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Investigation of E1 Strength in Coulomb Excitation of Light Nuclei

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Abstract

Contributions to the predominantly E2 Coulomb excitation of the first excited states of ^{6}Li, ^{7}Li, ^{10}B, ^{11}B, ^{12}C and ^{17}O due to virtual E1 transitions through intermediate states are calculated using for all states shell model wavefunctions of the lowest admissible configurations, obtained with a realistic Hamiltonian. When harmonic oscillator single-particle wavefunctions are used, the contributions can be calculated rigorously but are generally less than the experimental values. Increases due to use of Woods-Saxon wavefunctions are estimated in a semi-quantitative way. For ^{17}O, the additional increase due to admixtures from higher configurations in the wavefunctions is also considered.

1. Introduction

Coulomb excitation has been used to measure B(E2) values for transitions between the ground and excited states of nuclei, and also the quadrupole moments of ground and excited states, through the reorientation effect. In order to obtain reliable values, allowance has to be made for the contribution to Coulomb excitation due to the second-order process involving virtual E1 transitions through higher excited states, including the giant dipole resonance (GDR).

Various essentially equivalent parameters have been used to measure the strength of this E1 contribution. These include τ_{if}, S(E1) and k, which are defined in the next section. Recently there have been several measurements and calculations of τ_{if} for light nuclei, and particularly for ^{7}Li.

Many of the calculations for ^{7}Li have made use of an α+t cluster description of the ^{7}Li states involved (Smilansky *et al.* 1972; Mertelmeier and Hofmann 1986; Fatemian *et al.* 1986). Häusser *et al.* (1973) pointed out that significant contributions to τ_{if} could come from other than α+t channels, and Mertelmeier and Hofmann found that inclusion of a component in the scattering functions from ^{6}Li+n channels increased τ_{if} by about 50%. Kajino *et al.* (1988) estimated that the α+t channel contributes only one third of the total τ_{if}. Häusser *et al.* used a schematic model to obtain an upper limit on τ_{if}. Shell model calculations of τ_{if} for ^{7}Li have been made by Barker (1982*a*) and Gomez-Comacho and Nagarajan (1985). Barker assumed that all the E1 strength was concentrated at a single energy, taken to be the GDR energy, and then used the closure

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approximation to simplify the calculation.* Gomez-Comacho and Nagarajan used a very simplified Hamiltonian, with harmonic oscillator single-particle wavefunctions. In both of these calculations, the value of τ_{if} was renormalised by an appreciable factor obtained by making E2 matrix elements agree with experiment.

Here we perform a shell model calculation with a realistic Hamiltonian (van Hees and Glaudemans 1983, 1984) and wavefunctions restricted to the lowest configuration, and include the contribution from each intermediate dipole state individually without making use of the closure approximation. Initially harmonic oscillator single-particle wavefunctions are used, in which case the calculation can be made consistently and rigorously. Changes due to the use of the more realistic Woods-Saxon wavefunctions are estimated.

The above discussion concentrates on ^{7}Li, but similar calculations are possible for other 1p-shell nuclei. We here include ^{6}Li, ^{10}B, ^{11}B and ^{12}C, as well as ^{17}O. All of these except ^{11}B were considered previously by Barker (1982*a*, *b*). In each case an experimental value of τ_{if} is available, or a value of τ_{if} has been required in order to analyse experimental data. For the ^{6}Li and ^{7}Li cases, we calculate τ_{if} corresponding not only to excitation of the first excited state but also to elastic scattering.

In the particular case of ^{17}O, because of the simple nature of the wavefunctions of the ground and first excited states, it is possible to estimate the effect on τ_{if} of the introduction of admixtures from higher configurations, chosen to fit E2 matrix elements.

Compared with the earlier shell model calculations of Barker (1982*a*, *b*), we do not here make the closure approximation, which was assumed for simplicity, and compared with the calculations of Gomez-Comacho and Nagarajan (1985) for ^{6}Li and ^{7}Li, we use a Hamiltonian and single-particle wavefunctions that are more realistic, and we cover a wider range of nuclei.

Section 3 describes the calculation of τ_{if} in the approximation that the initial and final states belong to the lowest shell model configuration and the single-particle wavefunctions are harmonic oscillator. The changes in τ_{if} due to the use of Woods-Saxon wavefunctions are given in Section 4, and the additional contribution to τ_{if} for ^{17}O due to admixtures from higher configurations in Section 5. These and other calculations are discussed in Section 6.

2. Formulae and Previous Values of τ_{if}

The tensor moment of the electric polarisability, τ_{if}, for the transition from the state i to the state f, is defined by (Lopes *et al.* 1983; Weller *et al.* 1985)

$$\tau_{if} = \tfrac{8}{9}\pi(\tfrac{10}{3})^{\frac{1}{2}} \sum_{n} W(1\,1 J_i J_f, 2 J_n)\langle i \,||\, \mathcal{M}(E1) \,||\, n\rangle\langle n \,||\, \mathcal{M}(E1) \,||\, f\rangle/(E_n - E_i), \qquad (1)$$

where the sum is over all intermediate states n that are connected by the E1 operator to both i and f. The reduced matrix elements are as defined by de-Shalit and Talmi (1963). The relation between τ_{if} and the quantity S(E1)

* Contrary to the implication by Gomez-Comacho and Nagarajan (1985), Barker (1982*a*) included excitation of nucleons from both the 1s and 1p shells.

defined by Häusser *et al.* (1973) is

$$\tau_{if} = \tfrac{8}{9}\pi(\tfrac{10}{3})^{\frac{1}{2}} S(\mathrm{E1}). \tag{2}$$

The parameter k (Häusser *et al.* 1972), which is the ratio of the actual dipole contribution to that calculated from the hydrodynamic model as empirically modified by Levinger (1957), is given by (Barker 1982*a*)

$$k = X/X_0, \tag{3}$$

where

$$X = S(\mathrm{E1})/\langle \mathrm{i} \,||\, \mathcal{M}(\mathrm{E2}) \,||\, \mathrm{f}\rangle \tag{4}$$

and

$$X_0 = 0 \cdot 00058 A/Z \; e\,\mathrm{MeV}^{-1}. \tag{5}$$

One also has

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

In most measurements, the state i is the ground state (labelled 1) and f is the first excited state (labelled 2), so that values of τ_{12} are obtained. Elastic scattering of aligned ^{7}Li nuclei has also provided a value of τ_{11}. We use values of S(E1) and k corresponding to τ_{12} only. Although the signs of τ_{11} and of k are necessarily significant, the sign of τ_{12} is significant only if the relative sign of the wavefunctions of the states 1 and 2 is fixed in some way—we choose this so that $\langle 1 \,||\, \mathcal{M}(\mathrm{E2}) \,||\, 2\rangle > 0$.

Experimental and calculated values of τ_{11}, τ_{12}, S(E1) and k previously obtained for light nuclei are collected in Table 1. In each case the original value is given without parentheses, and derived values are in parentheses. The value of $B(\mathrm{E2};\, 1 \rightarrow 2)$ used to relate S(E1) and k is also given.

3. Calculation of τ_{if} using Harmonic Oscillator Wavefunctions

We initially calculate τ_{if} using harmonic oscillator single-particle wavefunctions. The calculation is described in detail for ^{7}Li, and only the significant differences are given for other nuclei. A summary of the input experimental data and the resultant calculated values is given in Table 2, for each of the nuclei considered.

(a) ^{7}Li

In the expression (1) for τ_{if}, the state i for ^{7}Li is the $J^\pi T = \frac{3}{2}^- \frac{1}{2}$ ground state and f is either the ground state or the $\frac{1}{2}^- \frac{1}{2}$ first excited state. The intermediate states n are the complete set of nonspurious states with $J^\pi = \frac{1}{2}^+$, $\frac{3}{2}^+$ and $\frac{5}{2}^+$ and $T = \frac{1}{2}$ and $\frac{3}{2}$.

Van Hees and Glaudemans (1983, 1984) performed $(0+1)\hbar\omega$ shell model calculations of the properties of both normal and non-normal parity states of all the 1p-shell nuclei, using a translationally invariant treatment that

Table 1. Parameter values for E1 strength in Coulomb excitation

Nucleus	τ_{11} (fm^3)	τ_{12} (fm^3)	S(E1) (e^2fm^{2}MeV^{-1})	k	B(E2;1→2) (e^2fm^4)	Ref.
		(a) Experimental				
^{6}Li		(0·28)	(0·038)	3·9	24	A
		(0·26)	0·036	(3·6)	24	B
		(0·19)	(0·026)	2·6	25·6	C
^{7}Li		(0·20)	(0·027)	3·5±0·7	8·3	D
		(0·21)	0·028±0·004	(3·6)	8·3	E
		(0·12–0·19)	0·017–0·026	(2·3–3·4)	7·2–7·9	F
		(0·15)	(0·020)	2·7±0·2	7·42	G
	0·23±0·06	0·23±0·06	(0·031)	(4·0)	8·3	H
	−0·269	0·289	(0·039)	(4·9)	8·90	I
^{10}B		(0·039)	(0·0054)	1·3±0·3	1·81	J
		(0·039)	(0·0053)	1·3±0·2	1·79	K
^{17}O		(0·18)	(0·025)	5·7±0·4	2·10	L
		(b) Calculated				
^{6}Li			0·034[a]			E
			0·0198[b,c]	1·94[b,c]	25·6	M
			0·010[c]			N
^{7}Li			0·014[d]		7·0	O,E
			<0·05[a]			E
			0·0177[b,c]	2·27[b,c]	8·3	M
			0·012[c]			N
	−0·017, −0·11[b,d]	≈ $\lvert\tau_{11}\rvert$				P
	−0·084, −0·066[d]	0·092, 0·031[d]				P
	−0·120[e]	≈ $\lvert\tau_{11}\rvert$				P
		≈0·05[d]				Q
^{10}B			0·0050[b,c]	1·22[b,c]	1·81	M
^{12}C			0·0057[b,c]	0·77[b,c]	38·8	M
^{17}O			0·0113[b]	2·59[b]	2·10	R
^{18}O				≈4	31·9	S

[A] Disdier *et al.* (1971). [B] Häusser *et al.* (1973), using data of Disdier *et al.* (1971). [C] Gemmeke *et al.* (1978). [D] Häusser *et al.* (1972). [E] Häusser *et al.* (1973). [F] Häusser *et al.* (1973), using data of Bamberger *et al.* (1972). [G] Vermeer *et al.* (1984*a*, *b*). [H] Weller *et al.* (1985). [I] Kajino *et al.* (1988). [J] Vermeer *et al.* (1982). [K] Vermeer *et al.* (1983*b*). [L] Kuehner *et al.* (1982). [M] Barker (1982*a*). [N] Gomez-Comacho and Nagarajan (1985). [O] Smilansky *et al.* (1972). [P] Mertelmeier (1985), Mertelmeier and Hofmann (1986). [Q] Fatemian *et al.* (1986). [R] Barker (1982*b*). [S] Barker (1982*c*).

