J. Phys. G: Nucl. Phys. **8** (1982) 679–685. Printed in Great Britain

Magnetic moments of $T=0$ states in $N=Z$ nuclei

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Received 10 November 1981

Abstract. Theoretical calculations for magnetic moments of $T=0$ collective states in $N=Z$ nuclei are compared with recent experimental results. The $J^{\pi}=2^{+}$ and 3^{-} states considered all have $0.49<g_{\text{theor}}<0.51$ if the wavefunctions have pure $T=0$ isospin. Isospin mixing increases the calculated ^{16}O 3^{-} g factor by 9%, consistent with experiment, but has a much smaller effect on the 2^{+} states. A reported negative experimental g factor for the 4^{+} state in ^{20}Ne is in complete disagreement with theory.

1. Introduction

It is well known (Kurath 1961, Sugimoto 1969, van Hienen and Glaudemans 1972, Zalm *et al* 1978, Raman *et al* 1978) that for nuclei where the states have good isospin the isoscalar magnetic moment $\mu_0(J)=\frac{1}{2}(\mu(J,T,T_z=T)+\mu(J,T,T_z=-T))$ can be directly related to the expectation value of the spin density

$$g_0(J)=\mu_0(J)/J=\tfrac{1}{2}+(\mu_p+\mu_n-\tfrac{1}{2})\langle S_3\rangle/J \tag{1}$$

where

$$\langle S_3\rangle=\left\langle J,M=J\left|\sum_i s_{3i}\right|J,M=J\right\rangle$$

and where μ_p and μ_n are the free-nucleon moments, $\mu_p+\mu_n=0.880$. For collective excitations built on $J^{\pi}=0^{+}$ ground states the angular momentum comes only from vibrational and rotational degrees of freedom and hence $\langle J_3\rangle=\langle L_3\rangle$ or $\langle S_3\rangle=0$, which immediately implies that $g_0=\frac{1}{2}$. The subject of this work concerns some interesting deviations from $g_0=\frac{1}{2}$ obtained from theoretical calculations of the effects due to other nuclear degrees of freedom and compared with recent experimental results. The experimental g factors given in table 1 have been obtained only recently due to new techniques which have been developed (Randolf *et al* 1973, van Middelkoop 1978) for measuring the moments of states with lifetimes of the order of 10^{-12} s.

For states with good isospin the contributions to the moments from the single-particle degrees of freedom can be easily understood in terms of the $\langle S_3\rangle$ expectation values. The single-particle values for $T=\frac{1}{2}$ states in odd–even nuclei are $\langle s_3\rangle=\frac{1}{2}$ for $j=l+\frac{1}{2}$ and $\langle s_3\rangle=-\frac{1}{2}+1/(2l+1)$ for $j=l-\frac{1}{2}$. The stretched two-particle configurations for $T=0$ states in odd–odd nuclei have $\langle S_3\rangle=2\langle s_3\rangle$. A good example of the latter case is the deuteron $J^{\pi}=1^{+}$ state which has $g_{\text{exp}}=0.8574$ or, from equation (1), $\langle S_3\rangle_{\text{exp}}=0.941$ which is quenched from the pure s-state value of $\langle S_3\rangle=1$ due to the d-state admixture. In fact, all one- and two-particle configurations have $\langle S_3\rangle_{\text{exp}}$ values which are systematically

Table 1. Experimental g factors and the $\langle S_3\rangle_{\mathrm{exp}}$ matrix elements deduced from equation (1) compared with $\Delta N=0$ shell-model predictions.

		g_{exp}	$g_{\Delta N=0}$	$\langle S_3\rangle_{\mathrm{exp}}$	$\langle S_3\rangle_{\Delta N=0}$
^{12}C	2^+		0.510		0.054
^{16}O	3^-	0.556(4)[a]	0.511	(0.44(3))[h]	0.085
^{20}Ne	2^+	0.54(4)[b]	0.510	0.21(21)	0.050
	4^+	−0.10(19)[c]	0.511	−6.3(20)	0.111
	6^+		0.522		0.355
	8^+		0.534		0.722
^{24}Mg	2^+	0.51(2)[b]	0.512	0.05(10)	0.063
	4^+		0.515		0.158
^{28}Si	2^+	0.56(9)[d]	0.513	0.3(5)	0.071
^{32}S	2^+	0.47(9)[e]	0.495	−0.2(5)	−0.024
^{36}Ar	2^+		0.491		−0.047
^{40}Ca	3^-	0.56(13)[f]		0.5(10)	
	5^-	0.54(10)[g]		0.5(13)	

