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IMAGINARY POTENTIALS FROM MANY-BODY THEORY

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Abstract: The imaginary parts of the single-nucleon, particle-hole (phonon) and particle-particle (deuteron) mass operators are considered. Expressions of a Fermi golden-rule form are derived for each case. Special emphasis is laid on many-body aspects and collectivity in the transition amplitudes. Implications for the calculation of optical potential for nuclei are discussed.

1. Introduction

Nuclear optical potentials are of continued theoretical and experimental interest. On the theoretical side, the calculation of in-medium cross-sections and the inclusion of collective phonon excitations is of prime interest. Though there exist 'standard' many-body techniques for calculation of optical potentials (mass operators) [1-5]), there have been recently new developments reported [6]). In this work, we continue along these lines, and derive new formulas for the optical potential of elastic nucleon and deuteron scattering. We treat the particle-hole (phonon) self-energy in a consistent manner. We concentrate on the formulation of the imaginary parts of the optical potentials. The corresponding real parts are in principle determined via dispersion relations. It should always be possible to write the diagonal elements of the imaginary potentials in a golden-rule form; their definite sign is then ensured. Expressions that cannot be arranged in such a form had been used when attempting to include the phonon excitations. Due to antisymmetrization, overcounting problems occur which are partially cured by subtraction procedures [1,2,7,4]). As a consequence the imaginary potential can at least in principle acquire a wrong sign (flux is produced rather than absorbed). In this work we show that the collectivity may be put into the vertex functions without loosing generality. The proper sign of the imaginary potential is then guaranteed, as the diagonal elements are indeed of the golden-rule form.

In sect. 2 we first treat the nucleon-nucleus potential in quite detail. In sect. 3 we generalize the results to the particle-particle and particle-hole self-energies. Finally in sect. 4 we confront our new formalism with commonly used procedures and give a critical assessment.

2. Single-Particle Potential

The nucleon-nucleus optical potential can be identified with the mass operator appearing in the Dyson equation for the single-particle Green's function

$$G = G^0 + G^0 MG, \tag{2.1}$$

where

$$G_{11'}^{t-t'} = -i\langle 0|T\{\psi_1(t)\psi_{1'}^\dagger(t')\}|0\rangle, \tag{2.2}$$

and G^0 is the Hartree-Fock propagator.

Being concerned with the imaginary part of the potential, we write [8,9]

$$\begin{aligned} \mathrm{Im}\, M_{11'}^{t-t'} &= \mathrm{Im}\, -i\langle 0|T\{j_1(t)j_{1'}^\dagger(t')\}|0\rangle_{irr} \\ &= -\frac{1}{2}\langle 0|[j_1(t),j_{1'}^\dagger(t')]|0\rangle_{irr}, \end{aligned} \tag{2.3}$$

with

$$j_1 = [\psi_1,\hat{V}]. \tag{2.4}$$

Explicitly the mass operator is

$$M_{11'}^{t-t'} = -\frac{i}{4} v_{1234} \langle 0|T\{(\psi_2^\dagger\psi_3\psi_4)_t(\psi_{4'}^\dagger\psi_{3'}^\dagger\psi_{2'})_{t'}\}|0\rangle_{irr}\, v_{3'4'1'2'}\ , \tag{2.5}$$

where v_{1234} is the antisymmetrized matrix element of bare interaction, and

$$\mathrm{Im}\, M_{11'}^{t-t'} = -\frac{1}{8} v_{1234} \langle 0|[(\psi_2^\dagger\psi_3\psi_4)_t,(\psi_{4'}^\dagger\psi_{3'}^\dagger\psi_{2'})_{t'}]|0\rangle_{irr}\, v_{3'4'1'2'}. \tag{2.6}$$

From eq. (2.5) we see that M is related to the irreducible 2p-1h (2h-1p) Green function,

$$R_{1231'2'3'} = -i\langle 0|T\{\psi_1{}^\dagger\psi_2\psi_3\psi_{3'}^\dagger\psi_{2'}^\dagger\psi_{1'}\}|0\rangle_{irr}\ , \tag{2.7}$$

and we further write

$$\mathrm{Im}\, M_{11'}^{t-t'} = -\frac{i}{8} v_{1234} \left(R^{>\,t-t'}_{2342'3'4'} + R^{<\,t-t'}_{2342'3'4'}\right) v_{3'4'1'2'}\ , \tag{2.8}$$

where

$$iR^{>\,t-t'}_{1231'2'3'} = \langle 0|(\psi_1{}^{\dagger}\psi_2\psi_3)_t(\psi^{\dagger}_{3'}\psi^{\dagger}_{2'}\psi_{1'})_{t'}|0\rangle_{irr} , \quad (2.9a)$$

and

$$-iR^{<\,t-t'}_{1231'2'3'} = \langle 0|(\psi^{\dagger}_{3'}\psi^{\dagger}_{2'}\psi_{1'})_{t'}(\psi_1{}^{\dagger}\psi_2\psi_3)_t|0\rangle_{irr} . \quad (2.9b)$$

Assuming a weak damping of 2p-1h (2h-1p) states, the functions (2.9) can be represented as

$$iR^{>\,t-t'}_{1231'2'3'} = 2\sum_L \bar{N}_L \Xi^L_{123}\Xi^{L*}_{1'2'3'} e^{-i\Omega_L(t - t')} , \quad (2.10a)$$

$$-iR^{<\,t-t'}_{1231'2'3'} = 2\sum_L N_L \Xi^L_{123}\Xi^{L*}_{1'2'3'} e^{-i\Omega_L(t - t')} . \quad (2.10b)$$

Here Ξ^L are the wavefunctions and Ω_L are the energies of the correlated states. From general thermodynamic considerations [10]), $\bar{N}_L = 0$ for $\Omega_L < \epsilon_F$ and $N_L = 0$ for $\Omega_L > \epsilon_F$, in the ground state of a system, where ϵ_F is Fermi energy. We can further assume that the states are normalized so that $\bar{N}_L = 1$ for $\Omega_L > \epsilon_F$ and $N_L = 1$ for $\Omega_L < \epsilon_F$. Then eqs. (2.9) and (2.10) yield the completeness relations

$$\sum_{\substack{L \\ \Omega_L > \epsilon_F}} \Xi^L_{123}\Xi^{L*}_{1'2'3'} = n_1\bar{n}_2\bar{n}_3 \; \delta_{11'}\tfrac{1}{2}(\delta_{22'}\delta_{33'} - \delta_{23'}\delta_{32'}) , \quad (2.11a)$$

$$\sum_{\substack{L \\ \Omega_L < \epsilon_F}} \Xi^L_{123}\Xi^{L*}_{1'2'3'} = \bar{n}_1 n_2 n_3 \; \delta_{11'}\tfrac{1}{2}(\delta_{22'}\delta_{33'} - \delta_{23'}\delta_{32'}) , \quad (2.11b)$$

where n are occupations of the single-particle states in the shell model, $\bar{n} = 1 - n$. Upon introducing a dual system of functions $\tilde{\Xi}$ such that

$$\tilde{\Xi}^{L*}_{123}\Xi^{L'}_{123} = \delta^{LL'}, \quad (2.12)$$

and summing eqs. (2.11) side by side, we can obtain the normalization condition

$$\tilde{\Xi}^{L*}_{123}(n_1\bar{n}_2\bar{n}_3 + \bar{n}_1 n_2 n_3)\tilde{\Xi}^{L'}_{123} = \delta^{LL'} . \qquad (2.13)$$

