Fast-phonon induced broadening of single-particle lines in excited nuclear matter

C. O. Dorso
Departamento de Física, Facultad de Ciencias Exactas y Naturales, Universidad de Buenos Aires, 1428 Buenos Aires, Argentina

E. S. Hernández*
Consejo Nacional de Investigaciones Científicas y Técnicas, Argentina
(Received 22 December 1982)

A modified kinetic equation for nuclear matter, whose generator of irreversible motion contains a contribution from coupling to a fast collective mode in addition to the usual kinetic kernel, is examined. The fermionic collision rates are evaluated in the relaxation-time approximation, assuming a finite amount of excitation is transferred from collective motion to the single-particle degrees of freedom in the form of a temperature increase. Both energy and angular dependence of the collision frequencies are discussed.

NUCLEAR STRUCTURE Modified kinetic equation; fermion-boson coupling; intrinsic kinetics; phonon-induced kinetics; total collision rate; resonance decay; finite temperature in nuclear matter; distortion of the Fermi sphere.

I. INTRODUCTION

The interplay between collective and intrinsic excitations in nuclear matter and finite nuclei is a current source of attraction for theoreticians in the field, especially in view of the original dynamical features that one can appreciate.[1-6] Complementarily to earlier studies on transport processes in heavy nuclei, where damping of the collective motion was especially examined,[7] these latest works make room for the consideration of either irreversible evolution. Both single-particle (s.p.) lifetimes and resonance broadenings could be extracted in a framework that treats both kinds of excitations on an equivalent footing.[5,6] In particular, in Ref. 6 (hereafter referred to as I) we have presented a detailed model where a modified kinetic equation for nucleonic states, that contains the averaged coupling to a harmonic collective mode, is derived in parallel to a master equation for the latter. In I we have examined the way in which the mode approaches thermal equilibrium while performing quantal brownian motion, described by the above referred master equation, in a static heat reservoir represented by nuclear matter at a finite temperature T. The stationarity of the fermionic environment is an exact consequence of the sharp-resonance hypothesis adopted to describe the propagation of the coupled system between successive interactions and it makes possible an analytic, exact diagonalization of the generator of irreversible harmonic motion.

In the present work we attempt some estimates of the s.p. broadenings or inverse relaxation times when the sharp-resonance hypothesis is abandoned. Such a study should precede a complete dynamical calculation involving the coupled fermion and harmonic systems, since it provides useful information about the initial and the final rate of s.p. decay and the subsequent distortion of the Fermi sphere in the course of the evolution. It is not our purpose to abound on details regarding the general theory since enough discussion has been devoted to it in I. We rather focus upon the examination of the modified kinetic equation and the computation of the fermion collision frequencies in the so-called relaxation time approximation. The notation is the same as in I and when needed we briefly recall the meaning of each symbol and concept, so as not to distract the reader. The convention to be used throughout this paper regarding density matrices is as follows: The same symbol ρ denotes either the fermion or the boson density, although lower case Latin indices n label oscillator states from zero to an upper bound N while Greek labels α and μ name, respectively, particle levels that participate in phonon creation and phonon annihilation; thus, $\epsilon_\alpha > \epsilon_\mu$. No confusion can ever arise since each meaning becomes clear in context.

In Sec. II we briefly review the appearance of the modified kinetic equation and discuss the balance of competing processes. The collision frequencies are examined in Sec. III from the formal viewpoint when the sharp-resonance approach is abandoned in favor of a broad, energy-dependent form factor that accounts for finite s.p. lifetimes. The generation of this form factor is traced to the fulfillment of the kinetic hypothesis. The numerical calculations and their interpretations compose the body of Sec. IV, while Sec. V contains the final summary.

II. THE MODIFIED KINETIC EQUATION

We have shown in I that the fermions in the heat reservoir surrounding the harmonic mode of interest obey, in the neighborhood of their thermal equilibrium, the linearized kinetic equation

$$\dot{\rho}(1)=[-i\mathscr{L}_{\rm kin}(1)+\mathscr{K}_{\rm kin}(1)+\mathscr{K}_{\rm osc}(1)]\rho(1) . \qquad (2.1)$$