[a] Bennett (1980).
[b] Horstman *et al* (1975).
[c] Speidel *et al* (1980).
[d] Eberhard *et al* (1975).
[e] Zalm *et al* (1979).
[f] Jain *et al* (1976).
[g] Hensler *et al* (1974).
[h] $\langle S_3\rangle=0.09(3)$ from equation (4) using the calculated $\delta g=0.044$ discussed in the text.

quenched relative to the single-particle estimate due to configuration mixing within major oscillator shells ($\Delta N=0$) (Brown and Wildenthal 1981 (unpublished), Wildenthal and Chung 1979) as well as higher-order mixing involving 2p–2h, 4p–4h, ... ($\Delta N\geqslant 2$) configurations (Shimizu *et al* 1974, Arima and Hyuga 1979). Thus for $T=0$ states, $\langle S_3\rangle$ is determined by the competition between the collective component which has $\langle S_3\rangle=0$ and the quenched two-particle component which has $-1<\langle S_3\rangle<+1$.

It is important to remember that equation (1) is not valid if the wavefunctions do not have good isospin or if extra-nucleon degrees of freedom are important. It will be shown below that isospin mixing effects are important for the ^{16}O $J^\pi=3^-$ state. The calculated meson-exchange corrections are small for the isoscalar moments (the one-pion-exchange current contributes only to the isovector magnetic moment operator) and will be ignored here (see Hyuga *et al* (1980) and table II of Raman *et al* (1978)).

Assuming that the wavefunctions have pure $T=0$ isospin, equation (1) has been used to extract $\langle S_3\rangle_{\mathrm{exp}}$ from the experimental g factors (see table 1). These are compared with shell-model calculations involving full configuration mixing within major oscillator shells $\langle S_3\rangle_{\Delta N=0}$: $(1\mathrm{p})^8$ with the Cohen–Kurath (1965) interaction for ^{12}C, $(1\mathrm{p})^{-1}$ $(1\mathrm{d},2\mathrm{s})^1$ with the Millener–Kurath (1970) interaction for ^{16}O and $(1\mathrm{d},2\mathrm{s})^n$ with the Chung–Wildenthal interactions (Chung 1976, Wildenthal 1977) for the sd-shell nuclei. For the 2^+ states considered $\langle S_3\rangle_{\Delta N=0}\leqslant 0.1$, and this is in agreement with experiment. However, except for ^{24}Mg, the experimental error bars are too large for any discriminating test of the theory and new measurements of the g factors with about 1% precision are needed.

The $\Delta N=0$ predictions for the ^{20}Ne ground-state band are interesting. In contrast to $\langle S_3\rangle_{\mathrm{SU3}}=0$ for all J values obtained with SU3 cluster wavefunctions (see, e.g., Strottman 1972), the $\Delta N=0$ shell-model values increase as J becomes larger and reach almost the stretched two-particle value for the 8^+ state. The experimental value (Speidel *et al* 1980) for the 4^+ state is completely inconsistent with these predictions and in fact has an $\langle S_3 \rangle$ value an order of magnitude larger than even a two-particle configuration. Confirmation of this experimental result is essential.

The value of $\langle S_3 \rangle_{\exp} = 0.44 \pm 0.03$ from the ^{16}O 3^- g-factor measurement at Oxford (Bennett 1980) is in fair agreement with that expected for the simplest shell-model configuration $\langle S_3 \rangle [(p_{1/2})^{-1}(d_{5/2})] = 0.33$. However, it is well known from the large $0^+ \rightarrow 3^- B(E3)$ value that this state is collective. Part of this collectivity comes out of the 1p–1h $(p)^{-1}$ $(sd)^1$ calculation which gives $\langle S_3 \rangle_{\Delta N=0} = 0.085$ and additional $\Delta N = 3$, 1p–1h and 3p–3h mixing which is needed to reproduce the B(E3) value might be expected to further reduce $\langle S_3 \rangle$. Thus $\langle S_3 \rangle_{\exp}$ for the ^{16}O 3^- state is in disagreement with theoretical expectations.

