
Nuclear Structure Physics



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Nuclear Structure Physics

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Preface

How neutrons and protons can be joined together to form a nuclei? At most how many neutrons and protons can be held together to form a nuclei? This fundamental riddle seeks to be answered yet. What we have learnt so far and has also been testified with extremely sophisticated experiments is that the nucleus of an atom is held together by the strong nuclear force that binds together protons and neutrons. Most importantly, this ‘glue’ – strong nuclear force, which is the strongest of the four fundamental forces – acts only between nuclear particles and over very short – typically nuclear – distances where this nuclear force is much more powerful than the electrostatic repulsion between protons. Also, this ‘glue’ is slightly stronger when acting between a proton and a neutron than between either pair of like particles. As a result, atoms are usually stable so long as the numbers of protons and neutrons are not too uneven. If this balance isn’t right, atoms can split apart through radioactive decay or nuclear fission. If an atom gets too heavy with neutrons or protons, extra nucleons simply won’t stick at all — not even for an instant. Nuclear physicists have long been trying to map out where this boundary of stability line – termed to be the ‘drip line’ – exists, as nuclei beyond this point are like oversized droplets that drip small fragments. Similarly, there are theoretical speculations that there could lie a region of very stable nuclei across the unstable ones at the extremes of mass, i.e. in the superheavy region of the nuclear landscape. Key to all these questions in knowing the exact form of nuclear force, this is yet to be figured out, making nuclear physics a forbidding subject, even for highly trained physicists. It is important to mention here that nuclear physics research began with the investigation of a handful of nuclei (few tens), but today, this number has grown many folds (few thousands) due to tremendous efforts and contributions made in last 100 years.

A systematic theoretical treatment of nuclei provides a simple lab of complex nuclei, which could provide access to those areas of nuclear landscape, which is still beyond experimental reach. In a way, it is an extraordinarily exciting time for researchers in nuclear structure physics due to the advance tools available to experimentalists as well as theoretical researchers. On the one hand, there has been a great progress in detector techniques helping in the detection of nuclei lying on the extremes of mass and isospin; on the other hand, involvement of supercomputing facilities has led to the development of ab initio-based theoretical approaches increasing the accuracy in the study of many body systems phenomenally. The book chapters cover many interesting topics to be seen through the window of nuclear structure. These include mean field model and recent parametrization, magic nuclei, superheavy nuclei, bubble nuclei, and also topics related to weak interaction. These chapters not only discuss the potential of the subject to impact other current research areas such as mapping the astrophysical p-process in nucleosynthesis, and the exotic one- and two-proton decay modes but also showcase the applicability and success of the different nuclear effective interaction parameters near the drip line, where hints for level reordering have already been seen, and where one can test the isospin dependence of the interaction. Overall, the chapters written by highly experienced and well-known researchers/experts from academia may serve as a standard, up-to-date research reference book for the postgraduate and graduate students, and researchers working in related fields.

We would like to convey our appreciation to all the contributors as well as our academic collaborators, who have been always a source of inspiration to work in this field of research.

Our special thanks are also due to Ms. Aastha Sharma, Ms. Shikha Garg, Mr. Varun Gopal, and the entire editorial team members from CRC publishers for their kind support and great efforts in bringing the book to fruition. Last but not the least, editors are sincerely thankful to their family members, whose unending support and affection has always been a driving force in making any professional assignment a success!

Amritanshu Shukla
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Editors



Dr. Amritanshu Shukla earned his Master's degree in Physics from University of Lucknow (1999) and his Ph.D. in Physics from IIT Kharagpur (2005), and did his postdoctoral work at Institute of Physics Bhubaneswar (under the Department of Atomic Energy, Govt. of India); University of North Carolina Chapel Hill, USA; University of Rome, Italy; and Physical Research Laboratory Ahmedabad (under the Department of Space, Govt. of India). Currently, he is an Associate Professor in Physics and also the Head of Division at Rajiv Gandhi Institute of Petroleum Technology (RGIPT) (set up through an Act of Parliament by the Ministry of

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His research interests include theoretical physics, nuclear physics, and physics of renewable energy systems. He has published about 100 research articles in various international journals and conference proceedings of national and international repute. Currently, he is working on the different problems in nuclear structure and nuclear reactions with a special focus on bubble nuclei, drip-line nuclei, and superheavy nuclei. He is also supervising several research students, actively involved in a number of sponsored research projects, and having several national as well as international active research collaborations from India and abroad on the topics of his research interests.



Prof. Suresh Kumar Patra is a well-known Indian Theoretical Nuclear Physicist. He earned his Ph.D. (Theoretical Nuclear Physics) from Institute of Physics, India. After his Ph.D., he has done postdoctoral studies at Tohoku University (Japan), J.W. Goethe University (Germany), Chung Yuan Christian University (Taiwan), University of Barcelona (Spain), and University of Surrey (U.K.). He also had short visits to other reputed universities for collaborative research work and to give lectures. Currently, he is a Professor at Institute of Physics, Bhubaneswar.

Prof. Patra has a wide-ranging expertise in various areas of nuclear physics, such as nuclear structure, nuclear reaction, and nuclear astrophysics. He has contributed considerably to the structure of exotic and superheavy nuclei, nucleus–nucleus reaction for both beta-stable and beta-unstable nuclei, and structure of neutron stars and gravitational waves from their merger. He has published about 170 research articles in international peer-reviewed journals, more than 13 book chapters, and about 160 contributions to various national/international symposiums/workshops and is the co-editor of several conference proceedings. He has also served as the organizer/advisory member of various national/international symposiums/workshops and the referee of various journals. He has guided more than 16 Ph.D. and 8 postdoctoral scholars.

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Magic Numbers of Cylindrical Symmetry

Andriana Martinou and Dennis Bonatsos

National Centre for Scientific Research “Demokritos”

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1.1 Introduction

Magic numbers were one of the first discoveries in nuclear structure [1]. It was realized that nuclei formed by particular numbers of protons (Z) and neutrons (N) were extremely stable. This experimental discovery led to the concept of a nuclear closed shell [2], in analogy to atomic closed shells. As a result, the shell model was introduced as the basic microscopic model for the description of atomic nuclei, its magic numbers for either protons or neutrons being 2, 8, 20, 28, 50, 82 [3–7]. The next magic number for neutrons is known to be 126 [8], while the one for protons remains elusive and is currently the subject of intensive research in superheavy elements [9–12].

At the infancy of nuclear structure, it has also been realized that excitation energies of the atomic nucleus can be arranged into bands [13] with increasing angular momentum J . Soon thereafter, it was realized that energy bands near closed shells would increase almost linearly with angular momentum, i.e. $E(J) = AJ$, therefore called vibrational or spherical, while energy bands near the middle of the nuclear valence shells would increase almost following the rigid rotator expression $E(J) = AJ(J+1)$, therefore called rotational or deformed [14]. From the macroscopic point of view, these bands have been described in the framework of the Bohr Hamiltonian [15], soon extended into the Bohr Mottelson model [16,17], in terms of the collective deformation parameters β and γ , of which the former describes the degree of deviation of the nuclear shape from sphericity, while the latter describes its departure from

spherical symmetry towards cylindrical symmetry realized through axial prolate shapes (elongated like the ball of American football) or axial oblate (pancake-like) shapes, or even towards purely triaxial shapes [18,19], in which the cylindrical symmetry is also broken. It should be noticed that the Bohr Mottelson model [15–17,20,21] involves five dimensions, the two collective variables β and γ plus the three Euler angles describing the orientation of the atomic nucleus in space. In addition to the angular momentum J , nuclear bands are additionally characterized by the projection of the angular momentum on the z -axis of the coordinate system attached to the nucleus [14,22], which is called the intrinsic set of coordinates. In the case of even nuclei, the symbol K is used for this quantum number. Bands with $K = 0$ contain the angular momenta $J = 0, 2, 4, 6, 8, \dots$. The lowest one is the band built on the ground state (which is always zero for even-even nuclei, i.e. nuclei with an even number of protons and an even number of neutrons), called the ground state band. For the rest of $K = 0$ bands, the term β bands is used. In contrast, bands with $K = 2$ contain the angular momenta $J = 2, 3, 4, 5, 6, \dots$, and are called γ bands [14,22].

