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GDR Contribution to Coulomb Excitation. II†
^{17}O

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Abstract

The contribution to the Coulomb excitation of the first excited state of ^{17}O due to virtual excitation of the giant dipole resonance (GDR) is calculated, using shell model wavefunctions for the ground and first excited states. A large value is obtained.

1. Introduction

In calculating the Coulomb excitation probability of the first excited state of a nucleus, the values of at least four quantities are needed. These are B(E2) for the E2 transition from the ground to the first excited state, the quadrupole moments Q and Q^* of the two states, and a parameter k measuring the strength of the contribution due to virtual excitation of the GDR. Values of k have usually been taken from estimates based on a hydrodynamic model, which is not necessarily reliable for light nuclei. Values of k obtained from shell model calculations in Part I (Barker 1982, present issue p. 291) agreed reasonably with experimental values for ^{6}Li, ^{7}Li and ^{10}B, but only after radial integrals had been renormalized in order to fit experimental B(E2) values in these nuclei. This renormalization was presumably necessary because higher configurations were neglected in the wavefunctions for the ground and first excited states. It is desirable to compare calculated and measured values of k for a light nucleus for which wavefunctions including higher configurations are available, so that such renormalization is not needed. A suitable case is ^{17}O.

There have been many studies of the wavefunctions of the ground and first excited states of ^{17}O, in which higher configurations (core excitations) have been included with the aim of fitting the values of Q and B(E2). In comparison with other light nuclei, ^{17}O has the advantages of being stable (so that Q has been measured accurately), of having a particle-stable first excited state with spin differing by two from that of the ground state (so that B(E2) can be obtained from a lifetime measurement), and of having the spin of the first excited state equal to $\frac{1}{2}$ (so that $Q^* = 0$). Further, the excitation energy of the first excited state is small (0·87 MeV), making the excitation probability sensitive to the value of k, and there is a large energy gap between the $\frac{1}{2}^+$ first excited state and higher states that might feed this $\frac{1}{2}^+$ state by virtual E2 excitation, so that their contribution can be neglected.

Since ^{17}O seems to be such a good test case, we here give a shell model calculation of the GDR contribution. A measurement of the Coulomb excitation of the first excited state of ^{17}O is in progress at Canberra (R. H. Spear, personal communication).

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2. Formulae

Notation and formulae are taken from Häusser *et al.* (1973) and Part I. The parameter k, which measures the strength of the GDR contribution, is defined by the ratio

$$k = X/X_0, \tag{1}$$

where

$$X = S(\mathrm{E1})/\langle \mathrm{i} \,\|\, \mathscr{M}(\mathrm{E2}) \,\|\, \mathrm{f} \rangle, \tag{2}$$

with

$$S(\mathrm{E1}) = \sum_n W(11I_\mathrm{i}\, I_\mathrm{f}, 2I_n)\langle \mathrm{i} \,\|\, \mathscr{M}(\mathrm{E1}) \,\|\, n \rangle \langle n \,\|\, \mathscr{M}(\mathrm{E1}) \,\|\, \mathrm{f} \rangle/(E_n - E_\mathrm{i}). \tag{3}$$

Here $|\mathrm{i}\rangle$ is the ^{17}O ground state with $I_\mathrm{i} = \frac{5}{2}$, and $|\mathrm{f}\rangle$ is the first excited state with $I_\mathrm{f} = \frac{1}{2}$. The unit X_0 in equation (1) is given by

$$X_0 = 0{\cdot}00058\, A/Z \;\; e\,\mathrm{MeV}^{-1} = 0{\cdot}00123 \;\; e\,\mathrm{MeV}^{-1}. \tag{4}$$

Contributions to $S(\mathrm{E1})$ come only from states $|n\rangle$ with $I_n = \frac{3}{2}$. We assume that all the E1 strength from the ground state to states of spin $\frac{3}{2}$ is concentrated in one eigenstate with energy E_g. Then (3) reduces to

$$S(\mathrm{E1}) = 5^{-\frac{1}{2}}(E_\mathrm{g} - E_\mathrm{i})^{-1}\langle \mathrm{i} \,\|\, \mathcal{O}^2 \,\|\, \mathrm{f} \rangle. \tag{5}$$