2. Effects of isospin mixing

It will now be shown that the discrepancy in ^{16}O mentioned above can be understood as an effect of isospin mixing. First the two-level mixing of the $T=0$ and 1 $J^\pi = 3^-$ configurations $[(p_{1/2})^{-1}(d_{5/2})]$ will be considered. In perturbation theory the magnetic moment of the lowest 3^- states is given by

$$\langle 3^- | \mu | 3^- \rangle = \langle 3^- T=0 | \mu | 3^- T=0 \rangle + \delta\mu \tag{2}$$

where

$$\delta\mu = -2 \frac{\langle 3^- T=0 | V_C | 3^- T=1 \rangle}{|\Delta E|} \langle 3^- T=1 | \mu | 3^- T=0 \rangle \tag{3}$$

and equation (1) must be modified to

$$g_0(J) = g(J) - \delta g = \tfrac{1}{2} + (\mu_p + \mu_n - 1)\langle S_3 \rangle / J \tag{4}$$

where $\delta g = \delta\mu / J$. For the derivation of the following formulae it is convenient to write the 1p–1h wavefunctions in proton–neutron formalism:

$$\begin{aligned}
|j_h^{-1} j_p, J, T=0, {}^{16}\mathrm{O}\rangle &= -(|(\pi j_h)^{-1}(\pi j_p)J\rangle - |(\nu j_h)^{-1}(\nu j_p)J\rangle)/\sqrt{2} \\
|j_h^{-1} j_p, J, T=1, {}^{16}\mathrm{F}\rangle &= |(\nu j_h)^{-1}(\pi j_p)J\rangle \\
|j_h^{-1} j_p, J, T=1, {}^{16}\mathrm{O}\rangle &= (|(\pi j_h)^{-1}(\pi j_p)J\rangle + |(\nu j_h)^{-1}(\nu j_p)J\rangle)/\sqrt{2} \\
|j_h^{-1} j_p, J, T=1, {}^{16}\mathrm{N}\rangle &= |(\pi j_h)^{-1}(\nu j_p)J\rangle.
\end{aligned} \tag{5}$$

One can show that

$$\begin{aligned}
\langle JT=1 | \mu | JT=0 \rangle
&= \left(\frac{4\pi}{3} \frac{J}{(J+1)(2J+1)} \right)^{1/2} \langle JT=1 \| \mathrm{M1} \| JT=0 \rangle \\
&= \frac{-1}{2(J+1)} \{ [J(J+1) + j_p(j_p+1) - j_h(j_h+1)] g_1(j_p) \\
&\quad + [J(J+1) + j_h(j_h+1) - j_p(j_p+1)] g_1(j_h) \} \mu_N
\end{aligned} \tag{6}$$

where $g_1(j) = \frac{1}{2}(g(\pi j) - g(\nu j))$ are the single-particle isovector g factors. The Schmidt values are $g_1(1\mathrm{d}_{5/2})=1.342$ and $g_1(1\mathrm{p}_{1/2})=-0.899$ and hence $\langle 3^- T=1|\mu|3^- T=0\rangle = -2.90\mu_\mathrm{N}$. Alternatively, this off-diagonal matrix element can be related to the $B(\mathrm{M}1)$ between these two states:

$$|\langle J\,T=1|\mu|J\,T=0\rangle| = \left(\frac{4\pi}{3}\frac{J}{J+1}\right)^{1/2}(B(\mathrm{M}1))^{1/2}. \tag{7}$$

The experimental value for the transition between the 13.26 MeV $T=1$ and 6.13 MeV $T=0$ 3^- states is $B(\mathrm{M}1)=(2.16\pm0.35)\mu_\mathrm{N}^2$ (Ajzenberg-Selove 1977, Gorodetzky *et al* 1968), which gives $|\langle 3^- T=1|\mu|3^- T=0\rangle| = (2.6\pm0.2)\mu_\mathrm{N}$ in fair agreement with the $(1\mathrm{p}_{1/2})^{-1}(1\mathrm{d}_{5/2})$ calculation given above.

