On the nature of the resonances observed in photonuclear reactions

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(Submitted 19 March 1985)
Yad. Fiz. **43**, 843–853 (April 1986)

The continuum states of the light nuclei ^{6}He, ^{6}Li, ^{7}Li, and ^{7}Be are investigated in a microscopic approach that takes account of the dynamics of the cluster and quadrupole collective degrees of freedom. It is shown that the interaction of these nuclei with electromagnetic radiation leads to the excitation of collective resonances with energies E higher than 20 MeV and widths Γ smaller than 1 MeV and giant quadrupole resonances with E in the range from 12 to 15 MeV and $\Gamma \sim 5$ MeV.

1. INTRODUCTION

The dipole, monopole, and quadrupole giant resonances occupy a special place among the resonance states of nuclei excited by electromagnetic radiation or charged particles. Interest in them was aroused by Migdal's prediction[1] of the existence of the giant dipole resonance, which stems from the collective oscillations of the protons, and the subsequent experimental discovery[2] of various manifestations of this resonance. Since then a large number of experimental and theoretical papers have been published in which the properties of the dipole, monopole, and quadrupole resonances are discussed.

The main problem of the theory of the giant resonances is to explain their nature and find out the principal factors determining their positions and widths. To what extent have the published theoretical models for the giant resonances been able to cope with this problem?

Two different approaches, which we shall conditionally call below single-particle and collective approaches, are used in the investigation of the giant multipole resonances. In the single-particle approach, which is possible in the random phase approximation (RPA) or in the Tamm–Dancoff approximation, the collective excitation of a given multipole order is a coherent superposition of a definite number of particle-like and hole-like excitations. The coefficients of this superposition are determined by the interaction in the particle–hole channel. We judge the degree of collectivization of the excitations in this approach from the electromagnetic-transition probabilities relating a given state to the ground state: the higher the transition probability is, the more strongly the level is collectivized. The position of a giant resonance in the single-particle approach is determined by the transition-probability spread over an energy range. The question of which modes of the collective motion form the giant resonance states remain open in this case. With the aid of an averaging procedure (see, for example, Ref. 3), we are able to reproduce fairly well the (γ,n)- and (γ,p)-photodisintegration cross sections. For a more accurate determination of the position and width of a giant resonance, use has been made in the last few years of the continuum states, which allow us to compute the width due to the emission of a nucleon from the nucleus. But this width turns out to be much smaller than the experimental width. To eliminate this discrepancy in the single-particle approach, account is taken of the excitations of the $2p \rightarrow 2h$ type, etc. The role of these states and other problems connected with the determination of the width of a giant resonance are dealt with in detail in Refs. 3–6.

In the collective approach we use a translationally invariant basis of many-particle functions that reproduce certain nucleonic correlations. This approach, in contrast to the single-particle approach, involves the use of those modes of the collective nucleon motion (the monopole, dipole, quadrupole, etc., modes) which form the giant resonance, and are excited as a result of the interaction of photons and charged particles with a nucleus. As is well known, the RPA is valid for systems executing small-amplitude collective oscillations. However, there are no such limitations in the variuos variants of the collective approach (which are discussed below).

The microscopic theory of the collective excitations of nucleon systems,[7-10] based on the idea that the collective modes of the motions play the dominant role in the formation of the giant resonances, correctly predicts the positions of these resonances, as well as the relatively large photodisintegration cross sections, but turns out to be incapable of reproducing the observed resonance widths.

All the known giant resonances lie in the continuous region of the spectrum of nucleon systems, high above the threshold for the disintegration of these systems via the various channels. Therefore, in order to compute the giant-resonance widths on the basis of the microscopic theory of collective excitations, we must include in the analysis, along with the collective modes, the modes connected with the decay of the nucleon system via the open channels. Thus far this step has not been taken, which clarifies the difficulty of the microscopic theory of collective excitations in accounting for the observed giant-resonance widths.

