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SHELLS EVOLUTION AND CORE EXCITATIONS IN SEMI-MAGIC NUCLEI*

Frédéric Nowacki

Institut Pluridisciplinaire Hubert Curien
23 rue du Loess, BP28, F-67037 Strasbourg Cedex, France

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Recent advances in Large Shell Model calculations allow now to treat extended valence spaces and more complete descriptions of (semi-)magic nuclei can be achieved with inclusion of core excitations. The interplay between shell evolution and core excitations in semi-magic nuclei will be illustrated for tin isotopic chains in the framework of Large Shell Model calculations. *pn* and *nn* monopole relative influence will be traced back on Effective Single Particle Energies and $B(E2)$'s.

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1. Introduction

Like other finite quantum many-body systems, atomic nuclei are characterized by shell structure but also residual correlations among its constituents. Close to stability, shell structure has been established for long now and show up in identified magic numbers (2, 8, 20, 28, 40, 50 ...) corresponding to spherical closed shell nuclei. On the opposite, when active nucleons can interact through residual interaction, correlations can develop leading to collective deformation or superfluidity. But correlations also manifest in low-lying spectrum of magic nuclei and the presence of deformed states of many particles-holes nature have been observed in double magic nuclei like ^{16}O, ^{40}Ca or ^{56}Ni.

One of the main interest actually in nuclear structure is to study how this competition between magic spherical configurations and deformed ones evolve far from stability with neutron excess and if shell closures or magic numbers observed around the stability line still persist. The case of ^{32}Mg and the disappearance of $N = 20$ shell closure [1] was a pioneering case of this new domain of nuclear structure investigations.

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In a recent paper, Otsuka *et al.* [2] proposed an appealing mechanism driving this shells evolution. This mechanism is related to one of the characteristic part of the nuclear force which is the tensor force. Its importance has been shown to be crucial to the description of many nuclear properties.

Starting from a schematic tensor of the form

$$V_{\mathrm{T}} = (\vec{\tau}_1 \cdot \vec{\tau}_2)([\vec{s_1}\vec{s_2}])^2 Y^2) f(r)$$

they derived a closure relation for the monopole part of this interaction:

$$(2j_> + 1)V^{\mathrm{T}}_{j_>,j'} + (2j_< + 1)V^{\mathrm{T}}_{j_<,j'} = 0\,.$$

Here $j_>$ and $j_<$ denote for $l + \frac{1}{2}$ and $l - \frac{1}{2}$ respectively, l being the orbital angular momenta. From this relation one can follow the evolution of spin-orbit partners $j_>$ and $j_<$ of a given proton or neutron fluid with the filling of the other fluid j': it is seen that $V^{\mathrm{T}}_{j_>,j'}$ and $V^{\mathrm{T}}_{j_<,j'}$ have opposite signs and that also the tensor interaction being of dipole type, it is attractive between shells $j_</j'_>$ and repulsive between shells $j_>/j'_>$.

2. Tin isotopes

In the particular case of the tin region, the tensor monopole effects develop and can be put in evidence, for example in the spectroscopy evolution between ^{91}Zr and ^{101}Sn.

In Fig. 1, the experimental low-lying spectrum of ^{91}Zr is shown. The corresponding spectrum of ^{101}Sn was obtained by a fitting procedure of single particle energies and two-body matrix elements in the $(g_{\frac{7}{2}}, d_{\frac{5}{2}}, d_{\frac{3}{2}}, s_{\frac{1}{2}}, h_{\frac{11}{2}})$ $(r4h)$ valence space on top of a ^{100}Sn core. The clear effect of the filling of the $g_{\frac{9}{2}}$ proton orbital causes the drop of the $\frac{7}{2}^+$ state from 1.88 to 0.2 MeV. This neutron state is of single particle nature and rises down due to the strong pn attraction between spin-orbit partners $V_{\frac{7}{2},\frac{9}{2}}$. As quoted in [2], the tensor mechanism predicts also repulsive interaction $V_{\frac{11}{2},\frac{9}{2}}$ as could be inferred from the rise of the $\frac{11}{2}^-$ between ^{91}Zr and ^{101}Sn. While the state in ^{101}Sn is evidently the $h_{\frac{11}{2}}$ single particle state, the one in ^{91}Zr is carrying only a fractional part of the spectroscopic factor strength [3] and has probably strong admixture of $[(p_{\frac{1}{2}})^{-1}(g_{\frac{9}{2}})^1]_\pi \otimes (d_{\frac{5}{2}})_\nu$ configuration (inferred from Gross–Frenkel interaction [4] in ^{89}Zr for example). The experimental determination of the location of the $\mathrm{h}_{\frac{11}{2}}$ single particle centroid should therefore be crucial to determine the nature and the strength of $V^{pn}_{h_{\frac{11}{2}} g_{\frac{9}{2}}}$ monopole interaction.

