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The spin degree of freedom in nuclei: levels and transitions in ^{58}Cu

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Abstract. Gamow-Teller (GT) transitions in the region of A=56 are described in terms of the coupling between isovector and isoscalar pairing phonons and Gamow-Teller excitations. The available experimental information is used to extract coupling constants and strength functions of addition (removal) isoscalar and isovector pairing phonons and GT phonons. The validity of the approach is tested by the calculation of intensities for GT transitions in ^{58}Cu.

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1. Introduction

The calculation of nuclear matrix elements of spin-isospin dependent operators is an esential component in the description of neutrino induced transitions and decays, like single and double beta decay. In this work we analyze the experimental information for transitions involving low spin and isospin states in nuclei near closed shells [1]. For the theoretical analysis we have adopted isoscalar and isovector pairing excitations and GT excitations as elementary degrees of freedom. In the present work we extend the analysis of Ref. [2] to include the spin. Therefore, we shall study the properties of those states which may be interpreted either as members of isoscalar pairing multiplets or as Gamow-Teller excitations of closed shells. This is in contrast with usual treatments of the isoscalar pairing in terms of an extension of the BCS formalism. This note is intented to show that the adopted formalism may indeed be applied to calculate the strenght distribution corresponding to spin-isospin excitation, for nuclei near close shells. The numerical results presented in this work may be taken as preliminary results. Final ones will be published elsewhere [3].

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2. Formalism

We take into account three boson degrees of freedom, labeled by the angular momentum I, the isospin T and the transfer number $\alpha = 0, \pm 1$. The quantum number ν orders the states with the three previous labels in common. We assume that the dominant coupling scheme is given by the isovector pairing phonons. On top of this basic structure we locate the I=1 phonons (either isoscalar pairing or Gamow-Teller). We include only those states which one may recognize from the empirical evidence in the region $52 \leq A \leq 60$. We define the following pairs of single-particle operators coupled to good angular momentum I and isospin T

$$\left[a^+_{j_1} a^+_{j_2}\right]^{I=0\,T=1}_{0T_z} \; ; \; \left[a^+_{j_1} a^+_{j_2}\right]^{10}_{M0} \; ; \; \left[a^+_{j_1} a_{j_2}\right]^{11}_{MT_z} , \tag{1}$$

where $a^+_j = b^+_j, c^+_j$ and $j = k, h$. We construct the operators

$$\begin{aligned} P^+_{IM,TT_z,1} &= f^{IT1}_{j_1 j_2} \, [a^+_{j_1} a^+_{j_2}]^{I\,T}_{MT_z} \\ P^+_{1M,1T_z,0} &= f^{110}_{j_1 j_2} \, [a^+_{j_1} a_{j_2}]^{1\,1}_{MT_z} \\ P^+_{IM,TT_z,-1} &= (-1)^{I+T+M+T_z} P_{I(-M),T(-T_z),1} \end{aligned} \tag{2}$$

Here

$$f^{01(\pm1)}_{j_1 j_2} = \delta_{j_1 j_2} \, \hat{j}_1 \; ; \qquad f^{101}_{j_1 j_2} = f^{110}_{j_1 j_2} = \frac{\langle j_1 || \sigma || j_2 \rangle}{\sqrt{3}} \tag{3}$$

and with them we construct separable residual interactions of the form

$$H^{IT\alpha} = -\frac{g^{IT\alpha}}{1+\delta_{\alpha 0}} \, P^+_{IM,TT_z,\alpha} P_{IM,TT_z,\alpha} \tag{4}$$

which are diagonalized through the usual RPA procedure. The coupling pairing strengths $g^{IT(\pm1)}$ are fixed from the lowest states of the systems with A=54 and 58 with I=0 and 1, while the value of Gamow Teller strength g^{110} is an educated guess based on the energetics of the GT mode. Thus, the calculation is parameter-free for all practical purposes.

3. Gamow-Teller transitions

In this section we discuss the matrix elements of the operator

$$Q_{1M} = \frac{\langle j' || \sigma || j \rangle}{\sqrt{3}} \, [b^+_{j'} c_j]^1_M \tag{5}$$

We illustrate the treatment through the reaction (^{3}He,t) on ^{58}Ni. Thus the target nucleus is represented by the one-phonon isovector pairing state $|0\rangle \equiv \Gamma^+_{00,1(-1),1;1}|\,\rangle$. The states of ^{58}Cu listed in Table 1 are populated in lower order of the NFT diagrammatic perturbation theory. For cases A-E the symbols κ_i in Table 1 stand for

graph	ν	energy	$(\kappa_i)_{bare}$	$(\kappa_i)_{ren}$
A	1	3.84	1.21	0.83
	2	5.32	0.35	0.22
B		5.81	0.23	0.16
C	1	10.84	0.016	0.026
D	1	10.84	0.172	0.286
E	1	10.84	0.158	0.257
F	1	9.34	-0.152	-0.256
G	1	9.34	0.444	0.540
H	1	9.34	0.021	0.033
I	1	8.59	-0.919	-0.134
J	1	8.85	0.778	-0.088
	2	12.13	0.191	0.250