[a] Schematic model. [b] Uses closure approximation. [c] Renormalised to fit B(E2). [d] Includes α+t channel only. [e] Includes α+t, ^{6}Li+n channels.

completely eliminated spurious states. We have repeated their calculation, using the Oxford-Buenos Aires-MSU shell model code (Brown *et al.* 1986) to obtain the energies and wavefunctions of all the states i, f and n. As discussed in Barker and Woods (1985), the energies of the non-normal parity ^{7}Li states calculated with the van Hees and Glaudemans interaction appear to be about 4 MeV too high, so we reduce all calculated values of $E_n - E_i$ by 4 MeV. The E1 matrix elements are calculated with harmonic oscillator single-particle wavefunctions, using for ^{7}Li a length parameter b_0 chosen to fit the measured rms charge radius r_{ch} = 2.35 fm (Ajzenberg-Selove 1984).

Table 2. Values of quantities relevant to calculated E1 strength

Nucleus	^{6}Li	^{7}Li	^{10}B	^{11}B	^{12}C	^{17}O
Shift in energy (MeV)	−1	−4	0	0	0	−4
r_{ch} (fm)	2·51	2·35	2·45	2·42	2·47	2·71
Q (e fm^2)[A]	(−0·064)	−4·06	(8·47)	4·065	(6)	(−2·578)
B(E2) (e^2 fm^4)[A]	25·6	(7·42)	4·18	(2·6)	7·76	2·101
$(1 \|\| \mathcal{M}(E2) \|\| 2)$ (e fm^2)	8·76	5·47	3·54	3·15	6·23	3·55
			(a) Harmonic oscillator			
b_0 (fm)	1·891	1·737	1·658	1·630	1·635	1·761
$\Delta e/e$	0·548	0·397	0·305	0·274	0·232	0·524
Q (e fm^2)	0·09		7·90		5·79	−3·25
B(E2) (e^2 fm^4)		7·48		2·47		
τ_{11} (fm^3)	−0·0062	−0·0599				
τ_{12} (fm^3)	0·0667	0·0566	0·0353	0·0357	0·0360	0·0500
S(E1) (e^2 fm^2 MeV^{-1})	0·00908	0·00771	0·00481	0·00486	0·00491	0·00681
k	0·89	1·04	1·17	1·21	0·68	1·56
			(b) Woods-Saxon			
r_0 (fm)	2·014	1·855	1·416	1·439	0·412	1·256
R (fm)	3·44	3·37	2·95	3·10	3·14	3·16
$\Delta e/e$	0·183	0·218	0·078	0·245	0·158	0·389
Q (e fm^2)	0·11		5·54		5·74	−2·76
B(E2) (e^2 fm^4)		7·32		3·86		
τ_{11} (fm^3)	−0·015	−0·125				
τ_{12} (fm^3)	0·186	0·115	0·066	0·038	0·034	0·125
S(E1) (e^2 fm^2 MeV^{-1})	0·0254	0·0156	0·0090	0·0051	0·0046	0·0170
k	2·5	2·1	2·2	1·27	0·63	3·9[B]

[A] Experimental values. Values not in parentheses used for deriving Δe.
[B] With inclusion of configuration mixing, $k = 5·3$ (see Section 5).

The E1 matrix elements calculated with harmonic oscillator wavefunctions satisfy various relations, which can be used to test the completeness of the set of intermediate states n. From Barker (1982*a*), equation (8) or (14) together with (5) and (18), one has

$$\sum_{nJ_nT_n} W(11J_iJ_f, 2J_n)\langle J_i^- \tfrac{1}{2} \|\| \mathcal{M}(E1) \|\| J_n^+ T_n\rangle\langle J_n^+ T_n \|\| \mathcal{M}(E1) \|\| J_f^- \tfrac{1}{2}\rangle = 0$$
$$(J_i = \tfrac{1}{2}, \tfrac{3}{2}; J_f = \tfrac{1}{2}, \tfrac{3}{2}). \qquad (7)$$

Because the operator $\mathcal{O}^1$ analogous to $\mathcal{O}^2$ (defined in equation (6) of Barker 1982*a*) is identically zero, one also has

$$\sum_{nJ_nT_n} W(11J_iJ_f, 1J_n)\langle J_i^- \tfrac{1}{2} \|\| \mathcal{M}(E1) \|\| J_n^+ T_n\rangle\langle J_n^+ T_n \|\| \mathcal{M}(E1) \|\| J_f^- \tfrac{1}{2}\rangle = 0$$
$$(J_i = \tfrac{1}{2}, \tfrac{3}{2}; J_f = \tfrac{1}{2}, \tfrac{3}{2}). \qquad (8)$$

From equations (7) and (8) for $J_i \neq J_f$, it follows that for each value of J_n

$$\sum_{nT_n} \langle J_i^- \tfrac{1}{2} \|\| \mathcal{M}(E1) \|\| J_n^+ T_n\rangle\langle J_n^+ T_n \|\| \mathcal{M}(E1) \|\| J_f^- \tfrac{1}{2}\rangle = 0 \qquad (J_i \neq J_f). \qquad (9)$$

The matrix elements also satisfy the sum rule

$$\sum_{nT_n} \langle J_i^- T_i \|\| \mathcal{M}(E1) \|\| J_n^+ T_n\rangle^2 = \frac{3}{8\pi}(2J_n + 1)\frac{NZ}{A} e^2 b_0^2. \qquad (10)$$

The values of τ_{if} obtained for ^{7}Li are

$$\tau_{11} = -0\cdot0599 \text{ fm}^3, \qquad \tau_{12} = 0\cdot0566 \text{ fm}^3. \tag{11}$$

The contributions from the individual intermediate states are shown in Fig. 1. From equation (2), one has

$$S(\text{E1}) = 0\cdot00771\ e^2 \text{ fm}^2 \text{ MeV}^{-1}. \tag{12}$$

In order to calculate the corresponding value of k, a value is required for the E2 matrix element $\langle \frac{3}{2}^- \frac{1}{2} \| \mathcal{M}(\text{E2}) \| \frac{1}{2}^- \frac{1}{2} \rangle$. It is well known that E2 matrix elements in 1p-shell nuclei calculated with wavefunctions belonging to the lowest configuration are greatly underestimated, and that effective charges are required in order to obtain agreement with experimental values. We choose the additional isoscalar charge in order to fit the measured quadrupole moment of ^{7}Li ground state, taken as $Q(\frac{3}{2}^-) = -4.06 \pm 0.08\ e \text{ fm}^2$ (Sundholm *et al.* 1984; Diercksen *et al.* 1988), giving $\Delta e = 0.397e$. In comparison, van Hees and Glaudemans (1984) found that $\Delta e = 0.35e$ gave good agreement with E2 moments and transition rates throughout the 1p shell. Then

$$\langle \tfrac{3}{2}^- \tfrac{1}{2} \| \mathcal{M}(\text{E2}) \| \tfrac{1}{2}^- \tfrac{1}{2} \rangle = 5\cdot47\ e \text{ fm}^2, \tag{13}$$

which corresponds to $B(\text{E2};\ \frac{3}{2}^- \frac{1}{2} \rightarrow \frac{1}{2}^- \frac{1}{2}) = 7.48\ e^2 \text{ fm}^4$. This value of B(E2) agrees well with the experimental value of $7.42 \pm 0.14\ e^2$ fm^4 obtained by Vermeer *et al.* (1984*a*, *b*). The above experimental values of Q and B(E2) are probably the most reliable; these and other values are given in Table 2 of Mertelmeier and Hofmann (1986). From equations (3)–(5), (12) and (13), one obtains

$$k = 1\cdot04. \tag{14}$$

(b) ^{6}Li

For ^{6}Li, the values of $E_n - E_i$ calculated with the van Hees and Glaudemans interaction are reduced by 1 MeV (Barker and Woods 1985). The measured charge radius is taken from Ajzenberg-Selove (1984), and Δe is determined by fitting the measured value $B(\text{E2};\ 1^+\ 0 \rightarrow 3^+0) = 25.6 \pm 1.6\ e^2$ fm^4 (Eigenbrod 1969; Ajzenberg-Selove 1984). The calculated value of $Q(1^+)$ is then 0.093 e fm^2, compared with the experimental value $-0.0644 \pm 0.0007\ e$ fm^2 (Ajzenberg-Selove 1984); the difference, including the different sign, is not very significant because of the small magnitude of $Q(1^+)$. Individual contributions to τ_{11} and τ_{12} are shown in Fig. 2, and the relevant values and results are given in Table 2.