Here, $\rho(1)$ is the reduced particle density for the s.p. labeled 1 and $\mathscr{L}_{\rm kin}(1)$ is the mean-field or time-dependent Hartree-Fock (HF) Liouvillian responsible for the reversible free flow. The label kin denotes that the given symbol is well recognized in kinetic theory. The real collisional kernels $\mathscr{K}_{\rm kin}(1)$, $\mathscr{K}_{\rm osc}(1)$ are, respectively, the collision term owing to the residual fermion-fermion interaction[6] and the frequency-broadening operator that bears the effect of coupling to the macroscopic coordinate. The kinetic contribution has been analytically constructed in Ref. 7 and a numerical study of the approximate eigenvalues has been performed in Ref. 8 for a range of nucleon energies and temperatures of the environment. Our interest here is to focus upon the collision kernel induced by the harmonic oscillation. This term originates in a single fermionic superoperator K_F whose structure is[6]

$$K_F=K_0\rho+K_1\rho_2^{(0)} , \tag{2.2}$$

where $\rho\equiv\rho(1)$, $\rho_2^{(0)}$ is the antisymmetrized two-fermion density, and K_0 and K_1 are one- and two-body superoperators, respectively. In the very close to equilibrium regime, where all density matrices are diagonal, the components of K_F are[6]

$$K_F=\sum_{\alpha\mu} |\lambda_{\alpha\mu}|^2F(\epsilon_{\alpha\mu})\{ |\mu\rangle\langle\mu|[(1-\rho_0)\rho_\mu(1-\rho_\alpha|\alpha\rangle\langle\alpha|)-(1-\rho_N)\rho_\alpha(1-\rho_\mu|\alpha\rangle\langle\alpha|)] \\ +|\alpha\rangle\langle\alpha|[(1-\rho_N)\rho_\alpha(1-\rho_\mu|\mu\rangle\langle\mu|)-(1-\rho_0)\rho_\mu(1-\rho_\alpha|\mu\rangle\langle\mu|)]\} , \tag{2.3}$$

where $F(\epsilon_{\alpha\mu})\equiv F_{\alpha\mu}$ denotes an energy-dependent form factor to be examined later.

We have shown as well in I that $\mathscr{K}_{\rm osc}$ represents the superoperator

$$\mathscr{K}_{\rm osc}(1)=K_0(1)+{\rm Tr}_2K_1(1,2)\rho_2^{(0)}(1,2) . \tag{2.4}$$

The procedure that gives rise to K_F has been indicated in I. This kernel has been constructed using a rather standard technique in nonequilibrium statistical mechanics, namely, a convenient truncation of a Bogoliubov-Born-Green-Kirkwood-Yvon (BBGKY) hierarchy. In the present case, the hierarchy is not equivalent to a Liouville equation for an N-body system that evolves in a reversible fashion with a Hermitian generator of the motion. Instead, what one constructs in the present case is a modified BBGKY hierarchy, starting from an equation of motion for the many fermion system that already contains, in its non-Hermitian generator of motion, the boson-averaged interaction with the collective mode. In such a frame, K_0 and K_1 are non-Hermitian superoperators that cannot be factorized as commutators; however, it is worthwhile to notice that $\mathscr{K}_{\rm osc}$ in Eq. (2.4) possesses the same structure as the mean-field, kinetic Liouvillian,[6]

$$\mathscr{L}_{\rm kin}(1)=[\mathscr{H}_{\rm HF}(1),\ ] \\ =[H_0(1)+{\rm Tr}_2H_1(1,2)\rho_2^{(0)}(1,2),\ ] . \tag{2.5}$$

We recall as well that in I, the boson-induced collisional terms have been given a diagrammatic interpretation that we will not reproduce here.

Now, using expression (2.3) we obtain

$$\mathscr{K}_{\rm osc}(1)\rho(1)=\sum_{\alpha\mu} |\lambda_{\alpha\mu}|^2F(\epsilon_{\alpha\mu})\{ |\mu\rangle\langle\mu|[(1-\rho_0)\rho_\mu(1-\rho_\alpha)-(1-\rho_N)\rho_\alpha(1-\rho_\mu)] \\ +|\alpha\rangle\langle\alpha|[(1-\rho_N)\rho_\alpha(1-\rho_\mu)-(1-\rho_0)\rho_\mu(1-\rho_\alpha)]\} . \tag{2.6}$$

