COLLECTIVE CORRELATIONS IN C^{12}†

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The strong coupling of the giant resonance to the surface vibrations in C^{12} results in the splitting of the single one-particle, one-hole, 1^- collective state into several components, thus improving the agreement between theory and experiment to a very large extent.

It is well known that pure single-particle transitions do not account for collective nuclear states. Elliot and Flowers[1] and others[2-3] showed that the inclusion of residual nuclear forces is necessary for the understanding of the excitation energy and dipole strength of the giant resonance. A number of detailed one-particle, one-hole (1p-1h) investigations followed.[4-7] For example, in the case of C^{12} one collective 1^- state has been obtained. However, the experimental photoabsorption cross section shows evidence for three to four major peaks in the region from 18 to 30 MeV which cannot be explained by simple 1p-1h calculations. On the other hand, the collective model for giant resonances explains such structures in a very natural way by the coupling of giant-dipole and surface-quadrupole degrees of freedom, as has been shown recently.[8-10]

It is our aim to demonstrate in this Letter the importance of such collective correlations in light nuclei. We therefore, firstly, formulate the idea of collective correlations in the particle-hole framework and, secondly, give results obtained for C^{12}.

The Hamiltonian of the dynamic collective theory is of the form[9]

$$H = H_D + H_Q + H_{DQ}, \tag{1}$$

where the first term, H_D, describes the giant resonances; H_Q describes the surface oscillations; and H_{DQ} is the interaction between the giant resonances and the surface oscillations. If H_{DQ} is neglected, the solution of (1) yields only one dipole state carrying all strength, and many states of the form one dipole phonon–many quadrupole phonons carrying no strength. If, however, H_{DQ} is taken into account, all these states mix and, as a result, several states with appreciable dipole strength occur.

The situation is similar in the particle-hole model, where the diagonalization of the residual force with 1p-1h states usually gives one collective dipole state. What is missing in such 1p-1h calculations is the interaction with the collective surface degrees of freedom (surface phonons) contained in $H_Q + H_{DQ}$.

The surface phonons might be thought of as complicated, coherent superpositions of 1p-1h, 2p-2h, 3p-3h, etc., configurations coupled to 0^+ and 2^+ in the particle-hole framework. Of course, if the complete particle-hole Hamiltonian, H_{ph}, is diagonalized in the full Hilbert space consisting of 1p-1h, 2p-2h, 3p-3h, etc., configurations, one obtains in principle an exact solution of the nuclear problem. However, such a procedure is both nontransparent and impossible to carry through because of the tremendous number of many-particle, many-hole configurations. Therefore, a more physical approach to the problem is necessary.

Guided by the collective Hamiltonian (1), we assume that from all many-particle, many-hole configurations only a few, namely, the collective ones describing surface modes, are essential for the structure of the giant resonances. We therefore propose the following Hamiltonian:

$$H = H_{\mathrm{ph}}^{(1)} + H_Q + H_{DQ}, \tag{2}$$

where $H_{\mathrm{ph}}^{(1)}$ shall describe the particle-hole Hamiltonian in the 1p-1h, 1^- subspace, and

$$H_Q = \tfrac{1}{2}\sqrt{5}B_2[\dot{\alpha}^{[2]}\times\dot{\alpha}^{[2]}]^{[0]} + \tfrac{1}{2}\sqrt{5}C_2[\alpha^{[2]}\times\alpha^{[2]}]^{[0]},$$

$$H_{DQ} = K_1[\alpha^{[2]}\times[\alpha^{[1]}\times\alpha^{[1]}]^{[2]}]^{[0]} + K_{20}[[\alpha^{[2]}\times\alpha^{[2]}]^{[0]}\times[\alpha^{[2]}\times\alpha^{[2]}]^{[0]}]^{[0]} + K_{22}[[\alpha^{[2]}\times\alpha^{[2]}]^{[2]}\times[\alpha^{[2]}\times\alpha^{[2]}]^{[2]}]^{[0]},$$

$$K_1 = -1.588C_1,\quad K_{20} = -0.708C_1,\quad K_{22} = -0.936C_1. \tag{3}$$

C_1 is calculated in the collective model.[10] $\alpha^{[2]}$ and $\alpha^{[1]}$ are the usual collective coordinates for surface quadrupoles and giant dipoles, respectively. The latter ones are expressed in terms of particle-hole operators by the require-