(c) ^{10}B

An adequate shell model description of the 1^+0 first excited state of ^{10}B is not obtained easily. Warburton *et al.* (1968) pointed out that there are significant discrepancies between experimental values of electromagnetic matrix elements involving this state and those calculated with the interactions

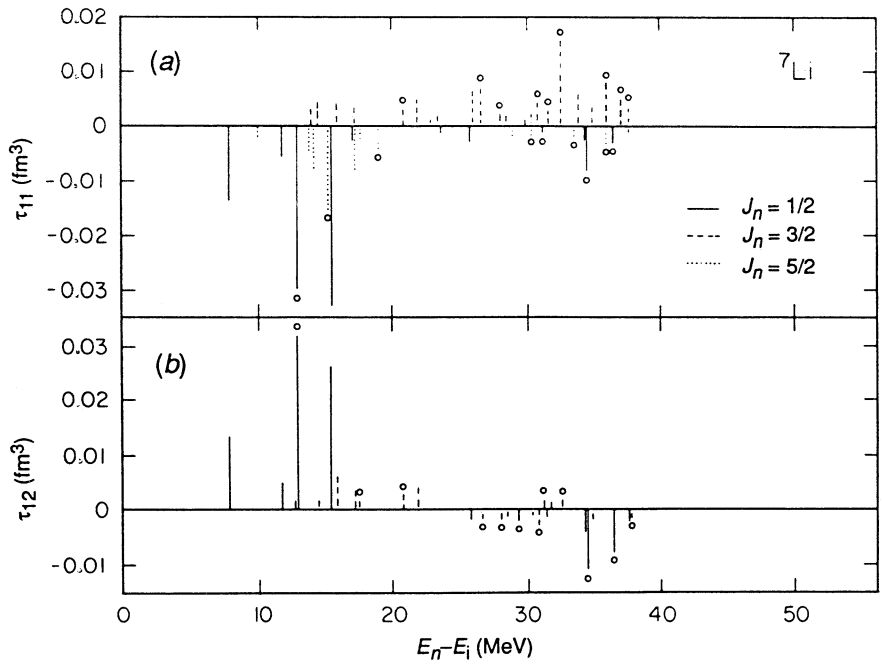

Fig. 1. Contributions from individual intermediate dipole states of ^{7}Li to (*a*) τ_{11} and (*b*) τ_{12}, calculated with harmonic oscillator wavefunctions. Contributions with magnitude less than 0·001 fm^3 are omitted. A circle denotes $T_n = \frac{3}{2}$, otherwise $T_n = \frac{1}{2}$.

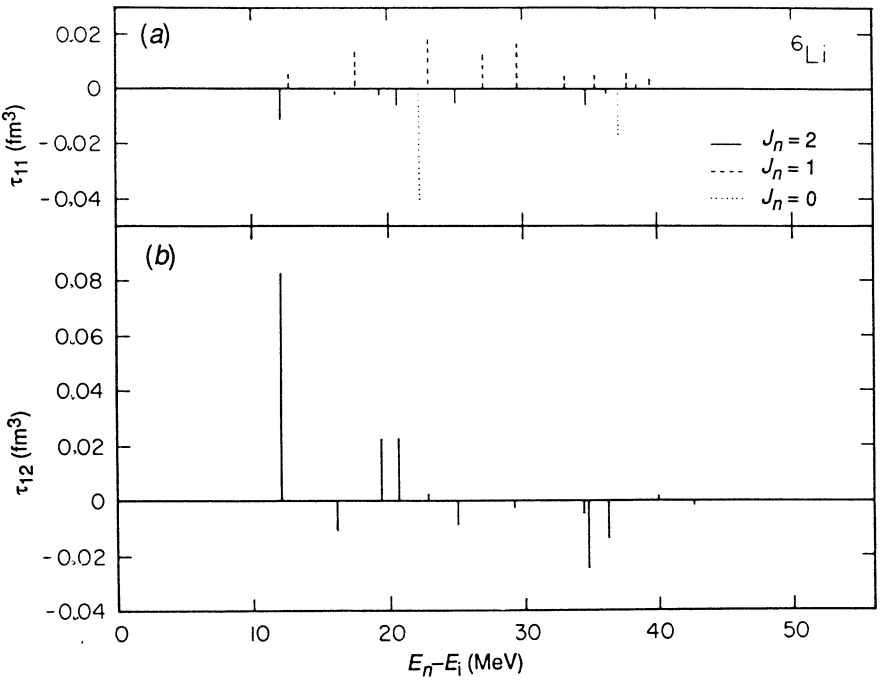

Fig. 2. Contributions from individual intermediate dipole states of ^{6}Li to (*a*) τ_{11} and (*b*) τ_{12}, calculated with harmonic oscillator wavefunctions. Contributions with magnitude less than 0.001 fm^3 are omitted. All states have $T_n = 1$.

of Cohen and Kurath (1965), and that much improved agreement can be obtained by mixing the lowest 1^+ 0 state with the next higher 1^+0 state. Kurath (1979) provided a qualitative understanding of this result, and Barker (1981) showed that such mixing could be obtained by modifying the values of some of Cohen and Kurath's interaction parameters. The van Hees and Glaudemans interaction automatically gives a satisfactory description of the two lowest 1^+0 states, the reason apparently being that although most of the two-particle matrix elements (written in the *LS*-coupling representation) have values close to those of Cohen and Kurath, particularly for the (6-16)2BME interaction, there is a marked difference in the value of $\langle LSJT|V|\bar{L}\bar{S}JT\rangle = \langle 1010|V|1010\rangle$. These values are (in MeV):

van Hees and Glaudemans	Cohen and Kurath (6–16)2BME	(8–16)2BME	(8–16)POT
4·11	0·07	0·05	−0·29

Kumar (1974) has argued that this matrix element should have a large positive value, about 2.5–3 MeV, in order to fit observed properties of the lowest $T = 2$ state of ^{8}Be, and on this basis Barker (1981) chose his parameter set III, which was used in the calculation of k for ^{10}B in Barker (1982*a*).

Thus we use the van Hees and Glaudemans wavefunctions for the 3^+0 ground state and 1^+0 first excited state of ^{10}B, and also for the 2^-1 intermediate states. The calculated energy of the lowest 2^-1 state is 7.73 MeV. Although there is some uncertainty and confusion regarding the $J^\pi T$ assignments for ^{10}B levels in the 7–8 MeV region (Ajzenberg-Selove 1984), it is believed that the lowest 2^- 1 state is at 7.79 MeV (Barker, to be published). Thus we do not modify the calculated values of $E_n - E_i$. The charge radius is taken from Ajzenberg-Selove (1979). With Δe chosen to fit the measured value $B(\text{E2};1^+0\rightarrow 3^+0) = 4.18\pm0.02\ e^2\ \text{fm}^4$ (Vermeer *et al.* 1983*b*; Ajzenberg-Selove 1984), one finds $Q(3^+0) = 7.90\ e\ \text{fm}^2$, compared with the experimental value of $8.47\pm0.06\ e\ \text{fm}^2$ (Ajzenberg-Selove 1984), and $Q(1^+0) = -0.61\ e\ \text{fm}^2$. Individual contributions to τ_{12} are shown in Fig. 3*a*.

(d) ^{11}B

The spins and parities of the states involved in the excitation of the first excited state of ^{11}B are the same as those in ^{7}Li. The energies of the lowest $\frac{1}{2}^+\frac{1}{2}$, $\frac{3}{2}^+\frac{1}{2}$, $\frac{5}{2}^+\frac{1}{2}$, $\frac{1}{2}^+\frac{3}{2}$ and $\frac{5}{2}^+\frac{3}{2}$ states calculated for the van Hees and Glaudemans interaction are 7.01, 9.20, 6.89, 12.97 and 14.75 MeV, compared with experimental values of 6.79, 7.98, 7.29, 12.56 and 14.34 MeV (Ajzenberg-Selove 1985). The states that have the largest E1 matrix elements with the ground state are a $\frac{3}{2}^+\frac{3}{2}$ state at 22.25 MeV and a $\frac{5}{2}^+\frac{3}{2}$ state at 25.36 MeV, while the giant dipole resonance is observed at 25–28 MeV (Ajzenberg-Selove 1985). We therefore use the calculated energies unchanged. The charge radius is given in Ajzenberg-Selove (1985). The value of Δe is chosen to fit the ground state quadrupole moment $Q(\frac{3}{2}^-\frac{1}{2}) = 4.065\pm0.026\ e\,\text{fm}^2$ (Ajzenberg-Selove 1985). Then $B(\text{E2};\ \frac{3}{2}^-\frac{1}{2}\rightarrow\frac{1}{2}^-\frac{1}{2}) = 2.47\ e^2\ \text{fm}^4$, which agrees satisfactorily with the only measured value of $2.6\pm0.4\ e^2\ \text{fm}^4$ (Fewell *et al.* 1980, 1984). Individual contributions to τ_{12} are shown in Fig. 3*b*.