The 3-body wavefunctions Ξ and the energies of the states can be searched for by solving the 3-body wave equation from ref. [11] (see also ref. [6]),

$$(\Omega_L + \varepsilon_1 - \varepsilon_2 - \varepsilon_3)\Xi^L_{123} - \tfrac{1}{2}(1 - n_2 - n_3)v_{232'3'}\Xi^L_{12'3'}$$
$$- (n_1 - n_2)v_{21'12'}\Xi^L_{1'2'3} - (n_1 - n_3)v_{31'13'}\Xi^L_{1'23'} = 0 , \qquad (2.14)$$

where ε are the shell-model particle-energies. On substituting eqs. (2.10) into (2.8) we find for the diagonal elements of the imaginary potential in the energy representation

$$\mathrm{Im}\, M = - \frac{\pi}{2} \sum_L |v_{1234}\Xi^L_{234}|^2(\bar{N}_L - N_L)\ \delta(\omega - \Omega_L) . \qquad (2.15)$$

Some further insight into eqs. (2.3), (2.6), and (2.15), may be gained by using the zero-temperature limit of finite-temperature theory to calculate the expectation values at the r.h.s. of (2.3) and (2.6). Two methods for calculating the expectation values exist. In appendix A we provide basic information on the method that yields chronological functions (such as in eqs. (2.2), (2.5), and (2.7)) in the amplitudes in the expressions for the expectation values. Details may be found e.g. in appendix C of ref. [12]. The other method [6] utilizes retarded functions in the amplitudes, and it has some advantages at finite temperatures and in nonequlibrium situations over the first method, as will be indicated. Unless explicitly stated, we use the more common chronological functions. We find

$$\mathrm{Im}\, M_{11'} = - \frac{i}{4} W_{1234} (G^<_{2'2}\, G^>_{33'}\, G^>_{44'}$$
$$+ G^>_{2'2}\, G^<_{33'}\, G^<_{44'}) W^*_{1'2'3'4'} + \cdots , \qquad (2.16)$$

where

$$G^{>}_{11'} = -i\langle 0|\psi_1\psi^{\dagger}_{1'}|0\rangle, \tag{2.17a}$$

and

$$G^{<}_{11'} = i\langle 0|\psi^{\dagger}_{1'}\psi_1|0\rangle. \tag{2.17b}$$

The vertex function can be written in terms of R as

$$W_{1234} = \frac{1}{2} v_{1567} R_{5672'3'4'} G^{-1}_{22'} G^{-1}_{3'3} G^{-1}_{4'4} . \tag{2.18}$$

The time-arguments of W are in general different from one another. The same holds for the right-hand time-arguments of R. The dots at the r.h.s. of (2.16) indicate contributions from the 3p-2h and 2h-3p states that are products of single-particle states, and such states with a still larger number of particle-hole pairs.

The result of the form (2.16) is valid at finite temperature when the expectation values that determine the functions in (2.3), (2.5), and (2.17), are taken with respect to the equilibrium density operator [6]. A complete contribution of the 2p-h and 2h-p states to the imaginary part of the mass operator is obtained, though, in that case, when the retarded rather than chronological amplitudes W are used [6]. The difference vanishes at zero temperature.

In the shell-model approximation we have

$$G^{>}_{11'} = -i\bar{n}_1 \delta_{11'} e^{-i\varepsilon_1(t_1 - t_1')}, \tag{2.19a}$$

and

$$G^{<}_{11'} = in_1 \delta_{11'} e^{-i\varepsilon_1(t_1 - t_1')}, \tag{2.19b}$$

where the ground-state occupation n is 1 for a state below Fermi surface, and 0 above; $\bar{n} = 1 - n$. On inserting (2.19) into (2.16), we find for the diagonal elements of imaginary potential in the energy representation,

$$\mathrm{Im}\, M \simeq -\frac{\pi}{2}|W_{1234}|^2\,(n_2\,\bar{n}_3\,\bar{n}_4 - \bar{n}_2\,n_3\,n_4)\;\delta(\omega + \varepsilon_2 - \varepsilon_3 - \varepsilon_4). \qquad (2.20)$$

We should, of course, be aware of the fact that the shell-model states in (2.20) may themselves have a decay width and/or have a reduced strength (Z-factors less than 1). We want to ignore such refinements here and rather concentrate on the vertex function W. With the help of the Watson-Fadeev series expansion of the 3-body propagator involved in (2.18) [refs. [9,11,6]], we can expand W in a series of pp and ph T-matrices. To lowest order W can be approximated by a pp T-matrix (a Brueckner T-matrix could be used, a Feynman-Galitski T-matrix, or any more sophisticated version thereof). Low-order terms of the expansion are shown in fig. 1, and the second equality indicated in the figure follows from an integral equation satisfied by the T-matrix (the dot stands for v_{1234}).

A subject that we want to particularily dwell on in this work, is how to incorporate collectivity from low-lying collective-states (including giant resonances) into the optical potential. For that purpose we add to the terms indicated in fig. 1, the next terms in the series, see fig. 2. The ph T-matrix satisfies an integral equation

$$\begin{aligned} T^{ph}_{1234} &= K^{ph}_{1234} - iK^{ph}_{123'4'}G_{3''3'}G_{4'4''}T^{ph}_{3''4''34} \\ &= K^{ph}_{1234} + K^{ph}_{123'4'}\Pi_{3'4'3''4''}K^{ph}_{3''4''34}\,, \end{aligned} \qquad (2.21)$$

that can be directly derived from the Bethe-Salpeter equation for the ph propagator [13,14],

$$\Pi = \Pi^0 + \Pi^0 K^{ph}\Pi\,, \qquad (2.22)$$

where explicitly

$$\Pi_{121'2'} = -i\langle 0|T\{\delta\rho_{21}\delta\rho^{\dagger}_{2'1'}\}|0\rangle \ , \tag{2.23}$$

$\delta\rho_{12} = \psi^{\dagger}_{2}\psi_{1} - \langle 0|\psi^{\dagger}_{2}\psi_{1}|0\rangle$, and K^{ph} is an irreducible kernel (mass operator in the ph channel). With this we can now write the expression for the vertex function indicated in fig. 2,

$$\begin{aligned} W_{1234}&(t - t') \\ &= T^{pp}_{1234}\delta(t - t') + T^{pp}_{12'3'4}\, e^{-i\varepsilon_4(t-t')}\, \Pi_{2'3'2''3''}(t - t')\, K^{ph}_{2''3''23} \\ &\quad + T^{pp}_{12'34'}\, e^{-i\varepsilon_3(t-t')}\, \Pi_{2'4'2''4''}(t - t')\, K^{ph}_{2''4''24} \ . \end{aligned} \tag{2.24}$$

Here we assume that T^{pp} and K^{ph} are energy independent (the dependencies on the relative times may be expressed with the δ functions). We choose for convenience the time representation since the forward- and backward-going contributions may be written concisely. Equation (2.24) together with (2.20) is the basic result of this work. Collectivity is accounted for via the density-density correlation function (2.23) which is available from extended RPA calculations [15,16]. For the interactions T^{pp} and K^{ph} there exist parametrizations based on Brueckner calculations and corresponding Landau parameters, respectively. Skyrme-like forces might be used at low energies. It should then not be too difficult to evaluate (2.24) numerically. The definite sign of Im M and correct Born limit are guaranteed in this case (see also discussion in sect. 4).