Symmetries that use the language and techniques of group theory [23–29] have also been used for the description of atomic nuclei [30–36]. It has been realized [37–42] that the Bohr Hamiltonian possesses a $U(5)$ symmetry, having an $O(5)$ subalgebra, which in turn has an $SO(3)$ subalgebra. The $SO(3)$ subalgebra is needed within all symmetries describing atomic nuclei, since the quantum number determining its irreducible representations is the angular momentum J . Therefore, $SO(3)$ has to be present, so that the labeling of energy levels by the angular momentum quantum number J is possible. $U(5)$ is counting the number of excitation quanta present, while the irreducible representations of $O(5)$ are characterized by the seniority quantum number [37–40] corresponding to the number of nucleon pairs coupled to nonzero angular momentum. The $O(5)$ symmetry imposes certain degeneracies in the spectrum. For example, the first excited $J = 2$ state has to be degenerated with the first $J = 4$ state, the first excited $J = 4$ and $J = 3$ states have to be degenerated to the first $J = 6$ state, and so on (see Table I and Figure 3 of Ref. [43] for further details). Since the seniority quantum number appears to be near closed shells [7], $O(5)$ is a symmetry expected to appear there.

For deformed nuclei, the $SU(3)$ symmetry, possessing an $SO(3)$ subalgebra, has been used. In addition to the three components of the angular momentum, forming the $SO(3)$ subalgebra, it also contains the five components of the quadrupole operator. Its relevance to nuclear spectra has been demonstrated by Elliott [44–47] and will be further discussed in Section 1.4.

An algebraic model incorporating $U(5)$ and $SU(3)$ as special subcases is the Interacting Boson Model [30–34,36] possessing an overall $U(6)$ symmetry. In addition to the $U(5)$ [48] and $SU(3)$ [49] subalgebras, it also contains an $O(6)$ [50] subalgebra, which in turn contains $O(5)$ and $SO(3)$ subalgebras. The $O(6)$ symmetry [50] turns out to be applicable to γ -unstable nuclei, i.e. to nuclei which can change easily their axial or triaxial character at the expense of no energy. Since the $O(6)$ symmetry contains the same $O(5)$ and $SO(3)$ subalgebras as the $U(5)$ symmetry, they exhibit the same set of degeneracies in their bands. In the Interacting Boson Model, the nucleus is approximated by a set of bosons corresponding to correlated valence proton pairs and valence neutron pairs, i.e. nucleon pairs outside closed shells. But symmetries appear also in the fermionic description of the nucleus [35], as we shall see below.

It should be noticed that the study of the three-dimensional harmonic oscillator (3D-HO) in spherical coordinates leads to shell model algebras possessing $SU(3)$ subalgebras [23,51,52]. We shall use the spectroscopic notation, in which angular momenta $J = 0, 1, 2, 3, 4, 5, 6, \dots$ correspond to the symbols s, p, d, f, g, h, i, $\dots$ respectively. The p shell closes a $U(3)$ algebra; the sd shell closes a $U(6)$ algebra; the pf, sdg, pfh, and sdgi shells close $U(10)$, $U(15)$, $U(21)$, and $U(28)$ algebras, respectively; and all of them possess $SU(3)$ and

SO(3) subalgebras [52]. The s shell in the beginning of this series possesses a U(1) algebra. Since each value of the angular momentum can accommodate $2(2J+1)$ nucleons of the same type (protons or neutrons), the factor of two due to the two possible orientations of spin, one can easily work out the magic numbers appearing in the 3D-HO, which are 2, 8, 20, 40, 70, 112, 168, ...

One can remark that the shell model magic numbers and the 3D-HO magic numbers are identical up to 20. This means that up to the sd shell, the SU(3) symmetry of the 3D-HO oscillator is present, as first realized by Elliott [44–47]. The agreement between the shell model magic numbers and the 3D-HO magic numbers is destroyed beyond the sd shell by the spin-orbit interaction, which is a relativistic effect which has to be introduced by hand in nonrelativistic shell model Hamiltonians in order to establish agreement with the experimentally observed magic numbers [3–7]. Various approximation methods have been developed over the years, attempting to reestablish the SU(3) symmetry beyond the sd shell. Pseudo-SU(3) [53–57] and quasi-SU(3) [58,59] were the earliest ones, followed recently by proxy-SU(3) [60–62], to be considered in more detail in Section 1.5. The search for approximate SU(3) symmetries beyond the sd shell has been stimulated by the tremendous amount of existing analytical and mathematical work in SU(3), which in this way would become available for use in the nuclear structure framework.

Following a different path, deformed nuclear shapes have been described since a very early stage within an elementary shell model, using a three-dimensional (3D) anisotropic harmonic oscillator with cylindrical symmetry (occurring when two of the three frequencies are equal to each other), to which the spin-orbit interaction has been added. This is the Nilsson model [63,64], which over the years has been extremely useful in classifying and understanding a huge bulk of experimental data [65,66]. The Nilsson Hamiltonian is diagonalized in a set of states which become exact eigenstates only at large deformations, called the asymptotic basis wave functions [67–69], which are characterized by the quantum numbers $K[\mathcal{N}n_z\Lambda]$, where $\mathcal{N}$ is the number of oscillator quanta, n_z is the number of quanta along the cylindrical symmetry axis, Λ and K are the projections of the orbital angular momentum and the total angular momentum, respectively, along the same axis. These quantum numbers remain rather good even at intermediate deformation values [64]. Ben Mottelson [66] has remarked that the asymptotic quantum numbers of the Nilsson model can be seen as a generalization of Elliott’s SU(3), applicable to heavy deformed nuclei.

A nontrivial question arising above is the following one: SU(3) is known to be the symmetry of the 3D isotropic harmonic oscillator [23,51]. How is one allowed to use it in the case of the deformed (anisotropic) harmonic oscillator with cylindrical symmetry [70]? This is a formidable mathematical problem, which has already been addressed in the literature in several different ways, based on the fact that the z -component of the angular momentum operators remains intact by the transition from spherical to cylindrical symmetry [71–76]. In particular, using quantum group techniques [77–79], it has been shown how the irreducible representations of the anisotropic HO with rational ratios of frequencies can be constructed, leaving the z -component of the angular momentum operator unchanged [74]. Furthermore, Smirnov and collaborators have shown how the isotropic HO can be transformed into the anisotropic oscillator through the use of a dilatation operator [80]. The resurrection of the SU(3) symmetry of the isotropic HO for large anisotropic deformations has been proved numerically in [81].

We have described above some basic nuclear shapes, namely, spherical, γ -unstable, axially deformed, and the symmetries corresponding to them, viz. U(5), O(6), and SU(3), respectively. There are also nuclei exhibiting behavior intermediate among these symmetries [82]. Considering various series of isotopes, one can see a gradual development from one symmetry into another. If at some point this transition becomes abrupt, we say that we have a shape/phase transition [83–86]. Such an abrupt change is seen in the transition

from vibrational (U(5)) to γ -unstable (O(6)) nuclei, called the E(5) critical point symmetry [87–91]. Another case appears in the transition from spherical (U(5)) to axially deformed prolate (SU(3)) nuclei, called the X(5) shape phase transition [92–95]. According to the Ehrenfest classification [84], E(5) is a second-order phase transition, while X(5) is a first-order one, i.e. it is more abrupt. As a consequence, it is much easier to locate experimental examples for X(5) than for E(5) (see the review articles [96–98] for specific experimental manifestations of these critical point symmetries). Extensive literature exists [86] on shape phase transitions, within both the Bohr collective framework and the Interacting Boson Model approach.

