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Shape coexistence in even-even Sn nuclei with mass number $A=110-120$

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Introduction

The atomic nucleus is a many-body quantum system composed of interacting fermions—protons and neutrons. Its structure and dynamics are governed by the nuclear force, which gives rise to a wide variety of phenomena across the nuclear chart. Among these, the emergence of coexisting shapes is one of the most intriguing. Deformed configurations can appear at low excitation energies as a consequence of nucleon-nucleon correlations, characterized by different spatial arrangements of nucleons, induced by specific terms of the nuclear interaction. This phenomenon, known as *shape coexistence*, is now recognized as a widespread feature of nuclei.

In this thesis, shape coexistence is investigated in the $Z = 50$ region of the nuclear chart. The Sn isotopic chain, which contains the largest number of stable isotopes, offers an ideal testing ground for systematic studies with intense stable beams. Evidence for shape coexistence in stable Sn isotopes has existed for more than fifty years, based on level-energy systematics and two-proton transfer cross sections. Further insights have been mostly obtained from studies of low-lying excited 0^+ states populated via Coulomb excitation. However, configuration mixing between these states gives rise to a complex scenario that remains poorly understood. Therefore, additional spectroscopic information on excited 0^+ states, together with state-of-the-art shell-model and mean-field calculations, is essential to shed light on the onset of shape coexistence in this region.

To address this issue, a series of experiments was conducted between 2021 and 2024 at IFIN-HH (Bucharest, Romania) and LNL (Legnaro, Italy), employing γ -ray spectroscopy and lifetime measurement techniques. At IFIN-HH, four experiments (denoted *Exp. 1-4*) were carried out with the ROSPHERE-SORCERER setup. Their primary aim was to study the population of low-lying 0^+ states in even-even $^{114-120}\text{Sn}$ isotopes via single- and multi-nucleon transfer reactions induced by ^{13}C and $^{16,18}\text{O}$ beams. The most favorable channels, namely one- and two-neutron transfer reactions, were then exploited for lifetime measurements using the Doppler Shift Attenuation and the Recoil Distance Doppler Shift method, thus allowing for the derivation of $B(E2)$ transition probabilities (or a limit) for several 0^+ states up to 0_4^+ , in even-even Sn isotopes with mass number $A=110-120$.

At LNL, a complementary experiment was performed using the $^{32}\text{S} + ^{110}\text{Cd}$ multi-nucleon transfer reaction to populate excited states in Cd and Sn isotopes near $A = 110$. Spectroscopic studies with thick and plunger targets were carried out. γ rays and light recoils were detected with the AGATA array and the PRISMA spectrometer, which also enabled inclusive cross-section measurements.

This thesis is organized into four main parts, comprising ten chapters in total.

- **Part I** introduces the physical context and experimental methods. Chapter 1 outlines the formalism of nuclear deformation and collective excitations, followed by an overview of global manifestations of shape coexistence within different theoretical frameworks. A state-of-the-art review of coexisting structures in Sn isotopes is also provided. Chapter 2 then describes the experimental methods employed, including heavy-ion transfer reactions, γ -ray spectroscopy, and lifetime measurements based on Doppler techniques.
- **Part II** presents the IFIN-HH experiments. Chapter 3 introduces the Tandem facility and the ROSPHERE-SORCERER setup, along with detailed accounts of the four experimental campaigns. Chapter 4 covers preliminary data treatment, including detector calibration, alignment, and the new implementation of the ROSPHERE simulation framework. Chapter 5 discusses the spectroscopic and lifetime analyses performed using the Doppler Shift Attenuation and Recoil Distance Doppler Shift methods.
- **Part III** reports on the AGATA-PRISMA experiment. Chapter 6 describes the Tandem accelerator and AGATA-PRISMA detection system, as well as the characteristics of the $^{32}\text{S} + ^{110}\text{Cd}$ reaction and the target configurations employed. Chapter 7 focuses on the preprocessing of PRISMA and AGATA data, including trajectory reconstruction and signal processing, and concludes with an evaluation of AGATA's performance in terms of linearity, resolution, and efficiency. Chapter 8 presents the cross-section and spectroscopic results and discusses the feasibility of lifetime measurements using the collected dataset.
- **Part IV** concludes the thesis. Chapter 9 summarizes the results of the entire experimental program, presenting the extracted $B(E2)$ and $\rho^2(E0)$ transition probabilities and comparing them with existing experimental data and theoretical predictions.

Part I

Context and Methods

Nuclear shape coexistence

In this Chapter, the relevant concepts for the description of the shape coexistence phenomenon will be introduced. In Section 1.1, the formalism of nuclear deformation and collective excitation of deformed nuclei will be presented. An overview of the shape coexistence phenomenon will be discussed in Section 1.2, covering its theoretical description, the experimental signature and selected evidences across the nuclear chart. Lastly, experimental findings and theoretical predictions in the Sn isotopic chain, subject of this work, will be treated in section 1.3.

1.1 Deformation and collective excitations

The concept of nuclear shape developed from the need to encompass collective properties of nuclei [1]. In an intrinsic reference frame, the shape of a nuclear system is defined by its mass and charge distributions, and can be expressed through a multipole expansion of the nuclear surface radius, providing a convenient parametrization of possible deformations. Although shape cannot be accessed directly in experimental studies conducted in the laboratory frame, it has proven to be a powerful concept, offering a unified description of the nucleus that connects the microscopic single-particle motion with macroscopic collective dynamics [2].

In this Section, an introduction to the formalism of nuclear deformation is given, together with a discussion of the collective excitations built upon deformed configurations.

1.1.1 Nuclear surface parametrization

In a three-dimensional space, any closed surface can be parametrized as an infinite series of spherical harmonics. Following the Bohr and Wheeler description of a nucleus as a liquid drop with homogeneous density [3], its radius R can be written as

$$R(\theta, \phi) = R_0 \left[1 + \sum_{\lambda\mu} \alpha_{\lambda\mu} Y_{\lambda\mu}(\theta, \phi) \right], \quad (1.1)$$

where $R_0 = 1.2A^{1/3}$ is the radius of a spherical nucleus of mass number A with the same uniform density, $\alpha_{\lambda,\mu}$ are the (complex) expansion coefficients, and $Y_{\lambda\mu}$ are the spherical harmonics of degree λ and order μ , respectively [4]. Since the term $\lambda = 0$ corresponds to a change of the nuclear volume and the term $\lambda = 1$ describes (for small deformations) the translation of the whole system, the expansion is restricted to $\lambda \geq 2$. The most important contribution comes from the quadrupole term ($\lambda = 2$) of the multipole expansion, which gives rise to ellipsoidal shapes. Among the five $\alpha_{2\mu}$ coefficients that characterize

the quadrupole term, only three determine the orientation of the liquid-drop nucleus in the space-fixed frame, corresponding to three Euler angles. In the body-fixed system, the five coefficients reduce to only two real and independent variables α_{20} and $a_{22} = a_{2-2}$, which, together with the three Euler angles, provide a complete description of the system. It is convenient to change notation and write α_{20} and a_{22} in terms of two variables β and γ through the expressions

$$\begin{aligned} a_{20} &= \beta_2 \cos \gamma, \\ a_{22} &= \frac{1}{\sqrt{2}} \beta_2 \sin \gamma. \end{aligned} \quad (1.2)$$

The β ($\beta \geq 0$) parameter defines the magnitude of the deformation, while γ quantifies deviations from axial symmetry. To ensure a unique correlation between parameters and quadrupole deformation, the parameter γ is restricted to the $0^\circ < \gamma < 60^\circ$ range. Similarly, the radii along the principal axes $k = 1, 2, 3$ of the body-fixed frame can be written as

$$R_k = R_0 + \sqrt{\frac{5}{4\pi}} R_0 \beta_2 \left(\gamma - \frac{2\pi k}{3} \right). \quad (1.3)$$

A value of $\gamma = 0^\circ$ identifies an axially-symmetric nucleus with a prolate shape, while $\gamma = 60^\circ$ corresponds to an axially-symmetric oblate shape. For intermediate values of γ , the axial symmetry is not conserved and the nucleus is referred as *triaxial*. Examples of quadrupole deformations are shown in Figure 1.1

It should be noted that, although these conventions are commonly used, other notations exist. Often, a $\beta < 0$ is considered to represent oblate shapes. Moreover, the description of ellipsoidal shapes in terms of β and γ parameter is valid in the limit of small deformations, a condition that will be shown to hold in the context of this thesis.