With these considerations, the kinetic equation (2.1) can be written as a gain-minus-loss evolution law for the matrix elements of $\rho(1)$,

$$\dot\rho_A=\dot\rho_A^{\rm kin}+|\lambda_{A,A-q}|^2F_{A,A-q}[(1-\rho_0)\rho_{A-q}(1-\rho_A)-(1-\rho_N)\rho_A(1-\rho_{A-q})] \\ +|\lambda_{A+q,A}|^2F_{A+q,A}[(1-\rho_N)\rho_{A+q}(1-\rho_A)-(1-\rho_0)\rho_A(1-\rho_{A+q})] , \tag{2.7}$$

where $A\pm q$ is a shorthand notation to indicate the participation of the s.p. level with momentum $k_A\pm q$, the latter being the phonon momentum. It is clear that ${\rm Tr}\rho(1)=cte$ and this equation can be understood as follows: Apart from the level decay associated with the residual two-body interaction, indicated by $\dot\rho_{\rm kin}$ (Boltzmann collisional derivative; see, for example, Refs. 7 and 8), phonon creation and annhilation events drive the irreversible evolution according to the usual gain-minus-loss pattern. The s.p. level denoted as $|A\rangle$ increases its population from states with momentum k_A+q (k_A-q) owing to creation (destruction) of phonons, the latter processes possessing an intrinsic weight proportional to the probability of nonexistence (existence) of quanta, namely $1-\rho_N$ $(1-\rho_0)$. We are using the notation of I, where N denotes the upper bound of the oscillator spectrum, 0 its ground state, and ρ_n the population of the state with n quanta, $n=0,\ldots,N$. On the other hand, we read in (2.6) that the s.p. level $|A\rangle$ depopulates towards those with momentum k_A+q (k_A-q) through phonon creation (destruction) with probability $1-\rho_0$ $(1-\rho_N)$. Equilibrium will be reached when populating and depopulating events take place at the same rate.

III. ESTIMATE OF THE KINETIC-PLUS-NONKINETIC COLLISION RATE

In this work we evaluate the modified collision frequencies in the spirit of Ref. 8. In that work, the kinetic collision frequencies for fermions in nuclear matter interacting through a residual force have been calculated as functions of s.p. energy for a range of equilibrium temperatures and a variety of realizations of the two-body interaction. The philosophy of these types of calculations consists of approximating the eigenvalue problem for the collision operator (that gives rise to the characteristic collision frequencies of the system) as

$$\mathscr{K}_{\rm kin}[\rho(1)]\approx -L_{\rm kin}\rho(1)\ , \tag{3.1}$$

where $L_{\rm kin}$ is the linearized operator for the loss processes, namely, those that depopulate the s.p. states. Each diagonal element of the matrix on the right-hand side of (3.1) is the so-called collision frequency, or inverse relaxation time, of the system for the given s.p. energy. This procedure provides a means of extracting orders of magnitude, qualitative functional dependence, and parametric trends of the eigenvalues of the collision operator.

Thus, in this work we select the loss amplitude in Eq. (2.7) and define the phonon-induced collision rate,

$$\begin{aligned}\nu_A = {} & |\lambda_{A,A-q}|^2 F_{A,A-q}(1-\rho_N)(1-\rho_{A-q}) \\ & + |\lambda_{A+q,A}|^2 F_{A+q,A}(1-\rho_0)(1-\rho_{A+q})\ .\end{aligned} \tag{3.2}$$

This quantity should provide an estimate of the s.p. width originated in the coupling to the harmonic mode, when evaluated within a model that yields reasonable figures for the fermion and boson densities. The approach selected here consists of fixing $\rho_{A\pm q}$ as a Fermi factor at a given temperature T and looking for the values of ρ_0 and ρ_N in a thermal canonical distribution for an oscillator with energy $\hbar\Omega$. In order to better fix the ideas, we will always have in mind a system that, at the origin of time, consists of cold ($T_0=0$), equilibrated nuclear matter and an excited mode, $\rho_n=\delta_{n,n_0}$, with $0<n_0<N$. We assume that the system evolves toward its overall equilibration, and at the end of time, the fermionic environment has reached a temperature $T_\infty=\sqrt{\hbar\Omega/a}$. Here a is the level density parameter in the Fermi gas model. Thus, $\rho_n=\exp(-\hbar\Omega/T_\infty)$ is approximately equal to $\delta_{n,0}$ for frequencies representing nuclear giant resonances (see I). We then realize that

$$\nu_A(\infty)\approx |\lambda_{A,A-q}|^2(1-\rho_{A-q})F_{A,A-q}\ . \tag{3.3}$$

This means that the major contribution to the thermally equilibrated line width arises from events in which the particle with momentum k_A "decays towards" a state with momentum k_{A-q} creating a phonon, this being the only permitted process in a phononless situation, i.e., $\rho_0\approx 1$.