We now want to make connection with formulas more commonly used for including the collectivity. To this end we carry out the following consideration. In the energy space we take the ph propagator in its spectral representation:

$$\Pi_{121'2'} = \sum_{\substack{\nu \\ \Omega_\nu > 0}} \frac{\chi^{\nu}_{12}\chi^{\nu *}_{1'2'}}{\omega - \Omega_\nu + i\Gamma_\nu/2} - \frac{\chi^{\nu *}_{21}\chi^{\nu}_{2'1'}}{\omega + \Omega_\nu - i\Gamma_\nu/2} . \tag{2.25}$$

The states in (2.25) may e.g. be calculated in the RPA approximation. The widths Γ_ν are supposed to come from coupling to the continuum states and from the coupling to the low-lying collective and single-particle states. Let us examine a term entering the mass operator, taking the first term at the r.h.s. of (2.25) and putting it into (2.21), and further into the second term at the r.h.s. of the equality in fig. 2. Using a shell-model approximation for the single-particle Green's functions and isolating few (collective) resonances in (2.25), we get a contribution to Im M, ignoring interference terms,

$$-\frac{1}{\pi} \sum_\nu{}' \sum_{234} \bar{n}_4 \frac{|\sum_{2'3'} T^{pp}_{12'3'4}\chi^{\nu}_{2'3'}|^2}{(\omega - \varepsilon_4 - \Omega_\nu)^2 + \Gamma_\nu^2/4} |\sum_{2''3''} \chi^{\nu *}_{2''3''}K^{ph}_{2''3''23}|^2 \, n_2\bar{n}_3$$

$$\times \, 2\pi\delta(\omega + \varepsilon_2 - \varepsilon_3 - \varepsilon_4) . \tag{2.26}$$

The prime on the sum indicates that the sum runs over the few states. The quantity

$$\sum_{23} |\sum_{2''3''} \chi^{\nu *}_{2''3''}K^{ph}_{2''3''23}|^2 \, n_2\bar{n}_3 \, 2\pi\delta(\omega + \varepsilon_2 - \varepsilon_3 - \varepsilon_4) \tag{2.27}$$

is recognized as the decay width of the state ν into the ph states. To the extent that these widths constitute most of the widths of the states, we get the following contribution to Im M

$$-2\frac{\pi}{2} \sum_{\substack{\nu 234 \\ \Omega_\nu > 0}}{}' |T^{pp}_{1234}\chi^{\nu}_{23}|^2 \, \bar{n}_4 \frac{\Gamma_\nu/2\pi}{(\omega - \varepsilon_4 - \Omega_\nu)^2 + \Gamma_\nu^2/4} + \dots . \tag{2.28}$$

The dots stand for an analogous term coming from the backward contribution involving the second term at the r.h.s. of (2.25). The overall factor of 2 in (2.28) accounts for the second term at the r.h.s. of equality indicated in fig. 2, and further the third term that yields an identical contribution to Im M as the second term.

Using the direct terms from the expansion of W indicated in fig. 2, we are now able to give a formula for Im M which resembles that in common use when dealing with collectivity [1,2]:

$$\begin{aligned}\mathrm{Im}\, M = & -\frac{\pi}{2}|T^{pp}_{1234}|^2(n_2\bar{n}_3\bar{n}_4 - \bar{n}_2 n_3 n_4)\delta(\omega + \varepsilon_2 - \varepsilon_3 - \varepsilon_4) \\ & - \pi \sum_{\substack{\nu \\ \Omega_\nu > 0}}{}' \Big\{ |T^{pp}_{1234}\chi^{\nu}_{23}|^2 \, \bar{n}_4 \frac{\Gamma_\nu/2\pi}{(\omega - \varepsilon_4 - \Omega_\nu)^2 + \Gamma^2_\nu/4} \\ & - |T^{pp}_{1234}\chi^{\nu *}_{32}|^2 \, n_4 \frac{\Gamma_\nu/2\pi}{(\omega - \varepsilon_4 + \Omega_\nu)^2 + \Gamma^2_\nu/4}\Big\} + \dots \, . \end{aligned} \qquad (2.29)$$

The dots in (2.29) indicate eventual interference terms and smooth background from the remainder in fig. 2.

We see that the contribution of collective vibrational states to the mass operator can either be put into the vertex function as in (2.24), or (as is usually done), following certain approximations, more directly as intermediate states as in (2.29) and as is displayed in fig. 3. We emphasize that the prime on the summation sign indicates the inclusion of isolated collective states only (see the more detailed discussion on this point in Sect. 4). Equation (2.29) is similar but not identical with the formulas in use. For instance note that expression (2.29) is of definite sign for ω below or above the Fermi energy.

The form (2.29) could have been obtained in another way. A resummation in the diagrammatic expansion can be used to obtain still a representation of the imaginary part of potential other than (2.16),

$$\mathrm{Im}\, M_{11'} = \frac{1}{4} F_{1234} \left(\Pi^{>}_{232'3'} G^{>}_{44'} + \Pi^{<}_{232'3'} G^{<}_{44'}\right) F^{*}_{1'2'3'4'} + \cdots \quad (2.30)$$

Here the vertex function is

$$F_{1234} = \frac{1}{2} v_{1567} R_{5672'3'4'} \Pi^{-1}_{2'3'23} G^{-1}_{4'4} , \quad (2.31)$$

and further

$$\Pi^{>}_{121'2'} = -i\langle 0|\delta\rho_{21}\delta\rho^{\dagger}_{2'1'}|0\rangle , \quad (2.32)$$

and

$$\Pi^{<}_{121'2'} = \Pi^{>}_{2'1'21} , \quad (2.33)$$

cf. (2.23). Inserting a complete set of states we obtain

$$\Pi^{>\,t-t'}_{121'2'} = -i\sum_{\nu} \chi^{\nu}_{12}\, \chi^{\nu *}_{1'2'}\, e^{-i\Omega_{\nu}(t - t')} , \quad (2.34a)$$

and

$$\Pi^{<\,t-t'}_{121'2'} = -i\sum_{\nu} \chi^{\nu *}_{21}\, \chi^{\nu}_{2'1'}\, e^{i\Omega_{\nu}(t - t')} . \quad (2.34b)$$

We consider the resummations that lead to (2.29) involving any of two ph pairs. We approximate the vertex-function F with the matrix T^{pp} and ignore the energy dependence of the latter. In consequence we find from (2.29) and (2.34) the contribution of a collective state ν to the diagonal element of imaginary potential in the energy representation:

$$-\pi|T^{pp}_{1234}\, \chi^{\nu}_{23}|^2\, \bar{n}_4 \delta(\omega - \Omega_{\nu} - \varepsilon_4)$$
$$+ \pi|T^{pp}_{1234}\, \chi^{\nu *}_{32}|^2\, n_4 \delta(\omega + \Omega_{\nu} - \varepsilon_4) . \quad (2.35)$$

For small values of Γ_ν this expression is the same as the one found in (2.29) where it adds to the incoherent contribution of Im M.

At finite temperature we have $n = 1/(\exp(\beta(\epsilon - \mu)) + 1)$, where β is temperature inverse. We use eq. (2.24) with the retarded propagator

$$\Pi^{+}_{121'2'} = \sum_{\nu} \frac{\chi^{\nu}_{12}\chi^{\nu*}_{1'2'}}{\omega - \Omega_\nu + i\Gamma_\nu/2} , \tag{2.36}$$

obtaining a contribution to Im M of the form (2.26). However, the quantity (2.27) does not represent now in general the width of a state, but instead the decay rate for $\Omega_\nu > 0$, and the production rate for $\Omega_\nu < 0$. The rates are related to the width with a factor $\bar{N}_\nu = \text{sgn}\,\Omega_\nu/(1 - \exp(-\beta\Omega_\nu))$. For the imaginary part of the mass operator, we get

$$\begin{aligned} \text{Im } M = & -\frac{\pi}{2}|T^{pp}_{1234}|^2(n_2\bar{n}_3\bar{n}_4 - \bar{n}_2 n_3 n_4)\delta(\omega + \epsilon_2 - \epsilon_3 - \epsilon_4) \\ & - \pi\sum_{\nu}{}' \Big\{ |T^{pp}_{1234}\chi^{\nu}_{23}|^2\, \bar{n}_4\bar{N}_\nu \frac{\Gamma_\nu/2\pi}{(\omega - \epsilon_4 - \Omega_\nu)^2 + \Gamma_\nu^2/4} \\ & - |T^{pp}_{1234}\chi^{\nu*}_{32}|^2\, n_4 N_\nu \frac{\Gamma_\nu/2\pi}{(\omega - \epsilon_4 + \Omega_\nu)^2 + \Gamma_\nu^2/4} \Big\} + \dots , \end{aligned} \tag{2.37}$$

with $N_\nu = \text{sgn}\,\Omega_\nu/(\exp(\beta\Omega_\nu) - 1)$.

