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NEW VISTAS IN SUPERDEFORMATION

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NEW VISTAS IN SUPERDEFORMATION

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1. INTRODUCTION

The issue of nuclear shape coexistence, or nuclear superdeformation dates back to the mid–fifties; see the recent review[1]. Already in 1956 Morinaga[2] interpreted the deformed band in ^{16}O built upon the $I^{\pi}=0^{+}$ state at 6.049 MeV in terms of deformed $4p$–$4h$ configuration[†]. Another early example of multiple shape coexistence is the spectrum of ^{40}Ca with two well-deformed bands associated with $4p$–$4h$ and $8p$–$8h$ states at 3.352 MeV and 5.213 MeV, respectively[4].

In heavy nuclei, three regions of superdeformed (SD) shapes have been established: fission isomers in actinides, high spin SD states around ^{152}Dy, and SD bands around ^{192}Hg. An impressive experimental and theoretical effort has been devoted to exploring the underlying physics. There is no doubt that these investigations have opened up a new exciting field of nuclear superdeformed spectroscopy.

In this short contribution I would like to concentate on several aspects of nuclear superdeformation. The area of interest is so huge, and the space available is so limited, that a dramatic selection of the material had to be made. Some issues, such as the structure of collective excitations built on SD states, new symmetries of SD states, new predictions, etc., have been covered in other contributions to this Conference (Dudek, Matsuyanagi, Skalski, Delaroche, Sugawara–Tanabe, and others).

2. HYPERDEFORMATIONS

Hyperdeformed nuclei, i.e., with quadrupole deformations significantly larger than $\beta_2=0.6$ are known or predicted in several mass regions.

Good examples of very elongated configurations can be found in light nuclei. For example, the hyperdeformed state in ^{12}C (three aligned alpha particles) built on the 0^{+} resonant state at 10.3 MeV becomes yrast already $I^{\pi}=4^{+}$. The calculated low–lying reflection–asymmetric hyperdeformed minimum ($\varepsilon_2=1$, $\varepsilon_3=0.3$) in ^{24}Mg[5] can be associated with the asymmetric ^{16}O+α+α (or ^{16}O+^{8}Be) structures or the symmetric hyperdeformed α+^{16}O+α states[6]. Other examples are the hyperdeformed states in ^{36}Ar (^{16}O+^{16}O+α),^{48}Cr (^{16}O+^{16}O+^{16}O) (see discussion in ref.[7]), or a six-α chain structure in ^{24}Mg reported by Wuosmaa *et al.*[8], and see the contribution by R.R. Betts.

In heavy nuclei the best hyperdeformed states are the so–called third minima in nuclei around ^{232}Th. In these nuclei the second saddle point is split leading to the weak reflection–asymmetric minimum with $\beta_2\sim0.85$, $\beta_3\sim0.35$[9–11]. Experimentally, the third minimum shows

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†For the $4\hbar\omega$ shell model calculations for ^{16}O, see ref.[3] and the contribution by W.C. Haxton.

up as an alternating-parity microstructure of resonances near the (n,f) fission thereshold[12]. Recent calculations based on the Gogny–HF model[13] or the Woods–Saxon model[14] predict the the third minima to be deeper than in the previous calculations based on the Nilsson model[11]. Figure 1 displays the Woods–Saxon potential energy surface for ^{232}Th. The heights of the