The selection of coupling matrix elements will be discussed in Sec. IV. A word of caution should be devoted to the energy form factor $F_{A\pm q,A}$ in Eq. (3.2). We recall that in I, this factor was given by

$$F_{A\pm q,A}=\pi\delta(\hbar\Omega-|\epsilon_{A\pm q}-\epsilon_A|)\ . \tag{3.4}$$

The origin of the δ function can be traced to the time integral of the propagator between two successive collisions between fermions and the collective mode. Strictly speaking, we should remind the reader that, in the formal description of irreversible dynamics for weakly coupled systems, one writes that the joint probability of a fermion and a boson configuration evolves in time as

$$\begin{aligned} i\frac{\partial}{\partial t}(\rho_B\rho_F)_t = {} & (\mathscr{L}_B+\mathscr{L}_F)_t(\rho_B\rho_F)_t \\ & -i\int_0^t d\tau L_{BF}U_c(\tau)L_{BF}(\rho_B\rho_F)_{t-\tau}\ ,\end{aligned} \tag{3.5}$$

where U_c is the so-called correlation, or intermediate, propagator. The time integral is extended up to infinity if it can be assumed that the lifetime of a correlation, τ_c, is a microscopic time scale related to the duration of a single collision and much shorter than the macroscopic observation time, t. One then writes the collision kernel,

$$\mathscr{K}_{BF}(\rho_B\rho_F)=i\int_0^\infty d\tau L_{BF}U_c(\tau)L_{BF}(\rho_B\rho_F)_{t-\tau}\ . \tag{3.6}$$

However, for practical calculations it is usual to replace the correlated (and complicated) propagator $U_c(\tau)$ by its unperturbed counterpart U_0. In other words,

$$U_c(\tau)\to U_0(\tau)=\exp(-i\mathscr{L}_B\tau)\exp(-i\mathscr{L}_F\tau)\ . \tag{3.7}$$

If the interaction L_{BF} is time independent, integration of (3.6) under the assumption (3.7) gives rise to a δ function in place of the form factor F. Now, it is worthwhile noticing that $U_0(\tau)$ in (3.7) is an infinitely long-lived kernel; thus, its introduction in (3.6) violates the previous hypothesis regarding the separation between time scales. This inconsistency is avoided if one selects a short-lived kernel as an approximation for $U_c(\tau)$, i.e.,

$$U_c(\tau)=\exp[-i(\mathscr{L}_B+\mathscr{L}_F)\tau]\exp(-\gamma\tau)\ , \tag{3.8}$$

where γ stands as a large frequency parameter whose inverse is a microscopic time τ_c. Integration of (3.6) with (3.8) gives rise to a Lorentzian form factor,

$$\begin{aligned}F(\Delta\epsilon) &= \mathrm{Re}\int_0^\infty d\tau\, U_c(\tau) \\ &= \frac{\hbar\gamma}{(\hbar\Omega-\Delta\epsilon)^2+(\hbar\gamma)^2}\ .\end{aligned} \tag{3.9}$$

We remark again that the validity of the kinetic assumption, and consequently the legitimacy of the equations of motion treated in this work, demands a *broad* Lorentzian form factor ($\gamma\gg\Omega$, $\Delta\epsilon/\hbar$).

On the other hand, it is necessary to observe that the use of approximation (3.7) gives rise to a peculiar behavior of the fermionic environment, however, making possible the calculations in I. Indeed, one can see in (3.2) that if the form factor is a δ kernel, the collision frequency vanishes except for those s.p. levels with momentum lying on either of two preferential planes, whose equation can be found from the simultaneous energy and momentum conservation. This has been discussed in I, where advantage has been taken of the fact that those fermion orbitals participating in collisions with the macroscopic object relax instantaneously. Thus, the heat reservoir can be considered to be in a steady state from the origin of time. Contrary to this extreme situation, the correct treatment of the kinetic assumption makes room for finite collision frequencies whose systematic analysis we pursue in the next section.

IV. NUMERICAL CALCULATIONS AND DISCUSSION

The computation of Eq. (3.2) demands the knowledge of an order of magnitude for the coupling constants $\lambda_{\alpha\mu}$ and the microscopic width γ in the form factor given in Eq. (3.9). The former are selected as independent of α,μ, and the value is taken from Broglia[10] as an estimate for isovector giant modes,

$$|\lambda| \approx \frac{130\pi \text{ MeV}}{A\langle r^2\rangle}b \ , \tag{4.1}$$

where b is proportional to a reduced electromagnetic transition probability.[11] This expression yields $|\lambda|$ between 2 and 5 MeV for nuclides with mass A larger than 100. We have done our calculations with $|\lambda| = \sqrt{10}$ MeV.