The main purpose of our paper was to carry out a consistent analysis of the collective degrees of freedom of a nucleus and the degrees of freedom of those channels via which the disintegration of the nucleus can occur. We limited ourselves to the consideration of p-shell nuclei and quadrupole resonances so as to keep the actual calculations sufficiently simple. As a result, we were able to compute for ^{6}He, ^{6}Li, and ^{7}Li the effective photodisintegration cross sections in the various partial states, as well as the corresponding cross sec-

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tions for radiative capture, and demonstrate that the large width of the giant quadrupole resonance of these nuclei is a direct consequence of the strong coupling between the collective degrees of freedom and the cluster degrees of freedom of the open disintegration channels for the nuclei in question. Finally, the approach developed by us leads to the conclusion that there exist narrow quadrupole resonances with energies of several tens of MeV. At such high energies of excitation of the collective degrees of freedom the coupling between these degrees of freedom and the cluster degrees of freedom is weak, and therefore the decay of the collective excitations via the open cluster channels is suppressed.

In Sec. 2 we indicate the principal steps in the development of the microscopic theory of collective excitations of nuclei. Then in Sec. 3 we present a new method of investigating the continuum states of a nucleon system. This method takes account of the interaction between the collective and cluster degrees of freedom of a nucleus, and allows us to solve the giant-resonance-width problem. In Sec. 4 we present and discuss the results obtained in numerical computations, and compare the theoretical results obtained by us with the known experimental data.

2. ON THE NATURE OF THE GIANT-RESONANCE WIDTHS

The presently well known work by Griffin and Wheeler[8] became the first stage in the development of a microscopic theory of collective excitations of nuclei. Griffin and Wheeler uncovered a wealth of resources available to the generator-coordinate method for the solution of the problems of the microscopic theory. As an example, they compute the position of the giant monopole resonance of the ^{16}O nucleus. We must, in investigating the quadrupole-mode excitations, supplement the generator-coordinate method by the procedure for projecting[11] a wave packet onto a state with a definite angular-momentum value. This has been done successfully by Abgrall, Caurier, and Bourotte-Bilwes. In the papers cited in Ref. 9 they compute the excitation energies of the various quadrupole modes of the magic nuclei ^{4}He, ^{16}O, and ^{40}Ca and of the ^{20}Ne nucleus.

A new independent line of development of the microscopic theory of collective excitations of nuclei was conceived, beginning with Simonov's work.[12] In place of the Hill–Wheeler equation[8] for the collective modes, dynamical equations in ordinary coordinate space were derived first for the monopole mode[13] and then for the quadrupole modes,[7,13,14] and it became clear that the interaction of the collective modes with the internal-motion modes of a nucleon system can be taken into account, using a basis of generalized hyperspherical functions.[7] The total number of quadrupole degrees of freedom responsible for the giant quadrupole resonance is, according to the microscopic theory, equal in the general case to nine,[7] and not five, as follows from the Bohr–Mottelson phenomenological model.

Moshinsky and Quesne's work[15] on the symplectic groups facilitated the bringing together of the two orginally independent lines of development of the microscopic theory. As Smirnov *et al.*[16] have shown, this work by Moshinsky and Quesne contains a natural—from the standpoint of the method of generalized hyperspherical functions[7]—classification of the basis states of teh multidimensional harmonic oscillator that form the quadrupole and monopole excitations of nuclei. Later on, Rosensteel and Rowe[17] proposed a symplectic model of the nucleus, and an Sp(2,*R*) symplectic model was realized for the first time for the light even-even nuclei in a series of papers by Arickx *et al.*[10,18] Then Smirnov, Okhrimenko, and three of the present authors (G. F. F., V. S. V., and L. L. C.), by expanding the nuclear wave function in a generalized Fourier series in terms of the complete oscillator basis of the Sp(6,*R*) or Sp(2,*R*) model, were able to reduce the dynamical equations of the microscopic theory of collective excitations to a system of algebraic equations.[19–21] In the process it was found that the matrix elements of the various operators used in the computation of the collective -excitation spectra of nuclei within the framework of the generator-coordinate method are the generating invariants of the matrix elements of the same operators, but in the basis functions of the symplectic model. This completed the unification of the various approaches—the generator-coordinate method, the method of generalized hyperspherical functions, and the algebraic methods assumed as the bases of the symplectic models—which had for a long time been developed independently.

Finally, it was recently found that the many-particle oscillator basis can equally well be used to investigate both the stationary and the continuum states of nucleon systems.[22,23] This enables us to expand the standard basis for the symplectic models, and include in it basis functions corresponding to the modes of the motions in the various types of cluster channels for nuclei.[24] In essence, this step should be regarded as the use of the ideas of the resonating group method for a more complete investigation of the properties of collective excitations and the solution at a new level of the giant-resonance width problem on the basis of the microscopic theory of collective excitations.