We will assume that the isospin mixing is due to the Coulomb interaction V_C between two protons. In the one-particle–one-hole model the off-diagonal matrix element $\langle T=0|V_\mathrm{C}|T=1\rangle$ can be related to $A=15$, 16 and 17 binding energies ($E=-BE$):

$$E(^{16}\mathrm{F}\,J,\,T=1) = \varepsilon(\pi j_\mathrm{p}) - \varepsilon(\nu j_\mathrm{h}) + \langle V_\mathrm{s}\rangle + E(^{16}\mathrm{O\ GS}) \tag{8a}$$

$$E(^{16}\mathrm{N}\,J,\,T=1) = \varepsilon(\nu j_\mathrm{p}) - \varepsilon(\pi j_\mathrm{h}) + \langle V_\mathrm{s}\rangle + E(^{16}\mathrm{O\ GS}) \tag{8b}$$

$$E(^{16}\mathrm{O}\,J,\,T=1) = \tfrac{1}{2}(\varepsilon(\pi j_\mathrm{p}) - \varepsilon(\nu j_\mathrm{h}) + \varepsilon(\nu j_\mathrm{p}) - \varepsilon(\pi j_\mathrm{h}) + \langle V_\mathrm{s}\rangle + \tfrac{1}{2}\langle V_\mathrm{C}\rangle + E(^{16}\mathrm{O\ GS}) \tag{8c}$$

$$\langle J,\,T=0|V_\mathrm{C}|J,\,T=1\rangle = -\tfrac{1}{2}(\Delta\varepsilon_1 + \langle V_\mathrm{C}\rangle) \tag{8d}$$

where

$$\Delta\varepsilon_1 = (\varepsilon(\pi j_\mathrm{p}) - \varepsilon(\nu j_\mathrm{p})) - (\varepsilon(\pi j_\mathrm{h}) - \varepsilon(\nu j_\mathrm{h}))$$

$$\varepsilon(j_\mathrm{p}) = E(A=17,\,T=\tfrac{1}{2}) - E(^{16}\mathrm{O\ GS}) \tag{8e}$$

and

$$\varepsilon(j_\mathrm{h}) = E(^{16}\mathrm{O\ GS}) - E(A=15,\,T=\tfrac{1}{2}). \tag{8f}$$

$\langle V_\mathrm{s}\rangle$ is the strong isospin-conserving (T_z-independent) particle–hole matrix element and $\langle V_\mathrm{C}\rangle$ is the Coulomb particle–hole matrix element between two protons. First equations (8*e*) and (8*f*) can be used to obtain $\varepsilon(j)$, then these are put into equations (8*a*) or (8*b*) to obtain $\langle V_\mathrm{s}\rangle$ and finally ε and $\langle V_\mathrm{s}\rangle$ are put into equation (8*c*) to obtain $\langle V_\mathrm{C}\rangle$. The numerical values obtained from the experimental binding energies (Wapstra and Bos 1977, Ajzenberg-Selove 1976, 1977) of $A=15$ $\frac{1}{2}^-$, $A=17$ $\frac{5}{2}^+$ and $A=16$ 3^- states are $\varepsilon(\pi\mathrm{d}_{5/2}) - \varepsilon(\nu\mathrm{d}_{5/2}) = 3.54$ MeV, $\varepsilon(\pi\mathrm{p}_{1/2}) - \varepsilon(\nu\mathrm{p}_{1/2}) = 3.54$ MeV, $\langle V_\mathrm{s}\rangle(^{16}\mathrm{F}) = 1.87$ MeV, $\langle V_\mathrm{s}\rangle(^{16}\mathrm{N}) = 1.95$ MeV and $\langle V_\mathrm{s}\rangle(^{16}\mathrm{O}) + \frac{1}{2}\langle V_\mathrm{C}\rangle = 1.73$ MeV. The values of $\langle V_\mathrm{s}\rangle$ obtained from ^{16}F and ^{16}N are inconsistent due to an effect which will be discussed below. The average value will be used to obtain $\langle V_\mathrm{C}\rangle = -0.36$ MeV. Notice that empirically $\Delta\varepsilon_1 \approx 0$ in equation (8*d*), leaving only $\langle V_\mathrm{C}\rangle$ to contribute to the isospin-mixing matrix element $\langle 3^- T=0|V_\mathrm{C}|3^- T=1\rangle = +0.18$ MeV.