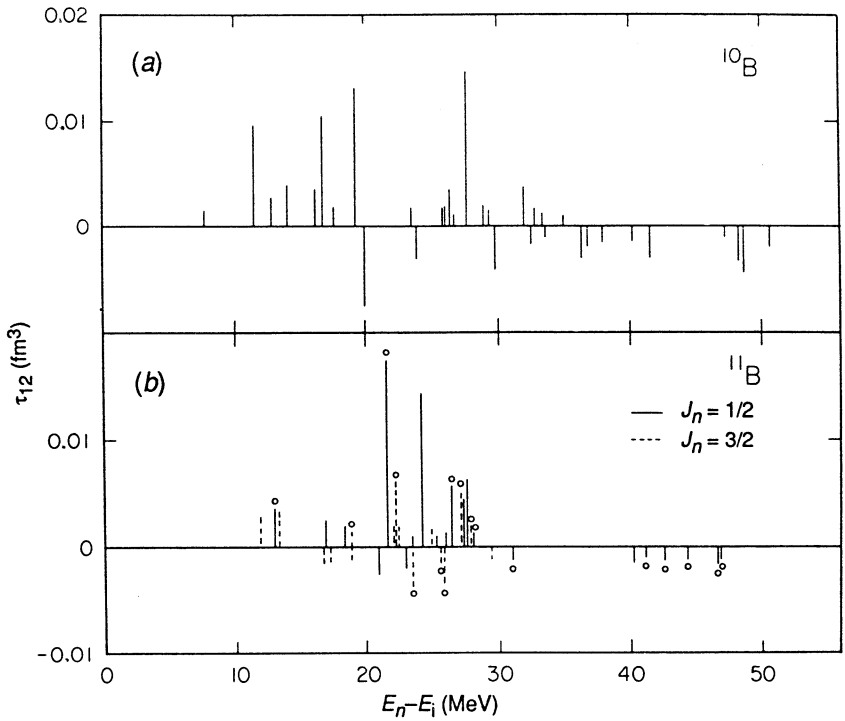

Fig. 3. Contributions to τ_{12} from individual intermediate dipole states of (*a*) ^{10}B and (*b*) ^{11}B, calculated with harmonic oscillator wavefunctions. Contributions with magnitude less than 0.001 fm^3 are omitted.

(e) ^{12}C

For the transition from the 0^+ 0 ground state of ^{12}C to the 2^+ 0 first excited state, the intermediate dipole states are the 1^-1 states. The van Hees and Glaudemans interaction predicts the lowest 1^-1 state at 18.26 MeV, compared with the experimental value 17.23 MeV (Ajzenberg-Selove 1985). The calculated levels associated with the GDR are those at 23.08, 24.99 and 26.00 MeV, while the experimental peaks are observed at 23.2 and 25.6 MeV (Ahrens *et al.* 1975; Ajzenberg-Selove 1985). We therefore use the calculated values of $E_n - E_i$ unchanged. The charge radius is taken from Ajzenberg-Selove (1985). The measured value $B(\mathrm{E2}; 2^+0 \rightarrow 0^+0) = 7.76 \pm 0.43\ e^2\ \mathrm{fm}^4$ (Ajzenberg-Selove 1985) gives Δe, from which one obtains $Q(2^+\ 0) = 5.79\ e\ \mathrm{fm}^2$, compared with the experimental value $6 \pm 3\ e\ \mathrm{fm}^2$ (Vermeer *et al.* 1983*a*; Ajzenberg-Selove 1985). Individual contributions to τ_{12} are shown in Fig. 4*a*.

(f) ^{17}O

The case of ^{17}O is exceptional, since it is not a 1p-shell nucleus and not all the interaction matrix elements required for a complete $(0+1)\hbar\omega$ calculation are provided by van Hees and Glaudemans. Nevertheless a reasonable calculation of τ_{12} is possible. In the spirit of the present section, we assume that the $\frac{5}{2}^+\frac{1}{2}$ ground and $\frac{1}{2}^+\frac{1}{2}$ first excited states of ^{17}O belong to the lowest configurations, namely $1s^4 1p^{12} 1d_{5/2}$ and $1s^4 1p^{12} 2s_{1/2}$ respectively. Relative to the ^{16}O

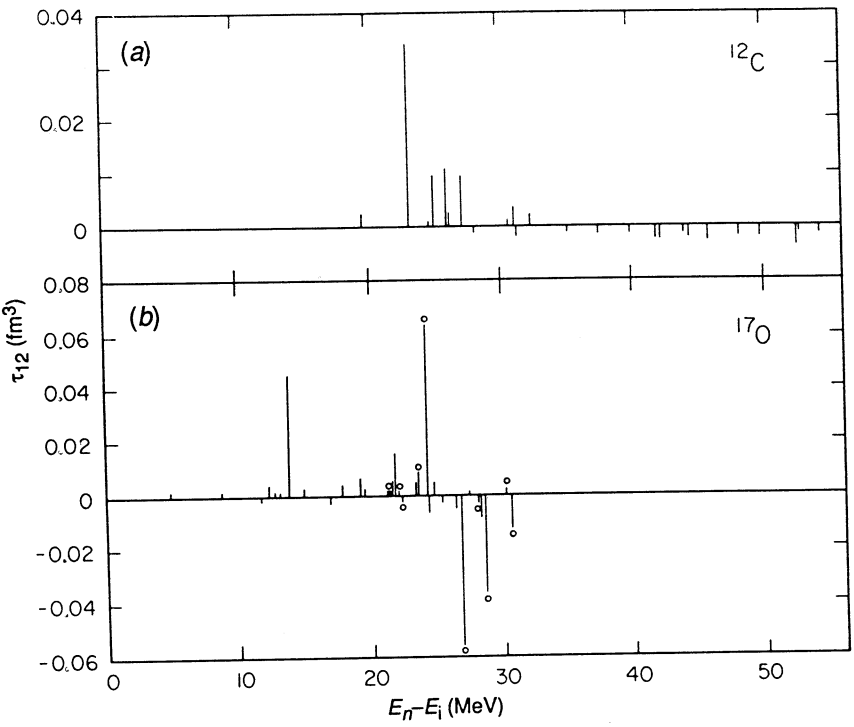

Fig. 4. Contributions to τ_{12} from individual intermediate dipole states of (*a*) ^{12}C and (*b*) ^{17}O, calculated with harmonic oscillator wavefunctions. Contributions with magnitude less than 0.001 fm^3 are omitted.

closed-shell ground state, these are one-particle (1p) states. The intermediate states are $\frac{3}{2}^-\,\frac{1}{2}$ and $\frac{3}{2}^-\,\frac{3}{2}$, and belong either to the 1p configuration $1s^4 1p^{12} 2p_{3/2}$, which contains only a single $\frac{3}{2}^-\,\frac{1}{2}$ state, or to the two-particle one-hole (2p–1h) configurations $1s^4 1p^{11}(2s,1d)^2$. These configurations contain three spurious states, all with $T = \frac{1}{2}$, which are eliminated by adding to the Hamiltonian (the SG16F interaction of Brown *et al.* 1986) a large multiple of the c.m. energy. The single-particle energies are adjusted to fit the observed energies of the pigmy and giant dipole resonances observed in $^{17}O(\gamma,n)^{16}O$. The pigmy resonance should be due essentially to the excitation of the valence neutron in the ^{17}O ground state from the $1d_{5/2}$ orbit to the $2p_{3/2}$, $1f_{7/2}$ or $1f_{5/2}$ orbits, while the GDR mainly involves excitations of nucleons from the p^{12} core, leading to 2p–1h states. The GDR is observed at about 23 MeV (Ajzenberg-Selove 1986). With the single-particle energies given by Brown *et al.*, the $\frac{3}{2}^-$ states with the largest E1 matrix elements to the ground state are at 25.74 MeV ($T = \frac{1}{2}$) and 28.23 MeV ($T = \frac{3}{2}$). Also the lowest $\frac{3}{2}^-\,\frac{1}{2}$ (mainly 2p–1h) and $\frac{3}{2}^-\,\frac{3}{2}$ states are calculated at 8.82 and 15.20 MeV compared with the observed values of 4.55 and 12.47 MeV. Thus we reduce the $2s_{1/2}$, $1d_{5/2}$ and $1d_{3/2}$ single-particle energies of Brown *et al.* by 4 MeV. Jury *et al.* (1985) interpreted their observations of the pigmy resonance in terms of $\frac{5}{2}^-$, $\frac{3}{2}^-$, $\frac{3}{2}^-$, $\frac{7}{2}^-$ and $\frac{7}{2}^-$ levels at 10.5, 13.0, 14.0, 16.6 and 21.0 MeV respectively. We therefore reduce the $2p_{3/2}$ single-particle energy by 1 MeV in order to locate the strength due

to the 1p state at about 13.5 MeV; this energy seems to be consistent with other evidence (Johnson 1973; Darden *et al.* 1973).

For the value b_0 = 1.761 fm chosen to fit the measured charge radius r_{ch} = 2.710±0.015 fm (Ajzenberg-Selove 1986), the calculated rms radius for the $d_{5/2}$ neutron is 3.33 fm, which is consistent with the value 3.31±0.12 fm measured by Hicks (1982). The relation corresponding to equation (7) gives

$$\sum_{nT_n} \langle \tfrac{5}{2}^+ \tfrac{1}{2} \| \mathcal{M}(\mathrm{E1}) \| \tfrac{3}{2}^-_n T_n \rangle \langle \tfrac{3}{2}^-_n T_n \| \mathcal{M}(\mathrm{E1}) \| \tfrac{1}{2}^+ \tfrac{1}{2} \rangle = 0 . \qquad (15)$$

The contributions to τ_{12} from individual dipole states are shown in Fig. 4*b*. The value of Δe that fits the experimental value $B(\mathrm{E2}; \tfrac{5}{2}^+ \tfrac{1}{2} \rightarrow \tfrac{1}{2}^+ \tfrac{1}{2})$ = 2.101±0.021 e^2 fm^4 (Ajzenberg-Selove 1986) gives $Q(\tfrac{5}{2}^+ \tfrac{1}{2})$ = −3.25 e fm^2, compared with the experimental value of −2.578 e fm^2 (Ajzenberg-Selove 1986).