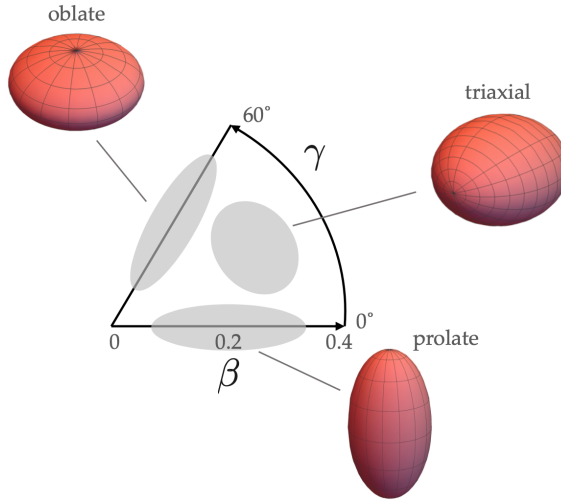


Figure 1.1: Quadrupole deformations ($\lambda = 2$) of the nuclear surface in the β - γ plane. Regions of the plane corresponding to prolate deformations ($\gamma = 0^\circ$), oblate deformations ($\gamma = 60^\circ$), and triaxial shapes are highlighted. The typical scale of 0.2 for the β parameter is also indicated. Courtesy of C. Porzio [5].

For axially-symmetric, quadrupole-deformed nuclei, two other important quantities are the moment of inertia and the quadrupole moment of the ellipsoidal shape, which

can be expressed in terms of β . Considering the moment of inertia of a rigid-body rotor $\mathcal{I} = \frac{2}{5}Mr^2$ and integrating over the nuclear surface of the ellipsoid, one obtains (to first order in β)

$$\mathcal{I} = \frac{2}{5}AMR_0^2(1 + 0.31\beta). \quad (1.4)$$

Since $R_0 \propto A^{1/3}$, the moment of inertia $\mathcal{I}$ is proportional to $A^{5/3}$ and, at leading order, has a linear dependence on β . At the second order, the quadrupole moment of the ellipsoidal nucleus is given by

$$Q_0 = \frac{3}{\sqrt{5}\pi}ZR_0^2\beta(1 + 0.16\beta). \quad (1.5)$$

1.1.2 Collective excitations of deformed nuclei

With the breaking of the rotational symmetry, nuclei can undergo collective rotations [6, 7]. In the case of axially-symmetric nuclei, these excitations are associated to the rotational angular momentum $\mathbf{R}$, which is perpendicular to the symmetry axis. Besides rotation, the nucleus is also characterized by an intrinsic angular momentum $\mathbf{J}_{\text{intr}}$, resulting from the microscopic configurations of the nucleons. Thus, the total angular momentum of the nucleus becomes $\mathbf{I} = \mathbf{J}_{\text{intr}} + \mathbf{R}$, with projection K along the symmetry axis. As $\mathbf{R}$ is orthogonal to the symmetry axis, only $\mathbf{J}_{\text{intr}}$ contributes to the K -projection [7]. If we consider the ground state of even-even nuclei, the intrinsic angular momentum vanishes ($\mathbf{J}_{\text{intr}} = 0$), resulting in a total projection of the angular momentum $K = 0$. As the work of this thesis focuses on even-even nuclei, the following discussion will be limited to the case with null ground-state spin.

In general, rotational motion can be superposed on intrinsic excitations. Among these, the common low-lying vibrational excitations of deformed nuclei are of quadrupole kind ($\lambda = 2$), in analogy with spherical nuclei [7]. Such vibrational modes are characterized by the projection K of the angular momentum along the symmetry axis, which can assume integer values from -2 to 2 . A twofold degeneracy persists for the $\pm K$ modes, corresponding to clockwise and counter-clockwise motions. Due to the symmetry of axially-deformed systems, the excitation modes are restricted to $K = 0$ or $K = 2$. Instead of describing an oscillation around the equilibrium shape, the $K = 1$ mode describes a rotation along an axis perpendicular to the symmetry one, and thus it is neglected. In the $K = 0$ mode, known as β vibration, the oscillation is along the symmetry axis (fluctuation in the quadrupole deformation β) and preserves the axial symmetry. On the other hand, $K = 2$ modes are called γ vibrations, as the motion involves the γ degree of freedom and breaks the axial symmetry.

Deformed systems can exhibit rotational bands superimposed to intrinsic vibrations, giving rise to β and γ bands in the excitation energy spectrum. In particular, the total angular momentum I will take on values $I = K, K + 1, K + 2, \dots$, except when $K = 0$, in which case only even values of I are allowed. Therefore, the excitation energy of rotational states, relative to the energy of the intrinsic excitation, can be expressed as

$$E_{\text{rot}}(I) = \frac{\hbar^2}{2\mathcal{I}}[I(I + 1) - K(K + 1)], \quad (1.6)$$

where $\mathcal{I}$ is the moment of inertia defined in Eq. (1.4).

Although successful in the description of low-lying excitations of several systems, the general applicability of this picture has been questioned. Experimental observation of large discrepancies with respect to the model predictions suggest the interplay of other effects, such as band mixing and γ softness [8, 9].

1.2 Shape coexistence

Shape coexistence is the phenomenon in which nuclear states associated to different intrinsic shapes appear at similar excitation energies. The concept of shapes emerging from the peculiarities of the microscopic wavefunctions was firstly suggested in 1956 by Morinaga [10], following the observation of a deformed 0^+ state at 6.05 MeV in ^{16}O , in contrast with the spherical nature of its ground state. Few years later, the fission shape isomers studied in the actinides were still considered “exotic rarity” [11]. Nowadays, after 60 years of extensive studies, shape coexistence is recognized as a ubiquitous feature across the nuclear chart [12–14]. Figure 1.2 shows the regions of the nuclear chart where shape coexistence was observed up to 2011 [12]. In particular, the $Z = 50$ region, which is the focus of this thesis, is highlighted and labeled with the letter J.

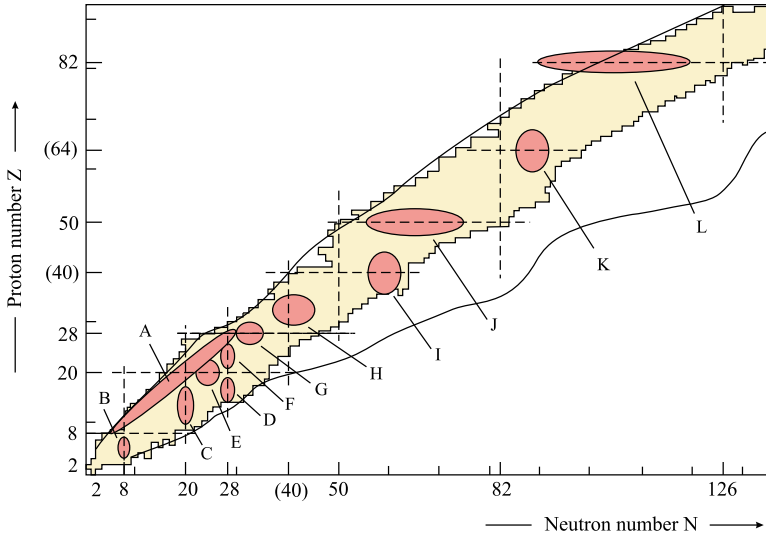


Figure 1.2: The regions of nuclear shape coexistence and their location with respect to nuclear magic numbers, as presented in Ref. [12]. The $Z = 50$ region, relevant for this work, is labeled with the letter J. Reprinted from Ref. [12].