A peculiar situation, which might be related to the concept of shape phase transitions [99,100], occurs in several even-even nuclei, in which the ground-state band is accompanied by another $K = 0$ band, which lies close in energy but possesses a radically different structure [101,102]. This effect, which also appears in odd nuclei [103], is called shape coexistence and is attracting recently wide interest. In particular, shape coexistence is known to appear in certain areas of the nuclear chart, while it seems to be absent in others (see Figure 8 of [102]). The arise of the additional coexisting $K = 0$ band is usually attributed to two-particle–two-hole (and, more generally, n -particle– n -hole) excitations across a nuclear closed shell [101,102]. Within the Bohr Hamiltonian approach, the two $K = 0$ coexisting bands are considered to live within two different minima of a sextic potential [104,105]. *However, no explanation has been provided yet for the borders of the regions, within which shape coexistence is observed.*

The main purpose of the present chapter is to show a path towards an explanation of shape coexistence in terms of a mechanism based on the two different sets of magic numbers (shell model, 3D-HO) mentioned above, taking advantage in parallel of the exact SU(3) symmetry present in the 3D-HO shells and the approximate proxy-SU(3) symmetry present in the shell model. In addition to predicting (free of any free parameters) the borders of the regions of coexistence, this method provides also specific predictions for regions of coexistence of a prolate with an oblate band, or of a prolate band with another prolate band of significantly different deformation. The collective model deformation parameters characterizing each band are also predicted by the theory, without involving any free parameters.

1.2 The Origin of QQ Interaction

A nucleus can exhibit collective features only if there is a kind of interaction between distant nucleons. So, there must be a long-range effective potential to unite the whole shell, or even different shells. The most important term of this long-range potential is proportional to the quadrupole interaction.

Suppose there is a kind of two-body central force. The Taylor expansion of it shall be [47]:

$$V^{(2)} = \sum_{i < i'} V\left(\frac{r_{ii'}}{\alpha}\right) = \sum_{i < i'} \left(\xi_0 + \xi_2 \frac{r_{ii'}^2}{\alpha^2} + \xi_4 \frac{r_{ii'}^4}{\alpha^4} + \dots \right), \quad (1.1)$$

where α is a range parameter.

The first term is dominant for a very long-range potential, which for an A-body problem is approximately $\xi_0[A(A-1)/2]$. The second term simply verifies that the average potential of all nucleons can be represented by an isotropic harmonic oscillator potential. The third term is [47]:

$$\sum_{i < i'} r_{ii'}^4 = \sum_{i < i'} \left[r_i^4 + r_{i'}^4 + \frac{8}{3} r_i^2 r_{i'}^2 - 4(r_i^2 + r_{i'}^2) r_i r_{i'} \cos \theta_{ii'} + \frac{4}{3} r_i^2 r_{i'}^2 P_2(\cos \theta_{ii'}) \right]. \quad (1.2)$$

The first three terms in the above expansion do not contribute any splitting in the energy degeneracy problem. Therefore, the only significant term is $\frac{4}{3} r_i^2 r_{i'}^2 P_2(\cos \theta_{ii'})$, where P_2 is the Legendre polynomial. This important term can be written as [47]:

$$r_i^2 r_{i'}^2 P_2(\cos \theta_{ii'}) = (r_i^2 Y_{20}(\theta_i \phi_i)) \cdot (r_{i'}^2 Y_{20}(\theta_{i'} \phi_{i'})), \quad (1.3)$$

where Y_{20} is the spherical harmonic with $l = 2, m = 0$.

But the spherical harmonics are related to the collective quadrupole moment. The collective operators have the following relation with the space variables [47]:

$$r_i^2 Y_{20}(\theta_i \phi_i) = \frac{b^2}{4} \sqrt{\frac{5}{\pi}} q_m^{(c)}, \quad (1.4)$$

where b is the oscillator length parameter with $b = \sqrt{\frac{\hbar}{m\omega}}$. The difference between q and Q is that the first stands for one nucleon, while the capital for the whole nucleus.

So, it happens that [47]:

$$r_i^2 r_{i'}^2 P_2(\cos \theta_{ii'}) = \left(\frac{b^2}{4} \sqrt{\frac{5}{\pi}} \right)^2 q_0^{(c)} q_0^{(c)}. \quad (1.5)$$

This result proves that any central potential contains a long-range effective QQ interaction.

1.3 The Nilsson Model

The Nilsson model is a microscopic nuclear model, in the spirit of shell model, applicable in deformed nuclei. The following will be a review of Eqs. (2)–(7) of [63]. The Hamiltonian for the nuclear coordinate system is:

$$H = H_0 + u_{ls} \hbar \omega_0(\delta) \Lambda \cdot \Sigma + u_{ll} \hbar \omega_0(\delta) \Lambda^2, \quad (1.6)$$

$$H_0 = -\frac{\hbar^2}{2m} \nabla'^2 + \frac{m}{2} (\omega_x^2 x'^2 + \omega_y^2 y'^2 + \omega_z^2 z'^2). \quad (1.7)$$

The frequencies are set to be:

$$\omega_x^2 = \omega_y^2 = \omega_0^2(\delta) \left(1 + \frac{2}{3} \delta \right), \quad (1.8)$$

$$\omega_z^2 = \omega_0^2(\delta) \left(1 - \frac{4}{3} \delta \right). \quad (1.9)$$

The volume conservation restriction leads to the relation:

$$\omega_0(\delta) = \tilde{\omega}_0 \left(1 - \frac{4}{3} \delta^2 - \frac{16}{27} \delta^3 \right)^{-1/6}. \quad (1.10)$$

The deformation parameter δ is connected with the deformation β of the Bohr and Mottelson model as:

$$\delta \approx \frac{3}{2} \sqrt{\frac{5}{4\pi}} \beta \approx 0.95 \beta. \quad (1.11)$$

The dimensionless coordinate system can be created by setting:

$$x = \sqrt{\frac{m\tilde{\omega}_0}{\hbar}}x', \quad y = \sqrt{\frac{m\tilde{\omega}_0}{\hbar}}y', \quad z = \sqrt{\frac{m\tilde{\omega}_0}{\hbar}}z'. \quad (1.12)$$

Then, the spatial part of the Hamiltonian becomes:

$$H_0 = \tilde{H}_0 + H_\delta, \quad (1.13)$$

$$\tilde{H}_0 = \hbar\omega_0(\delta)\frac{1}{2}(-\nabla^2 + r^2), \quad (1.14)$$

$$H_\delta = -\delta\hbar\omega_0(\delta)\frac{4}{3}\sqrt{\frac{\pi}{5}}r^2Y_{20}. \quad (1.15)$$

Therefore, the spatial part of the Hamiltonian can be interpreted as an isotropic 3D harmonic oscillator with frequency $\omega_0(\delta)$ plus a distortion to the isotropic field H_δ .

This distortion contains the term r^2Y_{20} , which (recall Eq. (1.4)) is the $q_0^{(c)}$ component of the quadrupole moment of the nucleon. So:

$$H_\delta = -\delta\hbar\omega_0(\delta)\frac{b^2}{3}q_0^{(c)}. \quad (1.16)$$

This term is the quadrupole interaction of the nucleon with the mean quadrupole field of all the others. It is noticeable that in the Nilsson model, only the q_0 component appears, while the other four ($q_{\pm 2}, q_{\pm 1}$) are neglected.

The non-vanishing matrix elements of H_δ are among the Nilsson states with [63]:

$$\Lambda = \Lambda', \quad \Sigma = \Sigma', \quad l = l', l' \pm 2, \quad \mathcal{N} = \mathcal{N}', \mathcal{N}' \pm 2, \quad (1.17)$$

where Λ, Σ are the projections of the angular momentum and spin, respectively; l is the angular momentum; and $\mathcal{N}$ is the total number of quanta. Obviously, the $q^{(c)}$ interplays with orbitals with the same parity, since they have $\mathcal{N}, \mathcal{N} \pm 2$.

The full Nilsson model Hamiltonian can therefore be written as:

$$H = \tilde{H}_0 + u_{ls}\hbar\omega_0(\delta)\Lambda \cdot \Sigma + u_{ll}\hbar\omega_0(\delta)\Lambda^2 - \delta\hbar\omega_0(\delta)\frac{4}{3}\sqrt{\frac{\pi}{5}}r^2Y_{20}. \quad (1.18)$$

1.3.1 The Nilsson Basis

The asymptotic wave functions of Nilsson orbitals are labeled usually as $K[\mathcal{N}n_z\Lambda]$ [64]. These wave functions occur when the deformation ϵ is large enough. An elegant method for the derivation of the asymptotic wave functions is presented in [64]. In the following, emphasis will be placed on an equivalent notation $|n_z r s \Sigma\rangle$ of the asymptotic wave functions, which highlights the number of quanta in the $x - y$ plane [64].