(g) Discussion

From Figs 1–4, it is seen that for each of the nuclei considered here the contributions to τ_{12} tend to be positive at low energies and negative at high energies, although this is more pronounced in some cases than in others. Because of relations such as (7) or (15), the low-energy contributions dominate and the resultant τ_{12} values are positive, leading to positive values of k. In ^{11}B and ^{12}C the main contributions come from the GDR region around 25 MeV, whereas for ^{6}Li and ^{7}Li, and to some extent ^{10}B and ^{17}O, they come from lower energies. As regards τ_{11}, the negative contributions at low energies dominate for ^{7}Li, but for ^{6}Li the significant contributions are widely distributed and largely cancel—this is presumably due to the mainly L = 0 nature of the ^{6}Li ground state.

4. Use of Woods-Saxon Wavefunctions

So far all E1 matrix elements have been calculated with harmonic oscillator single-particle wavefunctions. Long ago, Lane (1960) pointed out that this may lead to appreciable errors because the wavefunctions do not have the correct asymptotic form, and Barker (1961) showed that use of wavefunctions that do have the correct asymptotic form could greatly change the oscillator values of E1 matrix elements in ^{13}C. These calculations were put on a more rigorous basis by Barker and Ferdous (1980), for E1 transitions in both ^{13}C and ^{13}N. Millener *et al.* (1983) considered other cases of strong E1 transitions between low-lying levels of light nuclei, where use of Woods-Saxon wavefunctions increased B(E1) values calculated with harmonic oscillator wavefunctions by factors of up to 50. The increase is generally associated with an increase (by a factor of order 2–3) in the E1 single-particle matrix element (SPME) involving a 1p → 2s transition of a nucleon that is loosely bound, whose wavefunction therefore extends out to large distances. Similar increases can also be obtained when one of the nucleons is slightly unbound (Barker 1984).

We therefore consider the changes in our calculations produced by use of Woods-Saxon wavefunctions. In their calculations, Barker and Ferdous (1980) and Millener *et al.* (1983) expressed each E1 matrix element in terms of contributions from single-particle transitions corresponding to different core

states, so that the single-particle wavefunctions could be calculated with the appropriate binding energy. It is impracticable to do this for each of the E1 matrix elements that we need (the numbers range from 51 in ^{6}Li to 930 in ^{11}B). Consequently we proceed in an approximate way.

For each 1p-shell nucleus, the ground state and first excited state belong to the configuration $1s^4 1p^{A-4}$, and the intermediate dipole states to the configurations $1s^3 1p^{A-3}$ and $1s^4 1p^{A-5}(2s,1d)^1$, so that the E1 matrix elements involve SPME for the excitations $1s \rightarrow 1p$, $1p \rightarrow 2s$ and $1p \rightarrow 1d$. Since (p, 2p) experiments on 1p-shell nuclei suggest that the 1s protons are bound by over 30 MeV (Tyrén *et al.* 1966), the 1s wavefunctions do not extend to large distances and should be well represented by harmonic oscillator wavefunctions; we therefore assume that the $1s \rightarrow 1p$ SPME are unchanged.

The $1p \rightarrow 2s$ and $1p \rightarrow 1d$ SPME are calculated using wavefunctions for a central Woods-Saxon potential, with the conventional diffuseness (0.65 fm), a radius parameter r_0 ($R = r_0 (A-1)^{1/3}$) chosen to fit the measured charge radius r_{ch} (using formulae (10) and (11) of Millener *et al.* 1983, with the spectroscopic factors calculated for the van Hees and Glaudemans interaction), and depth appropriate to the particular intermediate state and core state. The separate contributions to the SPME are proportional to

$$\int_0^\infty r u_i(r)\, u_f(r)\, dr \Big/ \left(\int_0^\infty u_i^2(r)\, dr \int_0^\infty u_f^2(r)\, dr \right)^{\frac{1}{2}}, \tag{16}$$

with the proviso that if the state q (q=i, f) is unbound, then $\int_0^\infty u_q^2(r)\, dr$ in (16) is replaced by $(1 + \gamma_q^2 dS_q/dE) \int_0^a u_q^2(r)\, dr$, where a is the channel radius (Barker and Ferdous 1980). We choose a = 5.0 fm (cf. Barker and Ferdous 1980); smaller values of a give larger enhancements. Following the work of Holt *et al.* (1978), we use in (16) the resonant part of the continuum wavefunction with a complex asymptotic form corresponding to outgoing waves (see Barker and Ferdous, note added in proof), rather than the real asymptotic form given in equation (15) of Barker and Ferdous. Then the SPME are complex, and we take τ_{12} to be the magnitude of the expression in equation (1).

From a general survey of the dependence of the $1p \rightarrow 2s$ and $1p \rightarrow 1d$ SPME on the energies of the states involved, with the 1p state bound and the 2s or 1d state bound or unbound, as is appropriate for the cases of interest here, it appears that the greatest enhancements over the harmonic oscillator values are found when the 1p nucleon is loosely bound, and the 2s or 1d nucleon is slightly unbound. The enhancements are appreciable over a wide range of energies (many MeV), although reductions may be found for very unbound states. In general the enhancements are greater for $1p \rightarrow 2s$ than for $1p \rightarrow 1d$. An exceptional case concerns a loosely-bound s-wave neutron state, since the SPME approaches zero as the binding energy goes to zero (Barker 1984).

Because of the dependence of the SPME on the energies of the states concerned, and because the corresponding one-body density matrix elements (OBDME—see Millener *et al.* 1983) tend to be largest for states lying near the energy of the core+nucleon system, we expect the greatest enhancements to occur for the low-lying intermediate states, particularly if the initial and final states are not too tightly bound with respect to the core states with

which they have appreciable spectroscopic factors. This is illustrated in the particular cases that follow, where again ^{7}Li is discussed in the greatest detail. A slightly different procedure is used for ^{17}O. The results are summarised in Table 2.

For consistency, we also calculate the E2 matrix elements in a similar way using Woods-Saxon wavefunctions. This does not change the relative values of τ_{12} and k, because we make use of experimental values of the E2 matrix elements, but it does lead to changed values of the additional charges Δe required to fit these matrix elements.

(a) ^{7}Li

The van Hees and Glaudemans interaction, like other shell model interactions for light nuclei, predicts that only a few low-lying $A = 6$ core states have appreciable 1p spectroscopic factors for the $\frac{3}{2}^-\frac{1}{2}$ ground state and $\frac{1}{2}^-\frac{1}{2}$ first excited state of ^{7}Li, namely the lowest 1^+ 0, 3^+ 0, 0^+ 1 and 2^+ 1 states for $\frac{3}{2}^-\frac{1}{2}$, and the lowest 1^+ 0, 2^+ 0, 0^+ 1 and 2^+ 1 states for $\frac{1}{2}^-\frac{1}{2}$. These extend over the ^{7}Li excitation energy range from about 7 to 13 MeV (see Fig. 5). It is only the positive-parity ^{7}Li intermediate states with energies not too far from this region, and with loosely bound (or unbound) 2s or 1d nucleons, for which appreciable enhancements of the SPME might be expected. It is seen from Fig. 1 that this energy region contains the intermediate states giving the largest contributions to τ_{11} and τ_{12} , namely the $\frac{1}{2}^+\frac{1}{2}1$, $\frac{1}{2}^+\frac{3}{2}1$, $\frac{5}{2}^+\frac{3}{2}1$ and $\frac{1}{2}^+\frac{1}{2}4$ states (which are also shown on Fig. 5), and only for these do we consider possible changes in the SPME values.

For these four states, the quantity (16) calculated with Woods-Saxon wavefunctions, with $r_0 = 1.855$ fm ($R = 3.37$ fm) to fit r_{ch}, is enhanced over the harmonic oscillator values by factors ranging in magnitude up to 2.2 for the $1p \rightarrow 2s$ excitation, and up to 1.4 for $1p \rightarrow 1d$. The additional contributions to τ_{11} from these four states are -0.014, -0.028, -0.008 and -0.015 fm^3 respectively, and to τ_{12} approximately 0.017, 0.029, 0.0 and 0.012 fm^3, giving

$$\tau_{11} = -0\cdot125 \text{ fm}^3, \qquad \tau_{12} = 0\cdot115 \text{ fm}^3, \tag{17}$$

and

$$S(\text{E1}) = 0\cdot0156\, e^2 \text{ fm}^2 \text{ MeV}^{-1}, \qquad k = 2\cdot11. \tag{18}$$

It is seen from Table 2 that the value of Δe required to fit the experimental value of $Q(\frac{3}{2}^-)$ when Woods-Saxon wavefunctions are used is considerably smaller than that obtained using harmonic oscillator wavefunctions. This value of Δe gives $B(\text{E2};\frac{3}{2}^-\frac{1}{2} \rightarrow \frac{1}{2}^-\frac{1}{2}) = 7.32\ e^2$ fm^4.