The description of shape coexistence is challenging from both the theoretical and the experimental point of view [15], due to the complex interpretation of results in terms of nuclear shape and the variety of experimental tools needed to fully assess the shape of a nuclear state, respectively. In the following sections, the theoretical description of shape coexistence will be discussed in terms of the Shell model and the Mean Field approaches. Then, the experimental signatures of shape coexistence and few selected examples in different mass regions will be presented.

1.2.1 Theoretical interpretation

Over the past decade, there has been a remarkable improvement of theoretical calculations, thanks to both theoretical methods and upgraded computational facilities [16]. This led to systematic calculations over extended regions of the nuclear chart, performed with a single set of model parameter, thus extending the predictive power of models and guiding experimental efforts. In this context, two main approaches emerge as complementary ways to delineate the appearance of shape-coexisting structures in nuclei. The

first one is the configuration interaction Shell model, which lays its foundation on microscopic interactions tailored for specific orbital configurations. In the second approach, known as Mean Field, the many-body nuclear problem is solved in a self-consistent way, starting from the two-body nuclear Hamiltonian. Although the two approaches are conceptually different, it has been observed that they often converge to the same results, in cases when both of the calculations are computationally feasible [12].

Shell Model

In the nuclear shell model, the nucleus is depicted as a many-body quantum system of independent particles, moving in a mean field potential [17–19]. With the addition of a spin-orbit potential term, the model is able to reproduce the experimentally observed shell structure of nuclei, and, specifically, energy gaps between the eigenvalues of the Hamiltonian, i.e., the single particle energies. In fact, these are grouped based on the oscillator quantum number n , forming major *shells*. The spin-orbit term lifts the degeneracy on the orbital angular momentum l , leading to the *shell gaps* observed in correspondence to specific numbers of protons and neutrons, the spherical magic numbers 2, 8, 20, 28, 50, 82 and 126. Figure 1.3 depicts the single particle orbits of the Shell model, labeled according to the quantum numbers n , l and j , at the stability valley.

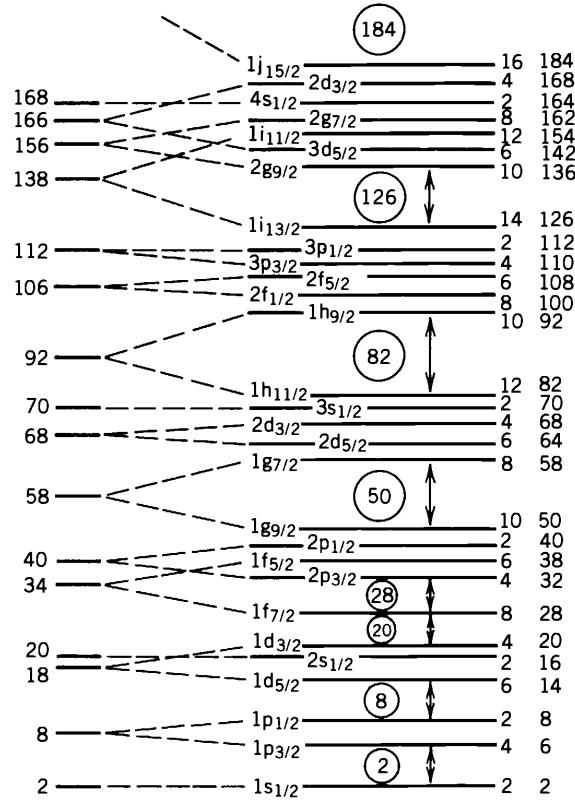


Figure 1.3: Shell model levels calculated with a Woods-Saxon potential plus a spin-orbit term. Adapted from Ref. [20]

In this primitive version of the shell model, also known as independent particle

model, only the single particle terms of the nuclear Hamiltonian are considered, such as the kinetic, the Wood-Saxon and the spin-orbit term. However, when two or more nucleons are outside a fully occupied shell (a *closed shell*), additional correlations and interaction between nucleons have to be taken into account with a residual interaction term, which cannot be absorbed by the single-particle potential [7]. As a consequence, the computational effort needed for diagonalizing the relevant Hamiltonian increases dramatically with the size of the shell. To keep the problem computationally treatable, the SM-CI (configuration interaction shell model) switches from an “exact” solution in the full Hilbert space to an approximate solution in a restricted space, using effective interactions [21, 22]. The latter encapsulate the interaction between single-particle configurations in partially occupied orbits (the *valence space*), which determine the behavior of the system, as opposed to inert nucleons in fully-occupied orbits (the *inert core*). The calculations are then carried out by diagonalizing the effective Hamiltonian in the basis of Slater determinants, built on the valence space up to a certain energy cut-off to keep the problem computationally tractable.

As argued by Dufor and Zuker in Ref. [23], any effective interaction can be split in two parts:

$$\mathcal{H} = \mathcal{H}_m + \mathcal{H}_M, \quad (1.7)$$

a monopole term $\mathcal{H}_m$ and a multipole term $\mathcal{H}_M$. The monopole part is obtained as the average over all the possible orientations of the magnetic substates of two interacting nucleons, hence independent on the quantum number m . It is responsible for the global saturation properties and the evolution of the spherical single-particle energies with the particle configuration in the valence space (also known as *shell evolution*) [22]. Indeed, the eigenvalues of an orbit j depend on the occupancies of all the other valence orbits j' . The multipole part contains terms depending on the angular momentum coupling and takes into account short-range correlations. Its main components, quadrupole and pairing, are responsible for inducing deformations, collective excitations and superfluidity [24]. Moreover, the structure of this component is considered to be of universal nature in the available effective interactions. While the multipole part of effective interactions can be reliably derived through Hamiltonian renormalization within the chosen valence space, the monopole component is often phenomenologically adjusted in order to reproduce the observed shell gaps [22]. In this context, three-body forces have proven to significantly affect single-particle energy, hence the monopole part of the Hamiltonian [25].

Of particular importance for shape coexistence and shell evolution is the tensor force [26, 27]. This non-central component of the nuclear interaction, and specifically its monopole term, is thought to govern two phenomenon: the variation of the shell structure nucleus as the N or Z change along a shell, and the change of shell structure within the same nucleus in systems with larger N/Z imbalance. These phenomena are referred as *Type I* and *Type II shell evolution*, respectively.

In Figure 1.4, the microscopic effects of the monopole part of the tensor force are depicted. Panel (a) shows the relevant neutron and proton orbitals. Proton orbitals are spin orbit partners $j_> = l + 1/2$ and $j_< = l - 1/2$, where l and $1/2$ are the orbital angular momentum and the spin, respectively. As neutrons occupy the $j'_>$ orbital, the effect of the tensor force is to reduce the spin-orbit splitting of the proton orbits. This is due to the attractive (repulsive) coupling of the monopole tensor force between $j'_>$ and $j_<$ ($j_>$) orbitals, which lowers (raises) the single-particle energy of the proton state. An illustrative example is shown in panels b) and c) of Figure 1.4. Since the monopole coupling increases linearly with the number of neutrons, this effect is particularly relevant for

