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Analogues of giant resonances and photoproduction of positive pions from ^{16}O

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Abstract. Using a pure configuration shell model, differential and total cross sections have been evaluated for the photoproduction of positive pions from ^{16}O leading to giant resonance states of various multipolarities in ^{16}N. The two-particle–two-hole (2p–2h) ground state correlations in ^{16}O are also taken into account. These 2p–2h correlations reduce the cross sections by about 30%. The total cross section for the 2^+ giant quadrupole state is found to have a large peak near the first pion–nucleon resonance region, making it thereby possible to experimentally identify it from other multipole resonance states by varying the incident photon energy.

1. Introduction

It has been recently pointed out (Überall 1969) that the inelastic nuclear pion photoproduction process

$$\gamma + \mathrm{N}(A, Z) \rightarrow \pi^{\pm} + \mathrm{N}(A, Z \mp 1) \tag{1}$$

is superior to the muon capture and radiative pion capture processes for a study of giant analogue resonances, since the momentum transfer can be varied in (1) by varying the angle at which the pion is observed. Kelly *et al.* (1969) have calculated the differential cross section for photoproduction of π^+ from ^{16}O, near threshold, leading to giant-resonance spin-flip states of various multipolarities in ^{16}N, on the basis of the generalized Goldhaber–Teller model assuming volume production and surface production of pions. We have earlier studied (Devanathan *et al.* 1967) the energy dependence of the cross section for photoproduction of π^+ from ^{16}O leading to the four bound states $J^p = 0^-, 1^-, 2^-, 3^-$ in ^{16}N studied by means of the transition amplitude obtained in the impulse approximation using configuration mixing particle–hole wave functions for the final nuclear states. In a recent paper (Srinivasa Rao and Devanathan 1970, to be referred to as I) we have shown that a better agreement between theory and experiment (Meyer *et al.* 1965) for π^+ photoproduction from ^{16}O leading to the four bound states of ^{16}N can be obtained by using the Kuo wave function with 'screening' correction for ^{16}N states together with the 'deformed' ground state wave function of ^{16}O. Here, we present a study of the differential and total cross sections for π^+ photoproduction of ^{16}O due to transitions to the giant multipole resonance states in ^{16}N treated in a pure configuration shell model, using the complete Chew *et al.* (1957, to be referred to as CGLN) transition amplitudes and taking into account the 2p–2h correlations in the ground state of ^{16}O. We compare our results with those of Kelly *et al.* (1969).

2. Theory

The free nucleon photoproduction amplitude has the general structure $(\boldsymbol{\sigma} . \boldsymbol{K} + L)$ where $\boldsymbol{K}$ and L are functions of momenta and energies of the incident photon and the outgoing pion, the polarization of the photon and the angle of the pion emission.

They also depend on the magnetic moments of the proton and the neutron. In our study, we choose K and L from the theory of CGLN. If we denote the initial ground state of ^{16}O as $|0^+, gs\rangle$ and the final nuclear state as $|f\rangle$, then the differential cross section we are interested in is given by

$$\frac{d\sigma}{d\Omega}(0^+ \rightarrow f) = (2\pi)^{-2}\mu\mu_0 \overline{\sum_{M_f}} |\langle f|T|0^+, gs\rangle|^2 \tag{2}$$

where T is the nuclear transition operator for reaction (1) and the bar over the sum denotes the average over photon polarizations. The details of the calculation can be found in earlier papers of Devanathan and Ramachandran (1963, 1965) and in Srinivasa Rao (1970).

The final nuclear states we are interested in are the dipole ($J^p = 0^-, 1^-, 2^-$) and the quadrupole ($J^p = 1^+, 2^+, 3^+$) giant resonances of the spin–isospin type. A great number of nuclear levels actually contribute to the photon absorption but they all cluster around the first excited state of the idealized giant resonance state. Hence, we sum over all states of a given J^p in our Independent Particle Model (IPM) study. Even though the transition cross sections to giant resonances are treated in a pure-configuration shell model, they are expected to remain the same when calculated in a realistic shell model with residual interactions (Kelly and Überall 1968).

