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Spin-isospin strength and spectral functions

G. A. Rijsdijk, W. J. W. Geurts, M. G. E. Brand, and K. Allaart

Department of Physics and Astronomy, Free University, De Boelelaan 1081, 1081 HV Amsterdam, The Netherlands

W. H. Dickhoff

Department of Physics, Washington University, St. Louis, Missouri 63130

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The issue of spin-isospin strength is addressed from the perspective of the information that is available about one-nucleon spectral functions. Both the effect of ground-state correlations and of the coherence by random-phase-approximation correlations are then incorporated in a natural way. The results for ^{48}Ca and ^{90}Zr are confronted with those obtained by earlier methods, which have been criticized recently. It is shown that ground-state correlations produce a large amount of strength in the first 20 MeV above the main peaks of the Gamow-Teller and spin-dipole resonances. They also yield for these nuclei with rather large neutron excess considerable (n, p) strength, thereby increasing the total (p, n) strength by 25–30%. In spite of this, the strength within the hitherto analyzed energy domains is not enhanced by these correlations. It is argued that a large amount of quasifree background that is assumed in the analysis of charge exchange reactions consists of broad distributions of genuine charge exchange strength.

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I. INTRODUCTION

The distribution of excitation strength over a certain range of excitation energies is governed by at least two distinct mechanisms. First, there are the coherence effects which may give rise to collective excitations, i.e., to concentration of the main portion of strength in a few states or resonances. Long-range attractive multipole forces give rise to low-lying collective isoscalar natural parity excitations, while for isovector excitations the collective state is often pushed upward in energy. As a consequence of this, the other states with the same spin and parity are then left with very little excitation strength in comparison with the independent particle shell model (IPSM) estimates. A well known example is the isovector dipole resonance for which a schematic model [1] illustrates these collective features. Indeed low-lying 1^- excitations are observed to be several orders of magnitude weaker than an IPSM estimate, while virtually all isovector dipole strength is absorbed by the giant resonances. For charge exchange Gamow-Teller (GT) and spin-dipole (SD) excitations a similar situation occurs, as was pointed out long ago by Ikeda *et al.* [2] and Ejiri *et al.* [3] on the basis of the observed hindrance factors in low-energy beta decay. These authors stressed the relation between (approximate) symmetries of the nuclear Hamiltonian, the widths of resonances, and the quenching of strength for low-lying states. An extreme example is the isospin symmetry which gives rise to a very narrow collective state, the isobar-analog resonance (IAS), and to low-energy Fermi transitions forbidden by several orders of magnitude. Spin-isospin symmetry is only approximate; consequently, the collective GT resonance (GTR) has a large width and low-energy GT beta decay is hin-

dered by typically one order of magnitude only [2]. For first-forbidden beta decay, i.e., spin-dipole strength, the situation is similar [3]. So on the basis of relatively weak beta decay strengths one could predict the existence of giant resonances which were observed much later in (p, n) and $(^3\text{He}, t)$ reactions [4–9].

A second mechanism which plays a role in the distribution of excitation strengths is the quasiparticle nature of the shell model [10]. This has been demonstrated nicely in recent $(e, e'p)$ experiments [11–15]. From these experiments one learns that the single-particle or single-hole states in the $A \pm 1$ nuclei, when A is a magic nucleus, should be understood as only, for about 60%, merely an addition or removal of a nucleon in a specific shell model orbit. The spectral function, i.e., the distribution of single-particle addition or removal strength, is fragmented over many states and over a wide energy range. This may be attributed to the coupling of the single-particle or single-hole configuration to $2p1h$ (two-particle-one-hole), $2h1p$, and more complicated configurations. For the description of these spectral functions the Dyson equation for the one-body Green function provides a most suitable tool [10,16]. Since the single-particle and single-hole states appear to have spectroscopic factors of only 50%–70% of the IPSM value, one may expect that particle-hole states which result from the combination of the particle and hole quantum numbers have excitation strengths that are only 25–50% of the corresponding strengths in the IPSM. This is indeed observed for states with a stretched coupling of angular momenta of the particle and hole [17–19]. The quenching by factors 2–4 relative to the IPSM is not noticed for isoscalar low-lying natural parity excitations, because it is overcompensated by the aforementioned collectiv-

ity produced by long-range forces, as described, e.g., in the random phase approximation (RPA). This implies, however, that the collective correlations are actually stronger than described in those RPA calculations with phenomenological forces which do fit the data, but do not properly account for the quasiparticle nature of the particle or hole states.

With the discovery of the spin-isospin resonances and the confrontation of the data [4–9], with largely model independent sum rules [20], the distribution of excitation strength within and beyond the experimental range of energies became an important issue. For instance, the observation of only about 60% of the sum rule strength for the GT resonance was ascribed to the coupling to the Δ_{33} excitation mode of the nucleon, as a consequence of which the missing strength had been moved a few hundred MeV further upward in excitation energy [21–24]. On the other hand, it was demonstrated that calculations which account for the coupling between 1p1h and 2p2h excitations [25–29] invariably predict a spread of a substantial amount of additional excitation strength over several tens of MeV above the resonance peaks. This result was also indicated by the RPA calculations of Ref. [30]. The conclusion from such calculations is then that the effect of the coupling to the Δ_{33} resonance is at most some 10% [28,31]. In other calculations [32,33] it has been shown that ground-state correlations increase the difference between the observed and calculated total strength by another 25% of the sum rule. However, this work was criticized recently by Van Neck *et al.* [34] who noted that these ground-state correlations, obtained with perturbation theory, have not been properly normalized. A “renormalization” [34] reduces their effect by typically a factor of 2.5. So the subject should be reconsidered from a theoretical point of view.

In this work, the spin-isospin distribution is addressed from the perspective of the present knowledge of spectral functions, which has been improved substantially by recent ($e, e'p$) experiments [11–15]. An outline of the theoretical background is presented in Sec. II, where a link between spectral functions and response functions is discussed as well as the relevant sum rules for spin-isospin strength. Details of the model and computational methods are given in Sec. III. The resulting spin-isospin strengths for ^{48}Ca are presented in Sec. IV and those for ^{90}Zr in Sec. V. Section VI contains a summary and conclusions.

II. THEORETICAL FRAMEWORK

A. A link between excitation strength and spectral functions

For the calculation of excitation strength distributions one may take as a starting point the so-called two-times polarization propagator [35]

$$\begin{aligned} iL_{\alpha\beta;\gamma\delta}(t-t') &= \langle \Psi_0 | T [c_\beta^\dagger(t)c_\alpha(t)c_\gamma^\dagger(t')c_\delta(t')] | \Psi_0 \rangle \\ &\quad - \langle \Psi_0 | c_\alpha(t)c_\beta^\dagger(t) | \Psi_0 \rangle \\ &\quad \times \langle \Psi_0 | c_\delta(t')c_\gamma^\dagger(t') | \Psi_0 \rangle. \end{aligned} \quad (1)$$

Here the greek indices refer to a single-particle basis and the operators c^\dagger and c are the particle creation and annihilation operators in the Heisenberg representation, e.g. (with $\hbar=1$),

$$c_\alpha(t) = e^{iHt}c_\alpha e^{-iHt}. \quad (2)$$

In (1) the state $|\Psi_0\rangle$ is the exact ground state of the interacting many-body system. This polarization propagator contains all essential information about one-body excitations of the system, to linear approximation in the probing field. This one notes by considering the Lehmann representation of (1) which is obtained by inserting a complete set of eigenstates of the Hamiltonian and subsequent Fourier transformation

$$\begin{aligned} L_{\alpha\beta;\gamma\delta}(\omega) &= \sum_{n \neq 0} \left[\frac{\langle \Psi_0 | c_\beta^\dagger c_\alpha | \Psi_n \rangle \langle \Psi_n | c_\gamma^\dagger c_\delta | \Psi_0 \rangle}{\omega - E^n + i\eta} \right] \\ &\quad - \sum_{m \neq 0} \left[\frac{\langle \Psi_0 | c_\gamma^\dagger c_\delta | \Psi_m \rangle \langle \Psi_m | c_\beta^\dagger c_\alpha | \Psi_0 \rangle}{\omega + E^m - i\eta} \right]. \end{aligned} \quad (3)$$

