(\varepsilon)$ (11). This is a smooth function of ε except in the region where $|\varepsilon|\approx|\varepsilon_n+W_s|$. The values of ε considered here, are far from these singularities since we take them close to the (unperturbed) s.p. energy corresponding to the first empty (full) level above (below) the Fermi level. In our calculation we take for ε the value that is obtained from the exact solution of the Dyson equation using the procedure of [6].

As long as the residual interaction is of short range ($a<b$) the function $\mu(r,\varepsilon)$ is seen to have most of its structure inside the harmonic oscillator well and approach unity outside (Fig. 2A). The most important effect of an increase in V_0 is through a change in Λ_s that diverges for values of V_0 where (9) has a root at zero energy. As V_0 approaches that phase transition region, the lowest RPA root W_1 tends to collect a larger fraction of the sum rule. Correspondingly, the oscillations of $\mu(r,\varepsilon)$ become larger. This behaviour is enhanced for larger a. All these results point to a correlation between an increasing importance of non local effects and a concentration of collective features in the RPA modes.

To discuss the spatial structure of the kernel and the fact that is an oscillating function of x let us consider two extreme cases. In the situation *a)* we assume that all particle (hole) vertex functions can be replaced by an average value $\bar{\Lambda}_p(\bar{\Lambda}_n)$ and that the only intermediate state to which the state ν is coupled, is built by the same state ν plus a boson of frequency W (Fig. 1B). In this case all particle (hole) levels are shifted by the same amount $F_{\nu_p}(F_{\nu_h})$:

$$\delta\varepsilon_{p(h)}=\varepsilon_{\nu_{p(h)}}-\varepsilon^0_{\nu_{p(h)}}=F_{\nu_{p(h)}}=\underset{(-)}{\pm}\frac{\bar{\Lambda}^2_{p(h)}}{\varepsilon_{\nu_{p(h)}}-(W+\varepsilon^0_{\nu_{p(h)}})}\approx\underset{(+)}{\mp}\frac{\bar{\Lambda}^2_{p(h)}}{W}. \tag{12}$$

The situation *b)* is quite the opposite. We assume that as a result of the coupling to the collective modes only the first empty and last full levels are affected, being shifted by the same amount but in