Thus in the two-level mixing approximation the matrix elements required for $\delta\mu$ can be obtained from experimental quantities and the sign can be deduced from the $(\mathrm{p}_{1/2})^{-1}(\mathrm{d}_{5/2})$ calculation. From equation (3) the result for mixing of the 13.26 MeV $T=1$ and 6.13 MeV $T=0$ 3^- states is ($|\Delta E| = 7.1$ MeV)

$$\delta\mu = \frac{-2(+0.18)(-2.6)}{7.1}\,\mu_\mathrm{N} = 0.13\mu_\mathrm{N}. \tag{9}$$

This, together with the $\Delta N=0$ value for the $T=0$ component $\langle 3^- T=0|\mu|3^- T=0\rangle = 1.53$, gives $\langle 3^-|\mu|3^-\rangle = 1.66\mu_N$ or $g(3^-)=0.555$, which is in remarkably good agreement with experiment.

Relationships between masses and isospin mixing were first used by Braithwaite *et al* (1972) to estimate isospin mixing for the ^{12}C 1^+ states and the relation given in their paper is equivalent to using only the matrix element $\langle V_s\rangle$ obtained from the neutron-rich nucleus (^{16}N in this case) and ignoring information about the proton-rich nucleus (^{16}F in this case). Similar relationships have been used since then (see, e.g., Sato and Zamick 1977, Shlomo and Wagner 1978).

The reason for the difference in $\langle V_s\rangle$ between ^{16}N and ^{16}F can be understood as an implicit effect of the Coulomb interaction in the model space (Lawson 1978). For ^{16}F the $1d_{5/2}$ proton single-particle wavefunction is bound by only 0.6 MeV while for ^{16}N the $1d_{5/2}$ neutron single-particle wavefunction is bound by 4.14 MeV. For a delta-function residual interaction the residual particle–hole interaction is proportional to the integral

$$\int R^2_{1p_{1/2}}(r) R^2_{1d_{5/2}}(r) r^2 \, dr$$

and the large spatial extent of the proton orbit in ^{16}F due to its small binding energy reduces the value of this integral compared with that for ^{16}N. The ^{16}F to ^{16}N ratio for this integral using Woods–Saxon wavefunctions is 0.90 compared with the empirical value of $1.87/1.95=0.96$ for the 3^- level (and $1.45/1.65=0.95$ for the 2^- level). (For the $(1p_{1/2})^{-1}$ $(2s_{1/2})$ $J^\pi=0^-$ and 1^- levels the calculation gives 0.76 compared with the empirical values of $0.65/0.90=0.72$ and $0.85/1.18=0.72$, respectively.) The discrepancies between theory and experiment may be due to the finite range of the residual interaction, but the effect is understood qualitatively. In this model $\langle V_s\rangle$ for the middle nucleus ^{16}O should be about the average of the values for ^{16}N and ^{16}F as we have assumed above.

In the limit of an infinitely long-range Coulomb interaction it is easy to see that $\Delta\varepsilon_1 = \langle V_c\rangle$ (≈ 0.36 MeV) and there would be no isospin mixing. $\Delta\varepsilon_1$ is nearly vanishing in this case partly because of the small binding energy of the $d_{5/2}$ orbit. In fact, for this reason the isospin matrix elements of the $(1p_{1/2})^{-1}$ $(2s_{1/2})$ $J^\pi=0^-$ and 1^- states should even be larger (about 0.40 MeV) since $\Delta\varepsilon_1 = -0.37$ MeV for the difference between the $1p_{1/2}$ and $2s_{1/2}$ displacement energies.

The effects due to more complicated structures for the lowest 3^- states as well as the effects due to isospin mixing with more highly excited $T=1$ 3^- states have been considered using the relation

$$\delta\mu = -2\sum_i \frac{\langle 3_1^- T=0|V_C|3_i^- T=1\rangle}{|\Delta E_i|}\langle 3_i^- T=1|\mu|3_1^- T=0\rangle.$$

The wavefunctions were obtained by allowing complete configuration mixing within the model space $(1p_{1/2}, 1d_{5/2}, 2s_{1/2})^4$ (ZBM) with the Reehal–Wildenthal (1973) interaction and a separate calculation within the model space $(1p)^{-1}$ $(1d, 2s)^1$ (PHSD) with the Millener–Kurath (1970) interaction.