Here the poles E^n are the excitation energies relative to the ground-state energy E^0 and $|\Psi_n\rangle$ the corresponding excited states. Note that both $|\Psi_0\rangle$ and $|\Psi_n\rangle$ may be of a much more complicated nature than just 0p0h and 1p1h states. For an arbitrary one-body operator

$$O = \sum_{\alpha\beta} \langle \alpha | O | \beta \rangle c_\alpha^\dagger c_\beta, \quad (4)$$

the excitation strength distribution may be expressed as

$$\begin{aligned} S_O(\omega) &= \sum_n |\langle \Psi_n | O | \Psi_0 \rangle|^2 \delta(\omega - E^n) \\ &= -\frac{1}{\pi} \text{Im} \sum_{\alpha\beta\gamma\delta} \langle \alpha | O | \beta \rangle^* L_{\alpha\beta;\gamma\delta}(\omega) \langle \gamma | O | \delta \rangle. \end{aligned} \quad (5)$$

The polarization operator may, at least formally, be calculated by a suitable splitting of the Hamiltonian which consists of a kinetic part and a two-body interaction

$$H = \sum_{\alpha\beta} T_{\alpha\beta} c_\alpha^\dagger c_\beta + \frac{1}{4} \sum_{\alpha\beta\gamma\delta} V_{\alpha\beta\gamma\delta} c_\alpha^\dagger c_\beta^\dagger c_\delta c_\gamma \quad (6)$$

into a part H_0 which has an uncorrelated eigenstate Φ_0 ,

$$H_0 |\Phi_0\rangle = E_0 |\Phi_0\rangle \quad (\text{note } E_0 \neq E^0), \quad (7)$$

and a residual part H_1 :

$$H = H_0 + H_1. \quad (8)$$

For H_0 one may take the kinetic energy plus a suitably chosen mean field, e.g., the Hartree-Fock (HF) field. One may then write the formal perturbation expansion with time-ordering operator T [35,36],

$$iL_{\alpha\beta;\gamma\delta}(t-t') = \sum_{m=0}^{\infty} \frac{(-i)^m}{m!} \int_{-\infty}^{\infty} dt_1 \cdots \int_{-\infty}^{\infty} dt_m \langle \Phi_0 | T [H_1(t_1) \cdots H_1(t_m) c_{\beta}^{\dagger}(t) c_{\alpha}(t) c_{\gamma}^{\dagger}(t') c_{\delta}(t')] | \Phi_0 \rangle, \quad (9)$$

which may be written out in Feynman diagrams.

When writing out Wick's theorem for the various contributions in (9) one may, in all orders of the perturbation H_1 , set apart those terms in which the operators c_{γ}^{\dagger} and c_{α} are contracted with operators of H_1 but not, via any chain of interactions, linked with the operators c_{β}^{\dagger} and c_{δ} . These so-called noninteracting quasiparticle terms then form the inhomogeneous term in the Bethe-Salpeter equation

$$L_{\alpha\beta;\gamma\delta}(t_1-t_2) = -ig_{\alpha\gamma}(t_1-t_2)g_{\delta\beta}(t_2-t_1) + \sum_{\mu\nu\kappa\lambda} \int \int \int \int dt_3 dt_4 dt_5 dt_6 \\ \times [-ig_{\alpha\mu}(t_1-t_3)g_{\nu\beta}(t_4-t_1)i\Gamma_{\mu\nu;\kappa\lambda}(t_3, t_4; t_5, t_6)L_{\kappa\lambda\gamma\delta}(t_5-t_2, t_6-t_2)], \quad (10)$$

in which the irreducible particle-hole interaction Γ is again an infinite perturbation series. This equation is graphically depicted in Fig. 1, in which the double lines are the one-body Green functions g . These are defined as [35,36]

$$g_{\alpha\beta}(t_1-t_2) = (-i)\langle \Psi_0 | T [c_{\alpha}(t_1) c_{\beta}^{\dagger}(t_2)] | \Psi_0 \rangle \quad (11)$$

and by again going to the Lehmann representation

$$g_{\alpha\beta}(\omega) = \sum_n \frac{\langle \Psi_0^A | a_{\alpha} | \Psi_n^{A+1} \rangle \langle \Psi_n^{A+1} | a_{\beta}^{\dagger} | \Psi_0^A \rangle}{\omega - (E_n^{A+1} - E_0^A) + i\eta} + \sum_m \frac{\langle \Psi_0^A | a_{\beta}^{\dagger} | \Psi_m^{A-1} \rangle \langle \Psi_m^{A-1} | a_{\alpha} | \Psi_0^A \rangle}{\omega - (E_0^A - E_m^{A-1}) - i\eta}, \quad (12)$$

one notes that it contains all the information for one-particle addition and removal spectroscopic factors, i.e., of the particle and hole spectral functions. Indeed, these spectral functions have been calculated recently for various nuclei [10,16] by solving the Dyson equation

$$g_{\alpha\beta}(\omega) = g_{\alpha\beta}^{(0)}(\omega) + \sum_{\gamma\delta} g_{\alpha\gamma}^{(0)}(\omega) \Sigma_{\gamma\delta}^*(\omega) g_{\delta\beta}(\omega), \quad (13)$$

in some approximation for the irreducible self-energy Σ^* .

The guiding idea of the present work is that the experimental information on spectral functions may be used to limit the uncertainties of the Green functions g and by Eq. (10) then also of the polarization propagator L . Of course Eq. (10) can still not be solved exactly, even if the $g_{\alpha\beta}(\omega)$ are known, since one has to adopt some approximation for the interaction Γ . The approximations applied in this work are $\Gamma = 0$ (Sec. IIB) and $\Gamma = V$ where V is a G -matrix interaction (Sec. IIC).

B. Dressed independent particle approximation: Fragmentation of strength

As the simplest approximation of Eq. (10) one may neglect the interaction Γ between the dressed particles completely and then with Eq. (11) obtain for the excitation strength (5) the expression

$$S_O(\omega) = -\frac{1}{\pi} \text{Im} \sum_{\alpha\beta\gamma\delta} \langle \alpha | O | \beta \rangle^* L_{\alpha\beta;\gamma\delta}^f(\omega) \langle \gamma | O | \delta \rangle, \quad (14)$$

with

$$L_{\alpha\beta;\gamma\delta}^f(\omega) = \int \frac{d\omega'}{2\pi i} g_{\alpha\gamma}(\omega + \omega') g_{\delta\beta}(\omega') \\ = \sum_{nm} \frac{\langle \Psi_0^A | a_{\alpha} | \Psi_n^{A+1} \rangle \langle \Psi_n^{A+1} | a_{\gamma}^{\dagger} | \Psi_0^A \rangle \langle \Psi_0^A | a_{\beta}^{\dagger} | \Psi_m^{A-1} \rangle \langle \Psi_m^{A-1} | a_{\delta} | \Psi_0^A \rangle}{\omega - (E_n^{A+1} - E_m^{A-1}) + i\eta} \\ - \sum_{nm} \frac{\langle \Psi_0^A | a_{\gamma}^{\dagger} | \Psi_m^{A-1} \rangle \langle \Psi_m^{A-1} | a_{\alpha} | \Psi_0^A \rangle \langle \Psi_0^A | a_{\delta} | \Psi_n^{A+1} \rangle \langle \Psi_n^{A+1} | a_{\beta}^{\dagger} | \Psi_0^A \rangle}{\omega - (E_m^{A-1} - E_n^{A+1}) - i\eta}. \quad (15)$$