The Nilsson problem is solved in the cylindrical coordinate system. The quanta on the $x - y$ plane are being created by the $R^\dagger, R, S^\dagger, S$ operators [106] which are exactly the same as the $a_+^\dagger, a_+, a_-^\dagger, a_-$ of Section 1.4:

$$R^\dagger = a_+^\dagger, \quad R = a_+, \quad S^\dagger = a_-^\dagger, \quad S = a_-. \quad (1.19)$$

It is valid that:

$$n_\perp = r + s = n_x + n_y, \quad (1.20)$$

where $R^\dagger R|r\rangle = r|r\rangle$ and $S^\dagger S|s\rangle = s|s\rangle$. The actions of the creation and annihilation operators on the relevant states are [107]:

$$\begin{aligned} R|r\rangle &= \sqrt{r}|r-1\rangle, & R^\dagger|r\rangle &= \sqrt{r+1}|r+1\rangle, & S|s\rangle &= \sqrt{s}|s-1\rangle, \\ S^\dagger|s\rangle &= \sqrt{s+1}|s+1\rangle. \end{aligned} \quad (1.21)$$

In addition, the absolute value of the projection of orbital angular momentum in the Nilsson model is [64]:

$$\Lambda = r - s. \quad (1.22)$$

Lastly, for the projection of total angular momentum K and the projection of spin Σ , the equation [64]:

$$K = \Lambda + \Sigma, \quad (1.23)$$

applies.

With these tools:

$$r + s = \mathcal{N} - n_z, \quad (1.24)$$

$$r - s = \Lambda, \quad (1.25)$$

$$\Sigma = K - \Lambda, \quad (1.26)$$

one can transform the Nilsson orbitals between the two notations:

$$K[\mathcal{N}n_z\Lambda] \rightarrow |n_z r s \Sigma\rangle. \quad (1.27)$$

Some transformations are given in Table 1.1.

TABLE 1.1

Nilsson Orbitals in Two Different Notations

Shell model	$K[\mathcal{N}n_z\Lambda]$	$ n_z r s \Sigma\rangle$
$1g_{9/2}$	$9/2[404]$	$ 040+\rangle$
	$7/2[413]$	$ 130+\rangle$
	$5/2[422]$	$ 220+\rangle$
	$3/2[431]$	$ 310+\rangle$
	$1/2[440]$	$ 400+\rangle$
$2p_{1/2}$	$1/2[301]$	$ 021-\rangle$
$1f_{5/2}$	$5/2[303]$	$ 030-\rangle$
	$3/2[301]$	$ 021+\rangle$
	$1/2[310]$	$ 111+\rangle$
$2p_{3/2}$	$3/2[312]$	$ 120-\rangle$
	$1/2[321]$	$ 210-\rangle$
$1f_{7/2}$	$7/2[303]$	$ 030+\rangle$
	$5/2[312]$	$ 120+\rangle$
	$3/2[321]$	$ 210+\rangle$
	$1/2[330]$	$ 300+\rangle$
$1d_{3/2}$	$3/2[202]$	$ 020-\rangle$
	$1/2[200]$	$ 011+\rangle$
$2s_{1/2}$	$1/2[211]$	$ 110-\rangle$
$1d_{5/2}$	$5/2[202]$	$ 020+\rangle$
	$3/2[211]$	$ 110+\rangle$
	$1/2[220]$	$ 200+\rangle$
$1p_{1/2}$	$1/2[101]$	$ 010-\rangle$
$1p_{3/2}$	$3/2[101]$	$ 010+\rangle$
	$1/2[110]$	$ 100+\rangle$
$1s_{1/2}$	$1/2[000]$	$ 000+\rangle$

1.4 The Elliott SU(3)

The majority of physicists nowadays are familiar with the nuclear shell model, but few have been engaged in the Elliott SU(3) symmetry, which was first introduced in [44–46]. *Briefly this symmetry is the fulfillment of the shell model. While the shell model treats single-particle states, the Elliott SU(3) teaches us, how to couple the valence nucleons towards the derivation of the collective nuclear states.* Moreover, the Elliott SU(3) has a remarkable and unique property: it is a fermionic, collective, nuclear model. As a consequence, the collective states emerge in accordance with the Pauli principle.

The nucleus consists of Z protons and N neutrons. The mass number is $A = Z + N$. All these nucleons create a mean field potential, which participates in the 3D isotropic, harmonic oscillator Hamiltonian:

$$H = \sum_{i=1}^A H_i = \sum_{i=1}^A \left(\frac{p_i^2}{2m} + \frac{1}{2} m \omega_i^2 r_i^2 \right). \quad (1.28)$$

The single-particle eigenstates of:

$$H_i = \frac{p_i^2}{2m} + \frac{1}{2} m \omega_i^2 r_i^2, \quad (1.29)$$

are labeled by the total number of oscillator quanta $\mathcal{N}_i$ in the three spacial dimensions.

The single-particle, isotropic, harmonic oscillator Hamiltonian can be transformed in the second quantization with the use of the quanta annihilation and creation operators in each Cartesian axis $a_j, a_j^\dagger$, with $j = x, y, z$. The Hamiltonian of the i^{th} particle in the x direction is $H_{x,i} = \frac{p_{x,i}^2}{2m} + \frac{1}{2} m \omega_i^2 x_i^2$. The eigenstates are $|n_x\rangle$. The action of the $a_x, a_x^\dagger$ on the eigenstates is [107]:

$$a_x |n_x\rangle = \sqrt{n_x} |n_x - 1\rangle, \quad a_x^\dagger |n_x\rangle = \sqrt{n_x + 1} |n_x + 1\rangle. \quad (1.30)$$

By changing the script x to y or z , one gets the relevant actions in the other two Cartesian directions. In addition, the number operators are:

$$a_j^\dagger a_j |n_j\rangle = n_j |n_j\rangle. \quad (1.31)$$

The above operators being boson (quanta) operators follow the well-known commutation relations:

$$[a_j^\dagger, a_{j'}^\dagger] = [a_j, a_{j'}] = 0, \quad [a_j, a_{j'}^\dagger] = \delta_{jj'}, \quad \text{for } j, j' = x, y, z. \quad (1.32)$$

Thus, the single-particle 3D Hamiltonian can be transformed as [107]:

$$H_i = \frac{p_i^2}{2m} + \frac{1}{2} m \omega^2 r_i^2 = \left(a_x a_x^\dagger + a_y a_y^\dagger + a_z a_z^\dagger + \frac{3}{2} \right) \hbar \omega_i, \quad (1.33)$$

with eigenvalues:

$$E_i = \left(n_x + n_y + n_z + \frac{3}{2} \right) \hbar \omega_i = \left(\mathcal{N}_i + \frac{3}{2} \right) \hbar \omega_i. \quad (1.34)$$

Now the problem can be turned in cylindrical coordinates, which are more suitable for axially deformed nuclei. In order to distinguish the $x - y$ plane from the z -axis, the left and right quanta operators are defined as [108]:

$$a_\pm = \frac{a_x \mp i a_y}{\sqrt{2}}, \quad (1.35)$$

$$a_\pm^\dagger = \frac{a_x^\dagger \pm i a_y^\dagger}{\sqrt{2}}. \quad (1.36)$$

The operators $a_+, a_+^\dagger$ are called right-hand operators, because they seem to destroy/create a circular quantum that rotates in right-hand side in the $x - y$ plane [108]. For the same reason, the other two operators are called left-hand operators. The operators in the z -axis are renamed:

$$a_0 = a_z, \quad a_0^\dagger = a_z^\dagger. \quad (1.37)$$

Now the Hamiltonian is [109]:

$$H_i = \hbar\omega_i \left(a_+^\dagger a_+ + a_-^\dagger a_- + a_0^\dagger a_0 + \frac{3}{2} \right). \quad (1.38)$$

With these ingredients, one can define eight operators which are the generators of SU(3) [109]:

$$l_0 = a_+^\dagger a_+ - a_-^\dagger a_-, \quad (1.39)$$

$$l_\pm = \mp(a_0^\dagger a_\mp - a_\mp^\dagger a_0), \quad (1.40)$$

$$q_{\pm 2} = -\sqrt{6}a_\pm^\dagger a_\mp, \quad (1.41)$$

$$q_{\pm 1} = \mp\sqrt{3}(a_0^\dagger a_\mp + a_\pm^\dagger a_0), \quad (1.42)$$

$$q_0 = 2a_0^\dagger a_0 - a_+^\dagger a_+ - a_-^\dagger a_-. \quad (1.43)$$

The first three operators ($l_0, l_\pm$) form the SU(2) algebra of angular momentum, while the last five are the five components of the single-particle quadrupole operator.