Recent experimental studies of $^{16}O(d, t)^{15}O$ and $^{16}O(d, {}^3He)^{15}N$ by Purser *et al.* (1969) have confirmed the existence of 'deformed' two-particle–two-hole (2p–2h) admixtures in the ground state wave function of ^{16}O, confirming thereby a prediction of Brown and Green (1966a,b). Assuming the ground state wave function of ^{16}O to contain 2p–2h components, Walker (1968) and Green *et al.* (1969, 1970) have shown that certain theoretical–experimental discrepancies for muon-capture and photodisintegration on ^{16}O can be resolved. We have shown in I that, without invoking the phenomenological surface production mechanism, a better agreement between theory and experiment for the photoproduction of π^+ from ^{16}O $(0^+, gs)$ leading to the four bound states of ^{16}N ($J^p = 0^-, 1^-, 2^-$ and 3^-) can be obtained by using the deformed ground state wave function of ^{16}O obtained by Purser *et al.* (1969). Therefore we wish to take into account the effect of these 2p–2h correlations in the ground state wave function of ^{16}O in the present study. As in I, we assume the ground state of ^{16}O to be approximated by the form:

$$|0^+, gs\rangle = \alpha|0p-0h\rangle + \beta|(1d_{5/2}{}^2)_{J=0,T=1}(1p_{1/2}{}^{-2})_{0,1}\rangle + \gamma|(2s_{1/2}{}^2)_{0,1}(1p_{1/2}{}^{-2})_{0,1}\rangle \tag{3}$$

where configurations with particles and holes separately coupled to (J, T) other than $J = 0$, $T = 1$ are neglected, based on the observation by Zamick (1965) that they lie much higher in energy and hence their coupling to the $|0p–0h\rangle$ state is much weaker. In table 1 are listed the values of the parameters α, β and γ determined by Purser *et al.* (1969) from an experimental study of the pick-up reactions $^{16}O(d, t)^{15}O$ and $^{16}O(d, {}^3He)^{15}N$.

In general, to be consistent with the choice of the ground state wave function of ^{16}O (equation (3)), the negative parity states of ^{16}N can be considered to be a combination of the usual particle–hole (1p–1h) states with three-particle–three-hole (3p–3h) configuration admixtures and the positive parity 1p–1h states can be considered to be mixed with 2p–2h states in ^{16}N, with two holes in the 1p shell and two particles in the 1d–2s shell. But, even in the case of the ground state and low-lying

Table 1. Ground state wave functions of ^{16}O

	Model	α	β	γ
I	(PS)†	1·00	0·00	0·00
II	(^{16}O(d, t)^{15}O expt)	0·87	0·26	0·27
III	(^{16}O(d, ^{3}He)^{15}N expt)	0·82	0·54	0·20

† (PS) denotes pure shell model wave function.

$T = 1$ negative parity states of ^{16}N, while Walker (1968) claims that the 3p–3h mixing can be as large as the 2p-2h mixing of the ground state, Green and Rho (1969) argue, on both theoretical and empirical grounds, that these 3p–3h admixtures are insignificant and they give the qualitative reason that $T = 0$ states (eg the ground state of ^{16}O) have a tendency to deform whereas $T = 1$ states (eg ^{16}N states) do not. Hence, in the absence of any conclusive theoretical or experimental confirmation about the relative importance of these higher particle–hole admixtures in the negative and positive parity states of ^{16}N, we choose to work in the framework of the simple IPM configuration for ^{16}N states.

The contribution from the 'deformed' 2p–2h components of the ground state wave function of ^{16}O to the matrix element $\langle f|T|0^+, gs\rangle$ will be nonzero if and only if the two particles (holes) of the 2p–2h initial state component have the same *nlj* quantum numbers as the particle (hole) of the 1p–1h final (IPM) state. The final IPM configurations of the giant multipole resonance states given in table 2 (see

Table 2. Cross sections for the reaction ^{16}O(γ, π^+)^{16}N when the final state is one of the excited states which decay by nucleon emission. The incident photon energy is 260 MeV