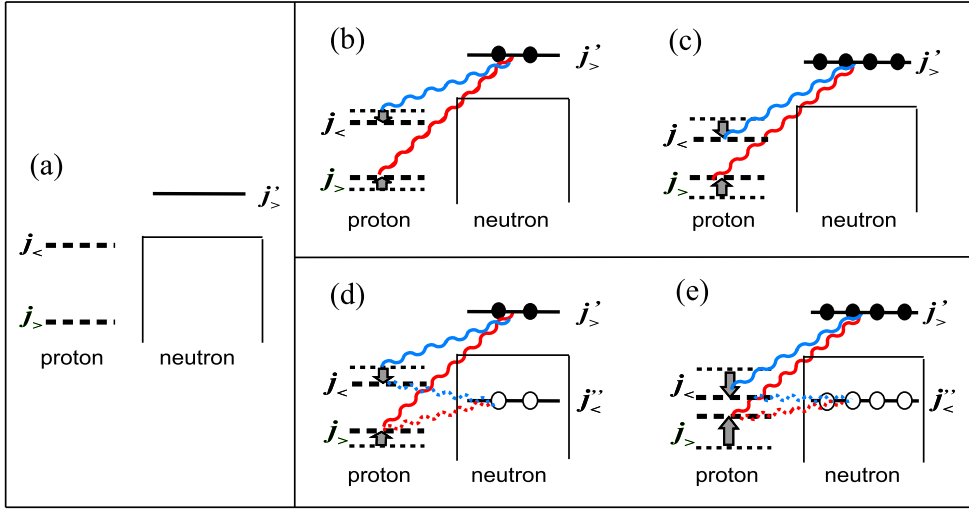


Figure 1.4: Schematic illustration of the microscopic mechanisms named Type I shell evolution (panels (b) and (c)), and Type II shell evolution (panel (d) and (e)), induced by the monopole component of tensor forces of the NN interaction. (wavy lines). Closed (open) circles denote neutron particles (holes). Reprinted from Ref. [26].

exotic nuclei, leading to the energy crossing of orbits, the disappearance of traditional magic numbers and the appearance of new ones [28–30]. These structural changes, originating from the change in the number of nucleons in the system, are referred as *Type I shell evolution*.

A similar effect on proton and neutron orbitals can be obtained by changing orbital configurations within the same nucleus. A schematic description of this phenomenon, defined as *Type II shell evolution*, is shown in panels d) and e) of Figure 1.4. In this case, the promotion of particles to higher lying orbits via particle-hole excitation leads to attractive and repulsive couplings due to the tensor force. If a neutron is excited from the $j_{<}''$ to the $j_{>}''$ orbit, the presence of a hole in the $j_{<}''$ orbit reinforces the reduction of the spin-orbit splitting given by the particle-orbit $j_{>}'$. In fact, the coupling between $j_{<}$ and $j_{<}''$ orbits is attractive, and depends linearly on the number of holes in $j_{<}''$. In contrast, the repulsive interaction between particles in $j_{>}$ and holes in $j_{<}''$ provides an additional increase of the single-particle energy of the $j_{>}$ orbital. Combining all these effects, the proton spin-orbit splitting is reduced even more than in the Type I shell evolution case. As the coupling is magnified by the number of holes and by specific orbital combinations, this phenomenon could induce significant structural changes in nuclei, especially at mid shell. In recent years, type II shell evolution turned out to be particularly relevant for the emergence of coexisting shapes. In the Ni chain, for example, experimental evidence of deformed configurations was predicted by Monte Carlo Shell Model (MCSM) calculations, ascribing their presence to the action of the monopole tensor force. [31–33]

So far we have mainly discussed about the monopole part of the nuclear force, but the multipole term covers a complementary and essential role in the appearance of the shape coexistence phenomena. Although the multipole term is not directly affected by the structural changes induced by the monopole term, its effects can be different for different single-particle energies. Indeed, the same quadrupole interaction could lead to different quadrupole deformations due to the shell evolution [26].

Mean field

As discussed previously, the Shell Model represents one possible solution to go beyond the independent particle model. An alternative approach to explicitly include nuclear correlations is provided by the mean-field methods [34]. These methods describe the motion of nucleons as independent (quasi)particles starting from effective, density-dependent interactions and using the self-consistent Hartree-Fock(-Bogoliubov) framework. In order to obtain single universal frame that is able to both reproduce global nuclear properties, such as radii and binding energies, throughout the nuclear chart and to provide reliable predictions in regions where experimental data is limited or unavailable (e.g., drip-lines), empirically-adapted Hamiltonians have been developed. Based on different assumptions, three different choices of parametrization are currently used: the non-relativistic Skyrme and Gogny forces [35, 36], and the relativistic mean field (RMF) models [37].

A fundamental aspect of mean-field calculations is the reliance on spontaneous breaking and restoration of the symmetries of the Hartree-Fock (HF) wavefunctions, with which correlations between nucleons are encompassed [38]. Indeed, the typical steps of these methods are the following [12]:

1. A nucleon-nucleon interaction is chosen, using parametrization fitted on the properties of spherical and deformed nuclei.
2. The self-consistent variational method is applied to derive the single-particle wavefunctions and occupation probabilities, using a density-dependent Hamiltonian functional which encodes pairing and multipole correlations, while breaking the symmetry of the original Hamiltonian.
3. The nuclear many-body wavefunction is obtained from the independent (quasi)particle states.
4. The broken symmetries of the intrinsic frame are enforced to restore the conserved quantum numbers by projecting the states over particle numbers N and Z , isospin T , and angular momentum J .

Hartree-Fock-Bogoliubov (HFB) or Hartree-Fock+Bardeen-Cooper-Schiffer (HF+BCS) equations are used to describe the many-body problem, including particle-particle pairing correlations in the calculations. By minimizing the HF energy with a constraint on a collective parameter q (generally the quadrupole deformation β), a set of wavefunctions $|\Phi(q)\rangle$ is obtained. After the symmetry restoration, q acts as a semiclassical parameters, and the eigenenergies of the system depend on it. This is schematically shown in Figure 1.5, where the energy surface $E(\beta)$ is depicted.

At this level, each potential energy surface (PES) is characterized by quantum numbers J and M , but no unique determination of the energy of a single state is provided. In order to properly reproduce excitation energy spectra, as in the experimental data, an additional step, beyond the static mean field approach, has to be made [39]. A possible solution is to use the Generator Coordinate Method (GCM) [4], which defines the wavefunction of a state k as the optimal superposition of projected mean-field states:

$$|\Psi_k\rangle = \sum_q f_k(q) |\Phi(q)\rangle, \quad (1.8)$$

where the parameter q acts as a collective generator coordinate, and $f(q)$ are the weights of the different wavefunctions. By applying the variational method, one determines the

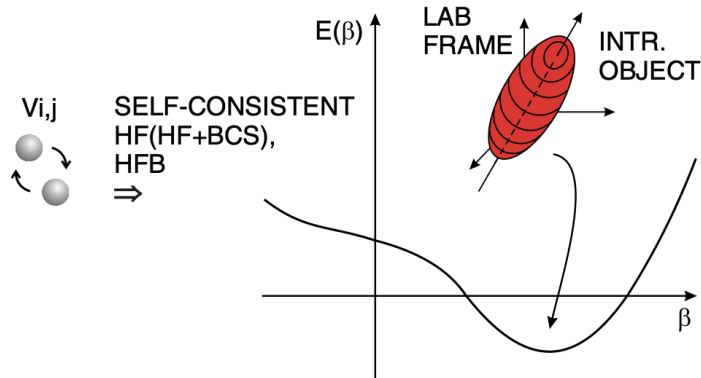


Figure 1.5: Schematic illustration of the energy surface $E(\beta)$ as a function of the collective parameter (quadrupole deformation) β , as can be obtained from self-consistent HFB calculations. Reprinted from [12].

weight functions f_k , the eigenstates $|\Psi_k\rangle$ and the energy eigenvalues E_k . Thus, for each J , a set of excited states, characterized by a mean parameter $\bar{q}$ (e.g., a mean deformation), is obtained from the configuration mixing of mean-field state. With this approach, the single-particle and collective nuclear dynamics, relevant for the shape coexistence phenomenon, are considered within a coherent quantum mechanical formulation [34].

The GCM method has been extensively used to study the shape coexistence phenomenon in a multitude of cases, employing both Skyrme- and Gogny-based interactions [40–42]. Alternative methods, such as the Time-Dependent Hartree-Fock-Bogoliubov (TDHFB) [43], the 5-dimensional Quadrupole Collective Hamiltonian (5DQCH) [44] and relativistic mean-field derivation of the Interacting Boson Model (IBM) Hamiltonian [45], are also widely used, addressing the evolution of the nuclear shape in large regions of the nuclear chart.