The eight of them close the commutation relations of the SU(3) algebra [109].

$$[l_0, l_\pm] = \pm l_\pm, \quad (1.44)$$

$$[l_+, l_-] = 2l_0, \quad (1.45)$$

$$[l_0, q_m] = m q_m, \quad (1.46)$$

$$[l_\pm, q_m] = \sqrt{6 - m(m \pm 1)} q_{m \pm 1}, \quad (1.47)$$

$$[q_0, q_{\pm 1}] = \pm 3\sqrt{3}l_\pm, \quad (1.48)$$

$$[q_1, q_{-1}] = -3l_0, \quad (1.49)$$

$$[q_2, q_{-2}] = 6l_0, \quad (1.50)$$

$$[q_{\pm 2}, q_{\mp 1}] = \pm 3\sqrt{2}l_\pm, \quad (1.51)$$

$$[q_0, q_{\pm 2}] = [q_{\pm 1}, q_{\pm 2}] = 0. \quad (1.52)$$

The commutator of two operators from the set (1.39)–(1.43) results to one operator from the set (1.39)–(1.43) again. Thus, the eight operators of (1.39)–(1.43) close the SU(3) algebra.

In simple words, the Elliott SU(3) consists of the three angular momentum operators and of the five components of the quadrupole operator. The eight generators of SU(3) are certain linear combinations of the boson operators $a_j^\dagger, a_j$. The 3D isotropic harmonic oscillator Hamiltonian is invariant under U(3) and SU(3) transformations. The $\frac{(\mathcal{N}+1)(\mathcal{N}+2)}{2}$ orbitals of a harmonic oscillator shell with $\mathcal{N}$ total number of quanta have SU(3) symmetry.

1.4.1 Derivation of the Highest Weight irrep

In the Elliott SU(3) model, the calculation of every observable is based on the Elliott quantum numbers (λ, μ) [44]. The highest weight (h.w.) irrep describes the ground-state properties of the nucleus. The method for the calculation of the h.w. (λ, μ) has been presented in steps at [110,111].

For the h.w. irreps, a simple method will be presented here.

1. The Cartesian states $|n_z, n_x, n_y\rangle$ are vectors of a $U\left(\frac{(\mathcal{N}+1)(\mathcal{N}+2)}{2}\right)$ algebra. The degenerate 3D isotropic harmonic oscillator orbitals are ordered as:

$$|n_z, n_x, n_y\rangle = |n, 0, 0\rangle, \quad |n-1, 1, 0\rangle, \quad |n-1, 0, 1\rangle, \quad \dots, \quad |0, 0, n\rangle. \quad (1.53)$$

For total quanta $\mathcal{N} = 2$, there are $\frac{(\mathcal{N}+1)(\mathcal{N}+2)}{2} = 6$ orbitals. One has to order these six orbitals as follows:

$$\begin{aligned} |1\rangle &= |2, 0, 0\rangle, & |2\rangle &= |1, 1, 0\rangle, & |3\rangle &= |1, 0, 1\rangle, & |4\rangle &= |0, 2, 0\rangle, & |5\rangle &= |0, 1, 1\rangle, \\ |6\rangle &= |0, 0, 2\rangle. \end{aligned} \quad (1.54)$$

2. Then, at most two protons or neutrons are placed in each orbital $|n_z, n_x, n_y\rangle$ with the order of Eq. (1.54). For instance, ${}^{31}_{12}\text{Mg}_{19}$ has $19 - 8 = 11$ neutrons in the 8–20 shell. Two neutrons are placed in each of the $|1\rangle, \dots, |5\rangle$ orbitals, while the last neutron occupies the $|6\rangle$.
3. Afterwards, the summations of the quanta in each Cartesian axis are calculated:

$$\sum_{i=1}^{Z_{val}, N_{val}} n_{zi}, \quad \sum_{i=1}^{Z_{val}, N_{val}} n_{xi}, \quad \sum_{i=1}^{Z_{val}, N_{val}} n_{yi}. \quad (1.55)$$

For the 11 neutrons in $U(6)$, the above expressions give 8, 8, 6, respectively.

4. The results of the summations are placed in decreasing order, and these are the quantum numbers $[f_1, f_2, f_3]$ of the U(3) algebra. For our example, $[f_1, f_2, f_3] = [8, 8, 6]$.
5. The λ and μ are:

$$\lambda = f_1 - f_2, \quad \mu = f_2 - f_3, \quad (1.56)$$

which leads to $(\lambda, \mu) = (0, 2)$ for 11 particles in U(6).

1.4.2 The Collective Operators

The single-particle operators can be extended to the many-particle ones, by summing over all valence particles [109]:

$$L = \sum_i^{A_{val}} l_i, \quad (1.57)$$

$$Q_m = \sum_i^{A_{val}} q_{im}. \quad (1.58)$$

The former one is the angular momentum, while the latter is the m component of the algebraic quadrupole moment for the whole shell. These many-particle operators satisfy the same commutation relations with their relevant single ones, and therefore, they close an SU(3) algebra.

The simplest SU(3) Hamiltonian of the many nucleon problem, which is suitable for a collective 0^+ state, is [112]:

$$H = \sum_i^A H_i - \frac{1}{2} \chi Q \cdot Q = H_0 - \frac{1}{2} \chi Q \cdot Q, \quad (1.59)$$

where χ is the strength of the algebraic quadrupole-quadrupole (QQ) interaction and if i and i' are two distinct valence nucleons:

$$QQ = \sum_{i,i'}^{A_{val}} \sum_{m=-2}^2 (-1)^m q_{mi} q_{-mi'} = \sum_{m=-2}^2 (-1)^m Q_m Q_{-m}. \quad (1.60)$$

The SU(3) algebra has a second- and a third-order Casimir operator [113]:

$$C_2 = \lambda^2 + \mu^2 + \lambda\mu + 3(\lambda + \mu), \quad (1.61)$$

$$C_3 = \frac{2}{9}(\lambda^3 - \mu^3) + \frac{1}{3}\lambda\mu(\lambda - \mu) + (\lambda + 2)(2\lambda + \mu). \quad (1.62)$$

The Casimir operators commute with all the generators of the algebra. The QQ interaction is measured via the C_2 as [113]:

$$QQ = 4C_2 - 3L^2. \quad (1.63)$$

The 3D isotropic harmonic oscillator Hamiltonian H_0 has dimensionless eigenvalues (i.e. $\hbar\omega_i = 1$):

$$N_0 = \sum_{i=1}^A \left(\mathcal{N}_i + \frac{3}{2} \right) \quad (1.64)$$

The strength χ is measured by the equation [114]:

$$\frac{\chi}{2} = \frac{\hbar\omega}{4N_0}, \quad (1.65)$$

where $\hbar\omega$ is measured in MeV.

The deformation variables β, γ follow the formulas [115]:

$$\beta^2 = \frac{4\pi}{5(A\bar{r}^2)^2} (C_2 + 3), \quad (1.66)$$

$$\gamma = \tan^{-1} \frac{\sqrt{3}(\mu + 1)}{2\lambda + \mu + 3}, \quad (1.67)$$

where A is the mass number and the dimensionless mean square radius is $\bar{r}^2 = 0.87^2 A^{1/3}$ [116]. The β^2 may be multiplied by the scaling factor $(A/S)^2$, where S is the size of the proton and neutron valence shell [61]:

$$\beta^2 = \frac{4\pi}{5(S\bar{r}^2)^2} (C_2 + 3). \quad (1.68)$$

1.4.3 The SU(3) $\rightarrow$ SU(2) $\times$ U(1) Decomposition

Most nuclei have quadrupole deformation. Such nuclei look like an ellipsoid with cylindrical symmetry. The cylinder is divided into the $x - y$ plane and the symmetry z -axis. If the ellipsoid has elongation across one axis (z), the shape is prolate, while if the elongation is across two axes (x, y), it is oblate.