Although the expressions (2.29) and (2.35) are more representative for what is common in evaluating the imaginary potential, we shall not advocate these expressions but rather stick to expression (2.24) where collectivity is put into the vertex, for reasons explained in sect. 4. First let us carry out similar considerations for the ph and pp mass operators.

3. Particle-Hole and Particle-Particle Mass Operators

We now make similar considerations for the particle-hole (phonon) and particle-particle (deuteron) mass operators.

The imaginary part of a particle-hole mass operator in a state ν may be written as [14])

$$\begin{aligned}
&\mathrm{Im}\, K^{ph\ \nu}(t - t') \\
&= \mathrm{Im}\, \chi^{\nu *}_{12}\, K^{ph\ t-t'}_{121'2'}\, \chi^{\nu}_{1'2'} \\
&= \mathrm{Im}\, -i\tilde{\chi}^{\nu *}_{12}\langle 0|T\{(\psi^{\dagger}_1 j_2 - j^{\dagger}_1\psi_2)_t(j^{\dagger}_{2'}\psi_{1'} - \psi^{\dagger}_{2'}j_{1'})_{t'}\}|0\rangle_{irr}\tilde{\chi}^{\nu}_{1'2'} \\
&= -\frac{1}{8}\,\tilde{\chi}^{\nu *}_{12}(v_{2356}\delta_{14} - v_{3415}\delta_{26}) \\
&\quad \times \langle 0|\{(\psi^{\dagger}_3\psi^{\dagger}_4\psi_5\psi_6)_t,(\psi^{\dagger}_{6'}\psi^{\dagger}_{5'}\psi_{4'}\psi_{3'})_{t'}\}|0\rangle_{irr} \\
&\quad \times (v_{5'6'2'3'}\delta_{1'4'} - v_{1'5'3'4'}\delta_{2'6'})\tilde{\chi}^{\nu}_{1'2'}\ .
\end{aligned} \tag{3.1}$$

Here $\tilde{\chi}$ is a wavefunction dual to χ,

$$\chi^{\nu}_{12} = \pm(n_1 - n_2)\tilde{\chi}^{\nu}_{12}\ , \tag{3.2}$$

where the upper sign refers to $\Omega_\nu > 0$ and lower to $\Omega_\nu < 0$, with the normalization set by $\tilde{\chi}^{\nu *}_{12}\chi^{\nu}_{12} = 1$. When ignoring the damping of the 2p-2h states, one can write

$$\begin{aligned}
&\langle 0|(\psi^{\dagger}_1\psi^{\dagger}_2\psi_3\psi_4)_t(\psi^{\dagger}_{4'}\psi^{\dagger}_{3'}\psi_{2'}\psi_{1'})_{t'}|0\rangle_{irr} \\
&= 4\sum_{\substack{J \\ \Omega_J>0}} \Phi^{J}_{1234}\Phi^{J*}_{1'2'3'4'}e^{-i\Omega_J(t - t')}\ ,
\end{aligned} \tag{3.3a}$$

$$\begin{aligned}
&\langle 0|(\psi^{\dagger}_{4'}\psi^{\dagger}_{3'}\psi_{2'}\psi_{1'})_{t'}(\psi^{\dagger}_1\psi^{\dagger}_2\psi_3\psi_4)_t|0\rangle_{irr} \\
&= 4\sum_{\substack{J \\ \Omega_J<0}} \Phi^{J}_{1234}\Phi^{J*}_{1'2'3'4'}e^{-i\Omega_J(t - t')}\ .
\end{aligned} \tag{3.3b}$$

To every wavefunction Φ^J there corresponds a wavefunction $\Phi^{J'}$, such that $\Omega_{J'} = \Omega_J$, and $\Phi^{J'}_{1234} = \Phi^{J*}_{4321}$. For the imaginary part of K^{ph} one then gets, in the energy representation,

$$\mathrm{Im}\, K^{ph\ \nu} = -\pi \sum_J |\tilde{\chi}^{\nu*}_{12}(v_{2356}\delta_{14} - v_{3415}\delta_{26})\Phi^J_{3456}|^2 (\bar{N}_J + N_J)\ \delta(\omega - \Omega_J)\ . \tag{3.4}$$

The wavefunctions χ might be searched for by solving a wave equation analogous to (2.14),

$$\begin{aligned}(\Omega_J &+ \varepsilon_1 + \varepsilon_2 - \varepsilon_3 - \varepsilon_4)\Phi^J_{1234} + \tfrac{1}{2}(1 - n_1 - n_2)v_{1'2'12}\Phi_{1'2'34} \\ &- \tfrac{1}{2}(1 - n_3 - n_4)v_{343'4'}\Phi^J_{123'4'} - (n_3 - n_1)v_{31'3'1}\Phi^J_{1'23'4} \\ &- (n_3 - n_2)v_{32'3'2}\Phi^J_{12'3'4} - (n_4 - n_1)v_{41'4'1}\Phi^J_{1'234'} \\ &- (n_4 - n_2)v_{42'4'2}\Phi^J_{12'34'} = 0\ .\end{aligned} \tag{3.5}$$

The completeness relation reads

$$\sum_{\substack{J \\ \Omega_J > 0}} \Phi^J_{1234}\Phi^{J*}_{1'2'3'4'} = n_1 n_2 \bar{n}_3 \bar{n}_4\ \tfrac{1}{4}(\delta_{11'}\delta_{22'} - \delta_{12'}\delta_{21'}) \times (\delta_{33'}\delta_{44'} - \delta_{34'}\delta_{43'}), \tag{3.6}$$

and on introducing a system of dual wavefunctions $\tilde{\Phi}$, the normalization condition can be written as

$$\tilde{\Phi}^{J*}_{1234}(n_1 n_2 \bar{n}_3 \bar{n}_4 - \bar{n}_1 \bar{n}_2 n_3 n_4)\tilde{\Phi}^{J'}_{1234} = \pm\delta^{JJ'}\ . \tag{3.7}$$

We can further expand diagrammatically the expectation value at the r.h.s. of (3.1) and apply a resummation as in the case of (2.6), to obtain

$$\begin{aligned}\mathrm{Im}\, K^{ph\ \nu} = -\frac{1}{8}\,\Gamma^{\nu}_{1234}\ (&G^{<}_{1'1}G^{<}_{2'2}G^{>}_{33'}G^{>}_{44'} \\ &+ G^{>}_{1'1}G^{>}_{2'2}G^{<}_{33'}G^{<}_{44'})\ \Gamma^{\nu}_{1'2'3'4'} + \cdots\ ,\end{aligned} \tag{3.8}$$

where the amplitude

$$\Gamma^{\nu}_{1234} = \frac{1}{2}\tilde{X}^{\nu*}_{56}\,(v_{61'3'4'}\delta_{52'} - v_{1'2'53'}\delta_{64'})\,Y_{1'2'3'4'1''2''3''4''}$$
$$\times\, G^{-1}_{11'}\,G^{-1}_{22'}\,G^{-1}_{3'3}\,G^{-1}_{4'4}\,, \qquad (3.9)$$

with the 2p-2h Green function

$$iY_{12341'2'3'4'} = \langle 0|T\{\psi^{\dagger}_{1}\psi^{\dagger}_{2}\psi_{3}\psi_{4}\psi^{\dagger}_{4'}\psi^{\dagger}_{3'}\psi_{2'}\psi_{1'}\}|0\rangle_{irr}\,. \qquad (3.10)$$