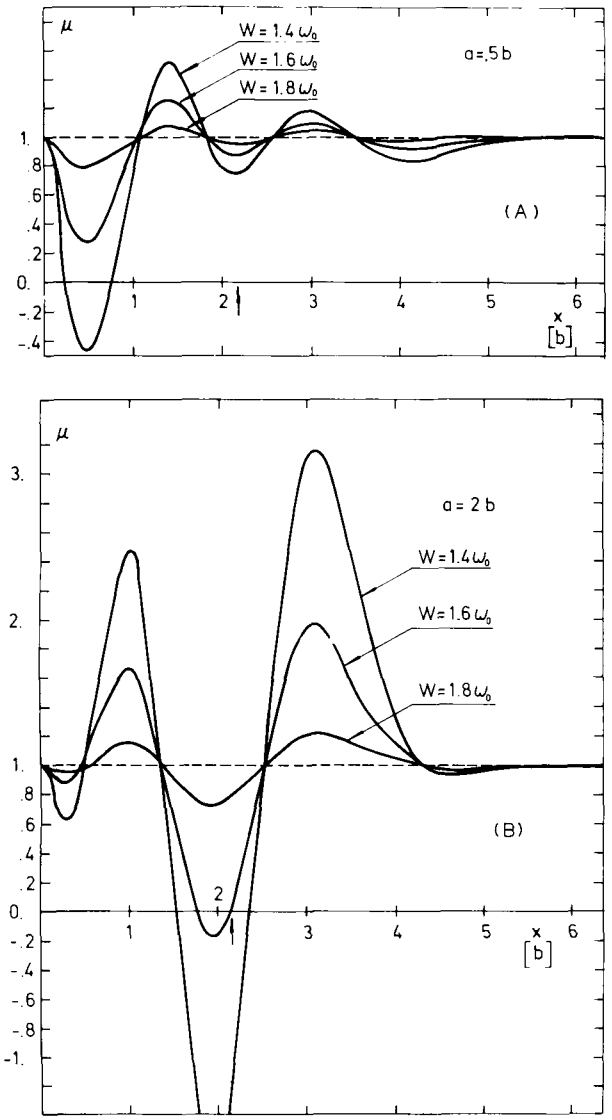

Fig. 2A and B. Effective mass ratio $\mu(x,\varepsilon)$ as a function of x for short **A** and long **B** range forces, the value W that labels each curve is the energy of the lowest *RPA* root (9) in units of the frequency W_0 of the central oscillator well. The calculation involves the boson space of Fig. 1. Fermions are assumed to fill four oscillator shells ($n_F=3$) and are allowed to move up to $n_{\max}=8$. The arrow on the abscissas indicates the "radius" of the system defined by the last occupied orbit

opposite directions

$$\delta\varepsilon_{\nu_F+1}=-\delta\varepsilon_{\nu_F}=F_{\nu_F}. \tag{13}$$

The corresponding kernels in the coordinate representation are:

$$F_a(x,x')=F_{\nu_p}\sum_{\nu'_p>\nu_F}\varphi^*_{\nu'_p}(x)\,\varphi_{\nu'_p}(x')+F_{\nu_h}\sum_{\nu'_h\leqq\nu_F}\psi^*_{\nu'_h}(x)\,\psi_{\nu_h}(x')$$
$$=F_{\nu_p}\,\delta(x-x')-(\delta\varepsilon_{\nu_p}-\delta\varepsilon_{\nu_h})\,\rho(x,x') \tag{14}$$

and

$$F_b(x,x')=F_{\nu_F}(\varphi^*_{\nu_F+1}(x)\,\varphi_{\nu_F+1}(x')-\varphi^*_{\nu_F}(x)\,\varphi_{\nu_F}(x'))$$
$$=F_{\nu_F}(\rho_{ph}(x,x')-\rho(x,x')). \tag{15}$$

In (14) and (15) $\rho(x,x')$ is the one body density matrix and $\rho_{ph}(x,x')$ corresponds to the density distribution obtained by promoting particles from the last filled shell to the first empty level. In (14) we notice an irrelevant local component proportional to $\delta(x-x')$ that shifts all (particle *and* hole) levels by the same amount. The non local term is given by $\rho(x,x')$ times the core polarization $\delta\hbar\omega_0$. In this case the behaviour of μ, follows that of the density, being roughly constant in the interior of the nucleus and approaching unity outside.

In a general situation it is not possible to define the average values $\bar{\Lambda}_p$ and $\bar{\Lambda}_h$ because levels close to the Fermi surface are affected in a different way than those far from it. We approach in this way to the second of the extreme pictures considered above. In (15) the non locality is no longer determined by ρ but rather by its (particle-hole) fluctuations. As a consequence μ presents the oscillatory behaviour that we observe in Fig. 2. This feature is a signature of a complicated type of non local effects that are originated in the coupling to the low energy vibrations of the many body system.

4. A Realistic Situation

We now proceed to the analysis of the realistic situation prevalent in the valence orbitals close to the Fermi surface of the ^{208}Pb core. As fas as non local effects are concerned, we follow the procedure of lumping them in the second moment of the kernel F, performing an average over all valence orbitals. We thus define the expectation value

$$\mu_\nu=\langle\nu|\,\mu(r,\varepsilon)\,|\nu\rangle=\alpha_\nu^2\left(1-\frac{m}{\hbar^2}F^{(2)}_{\nu\nu}(\varepsilon_\nu)\right) \tag{16}$$

(working in the state representations rather than in the coordinate representation).