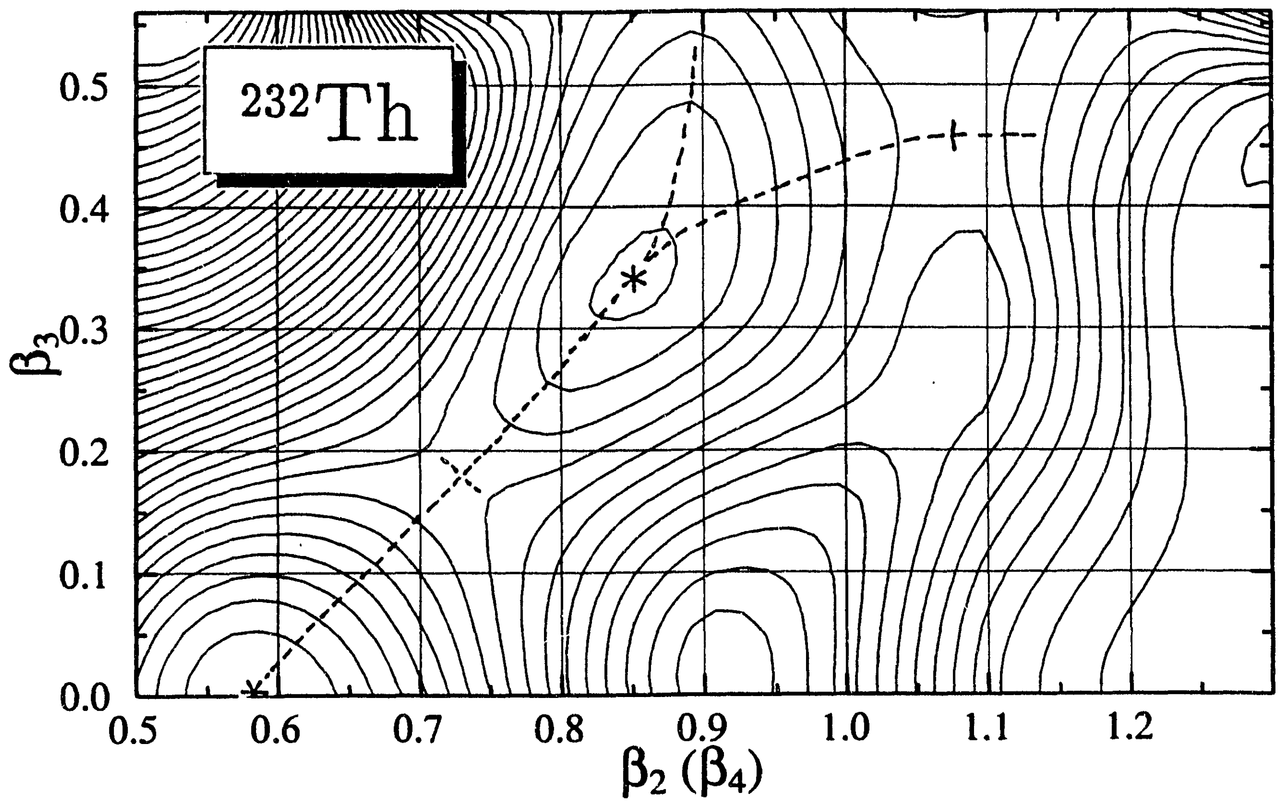

Fig. 1. The Woods–Saxon–Strutinsky total potential energy for ^{232}Th as a function of β_2 and β_3. At each (β_2, β_3) point the energy was minimized with respect to β_4–β_6. The distance between the contour lines is 0.5 MeV. (From ref.[14].)

second and third barrier (saddle point) are 1.4 MeV and 2.9 MeV, respectively. This should be compared with the value of ~400 keV from ref.[11].

For recent evidence for hyperdeformation in ^{152}Dy around I=80 $\hbar$; see the contribution by A. Galindo–Uribarri *et al.*.

3. SUPERDEFORMATIONS FAR FROM STABILITY

According to the model predictions (see, e.g., refs.[15, 16]) low–lying super and hyperdeformed configurations are expected in various mass regions, many of them being practically inaccessible with the present detector systems (too low cross sections), or inaccessible at all using combinations of stable beams and targets. Some of those "white spots" will, hopefully, be investigated in the future – thanks to the new–generation multidetector arrays (EUROGAM, GAMMASPHERE, EUROBALL), or exotic (radioactive) ion beam facilities currently being constructed in Europe, U.S.A., and Japan.