The denominator in (2) is related to $B(\mathrm{E2})$ by

$$B(\mathrm{E2}; \mathrm{i} \to \mathrm{f}) = (2I_\mathrm{i} + 1)^{-1}\, |\langle \mathrm{i} \,\|\, \mathscr{M}(\mathrm{E2}) \,\|\, \mathrm{f} \rangle|^2. \tag{6}$$

The wavefunctions of the ground and first excited states of ^{17}O, and of the analogue states in ^{17}F, have been discussed by many authors, mainly in connection with the problem of fitting the values of Q and $B(\mathrm{E2})$ for the two nuclei. We assume that they have the form (Barker 1964)

$$\begin{aligned} \Psi(M_T \tfrac{5}{2} M) = {} & \Psi([00\tfrac{5}{2}]M_T \tfrac{5}{2} M) + a_1\, \Psi([02\tfrac{1}{2}]M_T \tfrac{5}{2} M) + a_1'\, \Psi([12\tfrac{1}{2}]M_T \tfrac{5}{2} M) \\ & + a_2\, \Psi([02\tfrac{5}{2}]M_T \tfrac{5}{2} M) + a_2'\, \Psi([12\tfrac{5}{2}]M_T \tfrac{5}{2} M), \end{aligned} \tag{7a}$$

$$\Psi(M_T \tfrac{1}{2} M) = \Psi([00\tfrac{1}{2}]M_T \tfrac{1}{2} M) + a_3\, \Psi([02\tfrac{5}{2}]M_T \tfrac{1}{2} M) + a_3'\, \Psi([12\tfrac{5}{2}]M_T \tfrac{1}{2} M), \tag{7b}$$

where $M_T = \frac{1}{2}$ for ^{17}O ($-\frac{1}{2}$ for ^{17}F) and $[\overline{T}\overline{J}j]$ represents an $A = 16$ core state of isospin $\overline{T}$ and spin $\overline{J}$ coupled to an odd nucleon in the state nl_j (either $1\mathrm{d}_{5/2}$ or $2\mathrm{s}_{1/2}$). In Barker (1964), the terms in (7) involving a_i, a_i', which represent distortion of the ^{16}O core, were treated in first-order perturbation theory, the coefficients a_i, a_i' being of order 0·1–0·2. These distortion terms are necessary in order to explain the appreciable measured values of Q and $B(\mathrm{E2})$ for ^{17}O, since the zeroth-order terms vanish when the usual E2 operator (no recoil) is used. For ^{17}F the zeroth-order terms do not vanish, while the distortion contributions are smaller than in ^{17}O but are not negligible.† Recoil contributions, obtained by taking the origin of

† The formulae (25) in Barker (1964), which display these distortion contributions, are incorrect; the quantity $\langle r^2\rangle_{00}^{\frac{1}{2}}$ should in each case be replaced by 2·73 fm $\langle r^2\rangle_{00}^{\frac{1}{2}}$. Because $\langle r^2\rangle_{00}^{\frac{1}{2}} = 2{\cdot}64$ fm from (26), this change is numerically small and later results of that paper are not changed appreciably.

coordinates in the E2 operator as the centre of mass rather than the usual centre of the potential well, are relatively small. With the approximations made in Barker (1964), it was found that simultaneous fits to Q and B(E2) for ^{17}O could not be obtained with sets of expansion coefficients a_i, a_i' derived from any reasonable two-particle interaction. A similar result was found in more elaborate calculations, in which a distribution of E2 strength in the ^{16}O core was included by actual diagonalization of a given Hamiltonian (see e.g. Mavromatis and Singh 1969; Chung and Shin 1980). Other discussions of E2 matrix elements in ^{17}O and ^{17}F have been based on the concept of effective charges (Brown *et al.* 1977; Harvey 1978), but this does not seem to be a useful approach in the calculation of S(E1) (see Part I). We therefore evaluate S(E1) using wavefunctions of the form (7), working to first order in the coefficients a_i, a_i', which are adjusted to fit Q and B(E2).