1.2.2 Experimental signatures

Several experimental probes may reveal the presence of shape coexistence in atomic nuclei. The most relevant are: (i) level energies of low-lying states, (ii) electromagnetic transition probabilities, (iii) $E0$ transition strength, (iv) multipole moments, and (v) population mechanisms [13, 14]. Combining information from all of these observables is essential to determine the structure of states and provide a comprehensive description of shape evolution.

Level energies

A first indication of shape coexistence is given by the energy of low-lying excited states and by energy systematics. In spherical, even-even nuclei, the presence of low-lying 0^+ states may suggest the presence of deformed configurations. In particular, near closed shells (or sub-shells), the excitation energies of these states often display a parabolic-like behavior along the isotopic chain, which stems from the promotion of pairs of nucleons to orbitals lying above the shell gap (2p-2h excitations) [12–14]. Such excitations, known as *intruder* configurations, gain their binding energy from monopole, pairing and quadrupole correlations, the latter being mainly responsible for the parabolic trend.

The minimum of the parabola is generally found near or at mid-shell, where the proton-neutron quadrupole binding energy maximizes [46].

The appearance of a rotational band, built on top of excited 0^+ states, is another strong fingerprint of shape coexisting structures with enhanced deformation for nuclei with a spherical or weakly-deformed ground state. These excitation energy sequences have been frequently observed near closed shells, in different regions of the nuclear chart. The energy difference between rotational states can also give information about the moment of inertia of the system, hence its average deformation β , as discussed in Section 1.1.

Electromagnetic transition probabilities

In addition to the excitation energy spectrum, γ -ray spectroscopy information is crucial to establish the collectivity of rotational bands and the mixing between low-lying excited states (further details on electromagnetic transitions are provided in Section 2.2). In this context, electric quadrupole $E2$ transitions, and specifically their reduced transition strength $B(E2)$, are sensitive tools to reveal deformed structures. In fact, they probe the degree of shape mixing between an excited 0^+ state and lower-lying states in the main minima of Potential Energy Surfaces. Similar shapes correspond to similar wavefunctions and configurations, therefore, transitions between such states will be favored, resulting in $B(E2)$ values significantly larger than 1 W.u.. Conversely, transitions connecting excited states with strongly different shapes will be hindered, resulting in a small $B(E2)$ often below 1 W.u.. In a mean-field picture, these two states will reside in different minima of the Potential Energy Surface, separated by a sizable barrier which inhibits the electromagnetic decay. Such a case has been referred to as *shape isomerism*, highlighting the retardation of the transition due to the presence of a potential energy barrier between highly non-mixed configurations [14, 47, 48].

Experimentally, $B(E2)$ values are obtained from lifetime measurements or Coulomb excitation reactions. The former, being relevant for the work of this thesis, will be described in detail in Section 2.3.

Electromagnetic moments

As mentioned in 1.1, the intrinsic quadrupole moment Q_0 , defined in the principal-axis frame, is directly related to the deformation of the nuclear system (including its sign). Despite no direct measurement of Q_0 is possible, multi-step Coulomb excitation experiments provide an indirect way to extract its value from the spectroscopic quadrupole moment Q_s . This observable describes the apparent elongation of the charge distribution of a nuclear state in the laboratory frame [14], and, under the axially-symmetric model assumption, can be related to Q_0 through the expression [49–51]:

$$Q_s = \frac{3K^2 - I(I+1)}{(I+1)(2I+3)} Q_0, \quad (1.9)$$

where K denotes the projection of the nuclear spin I on the symmetry axis of the nucleus (see discussion of Section 1.1.2).

E0 transition strength

Electric monopole $E0$ transitions may serve as a powerful probe of the nature of shape coexistence. This decay mode, only for $\Delta J = 0$, proceeds through the emission of con-

version electrons and, if energetically possible, electron-positron pair emission (see Section 2.2.2 for further details). In the case of states with spin $J_i = J_f = 0$, monopole transitions are the only allowed decay mode, as the photon emission is forbidden by the angular momentum conservation. Nuclear structure information is carried by the dimensionless electric monopole strength [52], defined as

$$\rho(E0) = \frac{\langle f | m(E0) | i \rangle}{eR_0^2}, \quad (1.10)$$

where $m(E0)$ is the electric monopole operator and $R_0 = 1.2A^{1/3}$ fm is the typical radius of a spherical nucleus. Since the monopole operator is directly related to the mean-square charge radii, $\rho(E0)$ reveals shape changes of the system [13]. Specifically, in a two-state mixing assumption, this quantity is related to the difference of the mean-square charge radii $\langle r^2 \rangle$ of the two states according to the relation

$$\rho(E0) \simeq \alpha\beta\Delta\langle r^2 \rangle, \quad (1.11)$$

where α and β are the mixing coefficients between the two states [53]. As a result, large $\rho^2(E0)$ strengths may suggest significant differences in deformation between two strongly mixed states. On the contrary, very reduced $\rho^2(E0)$ values may point to either similar deformation or extreme cases of shape coexistence, where no mixing occurs. Electric monopole transition strengths offer information complementary to the $B(E2)$ values, thereby delineating different aspects of the shape coexistence phenomenon.

Population mechanisms

Besides γ -ray and electron spectroscopy, nuclear reaction cross sections provide valuable information about the wavefunction of shape-coexisting states.

One-nucleon transfer reactions are the optimal probe for the single-particle component of the wavefunction, populating states that have a significant overlap with the initial ground-state wavefunction. These reactions proved to be particularly efficient in populating 0^+ states, especially when a nucleon is added to a $J = 1/2$ ground state of a nucleus in the entrance channel [14].

Two-nucleon transfer reactions provide crucial information on the nuclear structure of the 0^+ excitations, specifically probing the pairing correlation of their wavefunctions. Typically, the $0_{gs}^+ \rightarrow 0_{gs}^+$ transition dominates, but unusually strong population of an excited 0^+ level may suggest substantial mixing with the ground state and indicate collective configurations of the nuclear wavefunctions. For instance, an enhancement of the two-proton transfer cross-section for low-lying, excited 0^+ states is considered an indicator of shape-coexistence, hinting at correlated proton-pair vibrations [13]. In the same way, two-neutron transfer processes are deemed as crucial probes for quantum shape phase transitions [54].

In general, for both the one- and two-nucleon transfer processes, detailed calculations that include both the nuclear structure and the reaction mechanism aspects are needed to gain deeper insight into the wavefunction compositions of the states. Due to the complexity and the large model spaces required to calculate the wavefunctions and the transfer amplitudes in medium-heavy systems, state-of-the-art Shell Model implementations have to be used, such as the MCSM [27].

An alternative approach to study excited 0^+ states involves β - and α -decay processes. When allowed by the selection rules, 0^+ states can be populated directly after the β electron (or positron emission). States with similar deformation to the parent nucleus will be

preferentially populated, while transitions to distinct configurations will be suppressed. In the same way, hindrance of energetically allowed α -decay transitions may highlight discrepancies in the microscopic structure of excited configurations. Over the past years, radioactive-decay studies of shape coexistence were preferentially applied in exotic regions of the nuclear chart, where the direct population of low-lying, off-yrast states is challenging.

Neutron capture represents an additional method for populating 0^+ states, characterized by its non-selective population of levels below the capture state. In this reaction, the final nucleus is always populated at the capture state, whose spin results from the coupling of the ground-state angular momentum J_i of the initial nucleus and the $1/2$ spin of the neutron. Thus, for low J_i of the target nucleus (e.g. $J^\pi = 1/2^+$), excited 0^+ states can be directly populated via $E1$ or $M1$ transition, and their spin can be unambiguously assigned with angular correlations. An example of this kind of studies is presented in Ref. [33] for the ^{64}Ni case.

1.2.3 Evidences across the nuclear chart

As introduced previously, shape coexistence is widespread across the nuclear chart, manifesting itself through a variety of phenomena. In the following, few selected examples in different mass regions will be briefly discussed based on historical context and relevance in the field.