The parameter γ possesses the dimensions of a frequency, and we assume its order of magnitude is

$$\gamma = \frac{v_F}{r_c} \tag{4.2}$$

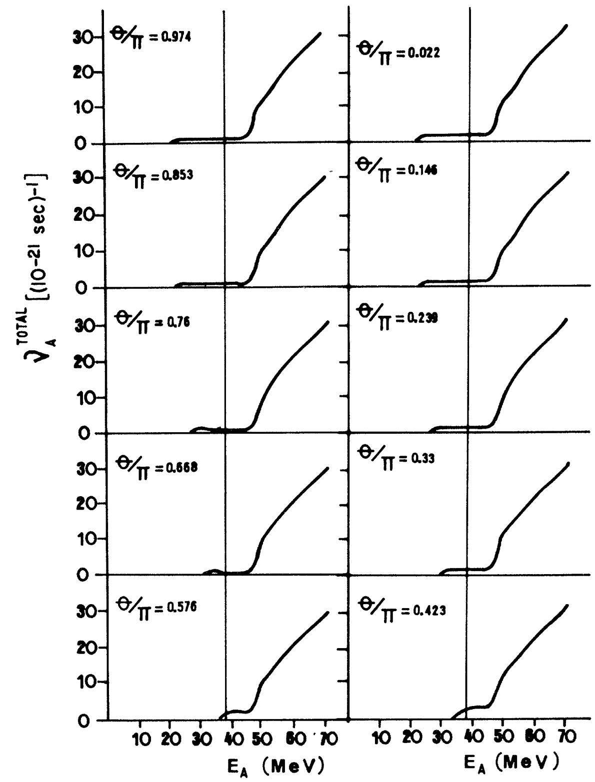

FIG. 2. The same as in Fig. 1, for a temperature $T = T_\infty = 1.3$ MeV.

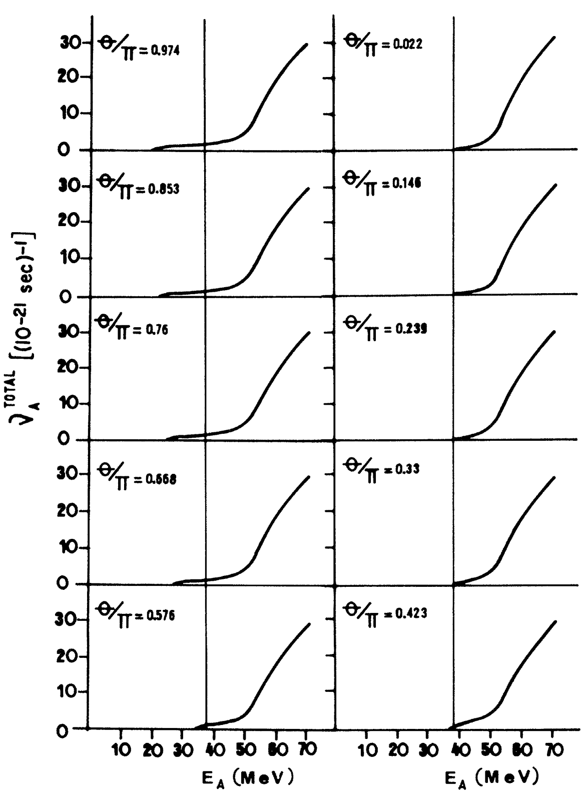

FIG. 1. The total collision frequency ν_A^{total} as a function of the single-particle energy ϵ_A for different angles calculated for a temperature $T = T_0 = 0$. The vertical line indicates the Fermi energy.

where v_F is the Fermi velocity, actually an average velocity for the fermions, and r_c is a correlation radius giving the range of the boson-fermion interaction. If we further suppose $r_c \approx r_0 = 1.12$ fm, and take into account that the Fermi momentum in nuclear matter is $k_F \approx 3/2r_0$, we get

$$\hbar\gamma \approx 3\frac{\hbar^2}{2mr_0^2} \ . \tag{4.3}$$

Then, $\hbar\gamma$ is three times the translational quantum whose value is 16.3 MeV. The latter is a typical resonant energy; we then obtain, as a result of this approximation, a Lorentzian width that permits the application of the kinetic hypothesis, since the lifetime γ^{-1} of the correlated propagator $U_c(\tau)$ is about one third of a characteristic oscillation period for the range of resonance frequencies under consideration.