Since $\mathcal{O}^2$ has one- and two-body terms, while the $\bar{J} = 2$, ^{16}O core states (with $\bar{T} = 0, 1$) are obtained by operating on the ^{16}O ground state with the E2 operators (taken as one-body operators by neglecting recoil), matrix elements between ^{17}O states of the lowest configuration would involve one-, two- and three-body operators. It is preferable to construct the $\bar{J} = 2$, ^{16}O core states explicitly, since then it is necessary to calculate matrix elements of only one- and two-body operators between the ^{17}O states based on these and the ^{17}O states of the lowest configurations.

Thus we take the ^{16}O ground state as the single state of the closed shell configuration $1s^4 1p^{12}$, and write the $\bar{J} = 2$, ^{16}O states relative to this as

$$\Phi_{\bar{T}2m} = \sum_{n=1}^{\infty} c_{1n}(\bar{T})\,|\,(1s^{-1}nd)\bar{T}02,02m) + \sum_{n=1}^{\infty} c_{2n}(\bar{T})\,|\,(1p^{-1}nf)\bar{T}02,02m)$$
$$+ \sum_{n=2}^{\infty} c_{3n}(\bar{T})\,|\,(1p^{-1}np)\bar{T}02,02m)\,, \qquad (8)$$

where values of the quantum numbers $TSL, M_T JM$ are given. The $c_{jn}(\bar{T})$ can easily be expressed in terms of one-body matrix elements such as the radial integral $\langle 1s \vdots r^2 \vdots nd\rangle$, where

$$\langle nl \vdots r^q \vdots n'l'\rangle = \int_0^\infty u_{nl}(r)\,u_{n'l'}(r)\,r^q r^2\,dr\,, \qquad \int_0^\infty u_{nl}^2(r)\,r^2\,dr = 1\,. \qquad (9a, b)$$

The matrix elements $\langle[00\,J]M_T J \,\|\, \mathcal{O}^2 \,\|\, [\bar{T}2j]M_T J'\rangle$ can be calculated, the coefficients of $c_{jn}(\bar{T})$ involving radial integrals such as $\langle 1p \vdots r \vdots nd\rangle$. By using relations such as

$$\sum_n \langle 1s \vdots r^2 \vdots nd\rangle\langle 1p \vdots r \vdots nd\rangle = \langle 1s \vdots r^3 \vdots 1p\rangle\,, \qquad (10)$$

one obtains

$$S(\mathrm{E1}) = \tfrac{96}{1445}\pi^{-1}(E_g - E_i)^{-1}e^2[\langle 1d \vdots r^2 \vdots 2s\rangle - 2\langle 1p \vdots r \vdots 1d\rangle\langle 1p \vdots r \vdots 2s\rangle$$
$$+ N_0^{-\frac{1}{2}}\{(\tfrac{1}{5}\sqrt{14})\langle 1p \vdots r \vdots 1d\rangle N_3\, a_1 - (\tfrac{1}{12}\sqrt{42})\langle 1p \vdots r \vdots 1d\rangle N_3\, a_1'$$
$$+ (\tfrac{145}{64}N_0 - \tfrac{145}{32}\langle 1s \vdots r \vdots 1p\rangle N_1 - \langle 1p \vdots r \vdots 2s\rangle N_3)a_2$$
$$+ \sqrt{3}(-\tfrac{17}{192}N_0 + \tfrac{17}{96}\langle 1s \vdots r \vdots 1p\rangle N_1 + \tfrac{5}{12}\langle 1p \vdots r \vdots 2s\rangle N_3)a_2'$$
$$+ \sqrt{3}(\tfrac{145}{192}N_0 - \tfrac{145}{96}\langle 1s \vdots r \vdots 1p\rangle N_1 - \tfrac{1}{5}\langle 1p \vdots r \vdots 1d\rangle N_2 - \tfrac{1}{30}\langle 1s \vdots r^2 \vdots 1d\rangle N_4)a_3$$
$$+ (-\tfrac{17}{192}N_0 + \tfrac{17}{96}\langle 1s \vdots r \vdots 1p\rangle N_1 + \tfrac{1}{4}\langle 1p \vdots r \vdots 1d\rangle N_2 - \tfrac{113}{960}\langle 1s \vdots r^2 \vdots 1d\rangle N_4)a_3'\}]\,, \qquad (11)$$