In the energy representation we get the contribution from 2p-2h states

$$\mathrm{Im}\,K^{ph\,\nu} = -\frac{\pi}{4}|\Gamma^{\nu}_{1234}|^{2}\,(n_{1}n_{2}\bar{n}_{3}\bar{n}_{4} + \bar{n}_{1}\bar{n}_{2}n_{3}n_{4})$$
$$\times\,\delta(\omega + \varepsilon_{1} + \varepsilon_{2} - \varepsilon_{3} - \varepsilon_{4})\,, \qquad (3.11)$$

where, to lowest order in T^{pp},

$$\Gamma^{\nu}_{1234} \simeq \tilde{X}^{\nu*}_{56}(T^{pp}_{6134}\,\delta_{52} - T^{pp}_{1253}\,\delta_{64} - T^{pp}_{6234}\,\delta_{51} - T^{pp}_{1254}\,\delta_{63})\,. \qquad (3.12)$$

The other important case is that of the deuteron (two-particle) mass operator K^{pp}, see also ref. [17]). With Ψ being a two-particle wavefunction the result analogous to (3.11) from retaining the transitions to 3p-h states, takes the form

$$\mathrm{Im}\,M^{pp} = \mathrm{Im}\,\Psi^{*}_{12}\,K^{pp}_{121'2'}\,\Psi_{1'2'}$$
$$= -\frac{\pi}{12}|\Lambda_{1234}|^{2}\,(n_{1}\bar{n}_{2}\bar{n}_{3}\bar{n}_{4} + \bar{n}_{1}n_{2}n_{3}n_{4})$$
$$\times\,\delta(\omega + \varepsilon_{1} - \varepsilon_{2} - \varepsilon_{3} - \varepsilon_{4})\,. \qquad (3.13)$$

Here the amplitude Λ is given by

$$\Lambda_{1234} = \frac{1}{2}\tilde{\Psi}^{*}_{56}\,(v_{1'62'3'}\delta_{54'} + v_{51'2'3'}\delta_{64'})$$
$$\times\,D_{1'2'3'4'1''2''3''4''}\,G^{-1}_{11''}\,G^{-1}_{2''2}\,G^{-1}_{3''3}\,G^{-1}_{4''4}\,, \qquad (3.14)$$

where $\tilde{\Psi}$ is a dual wavefunction, $\Psi_{12} = (1 - n_{1} - n_{2})\tilde{\Psi}_{12}$, and

$$iD_{12341'2'3'4'} = \langle 0|T\{\psi^{\dagger}_{1}\psi_{2}\psi_{3}\psi_{4}\psi^{\dagger}_{4'}\psi^{\dagger}_{3'}\psi^{\dagger}_{2'}\psi_{1'}\}|0\rangle_{irr}\,. \qquad (3.15)$$

It is important to note that into the vertices Γ, eq. (3.12), and Λ, eq. (3.14), enter the dual amplitudes with the result that integrations run over the whole rather than restricted phase space, even at zero temperature [see also ref. [14]].

4. Average Mass Operators of Nuclei

We now want to discuss in some detail the situation of nuclei. In nuclear physics essentially two different formulations are in use to calculate the imaginary part of the optical potential. The first one consists in calculating the expression (2.20) in infinite matter using the 2-body scattering matrix for W and going over to finite nuclei via the local density approximation [3,5,18].

The imaginary potential from the nuclear structure approach [1,2] takes the form

$$\begin{aligned} \mathrm{Im}\, M = &-\frac{\pi}{2}\sum_J |v_{1234}\Psi^J_{34}|^2\,(n_2\bar{N}_J - \bar{n}_2 N_J)\,\delta(\omega+\epsilon_2-\Omega_J) \\ &-\pi\sum_\nu |v_{1234}\chi^\nu_{23}|^2\,(\bar{N}_\nu\bar{n}_4 - N_\nu n_4)\,\delta(\omega-\Omega_\nu-\epsilon_4) \\ &+\pi|v_{1234}|^2\,(n_2\bar{n}_3\bar{n}_4 - \bar{n}_2 n_3 n_4)\,\delta(\omega+\epsilon_2-\epsilon_3-\epsilon_4)\ . \end{aligned} \qquad (4.1)$$

Here the sums run over <u>all</u> two-particle and particle-hole states and the last term at the r.h.s. is from subtraction of two second-order contributions to M that otherwise would be counted three times. When an effective interaction is used and the pp states are not collective, then the first term at the r.h.s. of (4.1) cancels one of the second-order contributions, leading to the result [2,7]

$$\mathrm{Im}\, M = -\pi\sum_\nu |v_{1234}\chi^\nu_{23}|^2\,(\bar{N}_\nu\bar{n}_4 - N_\nu n_4)\,\delta(\omega-\Omega_\nu-\epsilon_4)$$

$$+ \frac{\pi}{2}|v_{1234}|^2 \; (n_2\bar{n}_3\bar{n}_4 - \bar{n}_2 n_3 n_4) \; \delta(\omega + \varepsilon_2 - \varepsilon_3 - \varepsilon_4) \; , \qquad (4.2)$$

that is used in practise. The second term at the r.h.s. of (4.2) still serves to cancel half of the second-order contributions from the first term. There are inherent difficulties associated with (4.1) and (4.2) that we want to discuss. The correlation and polarization parts [cf. ref. [19]] of the imaginary potential should each have a definite sign, e.g. the correlation part should be positive definite as representing the particle absorption. This is not the case with the quantities in eqs. (4.1) and (4.2), as we explicitly demonstrate on a model in appendix B. This fact, however, seemed never to cause problems in actual calculations, and it was claimed [4,20-22] that the incoherent part is very small compared to the coherent part.
In general, this statement cannot be right for the following reason.
The incoherent part in (4.2) is up to a sign equal to the expression which is used in the local density approximation (LDA) to the imaginary potential [3,5]. These latter calculations give semi-quantitative agreement with experimental values. In ref. [23] it was shown that the LDA represents the correct average part of the corresponding quantal expression. It is in this sense that the second part of (4.2) can overall not be a small quantity. In the range of energies where collective states play an important role, it may be small; however, also there things can be more subtle. On examining the response function $\Pi = \Pi^0/(1 - K\Pi^0)$ in the interior of the nucleus corresponding to nuclear matter, one finds that the isoscalar interaction acting mostly in the surface is almost zero there. At the other hand, the repulsive isovector force $(S = 0,\ T = 1)$ is strong in the interior of the nucleus. So there Π is actually smaller than Π^0, and the subtraction in (4.2) possibly alters the sign of the imaginary part of mass operator. This is precisely what happened

in a study [24] involving one of the present authors. This never happened in a quantal calculation because of the possibly insufficiently averaged incoherent part that exhibits large fluctuations in energy (shell effects). Chance to look at energies for which the incoherent part is very small is then high. In fact the icoherent part should be averaged out to such an extent that no shell effects remain. This is what has been done in the aforementioned model study [24] in applying the Strutinsky smoothing procedure [23].