In this dressed independent particle approximation (DIPA) the effect of the fragmentation of the spectral functions is fully reflected, and so one expects a similar fragmentation of the excitation strengths over a very wide

range of energies. Still missing is the shift of strength by collective correlations caused by the residual interaction Γ . In fact, any approximation to Γ which does not fulfill the Baym-Kadanoff relation

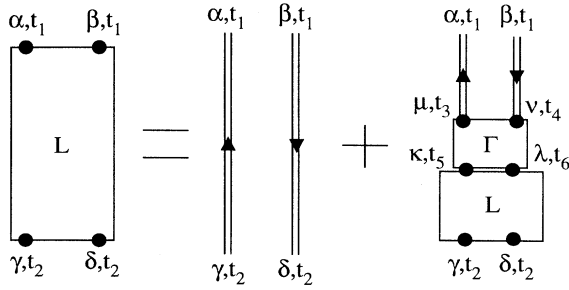


FIG. 1. Graphical representation of the Bethe-Salpeter equation (10) for the two-times polarization propagator. The double lines represent the fully dressed single-particle propagator (11).

$$\Gamma_{\alpha\beta\gamma\delta} = \frac{\delta\Sigma_{\alpha\beta}^*}{\delta g_{\gamma\delta}} \quad (16)$$

will in principle violate conservation laws [37,38]. Indeed, with a Σ^* in second order in the interaction [10], which shall be applied here, Eq. (16) is not fulfilled. For this reason it is necessary to check the sum rules of Sec. II E explicitly when Eq. (16) is not obeyed. If these rules are found to be satisfied to a good approximation, this will be considered as a positive indication for the usefulness of the method [39].

The physical interpretation of Eq. (14) is that it is similar in spirit to the independent particle shell model (IPSM), which is obtained by inserting a single-pole approximation for g , with pole strength unity. Now one has an independent dressed particle shell model. The dynamical interaction of the particle and the hole with the nuclear medium is fully included in the single-particle propagator $g(\omega)$, but the interaction between the quasi-particles is neglected. Note that due to the quasiparticle nature the response (14) contains not only ph and hp but also pp and hh contributions.

C. Dressed particle RPA (DRPA) with $\Gamma = V$

In order to account for the concentration of strength into collective states or resonances, the interaction Γ must be included. Here it is restricted to the lowest order approximation, i.e., $\Gamma = V$, where V is a G -matrix interaction. With this approximation the polarization propagator is now calculated by solving the equation

$$L_{\alpha\beta\gamma\delta}(\omega) = L_{\alpha\beta\gamma\delta}^f(\omega) + \sum_{\kappa\lambda\mu\nu} L_{\alpha\beta\kappa\lambda}^f(\omega) V_{\kappa\lambda\mu\nu} L_{\mu\nu\gamma\delta}(\omega) \quad (17)$$

and then calculating the excitation strength with Eq. (5). The approximation (17) will be referred to as the dressed particle RPA (DRPA), because if the free propagation of dressed particles in $L^f(\omega)$ is replaced by the free propagation of a normal particle or hole, one has the standard particle-hole RPA equation [40].

D. Comparison with extended RPA and other methods

From the presentation given here, it is clear that one should like to include terms of higher order in V into the effective interaction Γ that occurs in the Bethe-Salpeter equation (10). Such an effective interaction then becomes energy dependent, which involves a computational complexity outside the scope of the present work. In the extended RPA (ERPA) of Refs. [26–28] diagrams up to second order in the interaction have been included, but then a simple single-pole approximation of the single-particle propagator was used, instead of the solution of the Dyson equation (13), with a self-energy up to second order in V . The diagrams included in the ERPA are depicted in Fig. 2. This approach is somewhat complementary to the present one. In the ERPA the dynamical induced interaction causes a coupling between 1p1h and 2p2h states [cf. diagrams 2(b,c)], which is still absent in the present approach. On the other hand, the use of the full solution of the Dyson equation (13), even with self-energy up to second order, in Eq. (17), implies the treatment of particle-particle and hole-hole amplitudes on equal footing with particle-hole and hole-particle amplitudes. In the derivation of the ERPA equations [28] a restriction to only particle-hole and hole-particle amplitudes is imposed. In Refs. [26,27] pp and hh amplitudes were added in a perturbation approximation to the solution of the ERPA equations. So by comparing results of the ERPA of Ref. [28] and the present DRPA method one may get some feeling for the relevance of the induced interaction on one side and pp and hh contributions on the other side. Such a comparison will be made for the results of ^{48}Ca in Sec. IV and for the GTR of ^{90}Zr in Sec. V.

It should be mentioned that pp and hh amplitudes are

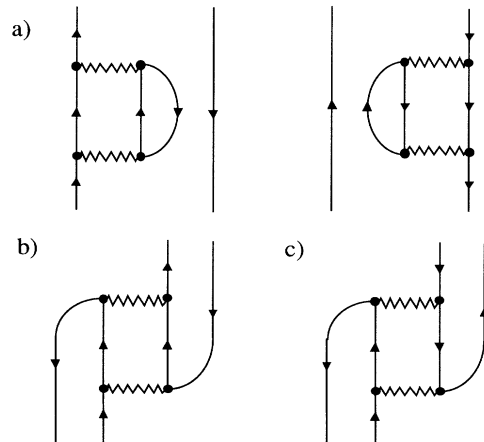


FIG. 2. Second-order diagrams for the polarization propagator that are included in the ERPA calculations of Refs. [26–29,31]. The lines represent single-particle (upward arrows) or hole (downward arrows) propagation in a mean field, e.g., HF field. The wiggly lines represent a G -matrix interaction.

also included in the calculations of Nishizaki *et al.* [33]. In this approach 2p2h amplitudes are used as ingredients. The method by which these 2p2h amplitudes are obtained has been criticized recently [34]. So the results of this method should be tested against other methods, which is one of the aims of the present work. Another method in which pp and hh amplitudes are treated on equal footing with ph amplitudes is the BCS quasiparticle RPA [41]. This method has the advantage that pairing correlations are treated fully in a nonperturbative way, but at the cost of particle-number conservation. As a consequence of this number nonconservation, spuriousity becomes a problem, which is still reasonably under control on the RPA level, but which becomes prohibitive for a further extension of the method [42,43]. For this reason this approach shall not be considered any further here.

E. Spin-isospin sum rules

When the nucleus is considered as composed of nucleons only, sum rules may be derived for the total excitation strength for a given multipole operator. Suppose that the strength for a (p, n) or (n, p) transition from the ground state $|\Psi_0\rangle$ to an excited state $|\Psi_n\rangle$ is given by the matrix element [20]

$$S_{\pm 1}^{LJ}(E^n) = \sum_{\mu} |\langle \Psi_{\mu}^n | \sum_{k=1}^A r^L(k) [Y_L(k) \sigma(k)]_{J\mu} \tau_{\pm 1} | \Psi_0 \rangle|^2. \quad (18)$$

Note that this expression differs by a factor of 2 from the usual one in the theory of beta decay [2,3] because of the isospin operator

$$\tau_{\pm 1} = \mp \frac{1}{\sqrt{2}} \tau_{\pm}, \quad \tau_0 = \tau_z. \quad (19)$$

Summing the strength (18) over all final states $|\Psi^n\rangle$ one may use closure and obtain the sum rule

$$S_{-1}^{LJ} - S_{+1}^{LJ} = \frac{2J+1}{2\pi} \langle \Psi_0 | \sum_{k=1}^A r^{2L}(k) \tau_z(k) | \Psi_0 \rangle, \quad (20)$$

with which the experimentally observed (p, n) strengths, for $L \neq 0$, are usually compared [6,44].