In (16) α_ν^2 is the residue of the Green's function at the one quasiparticle pole and ε_ν is the bare single particle energy that is closest to it. We first consider the energy dependent effects (E-mass) [9-11] as contained in the residue α_ν^2. In the present calculation we use the solution of the Dyson equation displayed in Fig. 1 A. The residues α_ν^2 can be calculated using the Ward identity that is equivalent to impose that the dressed fermions are eigenstates of the number of particles operator [19].

The collective modes of the ^{208}Pb are described by the RPA diagonalization of a residual multipole isovector and isoscalar particle hole interaction with multipolarity λ ($0<\lambda\leqq 9$). The strength of the isoscalar coupling constants are determined by the self consistency condition of requiring that the average over all nucleons of the single particle multipole field must be equal to the corresponding deformation amplitudes of the central potential [20]. The isovector strength is obtained from optical potential information [21]. This residual interaction provides a fairly good description of the ground state response function [1]. We refer the reader to [6] for further numerical details such as tables of excitation energies and $BE\lambda$ values obtained within this framework.

The single particle states are taken from the experimental evidence. This is certainly not a clean choice since the empirical values have already built in all possible renormalization processes. A choice free of this problem would imply a procedure such as that of [22]. This lengthy calculation is beyond the scope of the present paper and would not affect the main conclusions that are drawn from the present discussions that, as we shall see, are based on average estimates. We restrict our analysis to the neutron single particle space since the behaviour of protons is completely analogous adding no further insight into the problem. Nevertheless proton states are included in the RPA description of collective bosons. We refer the reader to [6] for further technical details.

The general features of the a_ν^2 that are relevant to the discussion can be summarized in the following comments:

a) The values of a_ν^2 throughout the last filled and first empty shell are roughly the same for all states.

b) An average value can be defined as:

$$S_{p(h)}=\frac{\sum_{\nu_{p(h)}}\Omega_\nu a_\nu^2}{\sum_{\nu_{p(h)}}\Omega_\nu} \tag{17}$$

Typical values of $S_{p(h)}$ 0.8 are obtained with a space of intermediate states that include particle-hole and pairing boson with $0\leqq\lambda\leqq 9$ and fermion states belonging to one major shell above and one below the Fermi surface.

c) Some individual values of a_ν^2 for a few states (notably the $j_{15/2}$) are sensible to the relative position of one and three quasiparticle states.

The signature of the coupling with strongly collective modes can be investigated more closely by breaking up the value of $S_{p(h)}$ into the contributions of each multipolarity involved in the calculation (Fig. 3). The low multipole modes, notably the 3^- (2.62 MeV) state, are seen to be the main responsibles for the departure from unity of $S_{p(h)}$.

These values for the residues a_ν^2 are almost the same as those of Ref. 14 (cf. Table 3 Ref. 14 and Fig. 6 Ref. 6 for numerical details) and slightly larger ($\sim 10\%$) than those of [7] in which is used a Landau-Migdal type of residual interaction with finite range corrections. In both these references no special consideration is made of the non local effects arising from the particle vibration coupling.

We now turn to consider these (k-mass) contained in the second moment of the kernel. The matrix elements of the second moment can be written in term of the matrix elements of the kernel F as:

$$F_{\nu\nu}^{(2)}(\varepsilon)=F_{ns_{\frac{1}{2}}}(\varepsilon)\sum_N(I_{n0}^{(0)}I_{N0}^{(2)}+I_{n0}^{(2)}I_{N0}^{(0)})\mathscr{I}_0^{(\nu)}(n,N)$$
$$-\tfrac{2}{3}(F_{np_{\frac{1}{2}}}(\varepsilon)+2F_{np}(\varepsilon))\sum_N I_{n1}^{(1)}I_{N1}^{(1)}\mathscr{I}_1^{(2)}(n,N) \tag{18}$$

where

$$I_{nl_{\frac{1}{2}}}^{(s)}=\int_0^\infty R_{nl}(r)r^{s+2}dr;$$

$$\mathscr{I}_l^{(\nu)}(n,N)=\int_0^\infty r^2|R_{n_\nu l_\nu}|^2R_{nl}R_{Nl}\,dr \tag{19}$$

being $R_{nl}(r)$ the radial wavefunction.