It is evident that only the excitations of the closed-channel modes, excitations which are, moreover, weakly coupled to the open-channel modes, can be narrow resonances. At relatively high energies only those channels in which the excitation energy gets to be uniformly distributed among all the nucleons, i.e., the channels of the collective degrees of freedom, are closed channels. But if the amplitude of the collective oscillations is not too large, the resonances of these channels can also be quite broad (as obtains in the case of the giant resonances), since under conditions of small-amplitude collective oscillations the collective-mode channels and the open nuclear-disintegration channels strongly overlap.

The elucidation of the nature of the giant-resonance widths on the basis of the microscopic theory of collective excitations amounts to the following. As a result of the interaction of the nuclei with the electromagnetic radiation or the charged particles (electrons, protons, etc.), the energy of the γ rays or charged particles is directly transferred to the collective (dipole, quadrupole, or monopole) degrees of freedom, resulting in the excitation of collective oscillations. The energy necessary for the excitation of specific collective oscillations can be estimated if the transfer of excitation energy from the collective oscillations to all the other modes of the nucleon system is ignored within the framework of the microscopic theory. This procedure, which has been repeatedly used for the monopole and quadrupole modes,[7,9,10,18,19,21,25] gives the correct positions of the monopole and quadrupole resonances, but ignores the problem of their width. The dissipation of the giant resonances is a consequence of the fact that the nucleus, having absorbed the energy transferred to it, goes over into a state lying high above the threshold for its disintegration via the various channels. The collective-mode channel is closed. Therefore, so long as the excitation energy remains localized on the collective degrees of freedom, the nucleus does not disintegrate. But this cannot last for long. The collective-degrees-of-freedom channel is coupled to other channels that are open at that excitation energy which has been absorbed by the collective channel. Moreover, this coupling is strong, and this ultimately establishes the large width of the giant resonances: the excitation energy of the nucleus is quickly transferred from the collective modes to the open-channel modes, and the nucleus immediately disintegrates.

This is roughly the interpretation of the giant resonances and their widths. It is corroborated by calculations carried out for the light nuclei. The principal, experimentally verifiable result of the microscopic theory is the conclusion that the giant resonance is a single state having a large widht. There may (it is not obligatory) occur in the general background of this state the peaks of a small number of narrow resonances, which then account for a certain portion of the sum rule. The possible origin of the narrow resonances has already been indicated above.

3. METHOD OF COMPUTING THE CONTINUUM WAVE FUNCTIONS

We shall consider only one cluster channel along with the collective isoscalar quadrupole excitations (the so-called longitudinal collective oscillations, *i.e.*, the oscillations along the direction made distinct by the vortex mode[21]). The simultaneous consideration of a large number of cluster channels is, in principle, possible, but it is expedient at the first stage to limit ourselves to the simplest situation so as to fully identify the effects stemming from the transfer of the excitation energy from the collective to the open-channel modes. The qualitative role of the neglected cluster channels is apparent. They are capable only of speeding up the dissipation of the excitation energy, which is initially concentrated only on the collective degrees of freedom. As to the collective isovector excitations, which are also ignored by us, they can form another giant resonance, the determination of whose position and width within the framework of the microscopic theory of collective excitations requires special calculations.

The results discussed below are, naturally, sensitive to some extent to the nucleon—nucleon potential chosen by us, since the microscopic theory essentially establishes the dependence of the position and width of the giant resonances on the parameters of the nucleon—nucleon potential. At the same time, it is well known that all the semirealistic nucleon–nucleon potentials used at present in the microscopic theory of light nuclei lead to at least qualitatively identical results. This gives us reason to limit ourselves in a first investigation to the consideration of one nucleon–nucleon interaction potential. As this potential, we chose the Brink-Boeker potential.

Finally, another nonobligatory simplification: In the computation of the photodisintegration and radiative-capture cross sections we used the long-wavelength approximation, which becomes invalid as the energy of the continuum states involved in the computation increases. We propose to investigate the degree of distortion introduced by the long-wave approximation in the course of the analysis of the continuum–continuum transitions.