For $L = 0$ one does not use the expressions (18) and (20) but rather

$$S_{\pm}^{\text{GT}}(E^n) = \sum_{\mu} |\langle \Psi_{\mu}^n | \sum_{k=1}^A \sigma_{\mu}(k) \tau_{\pm}(k) | \Psi_0 \rangle|^2, \quad (21)$$

which differs from the $L = 0$ substitution in (18) by the omission of Y_0 and by τ_{\pm} instead of $\tau_{\pm 1}$. So this GT sum rule reads

$$S_{-}^{\text{GT}} - S_{+}^{\text{GT}} = 3 \langle \Psi_0 | \sum_{k=1}^A \tau_z(k) | \Psi_0 \rangle = 3(N - Z). \quad (22)$$

As noted in Sec. II B these sum rules provide an important check on the theoretical methods.

III. COMPUTATIONAL FRAMEWORK

A. Model space, interactions, and single-particle energies

The calculations have been performed within a configuration space of harmonic oscillator basis functions with size parameters $b=2.00$ fm for ^{48}Ca and $b=2.12$ fm for ^{90}Zr . A model space including up to the $1h2f3p$ shell (21 shells) for ^{48}Ca and up to the $1i2g3d4s$ shell (28 shells) for ^{90}Zr was adopted. As an interaction within this space a parametrization [45] of a G matrix deduced from a Bonn meson-exchange potential [46] was used. Such a G matrix may be considered as a realistic effective interaction within a space which is large enough to incorporate collective long-range correlations, but too limited to allow for the effect of short-range correlations to be included [47].

For the single-particle energies ε_{α} a spacing between major shells of 16 MeV for ^{48}Ca and 14 MeV for ^{90}Zr was adopted. Those ε_{α} which are close to the Fermi level were fixed in such a way that energies of the states in adjacent odd nuclei with large spectroscopic factors are reproduced when the Dyson equation (13) is solved. This procedure was discussed in more detail in Ref. [10].

B. Solution of the Bethe-Salpeter equation and calculation of the strength function

For the dressed independent particle approximation (DIPA) of the polarization propagator (15), the response function (14) is obtained immediately with the solution of the Dyson equation (13) for the one-body Green function. In the discrete basis employed here, this one-body Green function is obtained as typically a few thousand poles and residues in Eq. (12) by the methods of Ref. [10]. In this way the response function (14) consists of typically 10^6 contributions, most of which are very small. In the actual calculations the number of poles in the calculated one-body Green function is reduced by grouping together poles in suitably chosen intervals such that the total sum of residues in the intervals is at least 0.01. The energy at which an effective pole appears is determined by the residue weighted average of the pole energies in the interval. For states close to the Fermi level which have their main portion of strength concentrated in one or more large fragments, the one-body Green function which enters Eq. (15) has typically some 20 poles, whereas the more fragmented states further away from the Fermi level are represented by up to 100 poles. Furthermore, the poles in the polarization propagator were provided with an imaginary part of 0.25–0.50 MeV. This not only simulates a certain mixing with more complicated configurations but also yields a continuous strength distribution which can be conveniently plotted. The result thereby becomes more comparable with experimental data, which are always smooth functions of energy due to finite energy resolution. This treatment of $L^f(\omega)$ with a certain energy smearing is also helpful for the calculation of the polarization propagator in the RPA which

is obtained by direct matrix inversion:

$$L(\omega) = [\{L^f(\omega)\}^{-1} - V]^{-1}. \quad (23)$$

Similar to $L^f(\omega)$, also $L(\omega)$ calculated in this way is continuous for real values of ω .

IV. GAMOW-TELLER AND SPIN-ISOSPIN MULTIPOLE STRENGTHS FOR ^{48}Ca

A. Gamow-Teller strength for ^{48}Ca

In the IPISM only two transitions contribute to the Gamow-Teller strength: The $1f_{7/2}$ neutron may transform into a $1f_{7/2}$ proton or into a $1f_{5/2}$ proton. These two possibilities are represented by the two thin dashed peaks in Fig. 3. In Figs. 3–5 as well in Tables I and II, the energy scale is with respect to the ^{48}Ca ground state. To obtain the excitation energy in ^{48}Sc , subtract 0.5 MeV, and for the excitation energy in ^{48}K , a shift of 11.3 MeV is required. The upper part of Fig. 3 illustrates clearly the results of fragmentation of the one-

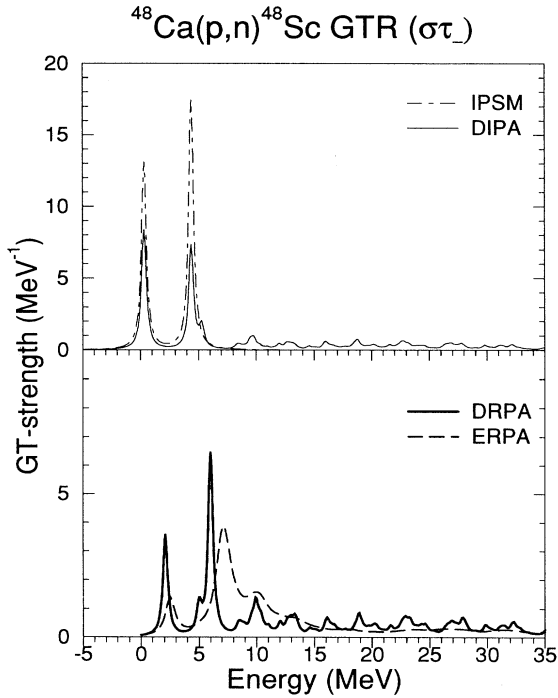


FIG. 3. The $^{48}\text{Ca}(p, n)$ Gamow-Teller response. The upper figure illustrates the fragmentation of strength when independent particles and holes (IPISM) are replaced by dressed particles and holes (DIPA) leading to the polarization propagator (15). With a G -matrix interaction between these dressed particles a further shift and fragmentation occurs (DRPA). The ERPA result of Ref. [28] is shown for comparison, illustrating the effect of induced interactions. Note the factor of 2 difference in scale between the upper and lower figures. The calculated energy is relative to the ^{48}Ca ground state. The ground state of ^{48}Sc is at 0.5 MeV.

body spectral functions, the inclusion of which leads to the thin solid distribution. The fragmentation effect is stronger for the $f_{7/2} \rightarrow f_{5/2}$ transition (the peak at 4.5 MeV) than for the $f_{7/2} \rightarrow f_{7/2}$ transition (the peak at 0.5 MeV), because the $1f_{7/2}$ proton shell is closer to the Fermi level than the $1f_{5/2}$ proton shell. If only this fragmentation effect were present, the excitation strength for the lowest state would be at least comparable to that of the next peak. However, as illustrated in the lower part of Fig. 3, the repulsive interaction causes interference, which brings about a strong reduction of the lowest-energy peak, the strength of which is partly transferred to the next one, now at about 6 MeV. Note that a large amount of strength appears as a wide distribution extending over an energy range of several tens of MeV. Experimentally the GT strength has been studied up to about 20 MeV [8,48]. This has been compared with various calculations [28,31,33] and with the sum rule value $3(N-Z)$. The global distribution obtained in the present calculation is indicated in Table I. First, it should be noted that the total (p, n) strength exceeds the sum rule value of 24 by roughly 25%, which is the size of the total (n, p) strength, also listed in Table I. This is entirely due to ground-state correlations, i.e., to particle-particle and hole-hole strengths. Note that in the ERPA of Ref. [28] such an enhancement does not occur, due to the absence of pp and hh contributions. The pp and hh amplitudes increase the destructive interference at low energy. As a consequence of this, the strength at low energy is less in the DRPA than in the ERPA while at energies above 20 MeV the DRPA strength is much larger than in the ERPA. This is clearly visible in Fig. 3. In the ERPA the strength distributions are smoother than in the DRPA, because the coupling of $1p1h$ to the multitude of $2p2h$ states is better accounted for, as discussed in Sec. IID. The induced interactions, of second order in the G matrix, which are included in the ERPA, also effectively enhance the repulsion in the isovector channel. This results in a stronger upward shift of the main GTR peak in the ERPA as compared to DRPA.