The algebra $SU(3)$ is decomposed into two subalgebras $SU(2)$ and $U(1)$ [45]. The former represents rotations in the $x - y$ plane, while the latter involves the z -axis. The $U(1)$ is characterized by the operator q_0 , which has eigenvalues $q_0 = 2n_z - n_\perp$, with $n = n_z + n_\perp$.

The $SU(2)$ algebra has the following generators [45]:

$$u_\pm = \mp \left(\frac{1}{2\sqrt{3}} \right) q_{\pm 2}, \quad u_0 = \frac{1}{2} l_0, \quad (1.69)$$

which satisfy the commutation relations:

$$[u_\pm, u_0] = \mp u_\pm, \quad [u_+, u_-] = -u_0. \quad (1.70)$$

The $SU(2)$ multiplets are characterized by a quantum number of j -type and one of m_j -type [113]. In the present algebraic structure, the m_j -type quantum number is the eigenvalue of u_0 and the j -type, let it be named Λ_E (Λ of Elliott), is

$$\Lambda_E = \frac{1}{2}(n_x + n_y) = \frac{1}{2}n_\perp. \quad (1.71)$$

Since:

$$u_0 = \frac{l_0}{2}, \quad (1.72)$$

and the possible values of u_0 are $-\Lambda_E, -\Lambda_E + 1, \dots, \Lambda_E$ obeying the angular momentum algebra, obviously:

$$l_0 = -2\Lambda_E, -2\Lambda_E + 2, \dots, 2\Lambda_E. \quad (1.73)$$

The basis of the $x - y$ plane is labeled as $|\Lambda_E, u_0\rangle$. These vectors with the use of (1.71) and (1.72) are more conveniently named $|n_\perp, l_0\rangle$.

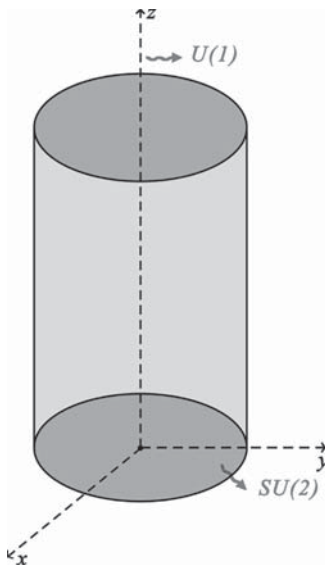
As a consequence, the $SU(2)$ subalgebra of the Elliott $SU(3)$ depends solely on the quanta on the $x - y$ plane. The physical quantity of the $SU(2)$ is the projection of the orbital angular momentum (l_0), while the $q_{\pm 2}$ are useful as ladder operators. The number of quanta in the z -axis is isolated in the $U(1)$ subalgebra (see Figure 1.1). This decomposition of the Elliott $SU(3)$ suits perfectly in nuclei with quadrupole deformation.

1.5 Proxy- $SU(3)$ Symmetry

A nucleus consists of Z protons and N neutrons. The nucleons interact mainly via the strong nuclear force. Since the exact formula for this force is not yet known, one can approach the result of all the inter-nucleon interactions as in Section 1.2. The main result of all the inter-nucleon interactions is the creation of a mean field potential, which is described by a 3D isotropic harmonic oscillator potential. This mean field creates nuclear shells, which consist of orbitals with a common number of oscillator quanta $\mathcal{N}$. The magic numbers of these shells are the 3D harmonic oscillator (HO) magic numbers: 2, 8, 20, 40, 70, 112, ...

Unfortunately the mean field is not the only interaction. There exists a single-particle spin-orbit (SO) term ls , which destroys the harmonic oscillator shells. The ls interaction forces the orbitals to move in energy. As an example, the 20–40 harmonic oscillator shell becomes 28–50. The mechanism of this procedure is the following:

- (a) The harmonic oscillator shell 20–40 consists of the shell model orbitals $1f_{7/2}, 2p_{3/2}, 1f_{5/2}, 2p_{1/2}$. The orbital $1g_{9/2}$ is part of the harmonic oscillator shell 40–70. The orbitals are listed in Table 1.2.


FIGURE 1.1

Cylindrical symmetry. The $SU(3)$ algebra is decomposed into an $SU(2)$, which describes the $x - y$ plane, and a $U(1)$, which refers to the z -axis. This decomposition suits perfectly to a prolate or oblate nuclear shape.

TABLE 1.2

Nilsson Orbitals for the Nucleons between 20 and 50

Shell model	$K[\mathcal{N}n_z\Lambda]$	$ n_z r s \Sigma\rangle$
$1g_{9/2}$	$9/2[404]$	$ 040+\rangle$
	$7/2[413]$	$ 130+\rangle$
	$5/2[422]$	$ 220+\rangle$
	$3/2[431]$	$ 310+\rangle$
	$1/2[440]$	$ 400+\rangle$
	$1/2[301]$	$ 021-\rangle$
$2p_{1/2}$	$5/2[303]$	$ 030-\rangle$
$1f_{5/2}$	$3/2[301]$	$ 021+\rangle$
	$1/2[310]$	$ 111+\rangle$
	$3/2[312]$	$ 120-\rangle$
$2p_{3/2}$	$1/2[321]$	$ 210-\rangle$
	$7/2[303]$	$ 030+\rangle$
$1f_{7/2}$	$5/2[312]$	$ 120+\rangle$
	$3/2[321]$	$ 210+\rangle$
	$1/2[330]$	$ 300+\rangle$

Due to the ls interaction, the $1f_{7/2}$ orbital is excluded from the shell under discussion, while the $1g_{9/2}$ orbital is included. Thus, the harmonic oscillator shell 20–40 becomes the nuclear shell 28–50. The 28–50 shell consists of some orbitals with $\mathcal{N} = 3$ and some orbitals with $\mathcal{N} = 4$ quanta [14]. Therefore, the Elliott $SU(3)$ symmetry is not valid in the 28–50 shell.

- (b) The ls interaction lowers the energy of the $1f_{7/2}$ orbital so much, that this orbital is no longer included in the shell under discussion.
- (c) Instead, the $1g_{9/2}$ from the 40–70 HO shell is lowered and becomes part of the shell under discussion.

To resume, the harmonic oscillator shell 20–40 loses the $1f_{7/2}$ orbital and gains instead the $1g_{9/2}$ orbital. This new shell lies between magic numbers 28 and 50.

Thus after, the ls interaction is applied, each harmonic oscillator shell with $\mathcal{N}$ quanta loses the orbital with total angular momentum $j = \mathcal{N} + \frac{1}{2}$ and gains the orbital with $\mathcal{N} + 1$ quanta and $j' = \mathcal{N} + 1 + \frac{1}{2}$.

Proxy-SU(3) was introduced in [60] in order to describe a nuclear shell which consists of some orbitals with $\mathcal{N}$ quanta and some orbitals with $\mathcal{N} + 1$ quanta. Nuclear shells above 28 protons or neutrons are of this type. Since such shells do not match with an exact SU(3) symmetry, an approximation is necessary in order to make the algebraic tools applicable. The proxy-SU(3) replacement of orbitals is justified in Table 1.3 and described in Table 1.4 for the 28–50 shell.

The proxy-SU(3) approximation affects only the orbitals with $\mathcal{N} + 1$ quanta (intruder orbitals), by erasing one quantum from the z -axis of each intruder orbital. The normal parity orbitals remain intact.

1.5.1 The Exact Symmetry Behind Proxy-SU(3)

Proxy-SU(3) has been very successful in predicting for various isotopic chains, without any parameters, the β and γ deformations of each nucleus, as well as the prolate over oblate dominance along with the prolate-oblate transition [61]. This success relies on the h.w. irrep and on the fact that there is an exact symmetry behind proxy-SU(3).