We have chosen the initial conditions to correspond to a one-phonon excited mode ($n_0 = 1$). The asymptotic equilibrium is selected at a temperature $T = 1.3$ MeV, this being the one to be reached in ^{208}Pb after absorption of an $\hbar\Omega = 13.8$ MeV phonon. The level density parameter a has been taken from Ref. 10. We have simulated the influence of low-energy collective modes upon particle dynamics by renormalizing the fermion inertial parameter as an energy-dependent, effective mass whose values have been extracted from Ref. 11. Our results are displayed in Figs. 1–6.

In Figs. 1 and 2 we show the total frequencies $\nu_A^{\text{total}} = \nu_A + \nu_A^{\text{kin}}$ as functions of the s.p. energy ϵ_A for dif-

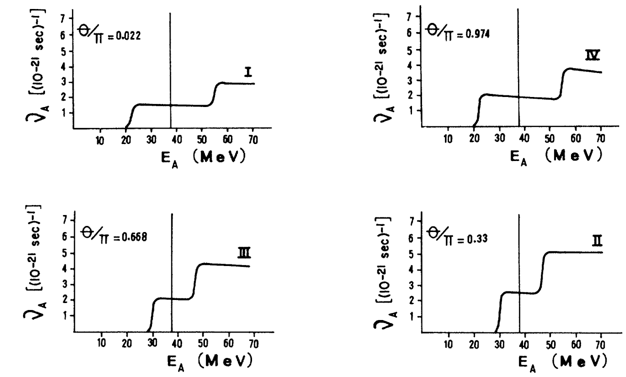

FIG. 3. The fermion-boson collisional frequency ν_A as a function of energy for a selected set of four angles. The temperature is zero and the correlation radius r_c has been chosen equal to $r_0 = 1.12$ fm.

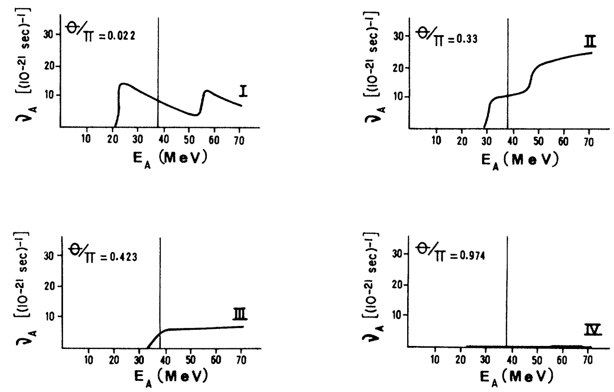

FIG. 5. The same as in Fig. 3, with a correlation radius $r_c = 10r_0$.

ferent angles, computed at $T_0 = 0$ (Fig. 1) and $T_\infty = 1.3$ MeV (Fig. 2). The kinetic frequencies $\nu_A^{\rm kin}$ have been calculated as indicated in Refs. 8 and 12. The fermion-boson collisional frequencies ν_A are shown in Figs. 3 and 4 for the same temperatures. The effect of the form-factor width is illustrated in Figs. 5 and 6, where the same calculations are presented, corresponding to a correlation radius $r_c = 10r_0$.

Let us first examine Fig. 1. We should remind the reader that at $T = 0$, the kinetic collision frequency vanishes for $\epsilon_A < \epsilon_F$ (see Refs. 8 and 12). In addition, the coupling frequency ν_A, whose expression is given in Eq. (3.2), consists of two step functions, $1 - \rho_{A-q}$ and $1 - \rho_{A+q}$, modulated by the Lorentzian form factors. Thus, ν_A deforms the kinetic collision frequency in two different fashions: (i) it provides a nonvanishing total rate below the Fermi level, and (ii) it changes the otherwise smooth slope of the increasing branch in the figures, insinuating a broad shoulder several MeV above the Fermi level. It is seen that as the angle increases from zero to $\pi/2$, the low energy plateau is displaced to the right (see as well Fig. 3), this tendency being reversed for back angles approaching π. The overall displacement is as large as 20 MeV, and we observe that for angles close to $\pi/2$, the collision rate is almost negligible below the Fermi energy. This means that the small forward momenta are more likely to participate in an event concerning the mode than those small momenta being almost perpendicular to the phonon momentum q. The same statement holds for the small backward momenta; this symmetry effect will be examined later in connection with Fig. 3.