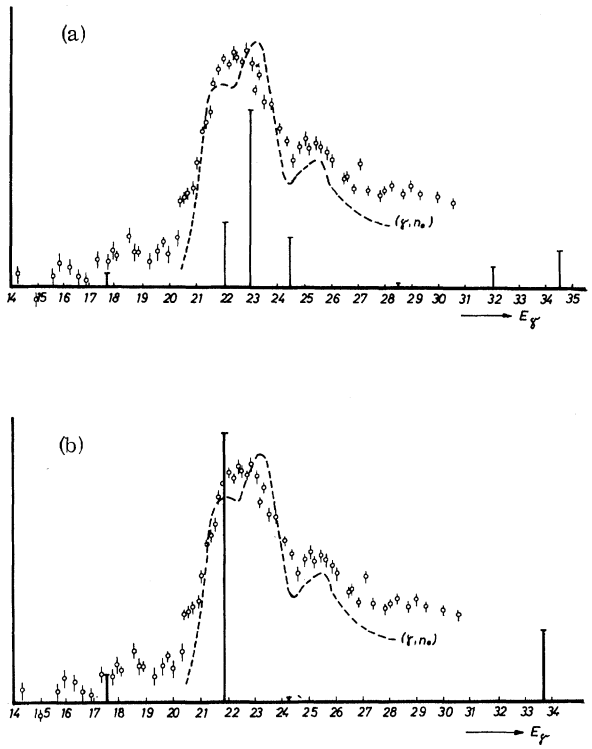

FIG. 1. Photoabsorption cross section of C^{12}. The experimental points give the total photoabsorption cross section [V. Shevchenko and N. Yudin, At. Energy Rev. 3, 3 (1965)]; the dashed line gives the (γ,n_0) cross section at 90° in arbitrary units [M. Danos and E. G. Fuller, Ann. Rev. Nucl. Sci. 15, 29 (1965); V. V. Verbinski and J. C. Courtney, Nucl. Phys. 73, 398 (1965)]. (a) The dipole strengths calculated with collective correlations, (b) with usual particle-hole calculations.

ment that the dipole operator has to be the same in the collective and particle-hole language:

$$D_{\mathrm{ph}}^{[1]} = D_{\mathrm{coll}}^{[1]}, \tag{4}$$

where

$$D_{\mathrm{ph}}^{[1]} = \left(\frac{4\pi}{3}\right)^{1/2} \sum_{\mu\nu} \langle \mu | \tau_{3/2} \cdot r Y^{[1]} | \nu \rangle a_{\mu}^{+} a_{\nu}, \tag{5}$$

and

$$D_{\mathrm{coll}}^{[1]} = M_0 (\alpha^{[1]} + M_1 [\alpha^{[1]} \times \alpha^{[2]}]^{[1]}).$$

M_0 and M_1 are constants.[9] Substituting $\alpha^{[1]}$ in H_{DQ} (3), we finally obtain[11]

$$\begin{aligned} H = {} & H_{\mathrm{ph}}^{(1)} + H_Q + \kappa_1 [\alpha^{[2]} \times [D_{\mathrm{ph}}^{[1]} \times D_{\mathrm{ph}}^{[1]}]^{[2]}]^{[0]} \\ & + \kappa_{20} [[\alpha^{[2]} \times \alpha^{[2]}]^{[0]} \times [D_{\mathrm{ph}}^{[1]} \times D_{\mathrm{ph}}^{[1]}]^{[0]}]^{[0]} \\ & + \kappa_{22} [[\alpha^{[2]} \times \alpha^{[2]}]^{[2]} \times [D_{\mathrm{ph}}^{[1]} \times D_{\mathrm{ph}}^{[1]}]^{[2]}]^{[0]}, \end{aligned} \tag{6}$$

where the coupling parameters are

$$\kappa_1 = -64K/(AR_0^2),$$

$$\kappa_{20} = -127.5K/(AR_0^2),$$

$$\kappa_{22} = -169K/(AR_0^2).$$

K is the symmetry energy of the Bethe-Weizsäcker mass formula, A is the atomic number, and R_0 is the equivalence radius. The last three terms in (6) describe the collective correlations between the 1p-1h states and the surface vibrations. They act like an additional interaction between the various 1p-1h, 1^- states.