Table 1. Model energies and GT matrix elements to excited states with good isospin projection T_z.

the matrix elements

$$\begin{aligned}
\kappa_A &= \langle \Gamma^+_{1M,00,1;\nu} | Q_{1M} | 0 \rangle \\
\kappa_B &= \frac{1}{\sqrt{2}} \langle [b^+_{\frac{3}{2}-} c^+_{\frac{1}{2}-}]^1_M - [b^+_{\frac{1}{2}-} c_{\frac{3}{2}-}]^1_M | Q_{1M} | 0 \rangle \\
\kappa_C, \kappa_D, \kappa_E &= \langle [\Gamma^+_{1,00,1;\nu} \Gamma^+_{1,10,0;1}]^1_M | Q_{1M} | 0 \rangle \qquad (6)
\end{aligned}$$

The isospin of the final state is well defined for these transitions. In the case of transitions F-H Table 1 includes the matrix elements

$$\begin{aligned}
\kappa_F &= \langle \Gamma^+_{1M,11,0;1} \Gamma^+_{00,1(-1),1;1} | Q_{1M} | 0 \rangle_F = -2 \, \langle \Gamma^+_{1M,10,0;1} \Gamma^+_{00,10,1;1} | Q_{1M} | 0 \rangle_F \\
\kappa_G &= 2 \, \langle \Gamma^+_{1M,10,0;1} \Gamma^+_{00,10,1;1} | Q_{1M} | 0 \rangle_G \qquad (7) \\
\kappa_H &= \langle \Gamma^+_{1M,1(-1),0;1} \Gamma^+_{00,11,1;1} | Q_{1M} | 0 \rangle_H = -2 \langle \Gamma^+_{1M,10,0;1} \Gamma^+_{00,10,1;1} | Q_{1M} | 0 \rangle_H
\end{aligned}$$

From (7) we extract the amplitudes for populating states with good isospin. The corresponding values obtained are listed in Table 2.

$$\kappa^T_{T_z=0} = \langle [\Gamma^+_{1M,1,0;1} \Gamma^+_{00,1,1;1}]^T_0 | Q_{1M} | 0 \rangle \qquad (8)$$

Transitions $I - J$ correspond to the matrix elements

$$\begin{aligned}
\kappa_I &= \langle \left(\Gamma^+_{1M,00,1;\nu} \Gamma^+_{00,11,-1;1} | 0 \rangle \right) | Q_{1M} | 0 \rangle \\
\kappa_J &= \langle \left(\Gamma^+_{1M,00,-1;1} \Gamma^+_{00,11,1;\nu} | 0 \rangle \right) | Q_{1M} | 0 \rangle \qquad (9)
\end{aligned}$$

T	κ_0^T	$\left(\kappa_0^T\right)_{bare}$	$\left(\kappa_0^T\right)_{ren}$
0	$\frac{\sqrt{3}}{2}(\kappa_F + \kappa_H - \kappa_G/3)$	-0.261	-0.373
1	$\frac{1}{\sqrt{2}}(\kappa_F - \kappa_H)$	-0.112	-0.209
2	$\frac{1}{\sqrt{6}}\kappa_G$	0.181	0.220

Table 2. The bare and renormalized amplitudes to states of good T arising from graphs F-H.

final state $\vert u\rangle$	ν	T	$E(i,A,T)$	$\vert\langle u\vert Q_{1M}\vert 0\rangle_{bare}\vert^2$	$\vert\langle u\vert Q_{1M}\vert 0\rangle_{ren}\vert^2$
$\Gamma^+_{1M,00,1;\nu}\vert\,\rangle$	1	0	1.50	1.42	0.69
	2	0	2.98	0.12	0.05
$\frac{1}{\sqrt{2}}\left([b^+_{\frac{3}{2}-}c^+_{\frac{1}{2}-}]^1_M - [b^+_{\frac{1}{2}-}c_{\frac{3}{2}-}]^1_M\right)\vert\,\rangle$	3.69	1	3.69	0.07	0.03
$[\Gamma^+_{1,1,0;1}\,\Gamma^+_{1,0,1;\nu}]^1_M\vert\,\rangle$	1	1	8.72	0.12	0.32
$[\Gamma^+_{1M,1,0;1}\,\Gamma^+_{00,1,1;\nu}]^T\vert\,\rangle$	1	0	5.50	0.79	1.96
	1	1	7.22	1.67	2.58
	1	2	10.67	0.96	0.40
$\Gamma^+_{1M,00,1;\nu}\,[\Gamma^+_{00,1,-1;1}\,\Gamma^+_{00,1,1;1}]^T\vert\,\rangle$	1	0	4.76	large	large
	10.67	1	6.48	-	-
		2	9.93	-	
$\Gamma^+_{1M,00,-1;1}\,[\Gamma^+_{00,1,1;\nu}\,\Gamma^+_{00,1,1;1}]^T\vert\,\rangle$	1	0	5.01	-	
		1	6.73	-	
		2	9.93	-	
				5.15	6.04