According to predictions of the mean field theory nuclei in the N=Z≃40 mass region favor 2:1 shapes. This tendency remains at high spins, and superdeformed configurations in nuclei around $^{82}_{38}Sr_{44}$ or $^{88}_{44}Ru_{44}$ are predicted to be yrast at spins ~30–40 ℏ[17–20]. Interestingly, cranked Skyrme Hartree–Fock calculations, cranked Woods–Saxon calculations, and cranked calculation based on the relativistic mean field theory give rather similar results. Since neutron and proton numbers are rather close, the neutron and proton SD configurations are similar, corresponding to the alignment of one or two $1h_{11/2}$ protons and neutrons. At very high spins also the lowest $1i_{13/2}$ orbitals become occupied. On the neutron rich side, the best prospects for superdeformation are expected to be in nuclei around $^{108}_{44}Ru_{64}$. Indeed, by combining the SD gaps at Z=44 and N=64 one obtains very favored SD structure which becomes yrast around I=35 ℏ[20, 21].

Very little is known about the very neutron–deficient Hg nuclei with N~96. The lightest system known from in–beam studies is ^{180}Hg[22]. The nucleus ^{178}Hg has lifetime τ~49 s whilst for ^{176}Hg τ drops to 34 ms (see ref.[23]), and ^{174}Hg is expected to be proton–unstable.

The potential energy surfaces for ^{170}Hg (proton unstable), ^{180}Hg (τ=5.9 s), ^{190}Hg (τ=20 m),