(b) ^{6}Li

The only $A = 5$ core states having non-zero p-wave spectroscopic factors for the 1^+0 ground state and 3^+0 first excited state of ^{6}Li are the $\frac{3}{2}^-\frac{1}{2}$ ground state and $\frac{1}{2}^-\frac{1}{2}$ first excited state (which we take at an excitation energy of 2.5 MeV in both ^{5}He and ^{5}Li—see Barker and Woods 1985). From Fig. 2, it is

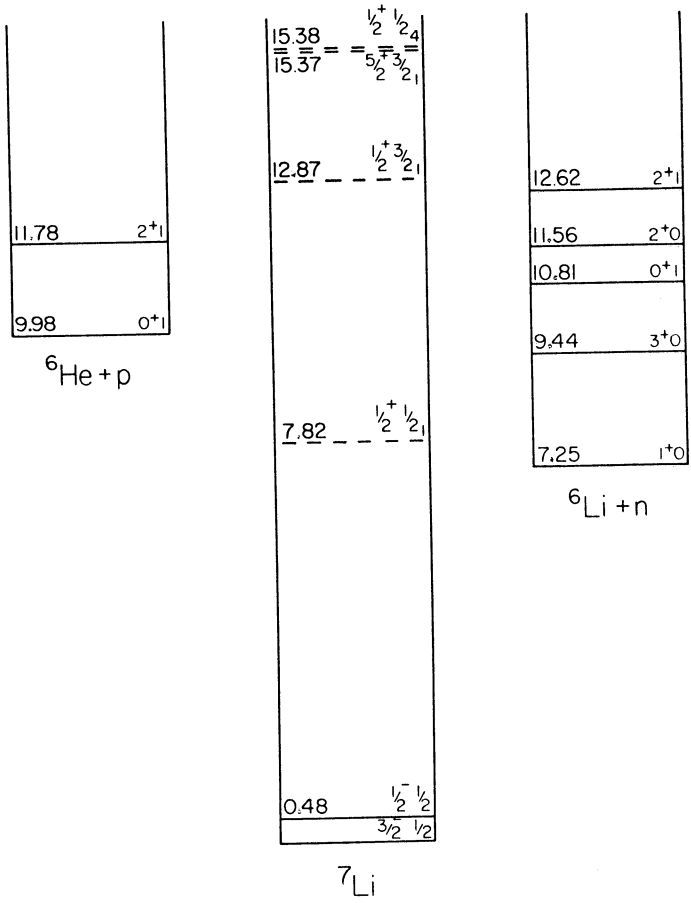

Fig. 5. Energy levels relevant to Coulomb excitation of ^{7}Li. All energies (in MeV) are given relative to the ^{7}Li ground state. Solid lines indicate experimental levels and energies (from Ajzenberg-Selove 1984), dashed lines indicate calculated intermediate dipole states.

seen that the intermediate state giving by far the largest contribution to τ_{12} is the 2^- 1_1 state at 12.07 MeV; we also consider changes for the next three 2^- 1 states at 16.24, 19.39 and 20.70 MeV, which give moderate contributions. For τ_{11}, in addition to these four 2^- 1 states, we consider changes in the contribution from the 1^- 1_2 state at 17.72 MeV; although the 0^-1_3 state at 22.31 MeV and the 1^- 1_4 state at 22.95 MeV give larger contributions, these come mainly from $1s \rightarrow 1p$ excitations, which we are assuming to be unchanged. The relevant states are shown in Fig. 6.

The enhancement of the quantity (16) for the $1p \rightarrow 2s$ excitations ranges up to about 1.7, and for $1p \rightarrow 1d$ up to 1.5. The additional contributions to τ_{11} from the 2^- $1_{1\text{-}4}$ and 1^- 1_2 states are −0.014, −0.001, 0.001, −0.002 and 0.008 fm^3 respectively, and to τ_{12} from 2^- $1_{1\text{-}4}$ are approximately 0.118, −0.003, −0.005 and 0.007 fm^3. The results are given in Table 2.

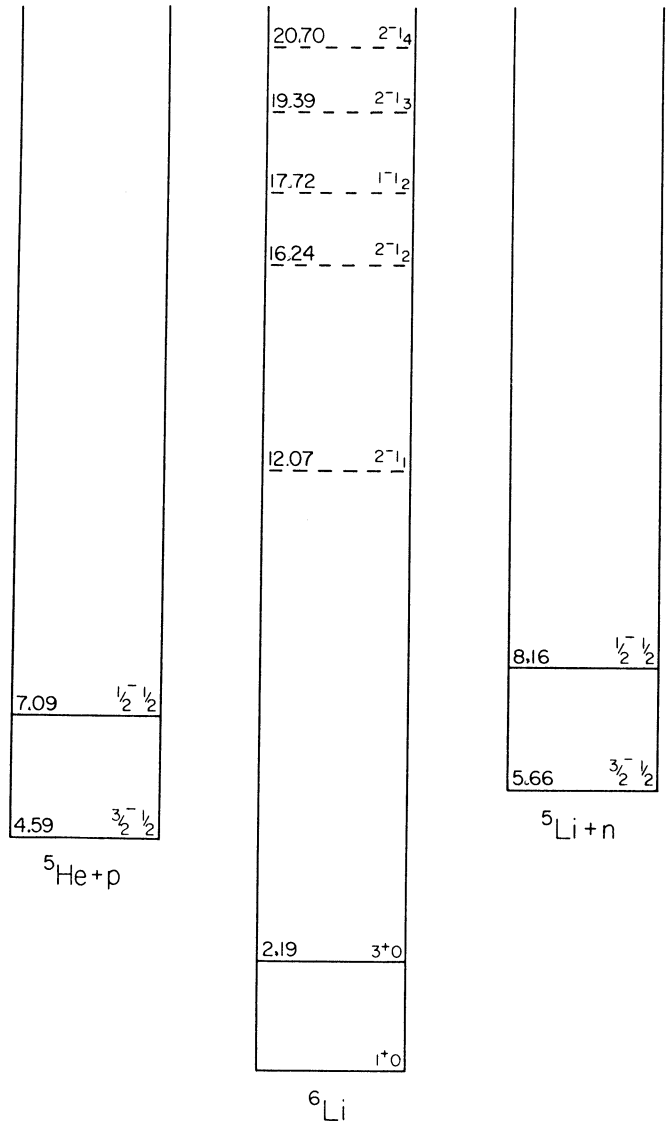

Fig. 6. Energy levels relevant to Coulomb excitation of ^{6}Li. See also Fig. 5 caption.

(c) ^{10}B

For the van Hees and Glaudemans interaction that we are using, the p-wave spectroscopic factors of the 3^+ 0 ground state and 1^+ 0 first excited state of ^{10}B are large for many $A = 9$ core states ($\frac{1}{2}^-\frac{1}{2}1\text{–}3$, $\frac{3}{2}^-\frac{1}{2}1\text{–}4$, $\frac{5}{2}^-\frac{1}{2}1,3,4$, $\frac{7}{2}^-\frac{1}{2}1,2$). The low-lying 2^- 1 intermediate states that give the main contributions to τ_{12} in general have small s- and d-wave spectroscopic factors except for the lowest state of each J, so that we include changes to the SPME's only for these core states. We consider changes only for the 2^-1_n states with $n = 1, 2, 6, 9$ and

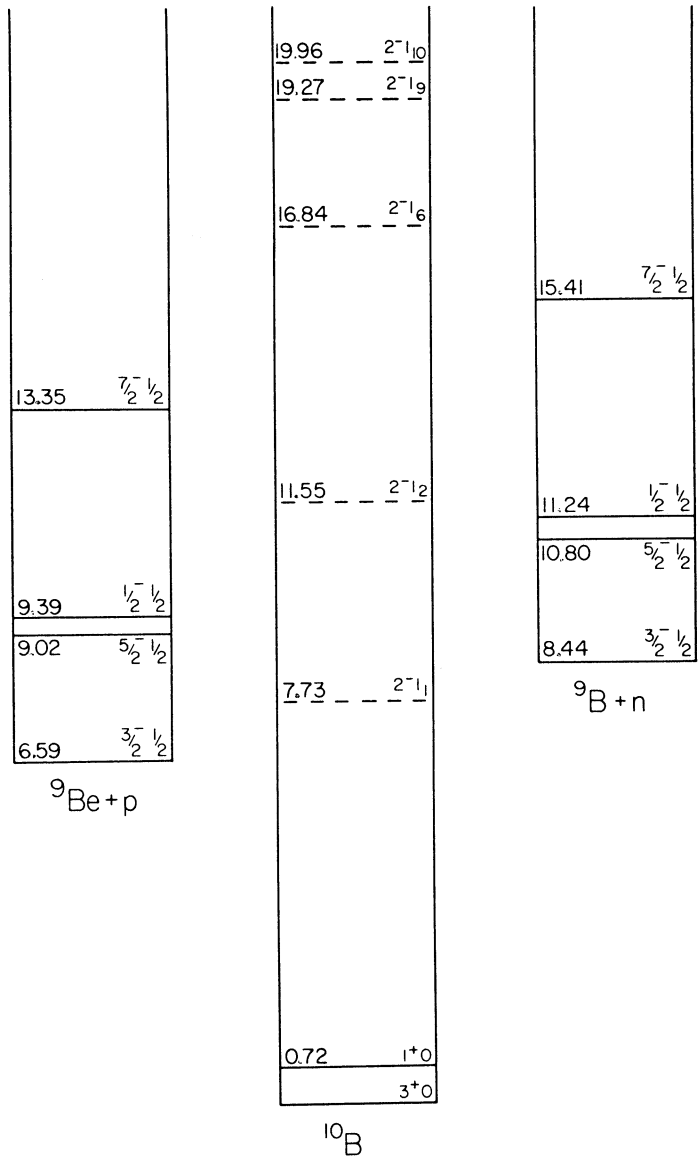

Fig. 7. Energy levels relevant to Coulomb excitation of ^{10}B. See also Fig. 5 caption.

10 (shown in Fig. 7), as these give the largest contributions to τ_{12} or seem likely to give large changes, due to appreciable cancellations in the harmonic oscillator values. The large contribution from the $n = 28$ state at 27.74 MeV (see Fig. 3*a*) is due mainly to $1s \rightarrow 1p$ excitations and to $1p_{3/2} \rightarrow 1d_{5/2}$ excitations with high-lying $A = 9$ core states. The additional contributions to τ_{12} from these five 2^- 1 states are approximately 0.010, 0.016, 0.005, 0.002 and -0.002 fm^2 respectively.