^{16}N state J^p	Configuration of ^{16}N state in IPM		σ (μb)
0^-	$(1p_{3/2}^{-1})$	$(1d_{3/2})$	3·531
1^-	$(1p_{1/2}^{-1})$	$(1d_{3/2})$	8·508
	$(1p_{3/2}^{-1})$	$(1d_{5/2})$	15·315
	$(1p_{3/2}^{-1})$	$(2s_{1/2})$	3·192
	$(1p_{3/2}^{-1})$	$(1d_{3/2})$	11·007
2^-	$(1p_{1/2}^{-1})$	$(1d_{3/2})$	7·755
	$(1p_{3/2}^{-1})$	$(1d_{5/2})$	10·560
	$(1p_{3/2}^{-1})$	$(2s_{1/2})$	5·753
	$(1p_{1/2}^{-1})$	$(1d_{3/2})$	4·609
1^+	$(1s_{1/2}^{-1})$	$(2s_{1/2})$	15·477
	$(1s_{1/2}^{-1})$	$(1d_{3/2})$	4·891
	$(1p_{1/2}^{-1})$	$(2p_{3/2})$	5·472
	$(1p_{1/2}^{-1})$	$(2p_{1/2})$	2·020
	$(1p_{3/2}^{-1})$	$(1f_{5/2})$	7·922
	$(1p_{3/2}^{-1})$	$(2p_{3/2})$	6·515
	$(1p_{3/2}^{-1})$	$(2p_{1/2})$	5·472
2^+	$(1s_{1/2}^{-1})$	$(1d_{5/2})$	10·578
	$(1s_{1/2}^{-1})$	$(1d_{3/2})$	11·213
	$(1p_{1/2}^{-1})$	$(1f_{5/2})$	9·415
	$(1p_{1/2}^{-1})$	$(2p_{3/2})$	4·379
	$(1p_{1/2}^{-1})$	$(1f_{7/2})$	16·139
	$(1p_{3/2}^{-1})$	$(1f_{5/2})$	8·759
	$(1p_{3/2}^{-1})$	$(2p_{3/2})$	5·038
	$(1p_{3/2}^{-1})$	$(2p_{1/2})$	4·379
3^+	$(1s_{1/2}^{-1})$	$(1d_{5/2})$	13·579
	$(1p_{1/2}^{-1})$	$(1f_{7/2})$	12·176
	$(1p_{1/2}^{-1})$	$(1f_{5/2})$	5·132
	$(1p_{3/2}^{-1})$	$(1f_{7/2})$	8·080
	$(1p_{3/2}^{-1})$	$(1f_{5/2})$	3·131
	$(1p_{3/2}^{-1})$	$(2p_{3/2})$	6·456

Note. The cross sections given here correspond to model I of table 1. To get the results for models II and III of table 1, these values have to be multiplied by 0·76 and 0·67, respectively (see text).

column 2) do not contain the states $(1p_{1/2}^{-1})(2s_{1/2})$ and $(1p_{1/2}^{-1})(1d_{5/2})$ since these are configurations which give rise to the four lowlying bound states of ^{16}N with $J^p = 0^-, 1^-, 2^-, 3^-$. Therefore, the cross sections to individual giant multipole resonance states should be multiplied only by the simple factor α^2, if we wish to take into account the effect of the 2p–2h correlations in the ground state wave function of ^{16}O. In other words

$$\left(\frac{d\sigma}{d\Omega}\right)_{\text{with corrl}} = \alpha^2 \left(\frac{d\sigma}{d\Omega}\right)_{\text{without corrl}} \tag{4}$$

where $(d\sigma/d\Omega)_{\text{without corrl}}$ is given by equation (2). Thus, it is clear that the inclusion of 2p–2h correlations in the ground state of ^{16}O reduces the cross sections to giant multipole resonance states by a factor of 0·76 for $\alpha = 0{\cdot}87$ or by a factor of 0·67 for $\alpha = 0{\cdot}82$.

3. Results and discussion

We use the value of the oscillator parameter, fitted to the electron scattering data (eg Hofstadter 1960)

$$b = 1{\cdot}76\,\text{fm}. \tag{5}$$

In table 2 are given the cross sections for the reaction $^{16}O(\gamma, \pi^+)^{16}N$ when the final state is any of the excited states whose wavefunctions are described in the simple IPM, for an incident photon energy of 260 MeV. The IPM configurations given in this table are ones which we expect would have contributed to an individual giant resonance state. These states of ^{16}N can decay by nucleon emission and hence these would not have contributed to the experimental cross sections of Meyer *et al.* (1965). These higher excited states contribute nearly thrice as much to the cross section as the contributions which arise from transitions to the low lying states of ^{16}N which are stable against nucleon emission (Devanathan *et al.* 1967).

Figure 1 shows the angular distribution for excitation of the individual giant multipole states at an incident photon energy of 200 MeV. We notice that for $\theta_{cm} > 60°$, the 1^+, 2^+ and 3^+ quadrupole state cross sections exceed the 1^- and 2^- dipole cross sections, so that varying θ_{cm} allows one to differentiate between the giant multipole resonances. This conclusion is in conformity with that drawn by Kelly *et al.* (1969) on the basis of a generalized Goldhaber–Teller model.