TABLE 1.3

The $1f_{7/2}$ Orbital has been Excluded from the Harmonic Oscillator Shell 20–40 due to the ls Interaction, while the $1g_{9/2}$ Orbital has Intruded in It

Shell Model	$ n_z r s \Sigma\rangle$	Shell Model	$ n_z r s \Sigma\rangle$
$1g_{9/2}$	$ 040+\rangle$	$1f_{7/2}$	X
	$ 130+\rangle$		$ 030+\rangle$
	$ 220+\rangle$		$ 120+\rangle$
	$ 310+\rangle$		$ 210+\rangle$
	$ 400+\rangle$		$ 300+\rangle$

The $1f_{7/2}$ and $1g_{9/2}$ orbitals consist of Nilsson orbitals with $\mathcal{N} = 3$ and $\mathcal{N} = 4$, respectively, as seen in Table 1.2. These orbitals are listed here using the $|n_z r s \Sigma\rangle$ notation. All of them have a common projection of spin Σ . Each of the orbitals of $1f_{7/2}$ has identical distribution of quanta in the $x-y$ plane ($|rs\rangle$) with an orbital of $1g_{9/2}$. Such orbitals appear in the same line in the table and are depicted in the same color. The orbitals with the same color differ by one quantum in the z -axis. Therefore, the operators $q_{\pm 2}, l_0$ of the SU(2) subalgebra of SU(3) (Section 1.4.3) give exactly the same eigenvalues when acting on the colored orbitals. As a consequence, the SU(2) operators cannot discriminate among the two sets of the colored orbitals [117].

TABLE 1.4

The Replacement of the Orbitals in Proxy-SU(3)

Shell model	Before		After	
	$K[\mathcal{N}n_z\Lambda]$	$ n_zrs\Sigma\rangle$	$K[\mathcal{N}n_z\Lambda]$	$ n_zrs\Sigma\rangle$
$1g_{9/2}$	9/2[404]	040+	X	X
	7/2[413]	130+	7/2[303]	030+
	5/2[422]	220+	5/2[312]	120+
	3/2[431]	310+	3/2[321]	210+
	1/2[440]	400+	1/2[330]	300+
$2p_{1/2}$	1/2[301]	021-		
$1f_{5/2}$	5/2[303]	030-		
	3/2[301]	021-		
	1/2[310]	111+		
$2p_{3/2}$	3/2[312]	120-		
	1/2[321]	210-		

One quantum from the z -axis is erased from each intruder orbital, while the n_x, n_y (or, equivalently, the r, s) of the intruders and the normal parity orbitals remain intact. The intruder with the highest projection of total angular momentum (K in Nilsson notation) is excluded. The new shell (after the approximation) is a complete harmonic oscillator shell; thus, it has an SU(3) symmetry.

The total wave function of a single-particle state is the product of a spatial and a spinor part. The spatial wave function is characterized by the number of quanta in the z -axis ($|n_z\rangle$) and in the $x - y$ plane ($|rs\rangle$). In a full harmonic oscillator shell with $\mathcal{N}$ quanta, the n_z gets values:

$$n_z = \mathcal{N}, \mathcal{N} - 1, \mathcal{N} - 2, \dots, 0. \quad (1.74)$$

The quanta in the $x - y$ plane are, respectively:

$$r + s = 0, 1, 2, \dots, \mathcal{N}, \quad (1.75)$$

with the restriction [64]:

$$r \geq s. \quad (1.76)$$

Thus, the following combinations are valid:

$$\begin{aligned} & (n_z = \mathcal{N}, \quad r = s = 0), \\ & (n_z = \mathcal{N} - 1, \quad r = 1, \quad s = 0), \\ & (n_z = \mathcal{N} - 2, \quad r = 2, \quad s = 0), \\ & (n_z = \mathcal{N} - 2, \quad r = 1, \quad s = 1), \\ & \dots \end{aligned}$$

As an example, the orbitals of the “Before” column of Table 1.4, excluding the 9/2[404] orbital, contain all the r, s values of a full harmonic oscillator shell with $\mathcal{N} = 3$ quanta.

Consequently, a nuclear shell, such as 28–50, 50–82, and 82–126, after excluding the Nilsson orbital with the maximum value of K , has the SU(2) subalgebra of the Elliott SU(3). The proxy-SU(3) approximation is not affecting the number of quanta in the $x - y$ plane. Thus, the orbitals of the proxy-SU(3) shell preserve the exact SU(2) symmetry.

1.6 Magic Numbers Below 28

In Section 1.5, it has been described how the spin-orbit interaction creates the nuclear shell 28–50. By the same procedure, the shells 50–82 and 82–126 are created. But a significant spin-orbit term in the Hamiltonian can create magic numbers below 28 particles, as described in Table 1.5, which we are going to call, following Ref. [8], SO-like magic numbers. The SO-like magic numbers, that will arise, are:

$$\text{SO-like magic numbers below 28: } 2, 6, 14. \quad (1.77)$$

Actually, the SO-like magic number 14 is already well established experimentally [8].

Clearly, the above SO-like magic numbers have not emerged in this work through the study of energy gaps, but through symmetry considerations. The sets of orbitals among the mentioned SO-like magic numbers, after excluding the Nilsson orbital with maximum K , preserve the $SU(2)$ symmetry of Elliott.

As seen in Table 1.5, the SO-like shells 6–14 and 14–28 have a proxy-SU(3) symmetry. These shells can be treated as in [60,61], to calculate the deformation variables (β, γ) . The shell 14–28 is useful for the neutrons of the Mg isotopes, while the 6–14 can be applied in the neutrons of the Be isotopes, to predict the parity inversion at ^{11}Be .

1.7 Magic Numbers $\rightarrow$ Shape Coexistence $\rightarrow$ Inversion of States

The Mg isotopic chain is the best example that demonstrates how the SO-like magic numbers below 28 may lead to right predictions about deformation β , shape coexistence and inversion

TABLE 1.5

A Significant Spin-Orbit Interaction may Create the SO-Like Shells 2–6 and 6–14

Shell model	Before $K[\mathcal{N}n_z\Lambda]$	After $K[\mathcal{N}n_z\Lambda]$	SO-like magic number
$1f_{7/2}$	7/2[303]	X	28
	5/2[312]	5/2[202]	
	3/2[321]	3/2[211]	
	1/2[330]	1/2[220]	
$1d_{3/2}$	3/2[202]	3/2[202]	14
	1/2[200]	1/2[200]	
$2s_{1/2}$	1/2[211]	1/2[211]	
$1d_{5/2}$	5/2[202]	X	14
	3/2[211]	3/2[101]	
	1/2[220]	1/2[110]	
$1p_{1/2}$	1/2[101]	1/2[101]	
$1p_{3/2}$	3/2[101]	X	6
	1/2[110]	1/2[000]	
$1s_{1/2}$	1/2[000]	1/2[000]	2

Proxy-SU(3) can be applied in these shells. The column “Before” refers to the original nuclear shell and the column “After” to the proxy-SU(3) shell [117].

of states. A bulk of experimental work [118–123] has established that the Mg isotopic chain exhibits both shape coexistence and inversion of states. A shell model mechanism that uses particle-hole excitations is the state-of-the-art theoretical approach for these phenomena in Mg [124–126].

Detailed calculations will be presented for $^{26-40}_{12}\text{Mg}_{14-28}$ within the proxy-SU(3) framework. In the following, a mechanism, which was initially proposed in [127], will be used to explain why:

- (a) $N = 20$ is not a magic number for the ground state of $^{32}_{12}\text{Mg}_{20}$,
- (b) shape coexistence occurs among neutron numbers $N = 18 - 20$ for this isotopic chain, and
- (c) inversion of states appears at $^{31}_{12}\text{Mg}_{19}$.

The isotonic chains with $N \approx 20$ possess values of the deformation β which exhibit a minimum at $Z = 14$ [128]. This provides experimental support for the findings of the previous section, in which 14 was found to be an SO-like magic number, for protons in the present case. The 12 protons of Mg lie in the 6–14 shell, and in proxy-SU(3) have $(\lambda, \mu)_p = (0, 0)$.