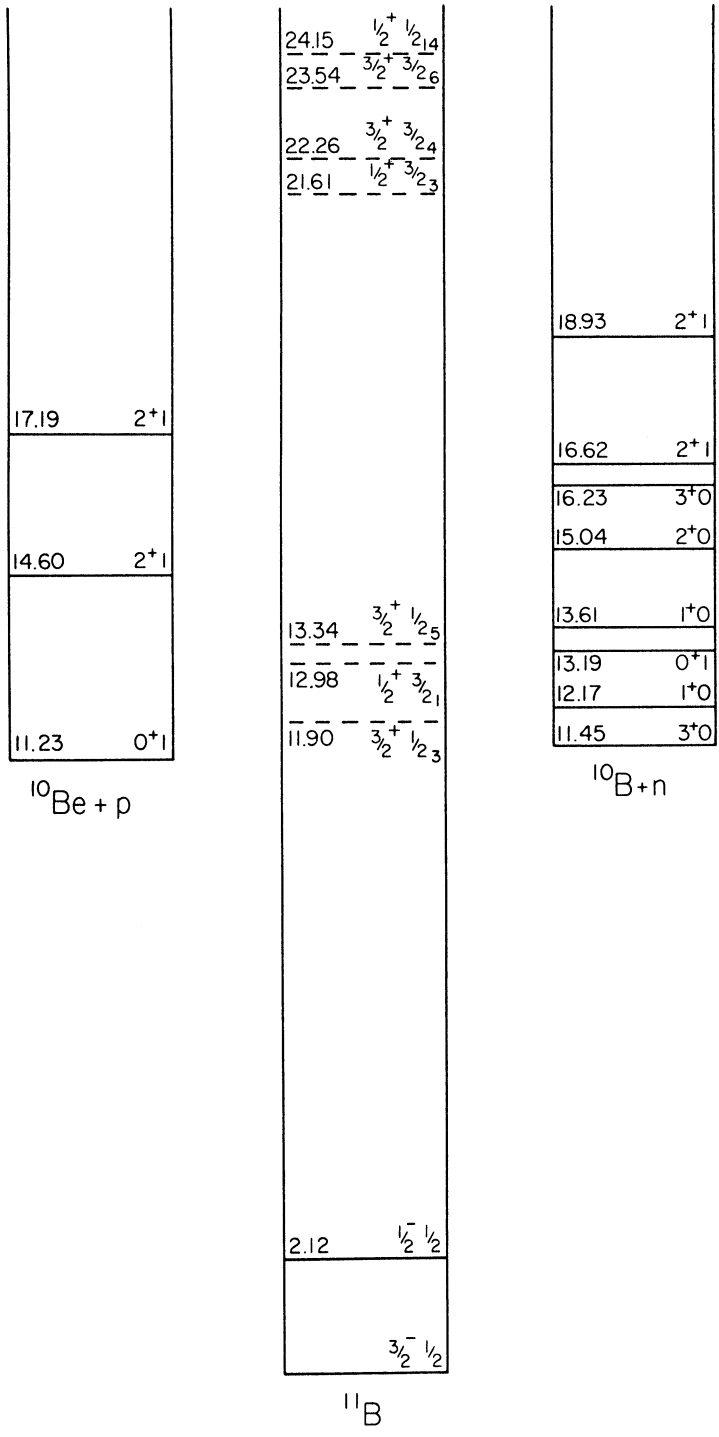

Fig. 8. Energy levels relevant to Coulomb excitation of ^{11}B. See also Fig. 5 caption.

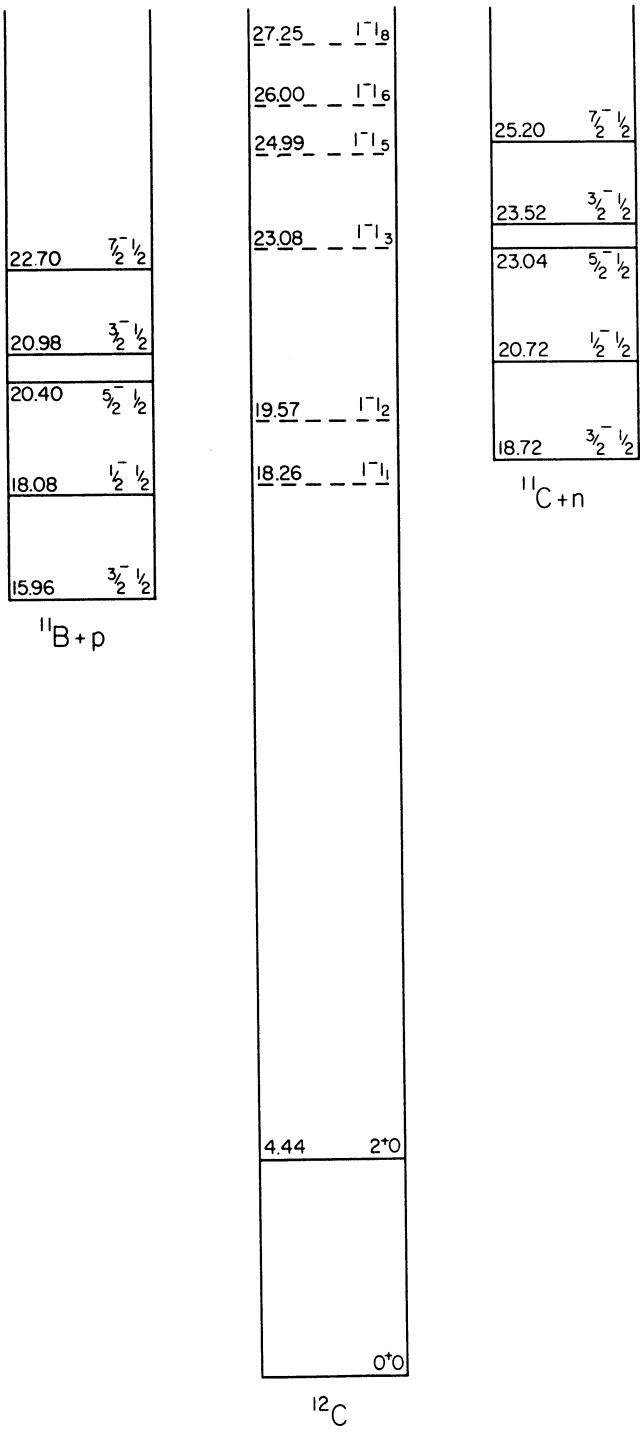

Fig. 9. Energy levels relevant to Coulomb excitation of ^{12}C. See also Fig. 5 caption.

(d) ^{11}B

The intermediate states that give large contributions to τ_{12} for ^{11}B, and have appreciable $1p \rightarrow 2s$ and $1p \rightarrow 1d$ excitations, are the $\frac{1}{2}^+\frac{1}{2}_4$, $\frac{1}{2}^+\frac{3}{2}_{1,3}$, $\frac{3}{2}^+\frac{1}{2}_{3,5}$ and $\frac{3}{2}^+\frac{3}{2}_{4,6}$ states. These are shown in Fig. 8 together with the ^{10}Be and ^{10}B core states that are included in the calculation. The additional contributions to τ_{12} from the seven intermediate states, in order of increasing excitation energy, are approximately 0.000, 0.002, 0.000, 0.000, 0.001, −0.001 and 0.000 fm^3 respectively.

(e) ^{12}C

Fig. 9 shows the 1^- 1 intermediate states for which possible changes are considered in the contributions to τ_{12} for ^{12}C, and the significant $A = 11$ core states that are included. The additional contributions to τ_{12} for these 1^- $1_{1-3,5,6,8}$ states are approximately 0.000, 0.001, 0.003, −0.001, −0.001 and −0.004 fm^3 respectively.

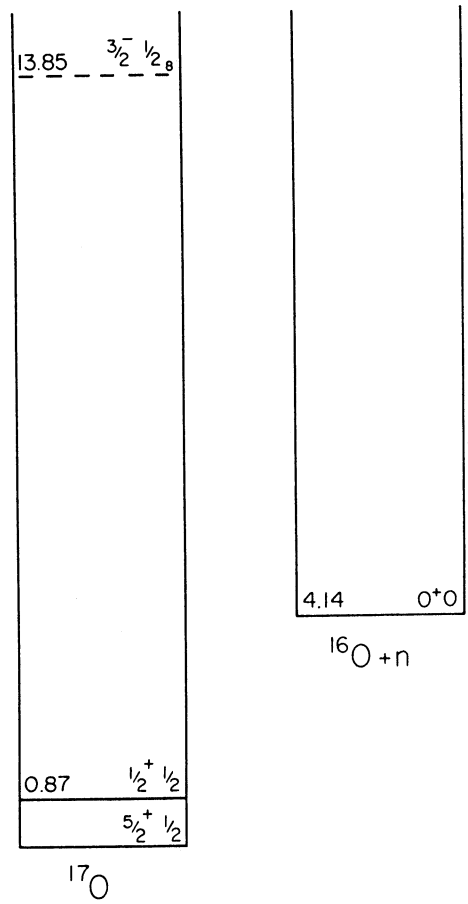

Fig. 10. Energy levels relevant to Coulomb excitation of ^{17}O. See also Fig. 5 caption.

(f) ^{17}O

The main contributions to τ_{12} for ^{17}O come from the four intermediate states $\frac{3}{2}^-\frac{1}{2}_8$ at 13.85 MeV and $\frac{3}{2}^-\frac{3}{2}_{10,11,14}$ at 24.24, 26.82 and 28.62 MeV. The $T = \frac{1}{2}$ state belongs mainly to the 1p configuration $1s^4 1p^{12} 2p_{3/2}$. The $T = \frac{3}{2}$ states are predominantly 2p–1h states built on the 1^-, $T = 1$ GDR in ^{16}O, and as such involve $1p \rightarrow (2s,1d)$ excitations in which the initial 1p nucleon has a binding energy of the order of 25 MeV. Consequently we assume that E1 transitions involving these $T = \frac{3}{2}$ states are well represented by harmonic oscillator values.