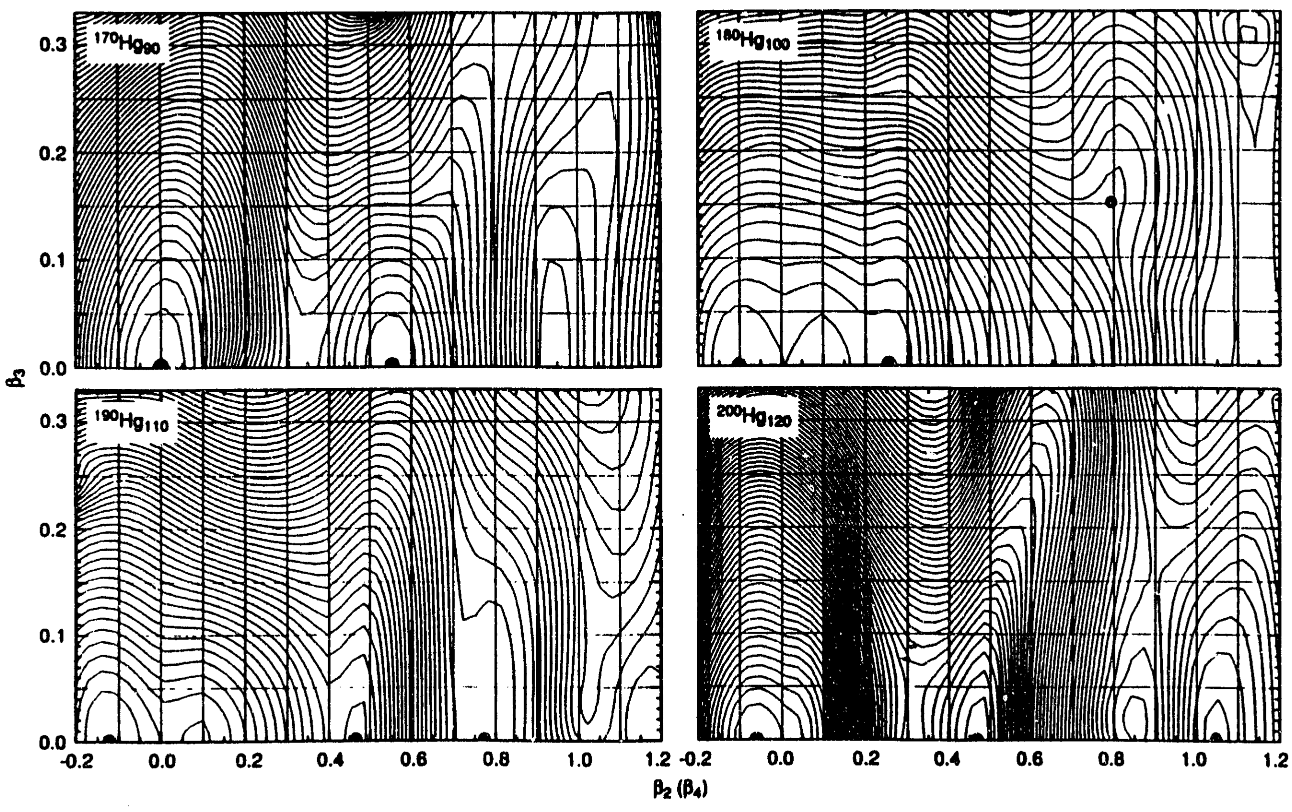

Fig. 2. The Woods–Saxon–Strutinsky total potential energy for ^{170,180,190,200}Hg as a function of β_2 and β_3. At each (β_2, β_3) point the energy was minimized with respect to β_4–β_8. The distance between the contour lines is 0.3 MeV.

^{200}Hg (stable), are shown in fig. 2. The SD minimum seen in ^{190,200}Hg disappears in ^{180}Hg[‡].

[‡]Note the presence of reflection–symmetric hyperdeformed minima in ^{190}Hg (β_2~0.8), ^{200}Hg (β_2~1.05), and

However, when decreasing neutron number the SD states reappear again; see the map for ^{170}Hg. Detailed calculations presented in fig. 3 indicate that the excitation pattern of low-deformation

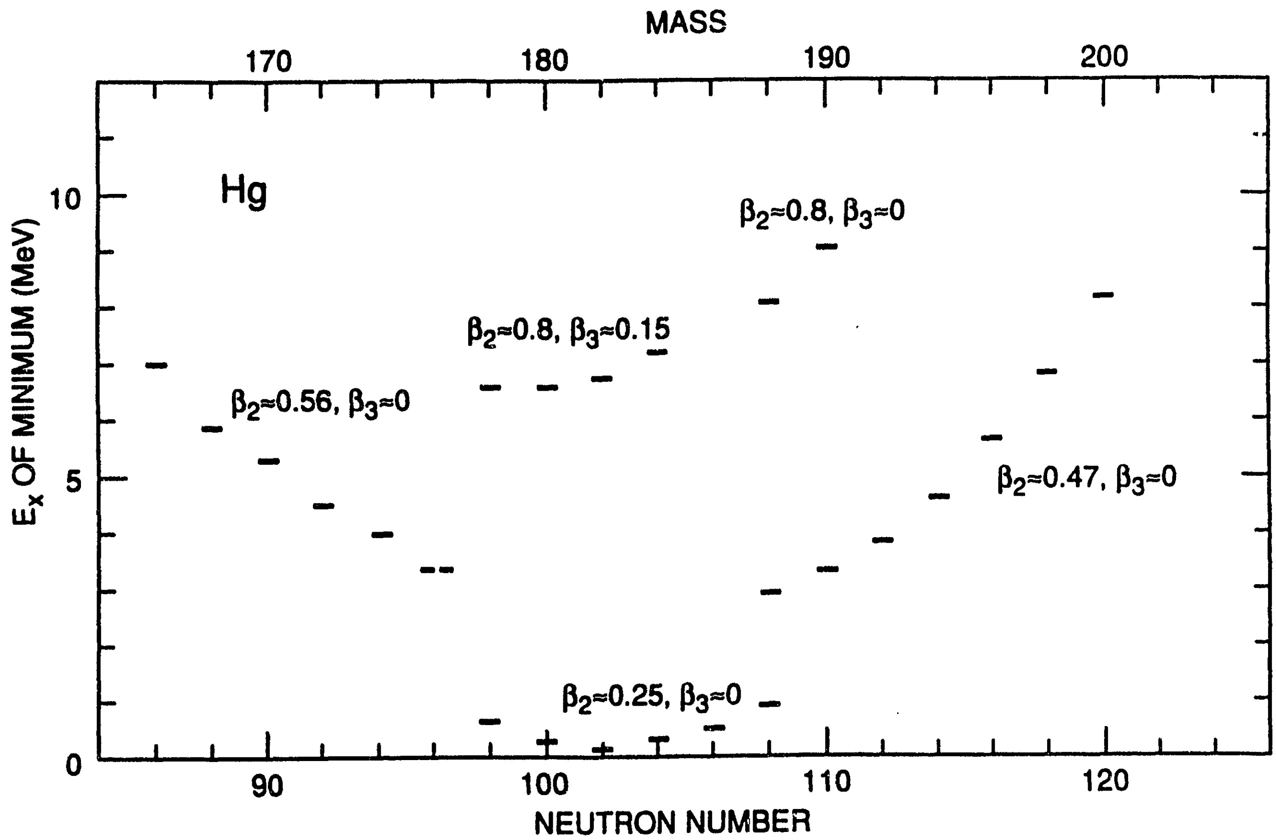

Fig. 3. Calculated energies of excited shape-coexisting states in even-even Hg isotopes with 86≤N≤120.