In Fig. 2 we observe a situation corresponding to the temperature T_∞ where the ground-state population ρ_0 is almost unity. This means the frequency ν_A is given by Eq. (3.3) and only one Fermi factor, $1 - \rho_{A-q}$, is visible. In this case, s.p. levels belonging to the Fermi sea participate in collisions only to the extent permitted by the range of the Lorentzian form factor, their contribution to the total rate being negligible. We see that the shoulder in Fig. 1 is here substituted by a smooth slope change, owing to the diffuseness of the Fermi surface at T_∞. The angular pattern displays an asymmetrical displacement of the lower bound of the curves towards smaller energies with increasing angle, opposite to the symmetry exhibited in Fig. 1.

In Figs. 3–6 we can appreciate more fully the effect of the energy-nonconserving kernel F. In Fig. 3, the form-factor width corresponds to a correlation length $r_c = r_0$. We clearly see the two step functions and their drift to-

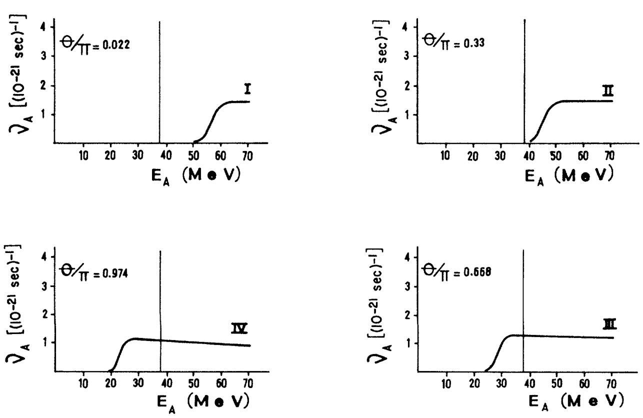

FIG. 4. The same as in Fig. 3, for $T = 1.3$ MeV.

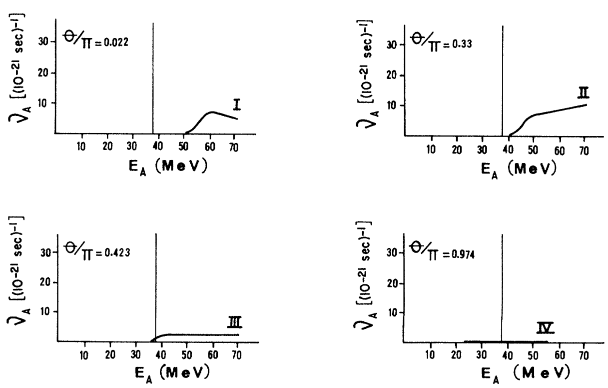

FIG. 6. The same as in Fig. 4, with a correlation radius $r_c = 10r_0$.

wards each other when θ increases from zero to $\pi/2$. Indeed, these two steps have almost collapsed for $\theta=0.42\pi$, and we realize, from Eq. (3.2), that $\theta=\pi/2$ is a transition point where $1-\rho_{A-q}$ and $1-\rho_{A+q}$ become interchanged. Since we are assuming an initial excitation of the oscillator where both ρ_0 and ρ_N are zero, we see that for a given momentum, ν_A would be a symmetric function of the angle with respect to $\pi/2$, except for the form factors F. This fact indicates to us that all differences between the θ and the $(\pi-\theta)$ patterns in Fig. 1 must be related to the Lorentzian kernels. In other words, at $T=0$, the rate at which a given momentum k_A lying in the plane perpendicular to q "decays towards" k_A+q, differs from the decay rate towards the momentum k_A-q in an amount corresponding to the respective Lorentzian weights.

Figure 4 displays the only existing step—considerably smoothed by the finite excitation of the environment—in the final condition, $\rho_0\approx 1$. We observe the leftwards drift with increasing angles and the presence of the Lorentz form factors clearly exhibited in the steady negative slope of the plateau.

Figures 5 and 6 were devised to illustrate to a better extent the effect of these form factors. The situation shown here lies closer to the limiting case[6] in which it merely consists of a δ kernel, since we have increased the correlation length by a factor of 10. We observe that the peaks in the collision rates are rather sharp for small angles (they tend to disappear as θ increases), giving rise to positive slopes related to peaks located at energies beyond the scale considered here. Furthermore, back angles possess very small collision frequencies, whose negative slopes in the region are associated with Lorentzian maxima lying far on the left of the plot.