where

$$N_0 = \langle 1\mathrm{s} \vdots r^4 \vdots 1\mathrm{s}\rangle + 3\langle 1\mathrm{p} \vdots r^4 \vdots 1\mathrm{p}\rangle - \tfrac{6}{5}\langle 1\mathrm{p} \vdots r^2 \vdots 1\mathrm{p}\rangle^2, \tag{12a}$$

$$N_1 = 2\langle 1\mathrm{s} \vdots r^3 \vdots 1\mathrm{p}\rangle - \langle 1\mathrm{p} \vdots r^2 \vdots 1\mathrm{p}\rangle\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle, \tag{12b}$$

$$N_2 = \langle 1\mathrm{p} \vdots r^3 \vdots 1\mathrm{d}\rangle - \langle 1\mathrm{p} \vdots r^2 \vdots 1\mathrm{p}\rangle\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle, \tag{12c}$$

$$N_3 = \langle 1\mathrm{p} \vdots r^3 \vdots 2\mathrm{s}\rangle - \langle 1\mathrm{p} \vdots r^2 \vdots 1\mathrm{p}\rangle\langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle, \tag{12d}$$

$$N_4 = \langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle - 2\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle. \tag{12e}$$

In the same notation we have

$$\begin{aligned} Q = {} & -\tfrac{32}{2023}e[\langle 1\mathrm{d} \vdots r^2 \vdots 1\mathrm{d}\rangle - \tfrac{7}{5}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle^2 \\ & + N_0^{-\frac{1}{2}}\{\surd 14(-\tfrac{271}{16}N_0 - \tfrac{9}{4}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 + \tfrac{7}{25}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_2 \\ & \qquad + \tfrac{3}{20}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle N_4 - \tfrac{3}{350}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle\langle 1\mathrm{p} \vdots r^3 \vdots 1\mathrm{d}\rangle)a_1 \\ & + \surd 42(\tfrac{85}{16}N_0 + \tfrac{17}{12}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 - \tfrac{7}{60}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_2 \\ & \qquad - \tfrac{11}{40}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle N_4 + \tfrac{1}{280}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle\langle 1\mathrm{p} \vdots r^3 \vdots 1\mathrm{d}\rangle + \tfrac{289}{80}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle^2)a_1' \\ & - \tfrac{1}{5}\langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle(7N_2 + 3\langle 1\mathrm{p} \vdots r^3 \vdots 1\mathrm{d}\rangle)a_2 \\ & + (\tfrac{1}{4}\surd\tfrac{1}{3})\langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle(7N_2 + 3\langle 1\mathrm{p} \vdots r^3 \vdots 1\mathrm{d}\rangle)a_2'\}], \end{aligned} \tag{13}$$

and $B(\mathrm{E2})$ is given by equation (6) with

$$\begin{aligned} \langle \mathrm{i} \| \mathscr{M}(\mathrm{E2}) \| \mathrm{f}\rangle = {} & \tfrac{24}{289}(6\pi)^{-\frac{1}{2}}e[\langle 1\mathrm{d} \vdots r^2 \vdots 2\mathrm{s}\rangle - 2\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle\langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle \\ & + N_0^{-\frac{1}{2}}\{(\tfrac{1}{5}\surd 14)\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_3\, a_1 - (\tfrac{1}{12}\surd 42)\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_3\, a_1' \\ & + (\tfrac{271}{8}N_0 + \tfrac{9}{2}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 - \langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle N_3)a_2 \\ & + \surd 3(-\tfrac{85}{8}N_0 - \tfrac{17}{6}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 + \tfrac{5}{12}\langle 1\mathrm{p} \vdots r \vdots 2\mathrm{s}\rangle N_3)a_2' \\ & + \surd 3(\tfrac{271}{24}N_0 + \tfrac{3}{2}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 - \tfrac{1}{5}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_2 \\ & \qquad - \tfrac{1}{30}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle N_4)a_3 \\ & + (-\tfrac{85}{8}N_0 - \tfrac{17}{6}\langle 1\mathrm{s} \vdots r \vdots 1\mathrm{p}\rangle N_1 + \tfrac{1}{4}\langle 1\mathrm{p} \vdots r \vdots 1\mathrm{d}\rangle N_2 \\ & \qquad + \tfrac{11}{60}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle N_4 - \tfrac{289}{120}\langle 1\mathrm{s} \vdots r^2 \vdots 1\mathrm{d}\rangle^2)a_3'\}]. \end{aligned} \tag{14}$$