The (n, p) strength has been experimentally investigated for various nuclei [49–52]. In the present calculations it is found distributed over a very large energy range of typically a hundred MeV. Part of this strength distribution is shown in Fig. 4. Note that the dressing of the single-particle propagators is the main origin of the calculated strength, which is only little influenced by

TABLE I. Gamow-Teller strength for ^{48}Ca with, in parentheses, the contribution of pp and hh amplitudes.

	DIPA	DRPA	ERPA [28]	Expt. [8,48]
(p, n) strength				
$E < 20$ MeV	18.2 (1.4)	14.5	17.8	14.4(± 2.4)
$20 < E < 40$ MeV	5.5 (2.7)	6.8	4.3	
$40 < E < 80$ MeV	5.9 (4.5)	7.2	1.8	
$E > 80$ MeV	1.1 (1.1)	1.5	0.1	
(n, p) strength	6.6 (6.6)	5.9	≈ 0	

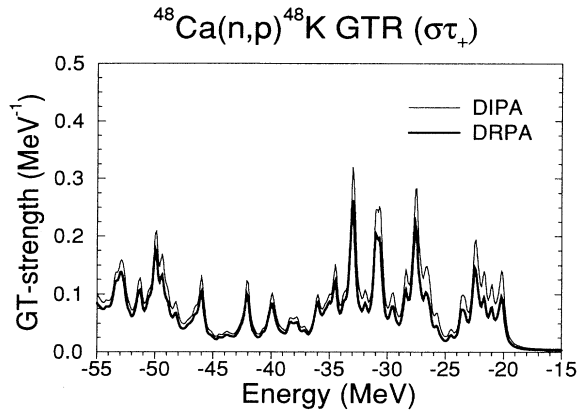


FIG. 4. Calculated $^{48}\text{Ca}(n,p)$ Gamow-Teller strength for the case of a dressed independent particle and hole (DIPA) and with a G matrix interacting between them (DRPA). The calculated energy is relative to the ^{48}Ca ground state. The ground-state energy of ^{48}K corresponds to -11.3 MeV.

the quasiparticle interaction Γ . Since the small sharp peaks will in reality be smeared out by mixing with more complicated configurations, one may conclude from Fig. 4 that this GT strength is roughly a constant function of energy with magnitude ≈ 0.1 MeV $^{-1}$. This amount of strength is comparable with the multipole analysis of the data for ^{40}Ca [52], where a total GT strength of 1.6 ± 0.1 was obtained within the first 15 MeV of the spectrum. Our prediction that much more strength is to be found at higher energies is also in agreement with the calculations of Adachi *et al.* [32].

The enhancement of the total (p,n) strength by pp and hh contributions in the present calculation is also of comparable size as to that found by Nishizaki *et al.* [33], although the method of those authors has been criticized by Van Neck *et al.* [34]. So it seems that, although certain ingredients of the method of Ref. [33] are not correct, the final strength distributions obtained are not unreasonable. However, the amount of strength below 20 MeV is in the present DRPA calculation only 60% of the sum rule, in accord with the data [8,48]. Our result is considerably lower than the 79% obtained in Ref. [33] or the 72% obtained in Ref. [31] with the standard RPA (SRPA) method. Therefore the conclusion [33] that approximately 25% quenching of the strength may be ascribed to the coupling to the Δ_{33} excitation of nucleons is not supported by the present results. To settle this issue it would be relevant to experimentally investigate the considerable (p,n) strength at higher energies as well as the non-negligible (n,p) strength. Another indication might be obtained from a comparison of similar quenching of higher multipole strengths, which are discussed next. In such excitations the Δ_{33} will play a minor role. This argument was put forward earlier in Ref. [30].

As a further theoretical investigation, it is desirable to include induced forces, of higher order in the G matrix, into the effective interaction. In this way also the Baym-Kadanoff prescription (16) should be better obeyed. These induced forces are included, in some ap-

proximation, in the ERPA [26–28] as well as the extended SRPA (ESRPA) [31,33] approaches which reproduce the resonance peak and the strength at lower energies better than the DRPA. However, the strength in the high-energy tails above the resonance peak, as well as the (n,p) spin-dipole strength that will be discussed next, are found to be almost equal in the DIPA and DRPA. So it seems that these are rather insensitive to the approximation made for Γ .

B. Spin-dipole strength for ^{48}Ca

It is well known that in nuclear beta decay also first-forbidden transitions, i.e., those with $\pi_i\pi_f = -1$ and $\Delta J = 0, 1$, or 2, are slower than predicted by the IPSM [53,54]. Figure 5 illustrates how this quenching of decay strength, or in the inverse reaction of excitation strength, is caused both by the fragmentation of the single-particle strength functions and by the repulsive spin-isospin force. The latter produces a destructive interference of transition amplitudes for the lowest states, shifting the strength to higher energies. Note that when only the fragmentation of the single-particle Green functions is included, strong transitions at low energies are still discernible, especially for 2^- and 0^- . When the interaction is switched on, these transition strengths almost disappear and very broad resonances are formed around 20–25 MeV for 0^- and 1^- and at slightly lower energies for 2^- . It should be remarked that in reality more smooth distributions than obtained with the DRPA are expected, as a consequence of coupling to a multitude of more complicated configurations in the giant resonance region. This is illustrated by the ERPA results, which are also displayed. The differences between the ERPA and DRPA results are similar as for the GT strength and have the same origins as discussed there. Note the stronger concentration into a response peak as a consequence of the stronger induced repulsive forces in the ERPA.

Although the maxima of the spin-dipole resonances are clearly below 30 MeV, it should be mentioned that almost half of the total (p,n) strength is calculated above 30 MeV. This feature is illustrated in Table II. This strength at higher energies may to a large extent be attributed to the fragmentation of the single-particle spectral functions. It is further increased by the interaction. Note that for these spin-dipole excitations there is a large amount of (n,p) strength already in the IPSM. As a consequence of this, in the DRPA calculations the (p,n) strength up to 30 MeV is typically 100% of the sum rule value, though it is only 60% of the total (p,n) strength. The strength above 30 MeV is predominantly due to ground state correlations, i.e., to pp and hh amplitudes.

V. CHARGE-EXCHANGE STRENGTHS FOR ^{90}Zr

A. Gamow-Teller (p,n) strength

In the reaction $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ the Gamow-Teller strength is observed [6,44,55] to be split in some low-