Let us, for the purpose of carrying out the planned program, represent the nuclear wave function in the form of an expansion in terms of two sets of many-particle oscillator functions:

$$\Psi=\sum_{\nu} C_{\nu}^{\rm col}|\nu,{\rm col}\rangle+\sum_{n} C_{n}^{\rm cl}\,|n,{\rm cl}\rangle. \qquad (1)$$

The first set of functions $\{|\nu,\mathrm{col}\rangle\}$ form the basis of an irreducible representation of the symplectic group $Sp(2,R)$. It is suited for the description of the collective quadrupole excitations of nuclei. If we limit ourselves to only this basis, we obtain the microscopic $Sp(2,R)$ model first realized by Arickx *et al.*[10,18] The second set of functions $\{|n,\mathrm{cl}\rangle\}$ form a basis of oscillator functions for the cluster model, or, in other words, of the resonating group method (RGM), which reproduces the fragment (cluster) motion in the considered channel. The coefficients $\{C_{\nu}^{\rm col}, C_{n}^{\rm cl}\}$ of expansion of the wave function Ψ in terms of the basis functions $\{|\nu,\mathrm{col}\rangle, |n,\mathrm{col}\}$ satisfy the system of linear algebraic equations

$$\sum_{\nu'}[\langle\nu,{\rm col}\,|\hat H|\,\nu',{\rm col}\rangle-E\delta_{\nu\nu'}]\,C_{\nu'}^{\rm col}+\sum_{n'}\,[\langle\nu,{\rm col}\,|\hat H|\,n',{\rm cl}\rangle-E\delta_{\nu n'}\lambda_{n'}]\,C_{n'}^{\rm cl}=0,$$

$$\sum_{\nu'}\,[\langle\, n,{\rm cl}\,|\hat H|\,\nu',{\rm col}\,\rangle-E\delta_{n\nu'}\lambda_{\nu'}]\,C_{\nu'}^{\rm col}+\sum_{n'}[\langle\, n,{\rm cl}\,|\hat H|\,n',{\rm cl}\,\rangle-E\delta_{nn'}]\,C_{n'}^{\rm cl}=0. \qquad (2)$$

Each of the set of functions $\{|\nu,\mathrm{col}\rangle\}$ and $\{|n,\mathrm{cl}\rangle\}$ used in an orthonormalized set, i.e.,

$$\langle\nu,{\rm col}|\nu',{\rm col}\rangle=\delta_{\nu\nu'},\quad \langle\, n,{\rm cl}\,|\,n',{\rm cl}\,\rangle=\delta_{nn'}.$$

But these sets are not biorthogonal. This means that the overlap integrals $\langle\nu,\mathrm{col}|n,\mathrm{cl}\rangle$ are nonzero. More precisely, they are not equal to zero when the functions $|\nu,\mathrm{col}\rangle$ and $|n,\mathrm{cl}\rangle$ belong to the same oscillator shell:

$$\langle\nu,{\rm col}|\,n,{\rm cl}\,\rangle=\lambda_n\delta_{\nu n}. \qquad (3)$$

The values of λ_n determine the coupling of one basis to the other, or, in other words, the coupling of the collective mode to the cluster mode. As the calculations show, the overlap integral $\lambda_0\equiv\langle 0,\mathrm{col}|0,\mathrm{cl}\rangle$ is equal to unity (i.e., the first functions of the two bases coincide), but as n increases, λ_n decreases monotonically, tending to zero as $n^{-(A-2)/2}$. From this it follows that the smaller the amplitude of the collective oscillations is, the more strongly they are coupled to the cluster mode. When the amplitude of the collective motions is large, the cluster and collective modes are practically independent, which preordains the existence of very narrow collective resonances with excitation energy in excess of 20 MeV.

We have previously proposed in a number of papers[19,20,26] a generating-function or generalized-coherent-state (GCS) algorithm that eliminates the main difficulties connected with the construction of the system of equations (2) and the computation in the $\{|\nu, \mathrm{col}\rangle\}$ and $\{|n,\mathrm{cl}\rangle\}$ bases of the matrix elements of the Hamiltonian operator $\hat{H}$ and the other operators required for the investigation of the various processes. In accordance with this algorithm, the GCS are constructed in the form of Slater determinants composed of some specially defined single-particle orbitals. In the case of the lightest p-shell nuclei, for which a single channel—the breakup of the nucleus into an α particle and an s-shell nucleus—is taken into account, the GCS should be constructed from orbitals of the form