Already in the 1950s, Hill and Wheeler suggested the possibility of having two well-separated minima in the potential energy surface of fissioning nuclear systems, which were described as liquid drops [55]. This picture was later confirmed in the 1960s with the discovery of fission isomers in the actinides. Specifically, an isomeric state in ^{242}Am was found at 3 MeV of excitation energy, which primarily decayed by spontaneous fission and showed hindered γ -decay probability [11]. Subsequently, this state was associated to the existence of an intermediate potential energy minimum at larger deformations, and the concept of shape coexistence (in particular shape isomerism) was proposed [14].

Based on these results, the actinides region was explored in detail, leading to the observation of extremely delayed γ -ray transitions from shape isomers in ^{236}U and ^{238}U [47, 48, 56]. The measured hindrances of the order of $10^7 - 10^9$ for $E2$ decays, expressed in terms of the hindrance factor $HF = (2J_f + 1)/B_\downarrow(\sigma L; J_i^\pi \rightarrow J_f^\pi)_{W.u.}$, are among the largest values observed in the entire nuclear chart.

In the following decades, evidences of shape coexistence gradually emerged in several mass regions. A breakthrough came in the 1970s from optical hyperfine-structure studies of mercury isotopes, where an anomalously large isotope shift between $A = 185$ and $A = 187$ was observed [57]. This was soon complemented by in-beam and decay-spectroscopy experiments, which revealed low-lying 0^+ states upon which rotational bands are built [58]. The phenomenon seen in the Hg isotopes raised the question about the survival of the $Z=82$ shell gap in the region, resulting in extensive studies of the Pb isotopes [12]. Similarly to the Hg case, a coexistence of spherical, prolate and oblate configurations was observed in Pb isotopes [60] via γ -ray and electron spectroscopy, as illustrated in Figure 1.6. The energy systematics of the bands built on top of the intruder 0^+ state showed the typical parabolic trend, with a minimum near the $N = 104$ midshell. Nowadays, neutron-deficient isotopes around the $Z = 82$ region are thought to constitute a proper island of shape coexistence, with one of the most extensive manifestations of this phenomenon.

Around the same period, first indications of shape coexistence in the $Z = 50$ re-

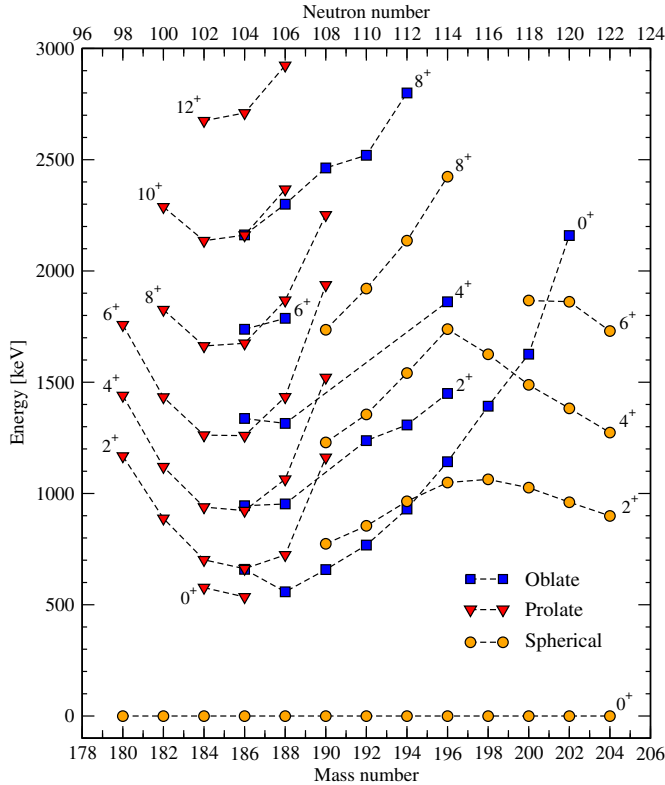


Figure 1.6: Adopted level energy systematics of the neutron-deficient even-mass Pb nuclei. Reprinted from [59].

gion appeared from collective excitations in even-mass Sn isotopes close to stability [61]. Intruder configurations were later observed in Te, Ru and Pd isotopes, using γ -ray spectroscopy and lifetime measurement techniques. More recently, cases of multiple shape coexistence were identified in Cd isotopes, showing similar features to neutron-deficient Pb region [13, 62]. Since the $Z = 50$ region and, in particular, Sn isotopes, are the main topic of this thesis, an up-to-date discussion of this region will be presented in Section 1.3.

Systematic investigations extended the picture to lighter systems, in particular the Kr, Sr, and Zr isotopes near $A \approx 100$, where strong shape competition was predicted by theoretical models [63]. On the proton-rich side, Kr isotopes were shown to exhibit nearly degenerate oblate and prolate configurations [64, 65].

Recently, extensive spectroscopy of the Ni isotopic chain revealed a complex scenario of multiple shape coexistence [14], already visible in $^{62,64}\text{Ni}$ while becoming particularly prominent across the $N = 40$ subshell closure. In $^{64,66}\text{Ni}$, the appearance of prolate-deformed configurations with significantly hindered γ decay branch was described in the MCSM framework as a manifestation of the type II shell evolution phenomenon [32, 33].

In lighter nuclei around $N = 20$, evidences of deformation in the ground-state configuration of the semi-magic ^{32}Mg were already available in the 1980s [13]. Such an observation paved the way to the identification of the so-called *island of inversion*, where

deformed intruder configurations dominate the low-lying structure of Na, Mg, and Al isotopes. Shell-model studies later clarified this as the inversion between normal spherical 0^+ states and neutron $2p - 2h$ intruder 0^+ states. On the other hand, competing configurations with different deformation were also found in light, $N \sim Z$ nuclei. [12, 66]. Examples include ^{24}Mg , where recent $E0$ strength measurements indicate a significant shape change between the ground state and first excited 0^+ state, and ^{28}Si , which is characterized an oblate ground state, a first excited 0^+ state with suggested prolate shape, and a possible superdeformed band at higher excitation energy [66].

Despite being known for more than 60 years, the shape coexistence phenomenon remains a central subject in nuclear-structure research, motivating continuous theoretical developments and experimental efforts throughout the Segrè chart.

1.3 The Sn isotopic chain

The Sn nuclei ($Z = 50$) form the longest shell-to-shell chain of semi-magic nuclei investigated in nuclear spectroscopy to date, offering stringent benchmarks to test theoretical models in both neutron-deficient and -rich regions, especially from the standpoint of systematic trends. Even-even Sn isotopes near the valley of stability are of particular interest for shape coexistence studies as well. In fact, shell and sub-shell closures tend to stabilize these systems into spherical configurations, while on the other side the residual interactions and correlations between protons and neutrons drive the system toward maximum collectivity. The balance between these two aspects may give rise to shape coexistence effects [12]. Moreover, multiple cases of shape coexistence have been identified in $^{186-190}\text{Pb}$ [59, 60, 67] and in cadmium isotopes [13, 62], which display structural features similar to those of the Sn nuclei. These findings open the possibility that spherical, oblate, and prolate configurations may coexist within individual isotopes of the Sn chain.

Experimental signatures of shape coexistence come from a multitude of observations. Firstly, a large number of 0^+ states were observed at low excitation energy, with spin-parity assignment firmly established through high-resolution transfer studies [14, 68]. Their energy systematics along the isotopic chain shows the typical parabolic intrusion of the 0_2^+ state also followed by the 2_2^+ state, as presented in Figure 1.7. The minimum of the parabola is centered at midshell, around $N = 66$. Such behavior is commonly interpreted as $2p - 2h$ proton excitations across the shell gap [12], that gain binding energy from quadrupole and tensor correlations. A similar pattern is observed for the 0_3^+ state, lying roughly 300 keV above the 0_2^+ , while a less pronounced trend is observed for the 0_4^+ state.