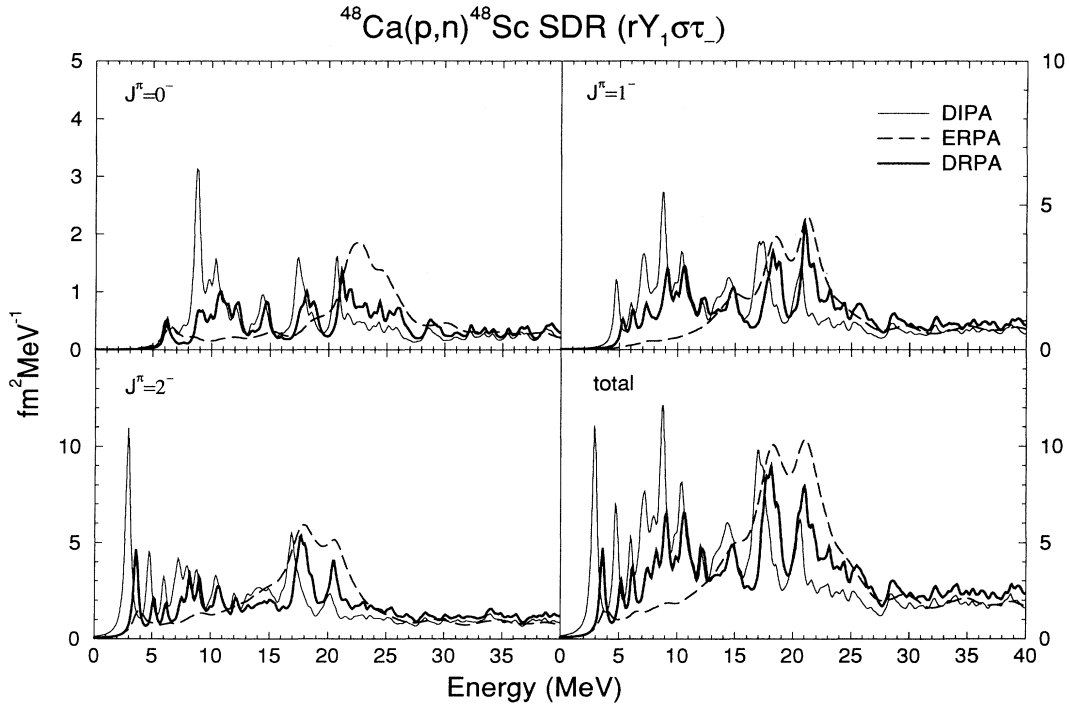


FIG. 5. Calculated $^{48}\text{Ca}(p,n)^{48}\text{Sc}$ SDR for $J^\pi = 0^-$, $J^\pi = 1^-$, $J^\pi = 2^-$, and the total $L=1$ strength. Plotted are the dressed independent particle approximation, Eq. (15) (DIPA), and the result with a G -matrix interaction between these dressed particles and/or holes, Eq. (17) (DRPA). For comparison also the result of the ERPA method [28] is shown. See caption of Fig. 3.

lying $T=4$ states, a giant GT resonance centered at 15.6 MeV ($E_x=8.7$ MeV in ^{90}Nb) with a full width at half maximum (FWHM) of 4.4 MeV, and a $T=5$ component at 20.3 MeV ($E_x=13.4$ MeV in ^{90}Nb) with a FWHM=3.1 MeV. The fraction of the GT sum rule (22) observed is $(61\pm 10)\%$ [56].

In the IPSM there are just two states which make up the GT strength, viz., with a neutron $g_{9/2}$ hole and a proton $g_{9/2}$ or $g_{7/2}$ particle. In the RPA the residual (G -matrix) interaction pushes these states upwards in energy. In [57] it was shown how in the ERPA a further shift of the higher peak and broadening due to coupling to 2p2h states is found. Up to 25 MeV ($E_x=18$ MeV in ^{90}Nb) the calculated integrated (p,n) strength is 99% of the GT sum rule in the RPA and 79% in the ERPA [57]. The results of the present approach are displayed in Fig. 6. In this and the following figures, as well as in Tables III and IV, the energy scale is with respect to the ^{90}Zr ground state. For the excitation energy in ^{90}Nb subtract 6.9 MeV; for excitation energy in ^{90}Y an energy shift of 1.5 MeV is required. The total strength is enhanced by the contribution of pp and hh amplitudes, but in spite of this the strength below 25 MeV is less than in the ERPA. A comparison of the ERPA and DRPA shows again that in the latter the fragmentation of strength is stronger as a consequence of a better treatment of the single-particle propagator. Considerable extra strength is then found above 25 MeV, and the total (p,n) strength is 36% larger than the GT sum rule.

A quantitative indication of the strength distribution is given in Table III. Note that the sum rule for the difference between total (p,n) and (n,p) strengths is very well obeyed by the calculations, in spite of a certain violation of the Baym-Kadanoff self-consistency prescription, Eq. (16) [37,38]. This lends support to the approach followed here, as such a good fulfillment of sum rules is found in all cases.

TABLE II. Charge-exchange spin-dipole strengths of ^{48}Ca (fm^2).

J^π		IPSM	DIPA	DRPA	ERPA [57]
0^-	$(p,n); E < 30$ MeV	21.0	14.8	12.4	14.6
	$E > 30$ MeV	0.0	6.8	9.3	6.2
	(n,p)	9.5	10.2	10.3	9.4
	Sum rule (20)	11.5	11.5	11.5	11.5
1^-	$(p,n); E < 30$ MeV	57	42	34	39
	$E > 30$ MeV	0	19	25	17
	(n,p)	23	26	25	22
	Sum rule (20)	34	34	34	34
2^-	$(p,n); E < 30$ MeV	76	59	52	60
	$E > 30$ MeV	0	27	35	19
	(n,p)	19	29	30	22
	Sum rule (20)	57	57	57	57

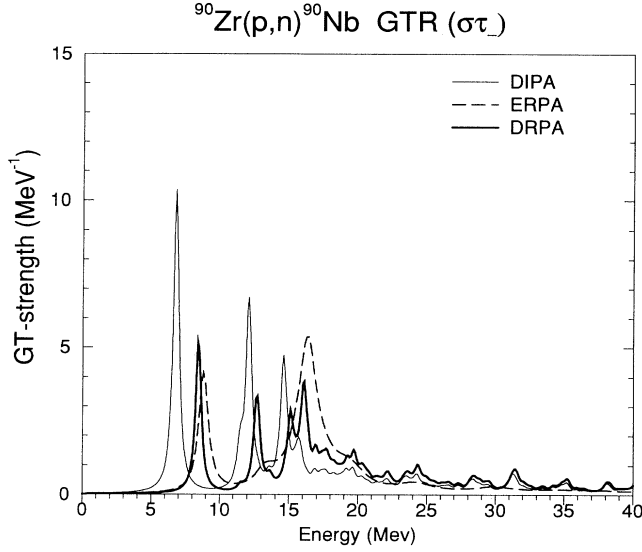


FIG. 6. The $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ Gamow-Teller response, for the case of a dressed independent particle and hole (DIPA) and with a G -matrix interaction between them (DRPA). For comparison, the result of the ERPA method, which includes all diagrams of Fig. 2, is also displayed. The calculated energy is relative to the ^{90}Zr ground state. The ^{90}Nb ground state is at 6.9 MeV.

B. Gamow-Teller (n, p) strength

The $^{90}\text{Zr}(n,p)^{90}\text{Y}$ reaction has been studied extensively [49,58,59]. A multipole decomposition was performed and considerable $L=0$ strength was identified up to 30 MeV excitation energy. This strength was interpreted [49] as a sum of GT and spin-isovector monopole (SIVM) strength. In the analysis [49] of the data it was assumed that the strength below 10 MeV is predomi-

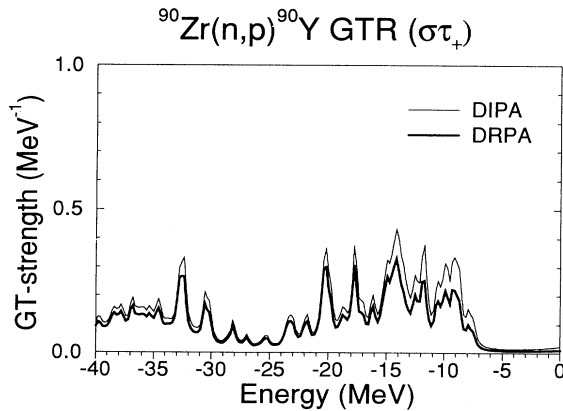


FIG. 7. Calculated $^{90}\text{Zr}(n,p)^{90}\text{Y}$ Gamow-Teller strength for the case of a dressed independent particle and hole (DIPA) and with a G -matrix interaction between them (DRPA). The calculated energy is relative to the ^{90}Zr ground state. The ^{90}Y ground state is at -1.5 MeV.

nantly GT strength whereas that above 10 MeV should be attributed to the SIVM excitation. This assumption was based on various RPA calculations [30,60–62] although in other calculations [63,64] an extension of the GT strength to higher energies was predicted. Indeed, in the present calculations the (n, p) GT strength is predicted to be spread over a larger energy range, essentially up to 100 MeV. These results are displayed in Fig. 7 and in Table III. If one assumes, as the other extreme, that all observed strength up to 20 MeV is merely GT strength, an upper limit of 12% of the GT sum rule is obtained [49,59]. In the present calculations the SIVM strength is mainly located around 20 MeV and upward, and so the main portion of the observed strength up to 20 MeV should indeed be GT strength. In Table III it is shown that the predicted amount of GT (n, p) strength between 20 and 40 MeV is almost as large as that below 20 MeV and that about half of the total GT (n, p) strength is predicted at still higher energies. This roughly constant strength as a function of energy is quite similar to that found for ^{48}Ca .