shape-coexisting configurations and SD states is symmetric with respect to the middle of the shell, i.e., N∼102. A similar situation has also been calculated for Pt and Pb isotopes.

A new region of hyperdeformed shapes has been predicted around ^{120}Ce[24]. These structures involve N=6 protonns and N=7 neutrons and are expected to become yrast around I=50 $\hbar$.

As discussed in the proposal for the Oak Ridge Radioactive Ion Beam Facility RIB[25], the new beams at RIB will provide the necessary tools to access high-spin states in the N=Z≃ 40 mass region, the region around ^{120}Ce, and around ^{176}Hg; see the contribution by J.D. Garrett.

4. EFFECTIVE INTERACTIONS, EFFECTIVE OPERATORS

The structure of single-particle states around the Fermi level in SD nuclei is significantly different from the pattern familiar from normal deformations. Indeed, the SD supershells consist of states with very different spatial character. Consequently, the commonly used effective and residual interactions are probably different in the superdeformed world.

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reflection-asymmetric hyperdeformed minima in ^{180}Hg (β_2∼0.8, β_3∼0.15).

It is well known that the low and high-frequency nuclear vibrations can be associated with the residual multipole–multipole interactions. Do the standard multipole–multipole forces describe correctly collective excitations at SD shapes? Sakamoto and Kishimoto[26, 27] have introduced the doubly-stretched multipole moments,

$$Q''_{\lambda K} = f_{\lambda K}(r'')Y_{\lambda K}(\Omega''), \quad x''_i \equiv \frac{\omega_i}{\omega_0}x_i, \ (i = 1, 2, 3), \tag{1}$$

where ω_i is the i–th oscillator frequency. The corresponding multipole–multipole doubly–stretched interaction can be viewed as an improved conventional multipole–multipole force. Firstly, it satisfies the nuclear–selfconsistency rigorously even if the system is deformed. Secondly, it yields the zero–energy RPA spurious modes, i.e., it automatically separates the translational (λ=1, T=0) and reorientation (λ=2, K=1, T=0) modes. Last but not least, for the doubly–stretched interaction the coupling between octupole and dipole modes disappears.

At normal deformations, the spatial difference between doubly–stretched and normal multipole–multipole forces is rather small. At large deformations, however, these interactions give markedly different predictions[26, 27]. For instance, the presence of large dipole moments (or enhanced B(E1) rates at SD shapes) is a direct consequence of the doubly–stretched octupole force. The K=0 and K=1 $(r^3Y_{3K})''$ operators are linear combinations of the ordinary octupole fields, r^3Y_{3K} and the compressional dipole fields, r^3Y_{1K}[28]. Again, if the doubly–stretched residual interaction is realised in nature, strong dipole transitions de-exciting superdeformed octupole states should be present. For more discussion, see the contributions by K. Matsuyanagi and J. Skalski.

The influence of large shape-elongations on the pairing field is still not well understood. For example, in the "doubly-magic" SD nucleus ^{152}Dy pairing is expected to play a minor role[29–32]. Indeed, due to the very low level density of single-particle states the superfluid-type correlations in this band are expected to be seriously quenched and mainly of a dynamical character. In ^{150}Gd the large increase of thy dynamical moment of inertia, $\Im^{(2)}$, in the lower part of SD band[33] has been interpreted as a *paired* band crossing associated with an alignment of the N=7 neutron pair[34, 31]. A similar crossing has been found in the first excited SD band in ^{149}Gd[35]. Another piece of experimental evidence suggesting the presence of pairing at SD shapes is a steady increase of $\Im^{(2)}$ in the SD bands in the A∼190 region, which can be attributed[36, 37] to the alignment of N=7 neutrons and N=6 protons. Calculations without pairing yield fairly constant moments of inertia.