$$\exp\left\{-\frac{1}{2}r_i^2-\frac{\beta}{1-\beta}(\mathbf{p}\mathbf{r}_i)^2+\frac{2(\mathbf{R}_k\mathbf{r}_i)}{1-\beta}-\frac{R_k^2}{1-\beta}\right\}, \quad (4)$$

where $\mathbf{r}_i$ is the radius vector of the ith particle, $\mathbf{R}_k = R_k\,\mathbf{p}$ is the cluster parameter indicating the position of the k th fragement ($k = 1,2$), and β is the deformation parameter of the nucleus. The unit vector $\mathbf{p}$ specifies the direction of the deformation and, simultaneously, the orientation in space of the vector $\mathbf{R} = \mathbf{R}_1 - \mathbf{R}_2$. The GCS constructed from the orbitals (4) will depend on $3A-3$ space coordinates (it is assumed that the motion of the center of mass has been eliminated), $3A$ spin and isospin variables, and four generator parameters: The three components of the vector $\mathbf{R}$ and the deformation parameter β. The expansion of this GCS [which expansion we shall denote by $\Phi(R,\beta,\mathbf{p})$] in powers of the generator parameters R and β generates the oscillator-function basis:

$$\Phi(R,\beta,\mathbf{p})=\sum_{n,\nu}\sum_{L,M}A_L^{n,\nu}R^{2n+f}\beta^{\nu}Y_{LM}^*(\mathbf{p})\,|n,\nu;LM\rangle, \quad (5)$$

where the $A_L^{n,\nu}$ are the structure constants and f is the minimally admissible—on the basis of the Pauli principle—power of the generator parameter R (for greater details, see Ref. 27). For p-shell nuclei, $f = A-4$ for the states of normal parity [$\pi = (-1)^A$], while for the states of anomalous parity [$\pi = (-1)^{A+1}$] we have $f = A-3$. It follows from the expansion (5) that, to separate out the oscillator function $|n,\nu;LM\rangle$, we must differentiate the GCS $\Phi(R,\beta,\mathbf{p})$ $2n+f$ times with respect to R and ν times with respect to β, let R and β go to zero, and then integrate over the angle variables of the unit vector $\mathbf{p}$ with weight $Y_{LM}(\mathbf{p})$. We can, using this same procedure, easily compute the matrix element of some operator $\hat{F}$ in the $\{|n,\nu;LM\rangle\}$ and $\{|n',\nu';L'M'\rangle\}$ bases if we know the generating matrix element of this operator, i.e., the matrix element

$$\langle\Phi(\tilde{R},\tilde{\beta},\mathbf{q})|\hat{F}|\Phi(R,\beta,\mathbf{p})\rangle = \sum_{n',\nu'}\sum_{n,\nu}\sum_{L',M'}\sum_{L,M}A_{L'}^{n',\nu'}A_L^{n,\nu}\tilde{R}^{2n'+f'}R^{2n+f}\tilde{\beta}^{\nu'}\beta^{\nu}$$
$$\times Y_{L'M'}(\mathbf{q})Y_{LM}^*(\mathbf{p})\langle n',\nu';L'M'|\hat{F}|n,\nu;LM\rangle. \quad (6)$$

The collective and cluster functions involved in the expansion (1) form only a small part of the complete function basis $\{|n,\nu;LM\rangle\}$. If we fix the parameter ν in the functions $|n,\nu\rangle \equiv |n,\nu;LM\rangle$, and set it equal to zero, then the resulting set of functions will coincide with the cluster basis $|n,\mathrm{cl}\rangle = |n,0\rangle$. Similarly, we can obtain the collective functions $|\nu,\mathrm{col}\rangle$ by setting n equal to zero: $|\nu,\mathrm{col}\rangle = |0,\nu\rangle$.

The general methods of constructing the GCS and computing the matrix elements of the various operators in the GCS basis, as well as the problems connected with the separation of the center-of-mass motion, are considered in detail in Ref. 26. The specific kinetic- and potential-energy-operator matrix elements needed for the investigation of the p-shell nuclei with $5 \leqslant A \leqslant 8$ nucleons are given in Ref. 28. The results of the computation of the collective -excitation spectra of the above-mentioned nuclei in the $\{|\nu, \mathrm{col}\rangle\}$ basis of the Sp(2,R) model are presented in Ref. 21, where the matrix elements of the operators $\hat{H}$ are also given in their explicit form. In Ref. 23 the main stages in the computation of the spectrum and wave functions in the cluster model are considered in detail. In Ref. 24 allowance is made for the monopole mode, along with the cluster mode. The above-listed papers contain all the details necessary for the investigation of the discrete and continuum states of the p-shell nuclei with allowance for the interaction of both the collective and cluster modes of the motion.