Table 3. The predicted population of final states in ^{58}Cu through the reaction (^{3}He,t) on ^{58}Ni. The predicted “true” excitation energies $E(i,58,T)$ are in MeV. Accidental degeneracies of energy denominators produced the results indicated as ” large”, and they have been omitted from the final sum

Matrix elements to states with good isospin are given by

$$\kappa^T_{I,J} = \langle 11; 1(-1); T0\rangle\,\kappa_{I,J} \tag{10}$$

The last two columns in Table 3 are obtained by squaring the sum of all amplitudes listed in Table 1 to the same final state $|u\rangle$. Only contributions larger than 0.01 are considered. The last column but one displays the bare matrix elements, while last one includes the renormalization through the Gamow-Teller phonon. This renormalization is taken into account in perturbation theory. The amplitudes to final states $[\Gamma^+_{1M,1,0;1}\,\Gamma^+_{00,1,1;\nu}]^T|\,\rangle$ include also the direct creation of the GT phonon

(GT giant resonance)

$$\langle\left(\left[\Gamma^{+}_{1M,1,0;1}\,\Gamma^{+}_{00,1,1;1}\rangle\right]^{T}\right)|\,(Q_{1M})_{coll}\,|0\rangle=-\frac{\Xi^{110}_{1}}{g^{110}}\langle\,11;1(-1);T0\rangle. \tag{11}$$

Our results are not restricted to cases in which the energies of the initial and final states are much smaller than the frequency of the giant resonance. In fact, we would like to find out how the sumrule is preserved and where the missing intensity lies. Thus we go beyond the usual prescription of effective charges. The details concerning the renormalization procedure are given in [3].

4. Ikeda's sumrule

For a two-particle state $|0\rangle$, Ikeda's sumrule reads

$$2=N-Z=|\langle a|Q_{1M}|0\rangle|^{2}-|\langle b|Q'_{1M}|0\rangle|^{2} \tag{12}$$

where the sum over M is omitted, the GT operator Q_{1M} is given in (5) and

$$Q'_{1M}=-\frac{\langle j'||\sigma||j\rangle}{\sqrt{3}}\,[c^{+}_{j'}b_{j}]^{1}_{M} \tag{13}$$

Let us verify the sumrule for the case of a pure two-neutron state $|0\rangle=[c^{+}_{k}c^{+}_{k}]^{0}|\,\rangle$, and assume that $|h\rangle$, the spin-orbit partner of $|k\rangle$, is filled. In this case,

$$\begin{aligned}
|a_1\rangle &= [b^{+}_{k}c^{+}_{k}]^{1}_{M}|\,\rangle\ ; \qquad |\langle a_1|Q_{1M}|0\rangle|^{2}=\frac{\langle k||\sigma||k\rangle^{2}}{3\hat{k}^{2}} \\
|a_2\rangle &= [b^{+}_{k}c_{h}]^{1}_{M}|0\rangle\ ; \qquad |\langle a_2|Q_{1M}|0\rangle|^{2}=\frac{\langle k||\sigma||h\rangle^{2}}{3} \\
|b,m'm\rangle &= c^{+}_{km}b_{h(m-M)}c^{+}_{km'}c^{+}_{k(-m')} \qquad (m\neq\pm m'\ ;\ m'>0) \\
|\langle b,m'm|Q'_{1M}|0\rangle|^{2} &= \frac{\langle k||\sigma||h\rangle^{2}\langle km;h(M-m);1M\rangle^{2}}{3\hat{k}^{2}}=\frac{\langle k||\sigma||h\rangle^{2}}{3}\left(1-\frac{1}{\hat{k}^{2}}\right)
\end{aligned} \tag{14}$$

Thus the sumrule is verified since

$$\frac{\langle k||\sigma||k\rangle^{2}\,+\,\langle h||\sigma||k\rangle^{2}}{3\hat{k}^{2}}=2 \tag{15}$$

The moral of this example is that the verification of the sumrule depends on the matrix elements of the operator Q'_{1M} (13) as much as on those of the original operator (5). The deviation from the leading order term in the last of eqs. (14) is a Pauli blocking effect on the giant resonance due to the existence of the two neutrons above the Fermi surface.

Within the NFT formalism, these effects are treated through the diagrammatic contributions.

5. Conclusions

In this note we have presented preliminary results concerning the treatment of spin-isospin transitions near close shell nuclei. We have adopted the formalism consisting of the simultaneous treatment of isoscalar pairing and GT vibrational states. The formalism was applied to a test case (transitions in ^{58}Cu) and the results show that the method is indeed feasible. The presence of accidental degeneracies may prevent for the inclusion of the complete set of diagrams prescribed by the NFT method, but the gross structure of the strength distribution, for spin excitations, is reproduced reasonably.

References

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