C. Spin-dipole (p, n) strength

In the reaction $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ the spin-dipole resonance is observed [6,44] around 21 MeV with a large width of about 8 MeV. Pion charge-exchange experiments [65] essentially confirm these data. The DRPA results, shown in Fig. 8, exhibit strength distributions around the same energy, but spread over a much wider range of energies. Possibly the DRPA method overestimates the width of this energy range; inclusion of higher-order terms in the interaction Γ may yield a further concentration of calculated strength. On the other hand, the analysis of the data may be biased by the predictions of RPA calculations, which always yield a concentration in too narrow resonances. It is clear [23] that in the early analysis of the data a large fraction of the widespread strength was subtracted as “background” and as such discarded in the interpretation of the giant resonance spectra. The present calculations indicate clearly that the $L=1$ strength must actually be spread out over a large domain between 10 and 40 MeV, with a long

TABLE III. Gamow-Teller strength for ^{90}Zr in % of the Ikeda sum rule $3(N-Z)$.

	ERPA [57]	DIPA	DRPA	Expt.
(p, n) strength				
$E < 25$ MeV	79	85	70	61(± 10) [56]
$25 < E < 40$ MeV	13	15	19	
$E > 40$ MeV	8	39	45	
(n, p) strength				
$E < 20$ MeV	≈ 0	12	9	< 12 [49,59]
$20 < E < 40$ MeV	≈ 0	6	5	
$E > 40$ MeV	≈ 0	22	21	

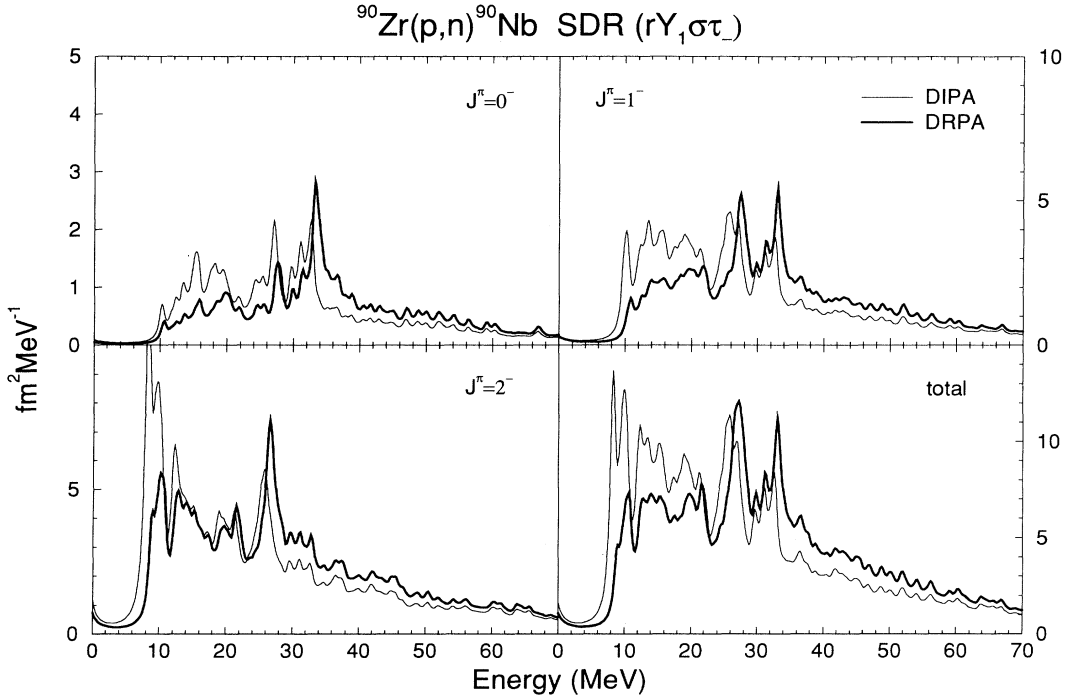


FIG. 8. Calculated $^{90}\text{Zr}(p,n)^{90}\text{Nb}$ spin-dipole response for $J^\pi = 0^-$, $J^\pi = 1^-$, $J^\pi = 2^-$, and the total $L=1$ strength. Results for a dressed independent particle and hole (DIPA) and with a G -matrix interaction between them (DRPA) are displayed. See caption of Fig. 6.

tail extending up to 100 MeV. In Table IV it is shown that the calculated strength in this tail is typically one-fifth of that below 40 MeV in the ERPA, while in the DRPA calculations this strength above 40 MeV becomes about half of that at lower energies. In Table IV also the (n,p) strength is listed. The sum rule (20) was checked by direct calculation of the radial matrix element on the right-hand side and found to be fulfilled within 2% in all cases. Note that for $L \neq 0$ the sum rule value is no helpful indication of the total (p,n) strength, as this strength exceeds the sum rule by 100%. This is not unexpected because the $L=1$ (n,p) strength is always of comparable size as (p,n) , even within the simplest independent particle model.

D. Spin-dipole (n,p) strength

In the multipole decomposition of the (n,p) data [49] the main portion of the $L=1$ strength is found as a broad structure up to about 20 MeV excitation energy with a tail extending beyond 40 MeV. In this analysis it was further assumed that peaks around 5 and 10 MeV represent the spin-dipole resonances, while most of the other $L=1$ strength from 10 MeV upward was attributed to a fitted quasifree background. The results of the present calculations are depicted in Fig. 9. The total $L=1$ strength indeed exhibits two concentrations of strength at about 5 and 12 MeV. The resonance structure seen at about 12 MeV is mainly due to 0^- strength while that at lower

energy is predominantly 1^- and 2^- strength. In the major part of the energy range these strengths of different J values overlap. Moreover, the calculation shows a large “background” tail, extending up to a 100 MeV. In view of this prediction it is likely that a considerable fraction of the fitted quasifree scattering strength in Refs. [49,65] is actually genuine spin-dipole strength. Therefore the spin-multipole strengths will in general be

TABLE IV. Charge-exchange spin-dipole strengths of ^{90}Zr (fm^2).