However, there are also many pieces of evidence that pairing correlations are extremely weak at large elongations. For instance, the moments of inertia in SD bands are very close to their rigid-body values and even a strongly reduced pairing field yields too strong quasiparticle alignment[37, 38]. Moreover, only very weak evidence for blocking effects in SD configurations has been found in the A∼190 region[39].

Let us consider the general pairing interaction

$$H_{pair} = -\sum_{i,j,k,l} G_{ij;kl}c_i^{+}c_{\bar{j}}^{+}c_{\bar{l}}c_k, \tag{2}$$

where $G_{ij;kl} =< i\bar{j}|v|k\bar{l} >_{AS}$ is the antisymmetrized matrix element of the two-body pairing interaction. The matrix elements $G_{ij;kl}$ have been calculated by many authors using various residual interactions like the delta force, the surface delta interaction, the Skyrme force, the finite-range D1S interaction, or a density dependent delta interaction[40–43] (see also refs.[44, 45]). It has been found that the pairing matrix elements are relatively enhanced for orbitals with similar values of $\langle n_z\rangle/\langle N\rangle$, i.e., orbitals with good angular overlap. In particular, it has been observed that the pairing matrix elements between the high-j intruder orbitals (such as [660]1/2 and [651]3/2) are rather large, as are those between high-j intruder states and the natural-parity orbitals with $j=N-1$ and similar Ω-values (such as [660]1/2 and [541]1/2). At normal deformations the single-particle unique-parity orbitals are relatively close to each other and to normal-parity states with similar spatial overlap. However, at large deformations states originating from completely different shells approach the Fermi level. These states are very weakly coupled through pairing interaction. Moreover, the "favored" coupling between unique-parity levels is diminished because of their large deformation splitting. In view of the above it is likely that the pairing correlation energy should decrease with deformation – an effect that is analogous to the fragmentation of pairing matrix elements caused by the Coriolis force.

A simple parametrization of the particle–particle interaction is usually based on the multipole expansion. Usually, only the monopole term (seniority pairing) is considered since for the delta force the $L=0$ component is about five times stronger than the $L=2$ (quadrupole) term[46]. At superdeformed shapes, however, most single–particle orbits close to the Fermi level are, on the average, of definite prolate character. That suggests that the quadrupole pairing interaction should play an important role. In particular, it is expected, that the inclusion of quadrupole pairing should shift the crossing frequency towards higher values[47] – a welcome effect in light of the recent data on SD band ^{192}Hg[38] (as shown in the contribution by R.V.F. Janssens). The previously obtained good fit for the dynamical moment of inertia[48] can easily be retained by adjusting the relative strengths of the $L=0$ and $L=2$ ($M=0,1,2$) components[47]. Preliminary results of calculations involving the quadrupole pairing interaction by R. Wyss[49] are very encouraging.

A new area of interest concerns spin polarization and the structure of the magnetic moment operator at large deformations. The measurement of the g factor of the fission isomer in ^{237}Pu[50] has put strict limits on the parameters of existing mean-field-based models: the value $g = -0.45(3)$ indicates that the parameters of the deformed single-particle potential exhibit a significant deformation dependence[51]. The intrinsic spin g-factors are usually reduced with respect to their free values to account for spin polarization of the core, i.e. $g_s=fg_s^{free}$, with f typically varying between 0.6 and 0.7 at normal deformations. However, for large elongations, the spin polarization may make contributions to the magnetic moments that cannot be simply accounted for in terms of the renormalization of the free values. Indeed, the spin-polarization effect in spin-unsaturated orbitals due to an odd particle is associated with the spin-dependent components of the effective nucleon-nucleon force, e.g., the spin-spin interaction[52]. The renormalization of g_s is, in the first order, due to the spin interaction between spin-orbit partners with $j=l\pm1$, and should exhibit a deformation dependence. Moreover, even within first-order perturbation theory, one can argue that spin- and isospin-dependent nuclear forces may lead to a renormalization of the orbital g_l factors and the presence of the tensor component[52, 53], which is expected to have a significant deformation dependence.