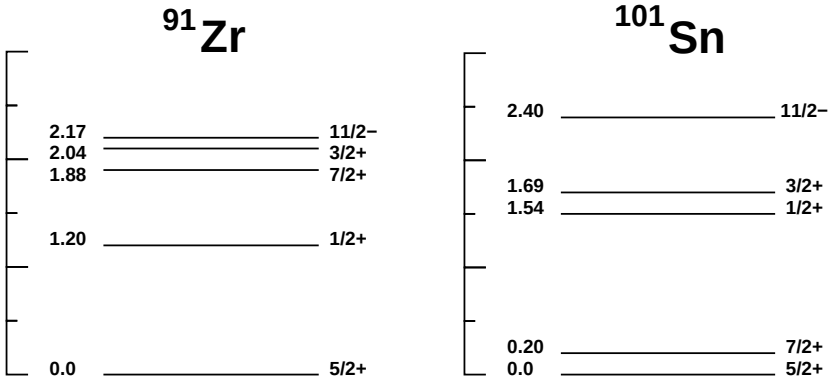

Fig. 1. Energy evolution of low-lying states between ^{91}Zr and ^{101}Sn.

At this point, it is interesting to notice that the spectroscopic properties all along the tin chain, from ^{101}Sn to ^{131}Sn (see Fig. 2), will be governed by the same interacting orbitals as the ones in the previous case from ^{91}Zr and ^{101}Sn. In particular, depending on the nature ($l \pm 1/2$) of the successive filled neutron orbitals, the interaction with $g_{\frac{9}{2}}$ protons of the core will induce variations of the proton $Z = 50$ gap.

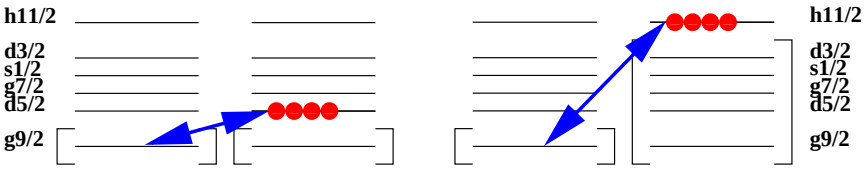

Fig. 2. Schematic filling evolution of neutron orbitals between ^{101}Sn and ^{131}Sn.

In left panel of Fig. 3 there are represented the $B(E2)$ values along all the isotopic tin chain. The discrete points are the experimental measurements while the continuous curves correspond to shell model calculations with different proton truncation levels ($0p0h$, $2p2h$, truncated and full $4p4h$). The calculations are performed in the *gds* valence space for protons and ($r4h$) for neutrons. Several points are to be commented from these curves:

- The experimental curves show an left/right asymmetry meaning enhanced transitions for the lightest tin isotopes. This is confirmed by the recent Rising measurement in ^{108}Sn [5].
- Shell model calculations presented here are able to reproduce the amplitudes of the transitions with standard polarization effective charge of 0.5. The need of core excitation is crucial and demonstrated by the large amplitude gain between $t = 0$ ($0p0h$) and $t = 4$ ($4p4h$) results.

- The convergence of calculations is somewhat reached for heavier tins (no strength gain between partial and full $4p4h$ calculations) while there is still some possible increase of the $E2$ strength for the lightest species, in particular due to neutron core excitation and non convergence of the $4p4h$ calculations. Such a gain should be sufficient to disymmetrize the curve, in agreement with the experimental trend.

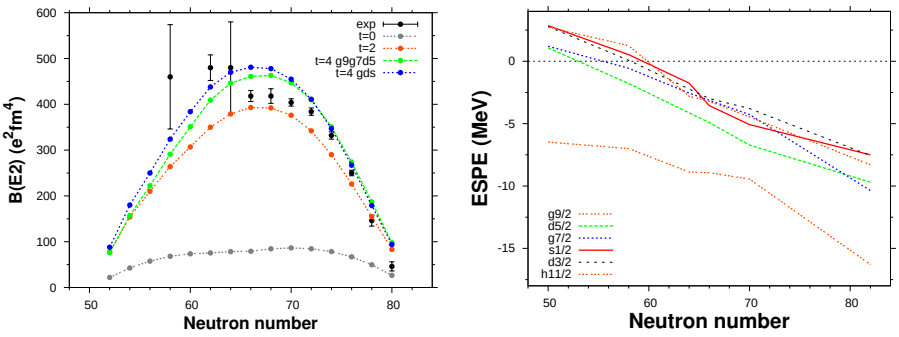

Fig. 3. Left: $B(E2)(2^+ \to 0^+)$ for Sn isotopes. Right: Proton ESPE for Sn isotopes.

It is interesting to put in correspondence previous panel with the right one of Fig. 3 where the underlying proton structure is illustrated with the evolution of the Effective Single Particle Energies [6] along the tin chain. The reduction of the proton gap at mid-shell explains the corresponding large amplitudes of the $B(E2)$'s. On the other hand, the $V^{pn}_{h\frac{11}{2}g\frac{9}{2}}$ monopole interaction (which could not be determined from ^{91}Zr spectrum) is needed to be sufficiently attractive (of the order of the $V^{pn}_{h\frac{11}{2}g\frac{7}{2}}$) to hinder the transitions at the end of the shell and close sufficiently ^{132}Sn. This appear to be in contradiction with the prediction of the tensor behaviour suggested in [2].

3. Conclusions

Large Shell Model calculations for the tin isotopic chain are performed for the calculations of electromagnetic $E2$ strength. Inclusion of core excitations allow to describe the amplitude of experimental values with standard polarization effective charges. In contradiction to the behaviour of a pure tensor force, the proton–neutron monopole interaction $V_{h\frac{11}{2}g\frac{9}{2}}$ need to be as attractive as $V_{h\frac{11}{2}g\frac{7}{2}}$ one.

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