We notice that (18) only involves the $l=0$ and $l=1$ components of F due to the selection rules of the

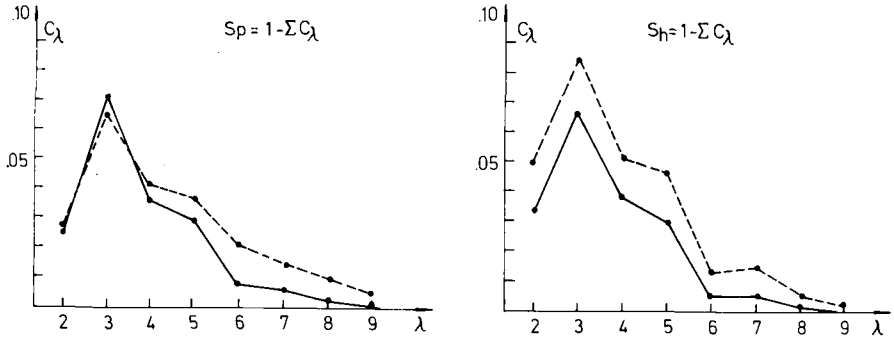

Fig. 3. Departure from unity of the average spectroscopic factor (17) as a function of the multipolarity of the collective (particle-hole) modes included in F. The magnitude $S_p(S_h)$ corresponds to the average over the first (last) empty (occupied) major shell. Full (dashed) lines are results obtained with only isoscalar (isoscalar plus isovector) bosons

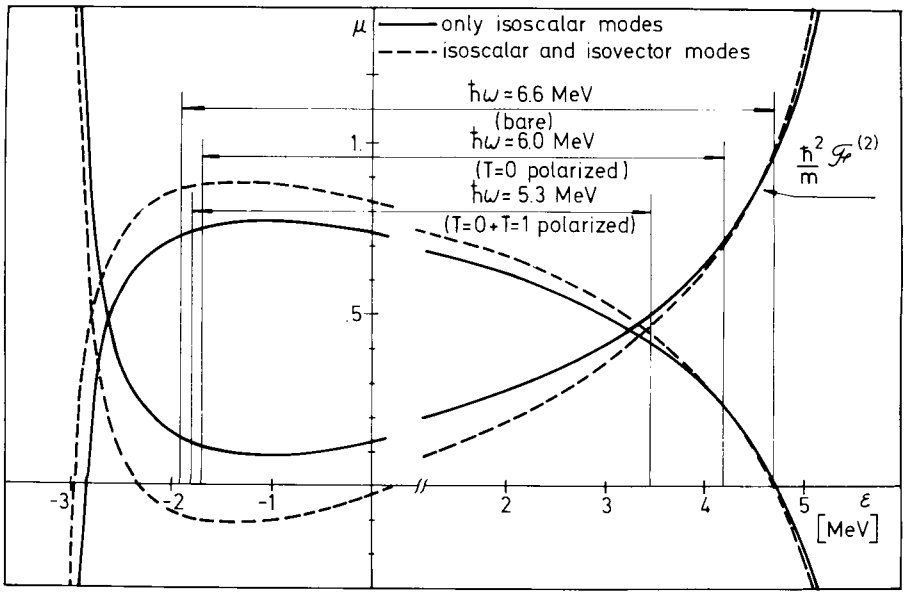

Fig. 4. Averaged second moment $F^{(2)}$ (20) and the corresponding effective mass ratio $\mu(\varepsilon)$ (including the $S_{p(h)}$ factors) as a function of the energy. The distance between the bare and polarized centroids of particle and hole levels are indicated in the upper part of the figure. Full and dashed lines have the same meaning as in Fig. 3