The inclusion of the collective, along with the cluster, degrees of freedom allows us to take account of the dynamical polarization of the clusters during their collision. It should be noted that there are other methods of taking account of the polarization. Thus, for example, in the papers cited in Ref. 29 the monopole polarization of one or each of the colliding clusters is taken into account. But such polarization has little effect on the results obtained, and does not allow us to identify the resonance states stemming from the collective motion of the nucleus.

4. COLLECTIVE RESONANCES AND PHOTODISINTEGRATION CROSS SECTIONS

Before proceeding to discuss the results of the theoretical calculations, let us note that, for the solution of the system (7), we took into account in the expansion (1) 25 functions of the collective basis and 100 functions of the cluster basis. This number of basis functions allows us to reproduce sufficiently correctly the wave functions of the nucleon systems in both the internal and the outer region. As has already been stated above, the interaction between the nucleons was modeled by the first variant of the Brink-Boeker potential.[30] The oscillator radius, which has the same value for the cluster and collective bases, was chosen on the basis of the requirement that the threshold for each of the investigated nuclei be a minimum. In view of the fact that the Coulomb interaction between the protons was ignored, and since the

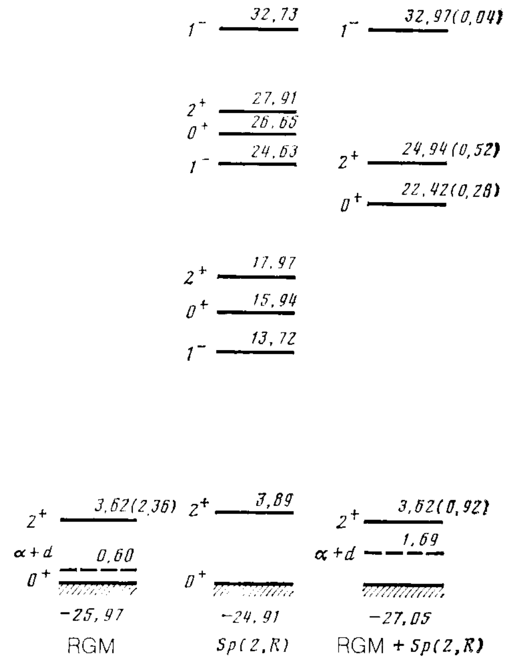

FIG. 1. Spectrum of the collective excitations and resonance states of the nucleus ^{6}He(^{6}Li). RGM: Calculation in the cluster basis; Sp(2,R): in the collective basis; RGM + Sp(2,R): calculation taking account of the coupling between the cluster and collective modes of the motion. The level energies and the resonance widths (numbers in the brackets) are given in MeV.

even and odd components of the Brink-Boeker potential are equal to each other, the threshold energies for the $\alpha + 2n$ and $\alpha + d$ (as well as $\alpha + t$ and $\alpha + {}^3$He) reactions and, consequently, the oscillator radii r_0 for the ^{6}He and ^{6}Li (^{7}Li and ^{7}Be) nuclei are equal. For this very reason, the results obtained for the nucleus ^{6}He(^{7}Li) also pertain to the nucleus ^{6}Li(^{7}Be).

Figure 1 shows the spectrum of the resonances and the collective excitations of the nucleus ^{6}He(^{6}Li), as obtained in the various approximations. As can be seen, the cluster basis (i.e., the RGM) allows us to describe the bound states and the resonance states (the 2^+ state in this case) stemming from the existence of the centrifugal barrier. With the aid of the collective functions [Sp(2,R)] we can obtain, besides the rotational excitations (the 2^+ state for ^{6}He), a set of the so-called vibrational excitations ($L^\pi = 0^+$ and 2^+ for ^{6}He), as well as states of anomalous parity ($L^\pi = 1^-$). The first 0^+ and 2^+ vibrational excitations of even-A nuclei and the 1^- and 3^- vibrational excitations of odd-A nuclei are usually identified in the microscopic theory of collective excitations with the monopole and quadrupole giant resonances, since they are connected with the ground state by large matrix elements of the monopole- and quadrupole-transition operators, and account for a significant fraction of the monopole and quadrupole energy-weighted sum rules (EWSR).