3. Calculated Values

For harmonic oscillator single-particle wavefunctions, the zeroth-order contribution to $S(\mathrm{E1})$ vanishes, for the same reason† as in 1p shell nuclei (see Part I). We evaluate the radial integrals using single-particle wavefunctions belonging to a real Woods–Saxon potential, with central and surface-peaked spin–orbit terms, plus the Coulomb potential of a uniformly charged sphere. Parameter values for such a potential appropriate to $^{17}\mathrm{O}$ were found by Brown *et al.* (1977) by fitting the measured $^{16}\mathrm{O}$ r.m.s. charge radius, giving $r_0 = 1\cdot324$ fm and $a = 0\cdot65$ fm. We adjust the central potential depth to fit binding energies of 4·144 and 3·273 MeV for the 1d and 2s neutron states respectively (Ajzenberg-Selove 1982), the spin–orbit

† This is not the same as the reason for the vanishing of $B(\mathrm{E2})$ in $^{17}\mathrm{O}$ in zeroth order, since the zeroth-order contribution to $S(\mathrm{E1})$ for $^{17}\mathrm{F}$ also vanishes.

strength for the 1d ($\equiv 1d_{5/2}$) state being $V_{ls} = 15$ MeV fm^2 as given by Brown *et al.* From the ^{16}O(e, e′p)^{15}N reaction, the 1s state is centred at $E_x \approx 41$ MeV, while the 1p strength is split between the $\frac{1}{2}^-$ ground state and the $\frac{3}{2}^-$ state at 6·32 MeV (Ajzenberg-Selove 1981). We therefore fit a 1s neutron binding energy of 56·7 MeV and, since we do not distinguish between $1p_{1/2}$ and $1p_{3/2}$ wavefunctions, we fit a 1p neutron binding energy of 19·8 MeV with $V_{ls} = 0$. Then we get

$$\begin{aligned}
\langle 1s\!:\!r^4\!:\!1s\rangle &= 23\cdot45\,\mathrm{fm}^4, & \langle 1s\!:\!r\!:\!1p\rangle &= 1\cdot917\,\mathrm{fm},\\
\langle 1s\!:\!r^3\!:\!1p\rangle &= 12\cdot87\,\mathrm{fm}^3, & \langle 1s\!:\!r^2\!:\!1d\rangle &= 4\cdot503\,\mathrm{fm}^2,\\
\langle 1p\!:\!r^2\!:\!1p\rangle &= 7\cdot520\,\mathrm{fm}^2, & \langle 1p\!:\!r^4\!:\!1p\rangle &= 82\cdot69\,\mathrm{fm}^4,\\
\langle 1p\!:\!r\!:\!1d\rangle &= 2\cdot715\,\mathrm{fm}, & \langle 1p\!:\!r^3\!:\!1d\rangle &= 32\cdot43\,\mathrm{fm}^3,\\
\langle 1p\!:\!r\!:\!2s\rangle &= -1\cdot586\,\mathrm{fm}, & \langle 1p\!:\!r^3\!:\!2s\rangle &= -30\cdot45\,\mathrm{fm}^3,\\
\langle 1d\!:\!r^2\!:\!1d\rangle &= 13\cdot31\,\mathrm{fm}^2, & \langle 1d\!:\!r^2\!:\!2s\rangle &= -14\cdot02\,\mathrm{fm}^2.
\end{aligned} \tag{15}$$