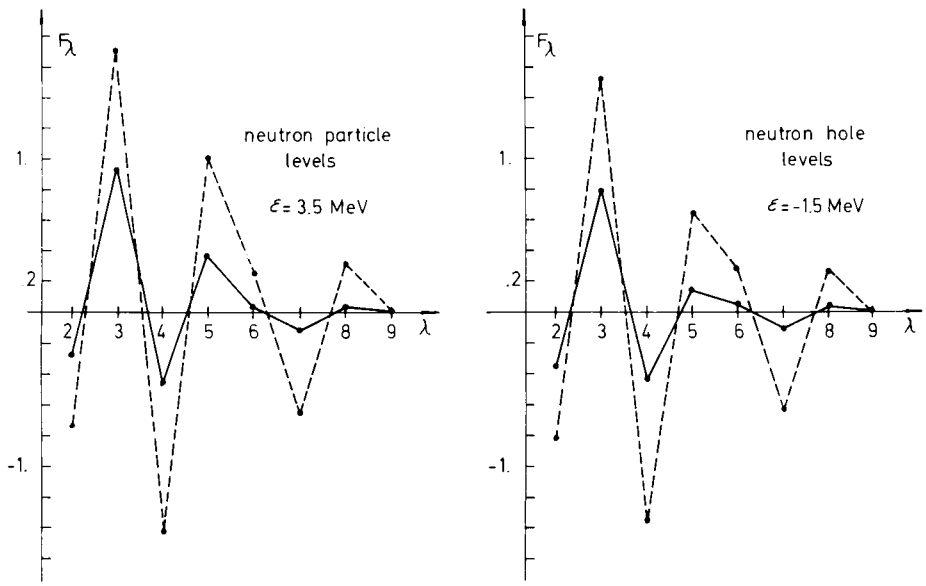

Fig. 5. Contributions to the second moment $F^{(2)}$ (20) for each multipolarity. The energy ε is chosen close to the centroids of (neutron) particle ($\varepsilon \approx 3.5$ MeV) and hole ($\varepsilon \approx -1.5$ MeV). Full and dashed lines have the same meaning as in Fig. 3

radial matrix elements. With these values (18), we thus define the average quantity

$$F^{(2)}(\varepsilon) = \frac{\sum_{\nu} \Omega_{\nu} F^{(2)}_{\nu\nu}(\varepsilon)}{\sum_{\nu} \Omega_{\nu}}. \tag{20}$$

This is to be evaluated at energies of the order of the centroids of the particle and hole s.p. levels. The function $F^{(2)}(\varepsilon)$ is shown in Fig. 4 together with the ratio $\mu(\varepsilon)$ including the factors $S_{p(h)}$.

The values of μ for hole states are seen to behave almost constantly in the energy range between the bare and polarized centroids with $0.70 \leqq \mu_{\text{holes}} \leqq 0.85$ this limits depending on whether the isovector modes are considered or not. The range of variation of $\mu_{\text{particles}}$ is instead larger in the corresponding energy range with values $\mu_{\text{particles}} \lesssim 0.45$. This is in direct relationship with the larger change of the centroid of particles levels as compared with that of holes.

These results of μ are somewhat smaller than those obtained in [23] in which an analytic expression of the kernel F is obtained through an averaging procedure over intermediate particle-boson states.

The non local effects accounted for in (20) can be inspected in more detail fixing the value of ε in the above mentioned regions and plotting separately the contributions of each multipolarity λ (Fig. 5). There