From the discussion it follows that the use of a calculational procedure for Im M which ensures the correct sign, is important. One of the purposes of the preceding sections was to show that such procedure is indeed possible. If, in particular, all collectivity is put into the vertex function W like in expression (2.24) the correct sign of Im M is guaranteed. Analogous considerations hold, of course, for the ph and pp mass operators studied in this paper.

5. Conclusion

In the paper we mainly addressed the very old question of how to put collectivity into the imaginary part of optical potentials. In our opinion the "factor 1/2 problem" was never really solved satisfactorily. We here give a novel prescription where no factor 1/2 ambiguity arises and which assures that correlation and polarization parts have a definite sign, i.e. only flux absorption and no flux creation is possible for elastic scattering. Our formulation resides in the fact that we could show that all collectivity can be put into the vertices, and thus a Fermi's golden rule form for the imaginary part of the optical potential can be applied under all cicumstances.

Our basic result is eq. (2.24) which also holds at finite temperature. It represents the first few terms when the vertex function is expanded in a Watson-Fadeev series. Direct numerical applications should be possible.

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Appendix A

We discuss here the computation of the expectation values such as at the r.h.s. of (2.3) and (2.6), or at the r.h.s. of (2.9).

From (2.3), we have

$$2\ \mathrm{Im}\ M_{11'}^{t-t'} = \langle 0|j_{1'}^{\dagger}(t')j_1(t)|0\rangle - \langle 0|j_1(t)j_{1'}^{\dagger}(t')|0\rangle. \tag{A.1}$$

For the first term at the r.h.s. of (A.1), we get

$$\begin{aligned}
&\langle 0|j_{1'}^{\dagger}(t')j_1(t)|0\rangle \\
&= \langle 0|U^I(-\infty,t')j_{1'}^{\dagger I}(t')U^I(t',\infty)U^I(\infty,t)j_1^I(t)U(t,\infty)|0\rangle \\
&= \langle 0|\left[U^I(\infty,t')j_{1'}^I(t')U^I(t',-\infty)\right]^{\dagger}U^I(\infty,t)j_1^I(t)U(t,\infty)|0\rangle \\
&= \langle 0|\left[T\left\{j_{1'}^I(t')\exp(-i\int_{-\infty}^{\infty}d\tau'\hat{V}^I(\tau'))\right\}\right]^{\dagger} \\
&\quad \times T\left\{j_1^I(t)\exp(-i\int_{-\infty}^{\infty}d\tau\hat{V}^I(\tau))\right\}|0\rangle.
\end{aligned} \tag{A.2}$$

When we expand the exponentials at the r.h.s. of (A.2) and apply a Wick decomposition, we obtain contractions of the interaction-picture operators of which both are associated with j, or with $j^{\dagger}$, or one with j and the other with $j^\dagger$. For both field operators associated with j, the contraction coincides with the function iG. For both field operators associated with $j^\dagger$, the contraction coincides with the complex conjugate function $(iG)^*$. The expansion of the exponential for $j^\dagger$ in (A.2) gives further also conjugate potential factors $(-iv)^*$. When one field-operator is associated with $j^\dagger$ and the other with j, then the contraction coincides either with $iG^>$ or $iG^<$, cf. (2.17). Disconnected diagrams merely give an overall factor $1 = \langle 0|U^I(-\infty,\infty)U^I(\infty,-\infty)|0\rangle$. Resummations can be carried in a standard manner, so as to eliminate the expectation values of interaction-picture operators in favor of the expectation values of Heisenberg-picture operators, $\langle 0|T\{\psi^I\psi^{\dagger I}\}|0\rangle \to \langle 0|T\{\psi^H\psi^{\dagger H}\}|0\rangle$, $\langle 0|\psi^I\psi^{\dagger I}|0\rangle \to \langle 0|\psi^H\psi^{\dagger H}|0\rangle$. Grouping together the terms with a minimum number of particle and hole density factors $G^{\gtrless}$, we get

$$\langle 0|j^\dagger_{1'}(t')j_1(t)|0\rangle_{irr}$$
$$= -\frac{i}{2} W_{1234}\, G^>_{2'2}\, G^<_{33'}\, G^<_{44'}\, W^*_{1'2'3'4'} + \cdots , \tag{A.3}$$

where dots indicate terms containing additional particle-hole factors $G^>G^<$, and the function W is given by (2.18).

The other expectation value in (A.1) can be treated similarily,

$$\langle 0|j_1(t)j^\dagger_{1'}(t')|0\rangle$$
$$= \langle 0|U^I(-\infty,\infty)|0\rangle\langle 0|U^I(\infty,t)j^I_1(t)U(t,\infty)|0\rangle$$
$$\times \langle 0|U^I(-\infty,t')j^{\dagger I}_{1'}(t')U^I(t',\infty)|0\rangle\langle 0|U^I(\infty,-\infty)|0\rangle$$
$$= \langle 0|T\{j^I_1(t)\exp(-i\int_{-\infty}^{\infty} d\tau\hat{V}^I(\tau))\}$$
$$\times \left(T\{j^I_{1'}(t')\exp(-i\int_{-\infty}^{\infty} d\tau'\hat{V}^I(\tau'))\}\right)^\dagger|0\rangle. \tag{A.4}$$

After carrying such steps as those leading from (A.2) to (A.3), an analogous result to (A.3) is obtained, with the densities of particles and holes interchanged,

$$\langle 0|j_1(t)j^\dagger_{1'}(t')|0\rangle_{irr} = \frac{i}{2} W_{1234}\, G^<_{2'2}\, G^>_{33'}\, G^>_{44'}\, W^*_{1'2'3'4'} + \cdots . \qquad (A.5)$$

Proceeding in an analogous manner as in the resummations leading to the interacting single-particle functions in the expansion of (A.2) or (A.4), one can carry a resummation that leads to the particle-hole function (2.32) or (2.33) for a chosen particle-hole pair, cf. (2.30).

Appendix B

We want to show explicitly in the case of the Lipkin model that formula (4.2) can lead to the wrong sign of Im M. The model [see e.g. [25]] consists of two equally degenerate levels (energy $\pm\varepsilon/2$) with the lower level filled in the noninteracting case. The interaction is of the monopole-monopole form and thus the hamiltonian reads

$$H = \varepsilon J_0 - \frac{V}{2}(J_+^2 + J_-^2) , \qquad (B.1)$$

with

$$J_0 = \sum_{m=1}^{\Omega} (c^\dagger_{1m}c_{1m} - c^\dagger_{0m}c_{0m}), \quad J_+ = \sum_{m=1}^{\Omega} c^\dagger_{1m}c_{0m} , \quad J_- = (J_+)^\dagger . \qquad (B.2)$$

We consider the single-particle Green's function $-i\langle T\{c_{0m}(t)c^\dagger_{0m}(t')\}\rangle$ with the corresponding mass operator

$$M^{t-t'} = -iV^2\langle T\{(c_{1m}J_-)_t(J_+c^\dagger_{1m})_{t'}\}\rangle_{irr} . \qquad (B.3)$$

Off-diagonal elements of M do not appear in the RPA approximation to the mass operator,

$$M = -iV^2\langle T\{c_{1m}c^{\dagger}_{1m}\}\rangle_0\langle T\{J_-J_+\}\rangle \; . \tag{B.4}$$

The index 0 at the expectation value indicates that the Green's function is calculated in zeroth order with respect to V. The noninteracting and interacting Green's functions are diagonal in m and further independent of m since the matrix element in H does not depend on m.