Possible changes due to the use of Woods-Saxon wavefunctions are therefore considered only in the E1 matrix elements involving the $\frac{3}{2}^-\frac{1}{2}_8$ state, and in these only in the parts coming from the ^{16}O core in its ground state. The relevant states and energies are shown in Fig. 10. Because the $2p_{3/2}$ neutron in the $\frac{3}{2}^-\frac{1}{2}_8$ state is unbound, we have to specify the value of the channel radius a in calculating the quantity (16); this is done as follows. We use a Woods-Saxon potential with central and spin-orbit components, with the radius and diffuseness the same for each component, and assume that the depths V_0 and V_{so} are the same for the even-parity states ($2s_{1/2}$, $1d_{5/2}$, $1d_{3/2}$). With the conventional value of the diffuseness, we choose V_0, V_{so} and R to fit the observed binding energies of the $2s_{1/2}$ and $1d_{5/2}$ states and the rms radius of the $1d_{5/2}$ neutron (Hicks 1982), and then choose a so that the $d_{3/2}$ resonance occurs at the appropriate energy. This is taken to be an excitation energy of 5.74 MeV (Johnson 1973), leading to $a = 3.85$ fm (the lowest possible energy for the $d_{3/2}$ state is 5.08 MeV, which gives $a = 4.27$ fm). We note that Johnson used $a = 3.86$ fm, whereas Hickey *et al.* (1974) and Holt *et al.* (1978) used $a \approx 4.92$ fm.

For the $2p_{3/2}$ state we use the same channel radius and the same Woods-Saxon potential parameters, except that V_0 is adjusted so that the $2p_{3/2}$ resonance corresponds to the energy of the $\frac{3}{2}^-\frac{1}{2}_8$ state. The resultant values of τ_{12}, $S(E1)$ and k are given in Table 2 (for $a = 4.27$ fm, these values would be reduced by about 8%).

5. Admixtures from Higher Configurations

Shell model calculations restricted to the lowest configuration fail to account for observed E2 matrix elements, as is illustrated by the nonzero values of Δe in Table 2. This failure is usually attributed to the neglect of higher configurations.

A calculation of the additional contribution to $S(E1)$ (or τ_{12}) for ^{17}O due to admixtures from higher configurations in the wavefunctions for the $\frac{5}{2}^+$ ground and $\frac{1}{2}^+$ first excited states has been made by Barker (1982*b*). The admixtures were treated in first-order perturbation theory, their coefficients in the wavefunctions being determined by fitting the observed values of $Q(\frac{5}{2}^+)$ and $B(E2; \frac{5}{2}^+ \rightarrow \frac{1}{2}^+)$, together with other reasonable restrictions. Woods-Saxon wavefunctions were used. The calculation of $S(E1)$ was made in the closure approximation, in which all the E1 strength to the $\frac{3}{2}^-$ states was assumed to be concentrated at the energy E_g; the value $E_g - E_i = 22$ MeV was estimated from the photonuclear cross sections for ^{17}O. The additional contribution to

S(E1) due to configuration mixing was found to be about 54% of the total value of 0.0113 e^2 fm^2 MeV^{-1}, *i.e.* 0.0061 e^2 fm^2 MeV^{-1}. Because of the perturbation treatment, we can here assume the same additional contribution to S(E1), and combine it with the value in Table 2 obtained for Woods-Saxon wavefunctions in the lowest configurations. This leads to total values for ^{17}O of τ_{12} = 0.170 fm^3, S(E1) = 0.0231 e^2 fm^2 MeV^1 and k = 5.3.

Such a calculation is practicable for ^{17}O because of the very simple wavefunctions of the ground and first excited states, if restricted to the lowest configurations. For the 1p-shell nuclei a similar calculation would be very complicated. Kurath (1959) used the method of generator coordinates to introduce configuration mixing of a collective type in 1p-shell nuclei, and found that the required enhancement of E2 transition probabilities could be obtained with small admixtures. The calculation of the corresponding enhancement of S(E1) would be considerably more difficult because it involves two one-body operators instead of one, and the same is true even in the closure approximation because both one- and two-body operators are involved. In a qualitative approach, an argument could be given similar to that used by Kurath when discussing the effect of centre-of-mass motion. This would suggest a relationship between the enhancements in E2 matrix elements and in S(E1) similar to that assumed in Barker (1982*a*) where the $\langle 1p \vdots r^2 \vdots 1p \rangle$ radial integrals were renormalised. From the arguments of Kurath, large enhancements would occur only for cases where the ground and first excited states belong to the same rotational band, and therefore not for ^{10}B or ^{11}B among the cases we consider. It may be noted that the concept of a state-independent effective charge does not apply very well for these two cases (for Woods-Saxon wavefunctions), as is seen from the values of Q and B(E2) in Table 2.

6. Discussion

The values of k calculated with Woods-Saxon wavefunctions, as given in Table 2, show considerable variation from nucleus to nucleus. These values should still be corrected for the effect of configuration mixing, the importance of which is indicated by the size of $\Delta e/e$. For ^{17}O, we estimate that configuration mixing increases k from 3.9 to 5.3, bringing it into good agreement with the experimental value of 5.7±0.4 (Table 1). Increases of k might also be expected in other cases, except perhaps for ^{10}B and ^{11}B where the ground and first excited states do not belong to the same rotational band. Thus k values for ^{6}Li and ^{7}Li of the order of 3, as suggested by experiment, do not seem to be unreasonable. For ^{10}B, our calculated value of k exceeds the experimental value. Fewell *et al.* (1980, 1984) assumed k = 1 for ^{11}B in deriving their experimental value of B(E2; $\frac{3}{2}^- \rightarrow \frac{1}{2}^-$), and Vermeer *et al.* (1983*a*) assumed k = 1 for ^{12}C in their measurement of $Q(2^+)$.

The calculations also indicate the regions of excitation energy that provide the main E1 contributions to Coulomb excitation in these nuclei. For Woods-Saxon wavefunctions, the harmonic oscillator distributions shown in Figs 1–4 should be modified by enhancement of the contributions from low-lying states. These distributions explain why the enhancements are larger for the nuclei with low-lying dipole states (^{6}Li, ^{7}Li, ^{10}B and ^{17}O) than for those where the main dipole states are at higher energies (^{11}B and ^{12}C). These calculations do not

support the use of the closure approximation with all the E1 strength taken at the GDR energy (Barker 1982*a*, *b*), except perhaps for ^{11}B and ^{12}C.

Almost all other calculations are for ^{7}Li (see Table 1). The main contributions to τ_{11} and τ_{12} for ^{7}Li, as shown in Fig. 1, come from $\frac{1}{2}^{+}$ intermediate states at excitation energies of about 8, 13 and 15 MeV (these states are expected to be very broad, so that their contributions to cross sections of photonuclear or other reactions would not be as distinctive as Fig. 1 might suggest). Nonnegligible contributions to τ_{11} and τ_{12} come from as high as 30–40 MeV, and these high-energy contributions tend to cancel those from low energy. Qualitatively similar results were obtained by Gomez-Comacho and Nagarajan (1985) in their shell model calculation for ^{7}Li (and similar remarks apply to their calculation for ^{6}Li). The α + t cluster calculations of Smilansky *et al.* (1972) and Mertelmeier and Hofman (1986) also found that the main contribution comes from $\frac{1}{2}^{+}$ intermediate states, with the peak contribution located about 1 MeV above the α+t threshold, *i.e.* at an excitation energy of about 3.5 MeV. Even with the inclusion of ^{6}Li+n channels in the scattering functions, Mertelmeier and Hofman found the peak at about the same energy, although the magnitude was increased by about 50% and the high-energy tail was appreciably larger. In our calculation, the individual contributions to τ_{if} are not attributed to the different breakup channels of ^{7}Li, although it is reasonable that the lowest-energy contributions should be associated with the α + t channel, which has the lowest threshold. Kajino *et al.* (1988) estimated that the α + t channel should contribute only about one third of the total value of τ_{11}, by using a relation (their equation 15) based on the assumption of the pure *LS*-coupling limit (for *both* the ^{7}Li ground state and the intermediate dipole states). In this limit, Kajino *et al.* gave

$$\langle \tfrac{3}{2}^{-} \| \mathcal{M}(\mathrm{E1}) \| \tfrac{5}{2}^{+} \rangle = 3 \langle \tfrac{3}{2}^{-} \| \mathcal{M}(\mathrm{E1}) \| \tfrac{3}{2}^{+} \rangle ;$$

however, this equation cannot be valid for all intermediate states because of the sum rule (10), and so equation (15) of Kajino *et al.* is open to question.

Although the present shell model calculation for ^{7}Li needs to be modified for the effect of configuration mixing, it does indicate that previous cluster calculations have not included all important clusters.

7. Summary

In the present shell model calculations of the strength k of the E1 contribution to Coulomb excitation for a selection of 1p-shell nuclei, we have restricted the wavefunctions of the ground and first excited states to the lowest configuration, and used harmonic oscillator or Woods-Saxon single-particle wavefunctions. The calculated k values show some correlation with experimental values, the variation from nucleus to nucleus being determined to a large extent by the energy region of the dipole states that give the main contributions to k. Accurate values of k cannot be expected, however, without the use of wavefunctions that include configuration mixing, chosen to fit observed E2 matrix elements. In the case of ^{17}O, the states belonging to the lowest configurations are sufficiently simple that such configuration mixing can be included, and good agreement with experiment is found for the resultant k value.

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