V. SUMMARY

In the present work we have illustrated the consequences, upon s.p. lifetimes or level widths, of the presence of an excited harmonic mode in nuclear matter. The coupling between fermions and phonons possesses a twofold effect; it damps away the collective motion in a manner that has been described in detail in I in the frame of a solvable model, and it broadens the s.p. levels in their mean field. Rather than solving the modified kinetic equation for the fermions simultaneously with the master equation for the mode, we have decided to perform a numerical analysis of the expected trends of the kinetic-plus-nonkinetic collision rates, in two different situations. The morphology of these examples schematically corresponds to (i) the initial condition of a dynamical process of resonance decay in cold nuclear matter, and (ii) the final or asymptotic state where the collective oscillation has released its energy, in the manner of a finite temperature T_∞, to the fermionic environment.

We have computed the collision frequencies in the usual approach, namely, the so-called relaxation time approximation, as functions of s.p. energy and angle between fermion and phonon momenta. It has been seen that the presence of a collective mode and a coupling device makes room for the participation, in the fermionic dynamics, of s.p. states lying rather deep in the Fermi sea. Such a participation would be forbidden in cold nuclear matter subject to a two-body, nucleon-nucleon effective interaction as the only source of relaxation, and highly inhibited even in the case of a finite temperature of the nuclear environment.[8,12] This effect is enhanced if a broad form factor is allowed to account for short-lived collisions between the fermions and the oscillation. In this sense, one could regard the coupling as responsible for an increase in the effective diffuseness of the Fermi surface, in addition to a distortion of the Fermi sphere in view of the angular dependence of the calculated figures.

The two examples discussed here provide estimates of the characteristic decay times of s.p. dynamics at the beginning and at the end, respectively, of the irreversible evolution of the overall system. A full numerical integration of the simultaneous equations of motion may give further insight into matters like the deformation of the Fermi sphere following the decay of a collective excitation. Calculations in this respect are currently being pursued and will be presented for publication.

ACKNOWLEDGMENTS

The authors are indebted to the Laboratorio de Física del Plasma, Universidad de Buenos Aires, for unlimited access to their computing facilities. One of us (E.S.H.) is grateful to the Departamento de Física, Universidad de Buenos Aires, for hospitality extended her as a nonstaff member. This work was performed under Grants No. 9413a/81 from CONICET, Argentina, and 10030106-019 from Subsecretaría de Ciencia y Tecnología of Argentina.

*Mailing address: Departamento de Física, Facultad de Ciencias Exactas y Naturales, Universidad de Buenos Aires, 1428 Buenos Aires, Argentina.

[1] S. Mukamel, U. Smilansky, D. H. E. Gross, K. Mohring, and M. J. Sobel, Nucl. Phys. A366, 339 (1981).
[2] N. Takigawa, K. Niita, Y. Okuhara, and S. Yoshida, Nucl. Phys. A371, 130 (1981).
[3] K. Niita and N. Takigawa, Nucl. Phys. A397, 141 (1983).
[4] D. H. E. Gross, K. Mohring, S. Mukamel, U. Smilansky, and M. J. Sobel, Nucl. Phys. A378, 375 (1982).
[5] J. Wambach, V. K. Mishra, and Li Chu-Hsia, Nucl. Phys. A380, 285 (1982).
[6] E. S. Hernández and C. O. Dorso Phys. Rev. C (in press).
[7] W. Nörenberg and H. A. Weidenmüller, *Lecture Notes in Physics, Vol. 5* (Springer, Berlin, 1976); L. G. Moretto and R. P. Schmitt, Rep. Prog. Phys. 44, 533 (1981).
[8] C. O. Dorso and E. S. Hernández, Phys. Rev. C 26, 528 (1982).
[9] E. S. Hernández and H. G. Solari, Nucl. Phys. A397, 115

(1983).
[10]R. Broglia, in *Nuclear Theory 1981, Proceedings of the Santa Barbara Conference,* edited by G. Z. Bertsch (World Scientific, Singapore, 1982), p. 93.
[11]A. Bohr and B. Mottelson, *Nuclear Structure* (Benjamin, New York, 1969), Vol. I.
[12]J. P. Blaizot and B. L. Friman, Nucl. Phys. A372, 69 (1981).
[13]M. T. Collins and J. J. Griffin, Nucl. Phys. A348, 63 (1980).