The $2p - 2h$ intruder interpretation is further supported by reaction data [13]. In particular, enhanced cross sections to excited 0^+ states (especially the 0_2^+) were measured in $^{110-116}\text{Cd}(^3\text{He},n)$ reactions [69], being in strong contrast with the weak population of such states (of the order of a few percent of the ground-state yield) found in the (p,t) two-neutron transfer reactions [70–73]. These cross sections, normalized to the ground-state values, are shown in Figure 1.8. On the other hand, larger two-neutron transfer cross sections were observed in (t,p) in ^{114}Sn of excited 0^+ states (with about 7-10 % of the ground-state yield) for the 0_2^+ and 0_3^+ , respectively [74]. Although the resolution of the two-proton transfer reaction was not sufficient to disentangle the 0_2^+ and 0_3^+ states, combined evidence from energy spectra and cross-section data indicate that the 0_2^+ state originates from correlated proton-pair excitations, whereas the 0_3^+ state has a distinct configuration.

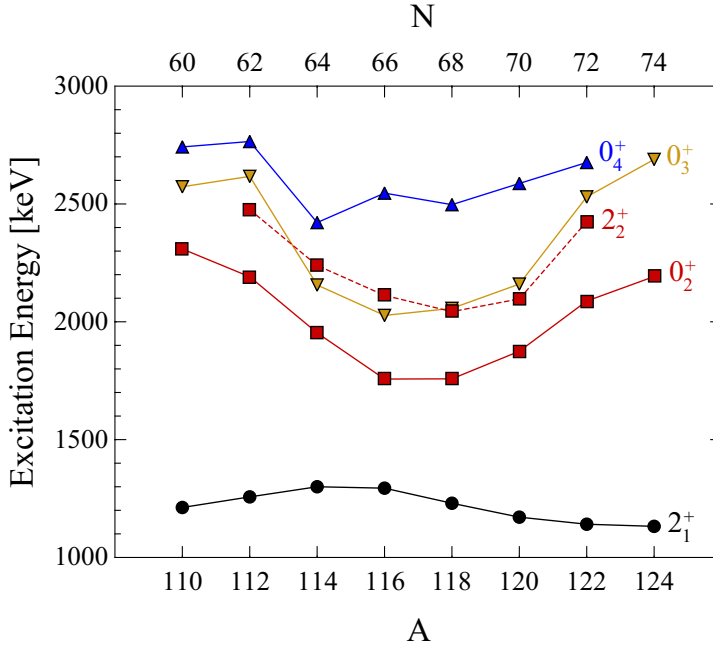


Figure 1.7: Energy systematics of the first 2^+ states (black), the intruder band built onto the 0_2^+ (red) and the 0_3^+ (yellow) and 0_4^+ (blue) states of Sn isotopes. Data taken from [68].

Additional evidence of coexisting shapes comes from $^{110-116}\text{Cd}(\alpha, 2n\gamma)$ reaction studies [61], where positive-parity rotational sequences, built on the 0_2^+ states were observed. These findings, later confirmed by fusion-evaporation experiments [75, 76], were interpreted in terms of proton $(g_{7/2})^2(g_{9/2})^{-2}$ excitations across the shell gap, which is expected to drive the deformation to a prolate shape. In the case of ^{112}Sn and ^{114}Sn , deformation parameters of $\beta \sim 0.18$ and $\beta \sim 0.26$, respectively, were obtained from lifetime measurements under the axially-symmetric rotor assumption [75, 76]. Since these bands tend to decay out at $J^\pi = 6^+$, only a fraction of the intensity feeds the band head. Although in most of these nuclei the 0_2^+ state is considered the band head, its assignment is still a matter of debate in ^{116}Sn . In fact, an indication for the 0_2^+ as band head comes from the sizable cross section in the $(^3\text{He}, n)$ transfer reaction, whereas recent β -decay [77] and (n, γ) [78] studies seem to favor the 0_3^+ according to lifetime measurements, in agreement with Interacting Boson Model (IBM) calculations. This re-assignment was originally proposed in a previous β -decay experiment [79], based on the observation of a 85-keV $2_2^+ \rightarrow 0_3^+$ transition with $B(E2; 2_2^+ \rightarrow 0_3^+) = 99.7(84)$ W.u., which is about twice the $B(E2; 2_2^+ \rightarrow 0_2^+) = 44(5)$ W.u.. Recent lifetime measurements in ^{114}Sn using the $(p, p'\gamma)$ reaction [80] opened the possibility of a similar reassignment of the intruder band-head, however, both the lack of direct lifetime measurement of the 0_3^+ state and the high level of mixing of 0_2^+ and 0_3^+ predicted by IBM calculations prevented a firm assignment.

In terms of transition probabilities, the 0_2^+ states share strong similarities among the Sn isotopes in the region. These states are characterized by enhanced $B(E2)$ transition probabilities towards the 2_1^+ state, with values between 10 and 20 W.u., suggesting a significant admixture with the ground state wavefunction [68]. Conversely, the

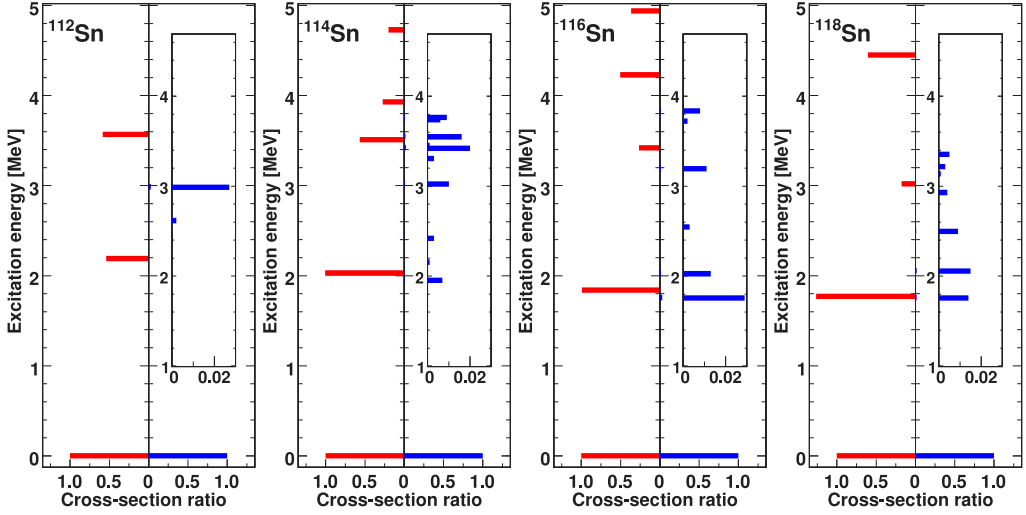


Figure 1.8: Cross sections for the population of excited 0^+ states normalized to those of the ground state. The bars on the left-hand sides of the plots represent the ratios observed in the $(^3\text{He}, n)$ reaction, and those on the right-hand sides are for the (p, t) reactions. The insets on the right-hand sides of the plots display a ratio scale expanded to 0.03 (3%). Data are taken from Refs. [69, 71–73]. Reprinted from [13].

$B(E2; 0_3^+ \rightarrow 2_1^+)$ was found to be just a fraction of the W.u. in $^{116,120}\text{Sn}$, which are the only two cases where the lifetime was measured [78, 81]. Moreover, the 0_3^+ states tend to decay via favored $E0$ transitions towards the 0_2^+ states, rather than the ground state. This behavior is indicative of a shape change and mixing between second and third 0^+ states. In ^{116}Sn and ^{120}Sn , a minimum of $\Delta\beta \gtrsim 0.22$ and $\Delta\beta \gtrsim 0.23$ between the 0_2^+ and 0_3^+ states is required, respectively, to reproduce the measured $\rho^2(E0)$ values [81, 82]. Similarly, the $10^3 \times \rho^2(E0) > 38$ observed for the $0_3^+ \rightarrow 0_2^+$ transition in ^{118}Sn , is evidence of strong mixing between the two states. Still, the complex structure of shape-coexisting states in the Sn calls for additional lifetime measurements and spectroscopic information to be definitively assessed. This is precisely the aim of this thesis.