The solutions of (6) were obtained by diagonalization in the basis consisting of products of 1p-1h states and phonons. The parameters $\hbar\omega_2 = \hbar(C_2/B_2)^{1/2}$ and $B_0 = (5\hbar\omega_2/2C_2)^{1/2}$ are taken from the low-energy spectrum of C^{12}.[12] The residual force of Gillet[4] has been used with a strength $V_0 = -35$ MeV. The results are shown in the upper part of Fig. 1.[13] For comparison we also have plotted the results of usual 1p-1h calculations. The nearly quantitative agreement of the three major peaks occurring at 22.1, 22.9, and 24.5 MeV with the experimental photoabsorption cross section demonstrates the importance of collective correlations in C^{12}. Similar results are expected for other nuclei.

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[11]Note that the collective correlations are completely different from the interaction of particle-hole states

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[12]Recent electron scattering experiments give $B(E2) = 44$ fm^4 and a transition radius of about 3.3 fm; thus β_0 is about 0.43. F. Gudden, private communication.

[13]The large vibrational amplitudes suggest that anharmonic terms might be of some importance. In fact, it is expected that they will improve the agreement between theory and experiment.

DOUBLE POLES AND NONEXPONENTIAL DECAYS IN ATOMIC PHYSICS*

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Atomic resonance fluorescence is examined as a scattering problem, and Goldberger-Watson double poles are found to occur in situations producible in the laboratory. Systematic study of this effect and associated nonexponential decay thus appears possible.

It was pointed out by Goldberger and Watson[1] that the evidence supporting purely exponential decay for every unstable particle is "far from convincing." These authors showed, in particular, that when the S matrix has multiple poles, the decay amplitude for the associated state becomes a polynomial in time multiplied by the usual exponential factor. Since then, several model situations have been discussed where double poles can occur.[2] In the present note physical situations with double poles are presented that can be produced and studied in the laboratory.

The transition matrix, or T matrix, under discussion is that for resonance scattering of light through more than one excited state. The calculational method and some results have been discussed elsewhere[3,4]; we now extend these to the case where the T matrix has a double pole. This approach[3] is similar to that of the Lee model[5] in that the calculation begins from a second-quantized Hamiltonian in which the unperturbed part gives the atomic energy levels and photon energies, and the interaction part describes transitions between the levels with photon emission and absorption. That such an approach is appropriate had been noted earlier also by Källén.[6]

Resonance flourescence through two excited levels is probably the case of most immediate experimental interest. This process is described by a T matrix of the following form (subscripts 1 and 2 refer to the two excited states; a single ground state is assumed, and the incident photon energy is ω):

$$T = \frac{g_1(\omega-\omega_2+i\tfrac{1}{2}\Gamma_2)f_1 + g_2(\omega-\omega_1+i\tfrac{1}{2}\Gamma_1)f_2 + g_2V_{21}f_1 + g_1V_{12}f_2}{(\omega-\omega_+)(\omega-\omega_-)}, \tag{1}$$

where

$$\omega_\pm = \tfrac{1}{2}(\omega_1+\omega_2-i\tfrac{1}{2}\Gamma_1-i\tfrac{1}{2}\Gamma_2) \pm \tfrac{1}{2}[(\omega_1-\omega_2-i\tfrac{1}{2}\Gamma_1+i\tfrac{1}{2}\Gamma_2)^2+4V_{12}V_{21}]^{1/2}, \tag{2}$$

ω_j is the energy ($\hbar=1$) of the excited state $|j\rangle$, Γ_j is the corresponding width, and V_{ij} is the matrix element of an external (or possibly internal) perturbation coupling the excited states. The f_i and g_i, in Franken's[7] notation, are abbreviations for the absorption and emission matrix elements, respectively, connecting the excited state $|i\rangle$ with the ground state.

A double pole in T requires $\omega_+ = \omega_-$. This