Values of the coefficients a_i, a_i' given in Barker (1964) were derived from a two-particle interaction, but they do not fit simultaneously the experimental values of both Q and B(E2) for ^{17}O. Here we choose the a_i, a_i' to fit Q and B(E2). Some further restrictions on the a_i, a_i' are needed, and we assume $a_i' = -a_i\frac{1}{2}\sqrt{3}(\Delta E/\Delta E')$, a relation found to be valid for a wide range of interactions (Barker 1964). Also from that paper we take $\Delta E/\Delta E' = 0\cdot61$, and $a_3 = 2a_2$. The latter assumption is not critical, since Q depends essentially only on the value of a_1 while B(E2) depends only on the combination $\sqrt{3}\,a_2 + a_3$ (this is clearly seen from the formulae (15) and (16) of Barker (1964), which neglect recoil). Also, S(E1) depends essentially only on a_1 and $\sqrt{3}\,a_2 + a_3$ (since the coefficients of a_2 and a_3 in (11) are each dominated by their first two terms, and the coefficients of a_2' and a_3' are relatively small). Then we fit $Q = -2\cdot578\,e\,\mathrm{fm}^2$ and $\tau_m(^{17}\mathrm{O}, \frac{1}{2}^+) = 258\cdot6$ ps (Ajzenberg-Selove 1982), which gives $B(\mathrm{E2}; \frac{5}{2}^+ \to \frac{1}{2}^+) = 2\cdot101\,e^2\,\mathrm{fm}^4$ and $\langle \mathrm{i}\,\|\,\mathscr{M}(\mathrm{E2})\,\|\,\mathrm{f}\rangle = -3\cdot550\,e\,\mathrm{fm}^2$, with

$$a_1 = -0\cdot134, \qquad \sqrt{3}\,a_2 + a_3 = -0\cdot495. \tag{16a, b}$$

Then equation (11) gives

$$S(\mathrm{E1}) = -0\cdot249(E_g - E_i)^{-1}\ e^2\,\mathrm{fm}^2. \tag{17}$$

As in the calculations for 1p shell nuclei in Part I, we assume that $E_g - E_i = \sigma_{-1}/\sigma_{-2}$, where σ_n is the nth moment of the photonuclear cross section. Total photonuclear cross sections (Ahrens *et al.* 1975) have not been measured for isotopic ^{17}O but only for oxygen of natural isotopic composition (99·8% ^{16}O); these give $\sigma_{-1}/\sigma_{-2} \approx 26$ MeV (for $E_{\gamma,\max} = 140$ MeV). Measurements of photoneutron cross sections (Jury *et al.* 1980), for $E_{\gamma,\max} \approx 40$ MeV, give $\sigma_{-1}/\sigma_{-2} \approx 20$ MeV for ^{17}O and 24–27 MeV for ^{16}O. Although the photoneutron and total photonuclear cross sections yield about the same values of σ_{-1}/σ_{-2} for ^{16}O, in which the proton threshold is 3·5 MeV below the neutron threshold, the same need not be true for ^{17}O, in which the proton threshold (at 13·8 MeV) is 9·6 MeV above the neutron threshold. We assume the photoproton and photoneutron cross sections to be about equal for

$E_\gamma \gtrsim 14$ MeV, and so estimate $\sigma_{-1}/\sigma_{-2} \approx 22$ MeV for the total photonucleon cross section. We therefore take $E_g - E_i = 22$ MeV. Then $S(E1) = -0\cdot0113\,e^2\,\mathrm{fm}^2\,\mathrm{MeV}^{-1}$, giving $X = 0\cdot00319\,e\,\mathrm{MeV}^{-1}$ and $k = 2\cdot59$.

4. Discussion

One of the main uncertainties in the calculation of the GDR contribution for 1p shell nuclei in Part I was due to the use of wavefunctions belonging entirely to the lowest shell model configuration, which necessitated the renormalization of the radial integrals in order to fit experimental B(E2) values. This is avoided in the present case by including terms in the ^{17}O wavefunctions belonging to higher configurations, and adjusting the coefficients of these to fit the experimental Q and B(E2) values (the contributions from the lowest configurations being almost negligible). These higher configuration components then contribute about 54% of S(E1) and therefore of the GDR contribution.

The other major uncertainty for 1p shell nuclei still remains, in the estimation of $E_g - E_i$. It is an approximation to take $E_g - E_i = \sigma_{-1}/\sigma_{-2}$, where the σ_n are taken from the photonuclear cross sections, since the value of $E_g - E_i$ should depend only on the location of the E1 strength to $\frac{3}{2}^-$ states of ^{17}O. Also the values of σ_n that we use are rather uncertain because total photonuclear cross sections are not available for ^{17}O.

From these calculations, it is expected that k for ^{17}O should be large, about $2\cdot6$. Such a large value is another reason, additional to those given in Section 1, why ^{17}O is a suitable nucleus for studying the GDR contribution to Coulomb excitation of low-lying excited states.

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