Table 3. Total cross section (in μb) due to photon induced transition to individual giant resonance states of ^{16}N for an incident photon energy of 200 MeV

^{16}N giant resonance state	Present IPM calculation			Calculation of Kelly *et al.*†	
J^p	$\alpha = 1{\cdot}0$	$\alpha = 0{\cdot}87$	$\alpha = 0{\cdot}82$	$a_1 = 0$	$a_1 = 2{\cdot}625$ fm
0^-	1·4	1·06	0·94	2·8	1·1
1^-	51·3	38·82	34·49	31·6	12·6
2^-	30·4	23·00	20·44	49·0	19·7
1^+	56·4	42·68	37·92	13·0	7·9
2^+	54·4	41·17	36·57	48·5	28·8
3^+	43·9	33·22	29·51	64·2	38·1

† The cutoff parameter $a_1 = 0$ corresponds to volume production and $a_1 = 2{\cdot}625$ fm corresponds to surface production of pions.

In table 3 we compare our results with those of Kelly *et al.* (1969) for an incident photon energy of 200 MeV. It is interesting to note the large differences which exist between the two model dependent calculations. We are unable to draw any conclusions due to the absence of experimental data but we notice that the 2p–2h

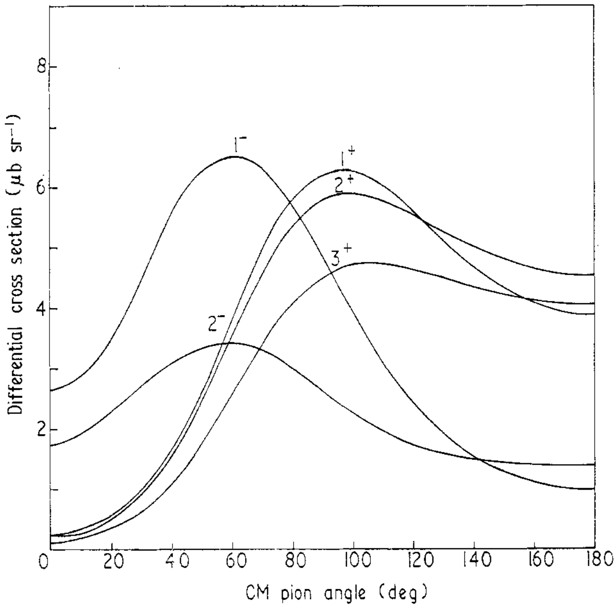

Figure 1. Photopion (π^+) angular distribution from ^{16}O for 200 MeV incident photon energy, with excitation to giant dipole and quadrupole resonance states of ^{16}N.

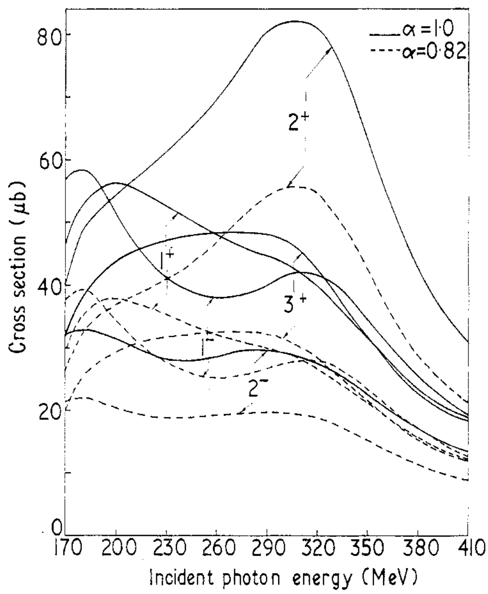

Figure 2. Energy dependence of the total cross section for π^+ photoproduction from ^{16}O due to excitation of the giant dipole and quadrupole resonance states of ^{16}N.

ground state correlations reduce the cross sections only by about 30% while the phenomenological surface production mechanism reduces the cross sections by almost a factor of two, as is to be expected from our earlier analysis (Devanathan *et al.* 1967) of $^{16}O\ (\gamma, \pi^+)^{16}N$ leading to the four bound states of ^{16}N.

In figure 2 we have plotted the energy dependence of the cross sections due to excitation of the multipole states. It is interesting to note that the 2^+ quadrupole state has a large peak at 320 MeV. It stands out from all other states, so that varying the incident photon energy allows us to differentiate it from other multipole resonance states.

In conclusion we wish to point out that since the cross sections due to excitations of the giant dipole and quadrupole states of ^{16}N in photoproduction of π^+ from ^{16}O are very much larger than those due to excitations to low-lying bound states of ^{16}N, it will be of value to measure these experimentally.

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