4. SPIN ASSIGNMENTS AND IDENTICAL BANDS

Parametrization of rotational spectra dates back to more than thirty years ago when Bohr and Mottelson used a phenomenological formula to characterize properties of rotational bands[54]. Later on, extended and revised versions of this early formalism were used to classify collective properties of the *ground state bands* in a wide range of nuclei[55, 56]. The problem of spectral fitting has recently been revisited in the context of the superdeformed bands in the Hg region[57]. Although a large number of high-spin superdeformed bands have been found around the doubly-magic SD ^{152}Dy and ^{192}Hg nuclei, their *absolute* spin assignments are so far lacking. This is because of the fact that the gamma rays that connect the SD bands to the known levels at lower spins have not been identified experimentally[§]. The correct angular momentum assignment has become a central issue after the discovery of *twinned bands* (i.e., bands having identical transition energies) in neighboring odd and even nuclei[58, 59].

In refs.[60, 61] a detailed analysis of the spin–fitting procedure has been made and several criteria have been introduced that are crucial for a meaningful spin assignment. The relative alignments of the twinned SD bands in the Hg region have been then obtained from a power-series expansion of angular momenta in terms of the measured transition energies.

Subsequent studies of such fitting procedures have pointed out[62, 63] that, because of the lack of knowledge about the low–lying transitions in the SD bands, the fitted spins are subject to uncertainties. In a recent ref.[64] it has been pointed out that possible presence of a non–zero initial alignment, i_o, can lead to a serious uncertainty in the absolute–spin determination of the *excited bands*. Since this quantity *cannot* be determined from a least-squares fitting procedure, the above fitting formalisms could yield incorrect spins, despite the superb quality of these fits as judged by their χ^2 values. The problem of non–zero initial alignment was recognized in refs.[65, 60]. However, the authors assumed a zero value as being the most plausible choice at superdeformed shapes where the deformation alignment (strong coupling) dominates. As briefly discussed in[62, 63], this problem seriously limits the applicability of the suggested fitting formalisms to the excited bands, which can potentially have non–zero initial alignments. Results of our extensive investigation of this problem in the normally–deformed excited bands where spins are known show that: (i) the fitting procedure oftentimes leads to wrong spin assignments even if γ–ray energies of all low-lying members of the band are known; (ii) there is no simple relationship between the angular momentum obtained from the optimization procedure and the experimentally determined spins; and (iii) the fitted spins are usually dependent on which data points are included in the fitting procedure.

Recently, several models have been proposed to explain identical SD bands, as well as the alignment patterns that have been obtained from phenomenological spin–fitting procedures, despite the problems alluded above. Among them are microscopic models based on the mean field approach (Hartree–Fock, Nilsson–Strutinsky) or the shell model (pseudo–SU_3 model, Fermion Dynamical Symmetry Model), formulas obtained from the symmetry limits of specific group theoretical models (Interacting Boson Model, supersymmetric schemes, models based on quantum groups), or other scenarios (triplet pairing). A critical review of several of the models

[§]As we have learnt during this Conference, such a connection has just been found in the nucleus ^{143}Eu, see the contribution by A. Ataç *et al.*

and scenarios that purport to explain the origin of the identical bands, or the patterns of the fitted alignments has also been presented in ref.[64]. It has been concluded that many of these formulations can be reduced to simple expressions for energies which strongly resemble the standard VMI model. However, a large majority of these models either explicitly *assume* identical moments of inertia, or *impose* an ad hoc symmetry on the model Hamiltonian to obtain the desired outcome of identical bands. Similarly, gross approximations made by some other models cast serious doubts on their ability to address such subtle effects as the constancy of moments of inertia. In spite of many efforts and new and interesting ideas and suggestions, none of the scenarios proposed so far addresses the fundamental question at hand, namely the microscopic origin of identical moments of inertia that have been observed in a wide range of normally–deformed and superdeformed nuclei. A more complete discussion is given in the contribution by C. Baktash.

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