J^π		DIPA	DRPA	ERPA
				[57]
0^-	$(p,n); E < 40$ MeV	31	26	32
	$E > 40$ MeV	12	17	9
	$(n,p); E < 40$ MeV	19	18	20
	$E > 40$ MeV	5	5	1
	Sum rule (20)	20	20	20
1^-	$(p,n); E < 40$ MeV	88	76	91
	$E > 40$ MeV	32	42	20
	$(n,p); E < 40$ MeV	47	46	50
	$E > 40$ MeV	13	13	3
	Sum rule (20)	59	59	59
2^-	$(p,n); E < 40$ MeV	126	115	134
	$E > 40$ MeV	44	54	19
	$(n,p); E < 40$ MeV	51	51	52
	$E > 40$ MeV	20	20	3
	Sum rule (20)	97	98	98

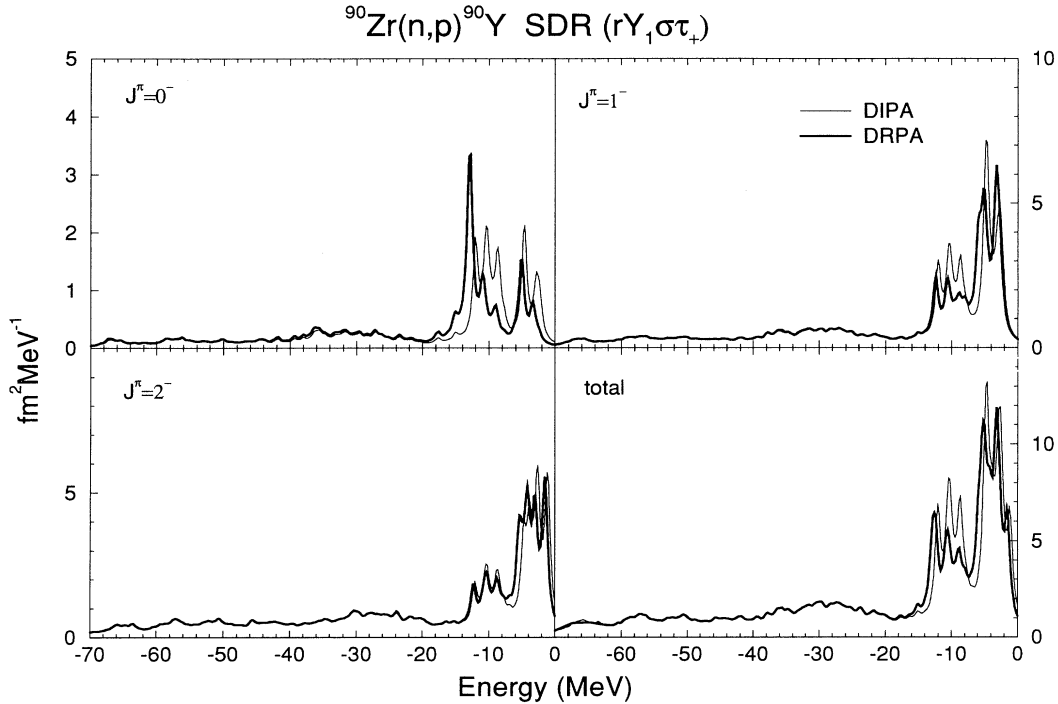


FIG. 9. Calculated $^{90}\text{Zr}(n,p)$ spin-dipole response for $J^\pi = 0^-$, $J^\pi = 1^-$, $J^\pi = 2^-$, and the total $L=1$ strength. See caption of Figs. 7 and 8.

underestimated when the data are analyzed as a sum of resonant peaks, with relatively small widths that are comparable to those of standard RPA calculations, and of a nonresonant quasifree background.

VI. SUMMARY AND CONCLUSIONS

The spin-isospin excitation strengths are computed by taking the fragmentation of single-particle strength as a starting point. For this purpose dressed single-particle propagators are obtained by solving the Dyson equation with self-energy up to second order in a G -matrix interaction. In this way one accounts qualitatively for the fragmentation of hole strength observed in recent $(e, e'p)$ experiments [11,15]. The fragmentation of single-particle strength is reflected in a similar fragmentation of the spin-isospin response, calculated with a polarization propagator which is just the product of independent, but dressed, single-particle propagators. This dressed independent particle approximation (DIPA) already yields excitation strength distributions which are quenched considerably at low energy in comparison with those in an independent particle shell model. A large fraction, several tens of percent, of the total strength is then found in large tails above the typical resonance energy region and extending up to a 100 MeV.

As a next step, the interaction is dealt with in a similar approximation as the RPA, but here it acts between the dressed single-particle propagators (DRPA). The already fragmented excitation strength is then pushed upward in

energy, in a similar way as this happens in the standard particle-hole RPA with a repulsive isovector particle-hole interaction. Although some resonance structure remains discernible, the strength distributions become broad with large tails at higher energies. In these calculations the consistency relation between self-energy and interaction [37,38], which is required by conservation laws, is not exactly fulfilled, but it turns out that sum rules for the total strength are very well fulfilled in all cases by the results. This supports the validity and quality of the present approach. On the other hand, it may be remarked that the strength distribution at low energies up to and including the resonance peak is not as well reproduced as in the ERPA or other approaches [31,33] which include induced forces in the effective interaction. However, the high-energy tails of the distributions are found to be rather insensitive to the effective interaction Γ . This part of the strength distribution seems in better agreement with the data [6,49,58] within the present approach than for earlier calculations.

In the DRPA formalism, particle-particle (pp) and hole-hole (hh) amplitudes, which reflect the importance of ground-state correlations in closed-shell nuclei, are treated on equal footing with particle-hole (ph) amplitudes. These pp and hh amplitudes enhance the total (p, n) Gamow-Teller (GT) strength by 25% in the calculation for ^{48}Ca and by 36% for ^{90}Zr . This confirms the calculations of Ref. [33]. It should be noted, however, that the strength within the experimentally analyzed energy domain is not enhanced, but rather decreased by the pp and hh amplitudes because of their destructive inter-

ference with ph amplitudes. The excess of strength is therefore mainly found above the resonance region. This role of pp and hh amplitudes is illustrated for the GT and spin-dipole charge-exchange strength for ^{48}Ca , by a comparison of DRPA results with those of the extended RPA (ERPA) approach [26–28]. In the latter the pp and hh amplitudes are not treated on equal footing with ph contributions, but discarded [28] or added as a perturbation [26,27]. On the other hand, the coupling of $1p1h$ to $2p2h$ states is treated more fully by the dynamical interaction in the ERPA, which leads to a smoother shape of the strength distributions and more pronounced resonance peaks. The DRPA and ERPA approaches may therefore be considered as complementary. As a further extension, the virtues of each could be combined by including a higher-order dynamical interaction within the here presented DRPA framework.

The large amounts of strength that are predicted in the high-energy tails of the distributions make it difficult to experimentally verify the GT sum rule with high accuracy. The unobserved strength in the tail of the (p, n) cross section is not necessarily compensated by the same amount of unobserved strength in the tail of the (n, p) cross section [66]. The (p, n) vs (n, p) asymmetry of the tails originates only partly from a (small) difference in occupation between proton and neutron orbits that are remote from the Fermi level. The main cause is the asymmetry of occupation around the Fermi level. The spectral functions of these valence orbits have long tails and thereby the proton-neutron asymmetry around the Fermi level is still contributing in the response at high energies. On the basis of a comparison of the DRPA calculations with the strengths deduced from the data [6,8,49,65] and keeping in mind that short-range correlations may reduce the calculated strengths by a 10–15% [47], there is no reason to invoke a substantial quenching of GT strength by coupling to the Δ_{33} resonance. The issue could be further clarified by an experimental determination of both spectral functions and response at higher energies.

The analysis of the experimental (p, n) as well as (n, p)

data has usually been performed by fitting the total cross section as a sum of resonance peaks with widths typically less than 10 MeV and a “background,” attributed to quasifree scattering [49,65]. The present calculations indicate that a substantial fraction of this “background” is genuine GT or spin-dipole strength. Of the observed small peaks on the top of the broad $L=1$ distributions, the lowest in energy should be attributed to 2^- strength. This has its maximum typically about a 5 MeV below that of the 0^- and 1^- distributions. This difference still reflects the underlying shell structure.

Summarizing, a formalism is applied that emphasizes the connection between excitation strength and spectral functions. The observed fragmentation of one-body addition and removal strengths implies a corresponding broad fragmentation of excitation strength, which is further redistributed by the effective quasiparticle interaction. It also implies the importance of pp and hh amplitudes, which increase the total excitation strength, although not within the experimentally investigated relatively low-energy domains. There is no reason to invoke a substantial quenching of GT strength by coupling to the Δ_{33} resonance. Inclusion of contributions of higher order in the G matrix into the effective interaction Γ is expected to improve the strength distribution at lower energies up to and including the resonance peaks. Such an extension is needed to reach more definite conclusions. Work along this line is in progress.

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