The neutrons of $^{26-40}_{12}\text{Mg}_{14-28}$ lie either in the SO-like shell 14–28 or in the harmonic oscillator shells 8–20 and 20–40. The h.w. SU(3) irreps for these shells are listed in Table 1.6. Since protons do not contribute at all, the deformation formula becomes:

$$\beta^2 = \frac{4\pi}{5(S_n \bar{r}^2)^2} (C_2 + 3), \quad (1.78)$$

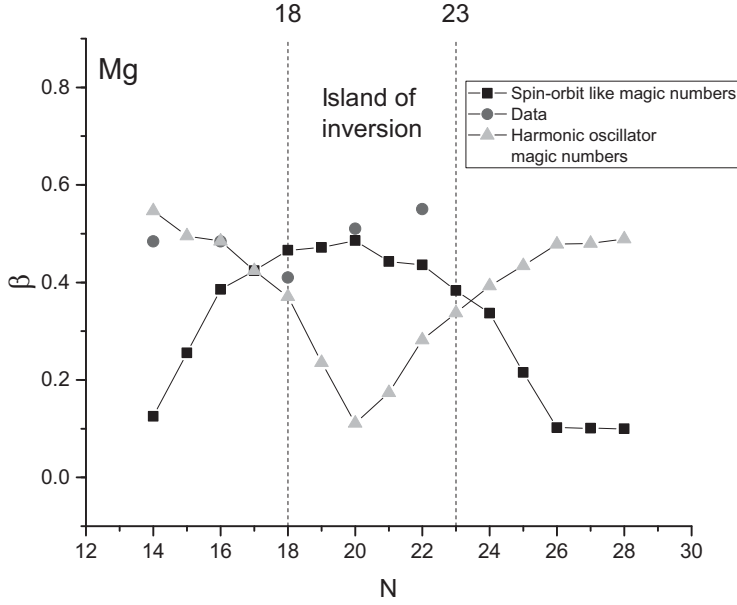
where $\bar{r}^2 = 0.87^2 N^{1/3}$ [116], $C_2 = \lambda_n^2 + \mu_n^2 + \lambda_n \mu_n + 3(\lambda_n + \mu_n)$ (where the subscript n stands for neutrons) and $S_n = 12$ for the 8–20 and the 14–28 proxy-SU(3) shell, while $S_n = 20$ for the 20–40 shell. The theoretical predictions for β for the Mg isotopic chain for the spin-orbit like magic numbers and the harmonic oscillator magic numbers are presented in Figure 1.2.

Shape coexistence is a phenomenon where the nucleus exhibits a ground-state band with a certain type of deformation (prolate for instance) and a slightly excited band with another type of deformation (oblate for example) [101–103]. The idea proposed in [117,127] is that shape coexistence is the consequence of two sets of magic numbers, namely the SO-like 6, 14, 28, 50, 82, 126 magic numbers and the harmonic oscillator 2, 8, 20, 40, 70, 112 magic numbers:

TABLE 1.6

The h.w. SU(3) Irreps [61] for 13 to 28 Nucleons, Listed in the Case in Which the SO-Like Magic Numbers 14 and 28 and the Proxy-SU(3) Symmetry Are Taken into Account, as well as in the Case in Which the Harmonic Oscillator Magic Numbers 8, 20, and 40 and the HO SU(3) Symmetry Are Considered

Even			Odd		
Particle Number	$(\lambda, \mu)_{SO}$	$(\lambda, \mu)_{ho}$	Particle Number	$(\lambda, \mu)_{SO}$	$(\lambda, \mu)_{ho}$
14	(0, 0)	(6, 0)	13	(0, 0)	(5, 1)
16	(4, 0)	(2, 4)	15	(2, 0)	(4, 2)
18	(4, 2)	(0, 4)	17	(4, 1)	(1, 4)
20	(6, 0)	(0, 0)	19	(5, 1)	(0, 2)
22	(2, 4)	(6, 0)	21	(4, 2)	(3, 0)
24	(0, 4)	(8, 2)	23	(1, 4)	(7, 1)
26	(0, 0)	(12, 0)	25	(0, 2)	(10, 1)
28	(0, 0)	(10, 4)	27	(0, 0)	(11, 2)

**FIGURE 1.2**

Shape coexistence begins when the deformation obtained with the harmonic oscillator magic numbers is less than the deformation calculated with the SO-like magic numbers, *i.e.* $\beta_{ho} \leq \beta_{SO}$ and usually stops at a harmonic oscillator magic number. The mechanism of [111,117, 127] predicts that the nuclei $^{30-32}_{12}\text{Mg}_{18-20}$ are candidates for shape coexistence. Data for β have been taken from [129]. The data indicate that the neutrons of the Mg isotopes with neutron number $N \leq 17$ follow the harmonic oscillator magic numbers and do not exhibit shape coexistence. As a consequence, when $N \leq 17$, there is a unique ground state obeying the rules of the 8–20 HO shell. But as soon as shape coexistence begins, $N > 17$, there may be two states low lying in energy. As shown in Eq. (1.82), in general the state with the higher deformation is the ground state. As seen from the comparison to the data [128], the ground state is derived from the SO-like shell 14–28. The excited state is derived in general from the coupling of the 8–20 HO shell with the 14–28 SO-like shell. To summarize, when $N > 17$ the ground state follows the 14–28 SO-like shell, while when $N < 17$ the ground state follows the 8–20 HO shell. This may be the cause of the inversion of states, which begins at $^{30}_{12}\text{Mg}_{18}$ [118].

When the nuclear deformation with the harmonic oscillator magic numbers becomes less than the deformation with the SO-like magic numbers:

$$\beta_{ho} \leq \beta_{SO}, \quad (1.79)$$

then shape coexistence is likely to appear.

The ground state of such nuclei derives from the SO-like magic numbers, because this shell possesses the maximum deformation and QQ interaction, as we can see using the Hamiltonian of Eq. (1.59). Indeed, schematically one has

$$H_{ho} = H_0 - \frac{\chi}{2} QQ_{ho}, \quad (1.80)$$

$$H_{SO} = H_0 - \frac{\chi}{2} QQ_{SO}, \quad (1.81)$$

$$QQ_{ho} \leq QQ_{SO} \Rightarrow H_{SO} \leq H_{ho}. \quad (1.82)$$

The excited $K = 0^+$ state of nuclei with shape coexistence derives from the coupling of the harmonic oscillator shell 8–20 with the SO-like shell 14–28. In the review articles [101,102], one can see that the experimentally known examples of shape coexistence in several series of isotopes appear below and stop at a harmonic oscillator magic number. For example, in the Hg series of isotopes, shape coexistence appears below and stops at $N = 112$, as seen in Figure 10 of Ref. [102]). In the present case, this means that shape coexistence should appear below and stop at $^{32}_{12}\text{Mg}_{20}$.

The inversion of states is a side effect of shape coexistence. One can divide the Mg isotopic chain in two halves:

- (a) the first half contains the isotopes with neutron numbers $N \leq 17$ with $\beta_{ho} \geq \beta_{SO}$,
and
- (b) the second contains the isotopes with $17 < N < 23$ with $\beta_{ho} < \beta_{SO}$.

The neutrons of the isotopes of the first half clearly follow the harmonic oscillator magic numbers 8–20 (see Figure 1.2). But the isotopes of the second half may exhibit shape coexistence. Such isotopes have two low-lying states: a ground-state and a slightly excited one.

While the ground state of Mg isotopes with $N \leq 17$ followed the 8–20 HO shell, now the ground state of isotopes with $N > 17$ follows the rules of the 14–28 SO-like shell. The inversion of states may be caused by a change of magic numbers.

The cause of the inversion of states may be hidden in the collective nuclear features instead of the single-particle ones. The Elliott SU(3) model contains all the necessary techniques, which can reveal the J^π of the ground state and of the excited state of even-odd nuclei, which lie within the islands of inversion.

1.8 Conclusions

A mechanism is proposed for shape coexistence and inversion of states [111,117,127]. The mechanism involves two sets of magic numbers: the spin-orbit (SO) like magic numbers 6, 14, 28, 50, 82, 126 and the harmonic oscillator magic numbers 2, 8, 20, 40, 70, 112. Shape coexistence arises when $\beta_{SO} \geq \beta_{ho}$ and stops at a harmonic oscillator closure. The state with the maximum deformation lies lower in energy. As a consequence, within the islands of shape coexistence shown in Figure 8 of Ref. [102], the ground state is derived by the SO-like magic numbers, while the coexisting excited 0^+ state is derived from the coupling of the harmonic oscillator with the SO-like magic numbers. The inversion of states is a side effect of shape coexistence. The proposed mechanism predicts without any parameters the borders of the islands of shape coexistence appearing at Figure 8 of [102]. This mechanism can be applied in every mass region to predict the borders of the islands of shape coexistence and all the relevant nuclear observables.

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