Table 1. Single particle energies of an odd neutron in the presence of the ^{208}Pb core

n j	[MeV]	n j	[MeV]
0 h 9/2	−5.17	1 g 9/2	1.70
1 f 7/2	−4.04	0 i 11/2	2.48
0 i 13/2	−3.33	0 j 13/2	3.13
2 p 3/2	−2.61	2 d 5/2	3.27
1 f 5/2	−2.27	3 s 1/2	3.74
2 p 1/2	−1.70	1 g 7/2	4.20
		2 d 3/2	4.24

are several points that are noteworthy. In the first place, we find the same feature displayed in Fig. 3 in the sense that the largest contributions are due to the lowest values of λ. In this case larger oscillations are found for even and odd values of λ, displaying in this fashion that non local effects can hardly be represented, in low energy, by the single number arising from the average of the expectation values of the second moment $F^{(2)}(\varepsilon)$. We thus recover the features concluded from our schematic model of Sect. 3. In the second place, even and odd multipolarities largely cancel their effects with each other giving a total correction that is small compared with each individual contribution. In the last place we note that the inclusion of the isovector modes have a significant contribution enhancing rather dramatically the pattern obtained with only $T=0$ bosons.

5. Discussion and Conclusions

In the preceeding sections we have analyzed the non local effects and the effective mass induced by the coupling of individual nucleons to collective nuclear surface oscillations.

The values of the spectroscopic factors as signature of energy dependent effects are discussed in the ^{208}Pb case. These are almost independent of the excitation energy of the s.p. levels and may present singular behaviours for individual cases under the form of a fragmentation of the one body transfer strength. This situation arises whenever two single particles energies ε_{ν_1} und ε_{ν_2} and one collective frequency W are tuned to each other in such a way that $\varepsilon_{\nu_1} \cong \varepsilon_{\nu_2} + W$ causing a vanishing energy denominator in F_{ν,ν_2}. From the physical point of view this kind of resonance implies that a nucleon in the level ν_1 has a large amplitude to change to level ν_2 as a consequence of being weakly coupled to the collective vibrational field. This feature is nevertheless taken into account by this treatment; since the kernel is evaluated summing up an infinite series of diagrams. In the Pb region we averaged the particle and hole spectroscopic factors to provide a number that can be associated to the whole of each of the two shells involved in the calculation*. If this is broken up into the contributions of each multipolarity, low values of λ and strong collective modes share the largest contributions.

The non local features of the kernel as measured by the second moment $F^{(2)}$, are studied within the framework of a schematic model and in the realistic situation of ^{208}Pb. In the ^{208}Pb case we define average values for particle and hole s.p. levels. In spite of the fact that we only consider averaged values the non local effects still show up as can be seen in the large oscillations shown in Fig. 5.

We find again that μ is particularly sensitive to the lowest multipoles with large cancellations between the contributions of even and odd λ. This feature is completely similar to that found in [6] studying the polarization of the ^{208}Pb core (see Fig. 5, Ref. 6). This resemblance is of course not casual. From the arguments at the end of section 3 it follows that if we limit our analysis to the average behaviour of particle and hole levels the absolute magnitude (rather than the range) of the non local effects in μ is given precisely by the core polarization $\delta\hbar\omega$. The coupling of single particles and collective modes do give rise to a compression of the single particle spectrum in the neighbourhood of the Fermi surface. As long as this is considered as an average effect it may be thought of as an energy-dependent redefinition of the mass of the valence nucleons. However the energy dependence and non local features show themselves very sensitive to the nature of the low energy nuclear collective motion. This influence may be indeed so important as to invalidate the truncation of the moment expansion to the second order. These features hamper any further refinement in the use of the concept of an effective mass at low energy for finite systems as a way of mocking up the coupling of s.p. and low energy collective vibrations. Thus, for instance a redefinition of the mass of nucleons moving in particular particle-hole configurations that coherently build up a collective mode, should be considered as an extremely crude approximation that perhaps amounts to not more than introducing an ad hoc free parameter. This only acquires physical meaning in the very restricted situations in which no strongly collective modes are present and three quasiparticle states remain uncoupled.

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* Thus ignoring the fragmentation of the $j_{15/2}$ and $i_{13}^{-1}{}_2$ [5-7].

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O. Civitarese
Departamento de Física, UNLP
C.C 67 (1900) La Plata, Argentina

R.P.J. Perazzo
S.L. Reich
M. Saraceno
Departamento de Física,
Comisión Nacional de Energía Atómica
1429 Buenos Aires, Argentina