But if, besides the collective mode, we also take account of the cluster moder [RGM + Sp(2,R)], then the above-mentioned vibrational collective excitations dissolve in the continuum, and the elastic-scattering phases in the corresponding cluster channel do not exhibit resonance behavior

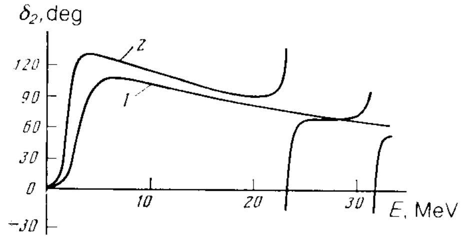

FIG. 2. D phase of elastic αd scattering, as computed in the cluster basis (RGM) (curve 1) and with allowance for the collective polarization [RGM + Sp(2,R)] (curve 2).

at those energies at which, according to the calculations carried out with the collective-function basis for a closed cluster channel, a giant resonance should occur. Consequently, the giant resonances cannot manifest themselves explicitly in interaction reactions involving light clusters, and we must find other ways of identifying them. The theoretical result obtained by us for the first 0^+ and 2^+ vibrational excitations is explained by the fact that the collective-oscillation amplitude in these excited states is relatively small, and therefore the coupling between the collective and cluster modes is a strong one. This leads to the immediate decay of the giant resonance via the cluster channel.

As to the higher 0^+ and 2^+ vibrational excitations, for them the collective-oscillation amplitude is typically large. The strength of the coupling between the collective and cluster modes in these excited states is smaller, and they do not get washed out upon the inclusion of the cluster basis, but appear as narrow resonances, manifesting themselves clearly on the energy-dependence curves for the elastic-scattering phases and the partial effective cross sections in the open cluster channel. The width of these collective resonances does not exceed 1 MeV, i.e., it is significantly smaller than the width of the 2^+ centrifugal resonance. A similar picture is observed in the odd-A nuclei ^{7}Li and ^{7}Be.

The plot, shown in Fig. 2, of the energy dependence (in the c.m. system) of the D phase δ_2 of elastic α-particle scattering by the deuteron illustrates the washing out in the continuum of the 2^+ quadrupole resonance, which, according to

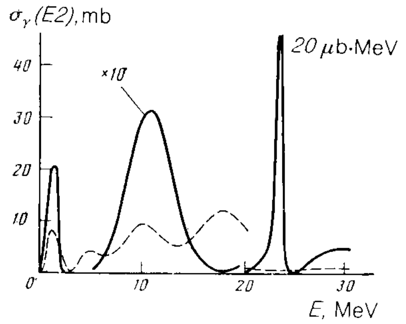

FIG. 3. Cross section $\sigma_\gamma(E2, 0^+ \to 2^+)$ for quadrupole photodisintegration of the ^{6}Li nucleus; E is the energy of the separating α and d fragments. Here and in the other figures the numbers beside the resonance peaks give the areas of the resonances. The continuous curve was computed in the RGM + Sp(2,R) basis; the dashed curve, in the RGM basis.

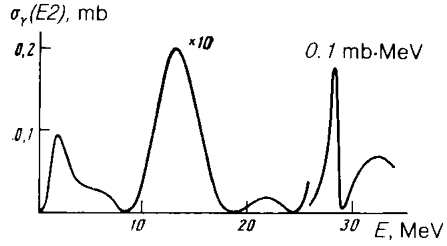

FIG. 4. Cross section $\sigma_\gamma(E2, 1^-\to3^-)$ for quadrupole photodisintegration of the ^{7}Li nucleus; E is the energy of the α and t fragments.

calculations carried out within the framework of the Sp(2,R) model, should occur in the 10–20-MeV energy region. In this region the phase δ_2 decreases monotonically, and exhibits resonance behavior only at $E = 23$ MeV, which corresponds to the excitation of the second 2^+ collective resonance.