The second-order contribution to M is given by

$$M^{(2)} = -iV^2\sum_{m'}\langle T\{c^{\dagger}_{0m}c_{0m'}\}\rangle_0[\langle T\{c_{1m}c^{\dagger}_{1m}\}\rangle_0\langle T\{c_{1m'}c^{\dagger}_{1m'}\}\rangle_0 - \langle T\{c_{1m}c^{\dagger}_{1m'}\}\rangle_0\langle T\{c_{1m'}c^{\dagger}_{1m}\}\rangle_0] \; . \tag{B.5}$$

We now use the diagonality and independence of m to obtain

$$M^{(2)} = -iV^2(\Omega - 1)\langle T\{c^{\dagger}_0c_0\}\rangle_0(\langle T\{c_1c^{\dagger}_1\}\rangle_0)^2 \; . \tag{B.6}$$

The usual approximation (4.2) to M corresponds then to

$$M \cong -iV^2[\langle T\{c_1c^{\dagger}_1\}\rangle_0\langle T\{J_-J_+\}\rangle - \langle T\{c^{\dagger}_0c_0\}\rangle_0(\langle T\{c_1c^{\dagger}_1\}\rangle_0)^2] \; , \tag{B.7}$$

where $\langle T\{J_-J_+\}\rangle$ should be calculated in the RPA approximation. To lowest order (B.7) is identical with (B.6) as it should.

With "1" being the particle index, only the polarization part of the mass operator survives. The Fourier transformation of the second and first term in the bracket at the r.h.s. of (B.7) gives, respectively,

$$\frac{i}{\omega - \frac{3}{2}\varepsilon + i\eta} \; , \tag{B.8}$$

and,

$$\frac{i|\langle 0|J_-|\nu\rangle|^2}{\omega - (E_\nu + \frac{\varepsilon}{2}) + i\eta} \; . \tag{B.9}$$

The RPA amplitudes $\langle 0|J_-|\nu\rangle$ and energies E_ν are easily calculated in this model:

$$|\langle 0|J_-|\nu\rangle|^2 = \frac{\Omega}{2}\left[1 + \frac{\varepsilon}{E_\nu}\right] , \quad E_\nu = \varepsilon\sqrt{1 - \chi^2} , \quad \chi = \frac{V}{\varepsilon}(\Omega - 1) . \tag{B.10}$$

Assuming a finite value of η as is usually done to obtain a smooth optical potential, we get for the imaginary part:

$$\mathrm{Im}\, M(\omega) = -V^2\left(|\langle 0|J_-|\nu\rangle|^2 \frac{\eta}{[\omega - (E_\nu + \frac{\varepsilon}{2})]^2 + \eta^2} - \frac{\eta}{(\omega - \frac{3}{2}\varepsilon)^2 + \eta^2}\right) . \tag{B.11}$$

The open parameters in (B.11) are χ, Ω, and η. One can see that for certain values of these parameters and ω the expression (B.11) can be positive which corresponds to nucleon production rather than absorption. This occurs e.g. when Ω is small, i.e. the RPA state is not very collective. This is precisely what can happen in a more realistic calculation where collective and less collective RPA states enter the complete spectrum. In fig. 4 we show the absorption rate for $\omega > 0$, $-2\mathrm{Im}\, M$, in units of ε, calculated using eq. (B.11) for $\Omega = 4$, $\chi = 0.88$, and $\eta/\varepsilon = 0.14$.

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Figure captions

Fig. 1 Approximation of the vertex function W by a scattering matrix T^{pp}.

Fig. 2 Approximation to the vertex function including ph correlations.

Fig. 3 Contribution to the imaginary part of the optical potential from the particle-hole coupling.

Fig. 4 Absorption rate -2ImM in the Lipkin model from eq. (B.11).