In both model spaces the sum was found to be dominated by more than 90% from the contribution from the lowest $T=1$ 3^- state. The theoretical $B(M1)$ values between the lowest $T=0$ and $T=1$ states are $1.25\mu_N^2$ for the PHSD model space and $2.30\,\mu_N^2$ for the ZBM model space compared with the $(1p_{1/2})^{-1}(1d_{5/2})$ value of $2.68\mu_N^2$ from equations (6) and (7) and the experimental value of $(2.16\pm0.35)\mu_N^2$. In the PHSD model space the M1 matrix element is small due to destructive interference between the large $(p_{1/2})^{-1}(d_{5/2})$ component and the relatively small $(p_{3/2})^{-1}(d_{5/2})$ and $(p_{3/2})^{-1}(d_{3/2})$ components. These results indicate that the Millener–Kurath interaction induces somewhat too large an admixture of the $1p_{3/2}$ and $1d_{3/2}$ orbits into the 3^- wavefunctions.

Recently the Oxford shell-model code has been extended to calculate two-body transition densities and two-body Coulomb matrix elements (Brown *et al* 1981, unpublished). The two-body Coulomb matrix elements were calculated with harmonic-oscillator wavefunctions and the single-particle energies were taken as adjustable parameters to fit the $A=15$ and $A=17$ displacement energies. The isospin-mixing matrix element was calculated to be 0.15 MeV in the ZBM model space and 0.12 MeV in the PHSD space, to be compared with the $(1p_{1/2})^{-1}(1d_{5/2})$ value obtained above of 0.18 MeV. Although the calculations can be criticised because harmonic-oscillator wavefunctions were used for the two-body Coulomb matrix elements, the reductions relative to 0.18 MeV are expected because of the more complex structure of the 3^- states in these model spaces. The moment correction becomes $\delta\mu=0.09\mu_N$ when a value of 0.12 MeV is used for the matrix element V_C in equation (3) (together with the experimental off-diagonal M1 matrix element), which is still in fair agreement with the experimental value of $\delta\mu=0.135\pm0.012$.

The quantity $\delta\mu$ is rather large for the 3^- state in ^{16}O because of the strong $(T=1)\rightarrow(T=0)$ M1 strength of (1.2 ± 0.2) Wu, the relatively small gap (7.1 MeV) between the states and the large isospin-mixing matrix element. As discussed above, normally the isospin matrix element would be smaller because of a cancellation between the terms involving $\Delta\varepsilon_1$ and $\langle V_C\rangle$ in equation (8*d*). For example, in ^{12}C $\Delta\varepsilon_1=0.24$ MeV from the $A=11$ and 13 binding energies and $\langle V_C\rangle$ should have about the same value of -0.36 MeV and thus $\langle T=0|V_C|T=1\rangle\approx0.06$. Also the $(T=1)\rightarrow(T=0)$ strength is smaller ((0.34 ± 0.06) Wu for ^{12}C) and $\Delta\varepsilon$ is larger (11.7 MeV) and hence $|\delta g|\approx0.007$ for the 2^+ state in ^{12}C. From the experimental properties of the 4^+ $T=1$ state in ^{20}Ne (Fifield *et al* 1980), equation (9) takes the form $|\delta\mu|=2(0.10)(1.7)/(6.8)=0.050$ or $|\delta g|=0.012$ (an isospin-mixing matrix element of 100 keV was assumed), which is far too small to account for the experimental results reported by Speidel *et al* (1980). Isospin-mixing effects on other positive-parity states should be similar to these two examples. For the 3^- and 5^- states in ^{40}Ca the isospin-mixing effects should again be large since $\Delta\varepsilon_1\simeq0$, but the present experimental errors are an order of magnitude too large to be sensitive to this effect.