Let us now consider the energy dependence of the photodisintegration cross sections σ_γ for the nuclei ^{6}He(^{6}Li) and ^{7}Li(^{7}Be). Since the cross sections σ_c for radiative capture are related to the cross sections σ_γ by the conditions for detailed balance, our conclusions will equally well hold true for the cross sections σ_c. A characteristic of the computed σ_γ cross sections (Figs. 3–6) is the presence of resonance peaks on them. Some of them pertain to the narrow collective resonances excitable by γ-rays. The radiative width of these resonances is equal to their α-decay width. Moreover, the photodisintegration cross sections have sharp peaks in the near-thershold region. These peaks occur as a result of the fact that the continuum wave functions have, when their energies exceed the threshold energy by small amounts, a structure similar to the structure of the wave function of the ground state, which lies not far below the threshold, and the electromagnetic-transition-operator matrix elements connecting the ground state to the above-the-threshold continuum states are really large. In sum, this is what gives rise to the above-the-threshold resonances. A detailed discussion of the above-the-threshold monopole resonances can be found in Ref. 31.

Finally, for ^{6}He, ^{6}Li, ^{7}Li, and ^{7}Be we found a quadrupole-resonance peak of width ~ 5 MeV in the 12–15-MeV excitation-energy region. It accounts for more than 15% of

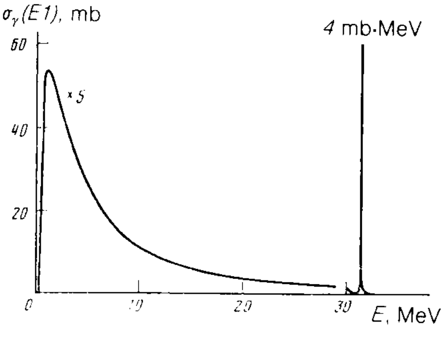

FIG. 5. Cross section $\sigma_\gamma(E1, 0^+\to1^-)$ for the dipole photodisintegration of the ^{6}He nucleus; E is the energy of the α and $2n$ fragments.

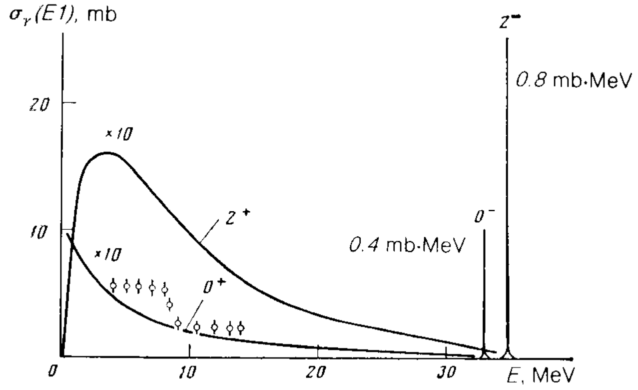

FIG. 6. Cross sections $\sigma_\gamma(E1, 1^-\to2^+)$ and $\sigma_\gamma(E1, 1^-\to0^+)$ for dipole photodisintegration of the ^{7}Li nucleus; E is the energy of the α and t fragments. The points correspond to experimental data taken from Ref. 31.

the energy-weighted quadrupole sum rule, and can be interpreted as a giant quadrupole resonance, since it is located at the place where an Sp(2,R) collective model calculation shows such a resonance should occur.

As a result of the washing out of the collective quadrupole excitation in the continuum, it is the quadrupole transition from the ground state of the nucleus into a continuous series of states grouped around some center, and not the transition into a single state, that has a large matrix element. The peak in the photodisintegration cross section separates out these states from among the other continuum states.

Thus, the photonuclear reactions are the most convenient and natural means of detecting the giant resonances in the light nuclei.

The experimental situation, as obtains at present, does not allow a complete verification of the quantitative results of the theory. Only Ref. 32, which reports the results of measurements of the effective differential cross section for the reaction

$$^7\text{Li}(\gamma, \alpha)^3\text{H} \qquad (7)$$

in the range of γ-ray energies from 5 to 50 MeV, has a direct bearing on the theoretical results obtained by us. The theoretical estimates of the effective cross section for the reaction (7) yield a value that is greater than the measured value by more than a factor of two. At least two factors could have contributed to the overestimation of the theoretical cross section. First, we ignore in the calculations the Coulomb interaction between the clusters (it suppresses the relative-motion wave function for the clusters at small intercluster separations). Second, we ignore the (^{6}Li + n) ^{7}Li-disintegration channel, which opens at relatively low energies.

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Translated by A. K. Agyei