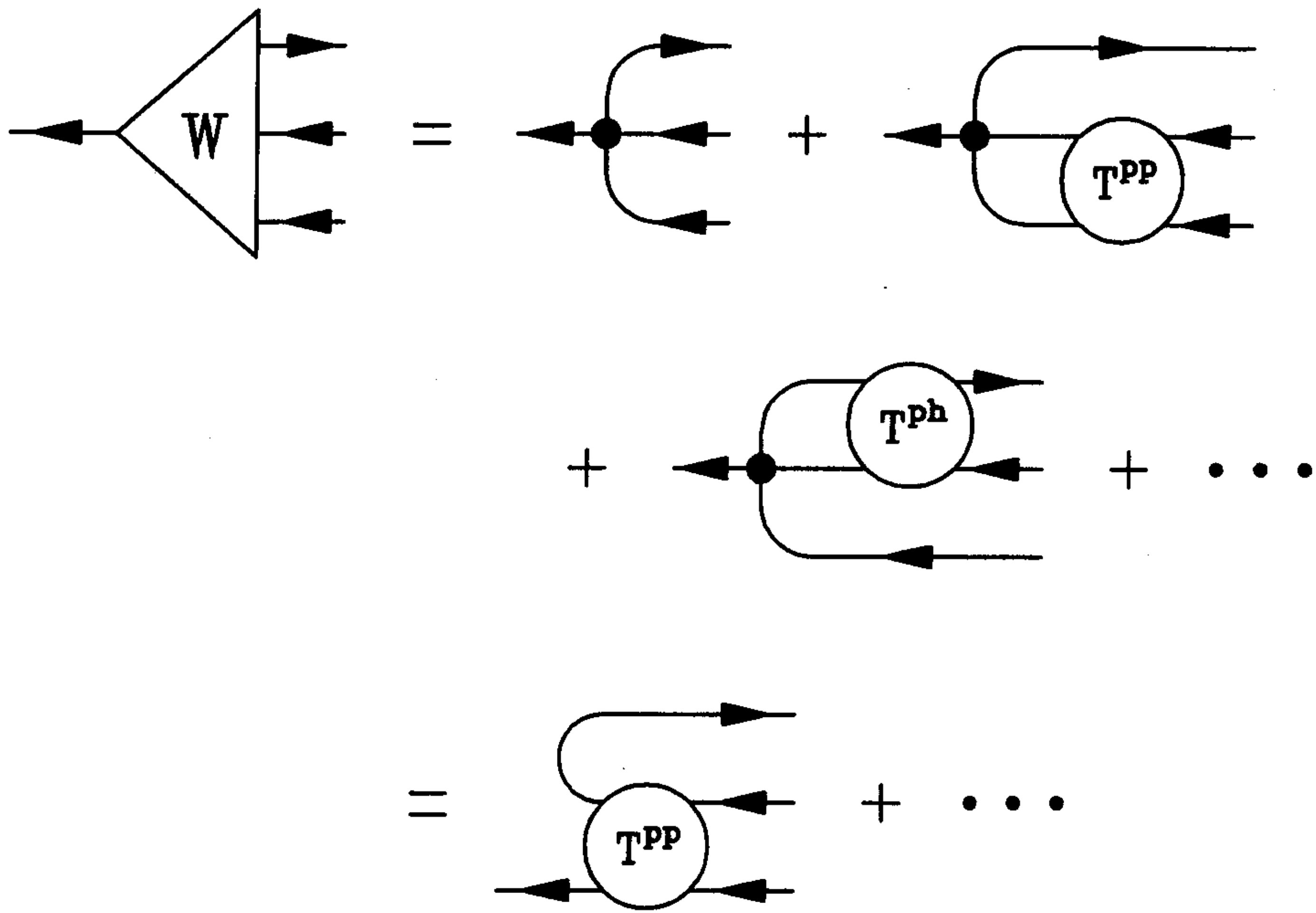

Fig. 1

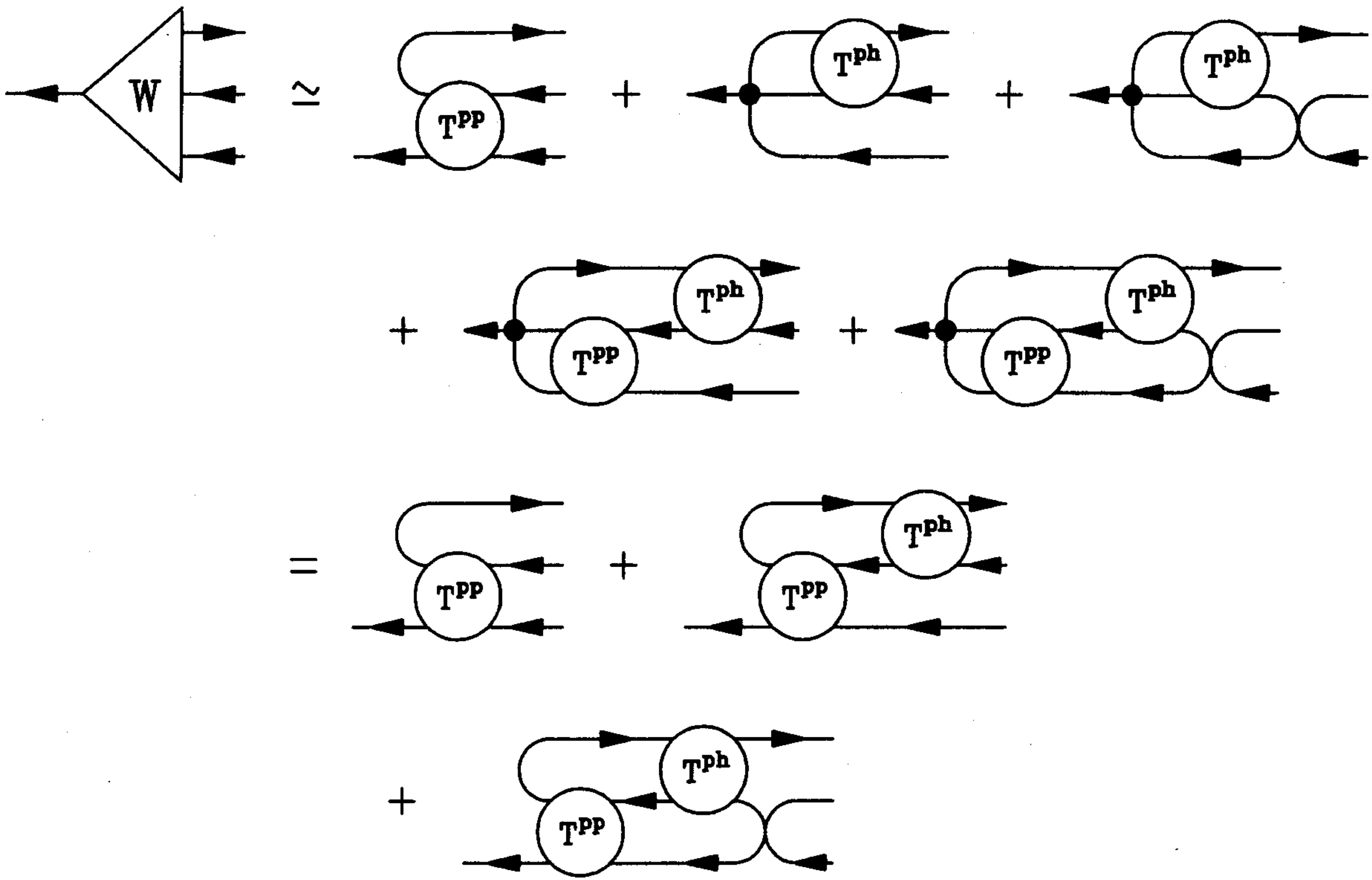

Fig. 2

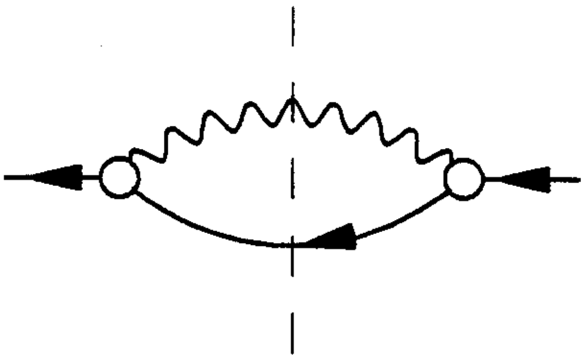

Fig. 3

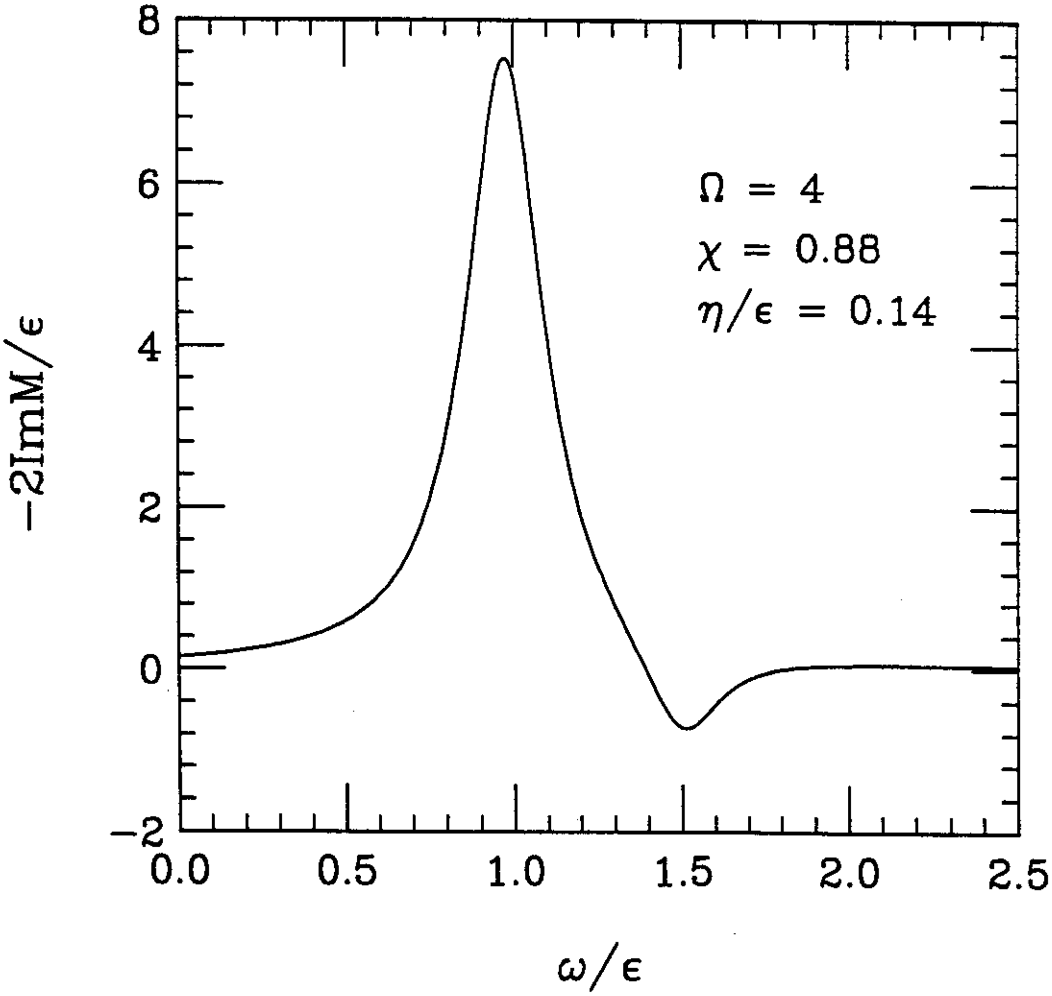

Fig. 4