3. Conclusions

In conclusion, shell-model calculations for the collective 2^+ and 3^- states in $N=Z$ nuclei give small values for $\langle S_3\rangle$ consistent with the collective-model assumption that $\langle S_3\rangle=0$. Since the meson-exchange corrections are small for the isoscalar magnetic operator it is expected that $g=\frac{1}{2}$ for the collective states in $N=Z$ nuclei. Experimental results for 2^+ states are in agreement with this expectation, but at present the experimental errors are too large to offer a discriminating test of the theory. The g factor of the ^{16}O 3^- state at 6.13 MeV excitation has been measured an order of magnitude more precisely than any other collective state and the experimental g factor is 10% larger than $g=\frac{1}{2}$. This deviation is found to be due to isospin mixing with the $T=1$ 3^- state at 13.26 MeV excitation, which is important because of the large $(3^-, T=1)\rightarrow(3^-, T=0)$ M1 transition strength and the large isospin-mixing matrix element. For the positive-parity states the isospin mixing effects are estimated to be smaller, of the order of $|\delta g|\leqslant0.01$. The calculated values for $\langle S_3\rangle$ in the ground-state band of ^{20}Ne increase with increasing J up to $\langle S_3\rangle=0.72$ for the 8^+ state. The shell-model g factor is then $g_{\Delta N=0}(8^+)=0.534$ compared with the collective value of $\frac{1}{2}$, which would be interesting to confirm experimentally. The measured value of

$g(^{20}\mathrm{Ne}, 4^+) = -0.10 \pm 0.19$ (Speidel *et al* 1980) is completely inconsistent with existing theory and a confirmation of this result is essential.

References

Ajzenberg-Selove 1976 *Nucl. Phys.* A **268** 1
—— 1977 *Nucl. Phys.* A **281** 1
Arima A and Hyuga H 1979 *Mesons in Nuclei* ed M Rho and D Wilkinson (Amsterdam: North-Holland) vol II p 683
Braithwaite W J, Bussoletti J E, Cecil F E and Garvey G T 1972 *Phys. Rev. Lett.* **29** 376
Bennett D W 1980 *Thesis* University of Oxford
Chung W 1976 *Thesis* Michigan State University
Cohen S and Kurath D 1965 *Nucl. Phys.* **73** 1
Eberhardt J L *et al* 1975 *Nucl. Phys.* A **244** 1
Fifield L K *et al* 1980 *Nucl. Phys.* A **334** 109
Gorodetzky S *et al* 1968 *Nucl. Phys.* A**113** 221
Hensler R *et al* 1974 *Phys. Rev.* C **10** 919
van Hienen J F A and Glaudemans P W M 1972 *Phys. Lett.* **42B** 301
Horstman R E *et al* 1975 *Nucl. Phys.* A **248** 291
Hyuga H, Arima A and Shimizu K 1980 *Nucl. Phys.* A **336** 363 and references therein
Jain H C *et al* 1976 *Phys. Rev.* C **14** 2013
Kurath D 1961 *Phys. Rev.* **124** 552
Lawson R D 1978 *Phys. Lett.* **78B** 371
van Middelkoop G 1978 *Hyperfine Interactions* **4** 238 and references therein
Millener D J and Kurath D 1970 *Nucl. Phys.* A **255** 315
Raman S, Houser C A, Walkiewicz T A and Towner I S 1978 *Atomic Data and Nuclear Data Tables* **21** 567
Randolph W L *et al* 1973 *Phys. Lett.* **44B** 36
Reehal B S and Wildenthal B H 1973 *Particles and Nuclei* **6** 137
Sato H and Zamick L 1977 *Phys. Lett.* **70B** 285
Shimizu K, Ichimura M and Arima A 1974 *Nucl. Phys.* A **226** 282
Shlomo S and Wagner G J 1978 *Z. Phys.* A **285** 283
Speidel K-H *et al* 1980 *Phys. Lett.* **92B** 289
Strottman D 1972 *Phys. Lett.* **39B** 457
Sugimoto K 1969 *Phys. Rev.* **182** 1051
Wapstra A H and Bos K 1977 *Atomic Data and Nuclear Data Tables* **19** 177
Wildenthal B H 1977 *Elementary Modes of Excitation in Nuclei* ed R Broglia and A Bohr (Rome: Italian Physical Society)
Wildenthal B H and Chung W 1979 *Mesons in Nuclei* ed M Rho and D Wilkinson (Amsterdam: North-Holland) vol II p 721
Zalm P C, van Hienen J F A and Glaudemans P W M 1978 *Z. Phys.* A **287** 255
—— 1979 *Nucl. Phys.* A **315** 133