From a theoretical perspective, Hartree–Fock–Bogoliubov (HFB) calculations of the potential energy surface (PES) using the interaction of Togashi et al. [83] predict a dominant spherical minimum, together with the emergence of a prolate minimum in $^{112,114,116}\text{Sn}$. This prolate minimum weakens in ^{118}Sn and vanishes for ^{120}Sn and heavier isotopes, as well as for nuclei lighter than $A = 110$. A similar, though less pronounced, coexistence scenario is obtained from HFB calculations employing the Gogny force [84, 85]. In contrast, macroscopic–microscopic mean-field calculations by Moeller et al. [86] predict only a spherical minimum. The resulting PES for $^{110-120}\text{Sn}$ from these different theoretical approaches are displayed in Figure 1.9.

Microscopically, Sn isotopes present a similar orbital configuration to that of Ni. An illustration of the orbitals involved in this process, compared with the Ni case, is shown in Figure 1.10. In Ni isotopes, the spontaneous symmetry breaking induced by the proton–neutron interaction was interpreted in terms of the type II shell evolution framework. The monopole term of the tensor force was responsible for lowering the proton $f_{5/2}$ orbital, hence reducing the $Z = 28$ shell gap and favoring proton excitations. This led to the appearance of a deformed minimum. A similar mechanism has

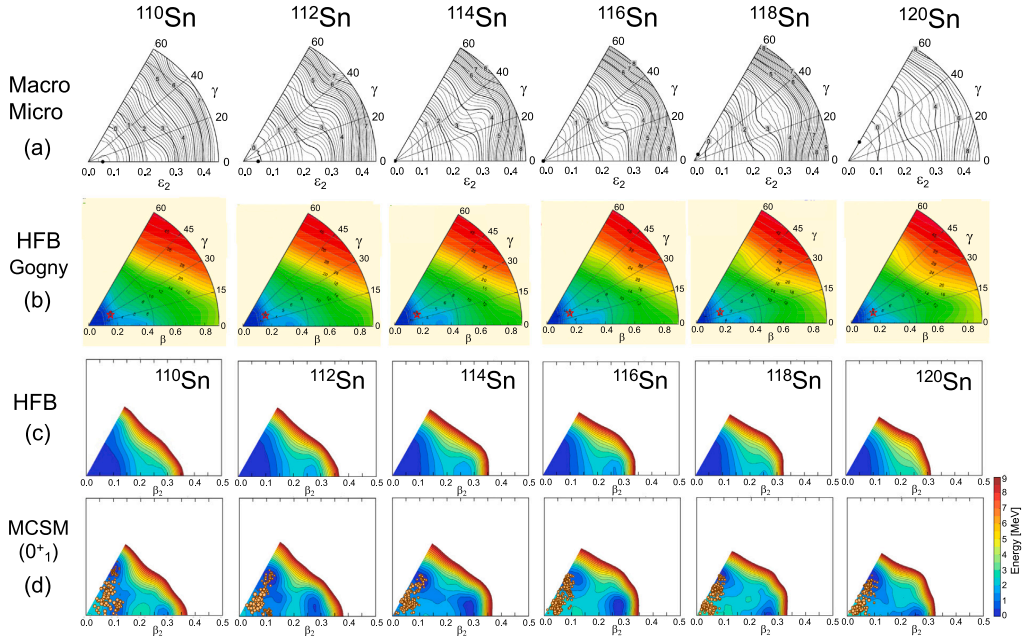


Figure 1.9: Potential energy surfaces for even-even $^{110-120}\text{Sn}$ obtained from different calculations: (a) macroscopic–microscopic mean field model of Moeller et al. [86], (b) Hartree–Fock–Bogoliubov calculations based on the Gogny force, by S. Hilaire and M. Girod [84, 85], (c) Hartree–Fock–Bogoliubov calculations with the interaction of Togashi et al. [83], and (d) Monte Carlo Shell Model calculations by Togashi et al. [83], with circles indicating the main component of the 0_1^+ ground state. Energies (in MeV) are given on the contours in (a)–(b), and by common color scale in (c)–(d). Reprinted from [14].

been proposed for Sn. Neutron excitations in the $h_{11/2}$ intruder orbital would reduce the splitting between the proton $g_{7/2}$ and $g_{9/2}$ orbitals, facilitating the promotion of protons across the $Z = 50$ gap. Based on this, deformed configurations lying in prolate-deformed secondary minima may appear. To date, no experimental evidence has been reported for such type of prolate excitations among the 0_2^+ and 0_3^+ states, therefore, higher-lying levels might be involved (e.g., the 0_4^+ state), as in the Ni case.

To clarify the nature of the coexisting deformed structures in Sn isotopes around $A = 110 - 120$, detailed spectroscopic studies of excited 0^+ states are required, with particular emphasis on lifetime measurements. So far, information on the lifetimes of 0^+ states is scarce, being mostly limited to the 0_2^+ state and, in a few cases, the 0_3^+ state. The available data originate primarily from Coulomb excitation experiments performed in the 1980s by Backlin et al. [87], as well as from fast-timing measurements in β -decay and (n, γ) studies [78, 81]. Complementary direct lifetime measurements could therefore reveal shape-isomer-like structures associated with secondary minima of the PES, while simultaneously providing crucial input toward a more comprehensive understanding of shape coexistence along the Sn isotopic chain.

1.3.1 The experimental campaign at LNL and IFIN-HH

The main aim of this work is to provide further experimental information on the low-lying excited 0^+ states, up to the 0_4^+ , in even-even Sn isotopes near the valley of stability.

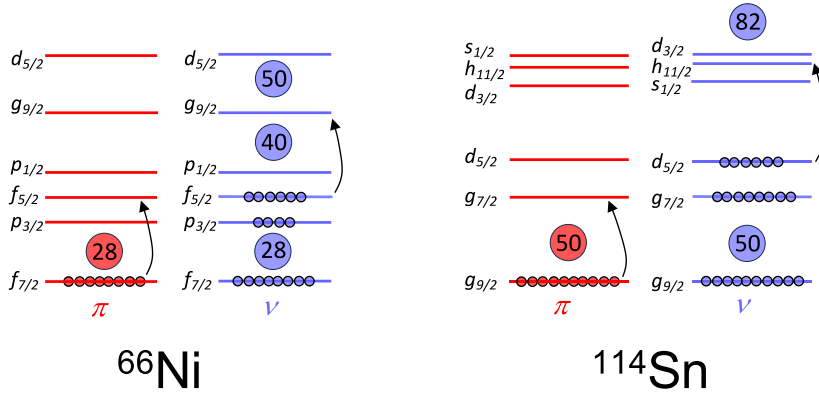


Figure 1.10: Single-particle orbitals of ^{66}Ni and ^{114}Sn . Arrows indicate the excitation of neutrons which favour proton excitations across the shell gap under the action of the tensor force. Reprinted from [14].

For this purpose, a series of experiments was performed at IFIN-HH and LNL, which focused on a spectroscopic investigation of 0^+ states in $^{110-120}\text{Sn}$ via transfer reactions and lifetime measurement techniques. In the experiments performed at IFIN-HH, several transfer reactions were tested to probe the population of the excited states of interest. One- and two-neutron transfer reactions were later employed for lifetime measurements with the DSA and RDDS methods. An alternative approach was chosen for the experiments at LNL. A large set of nuclei were populated in the Cd-Sn region via a multi-nucleon transfer reaction. Spectroscopic study and lifetime analysis, when possible, were carried out. The results of these campaigns are presented in Chapters 5 and 8.

Experimental Approach

This Chapter briefly reviews the experimental techniques used in this work. Section 2.1 will provide an overview on heavy-ion transfer reactions, which were used in the experiments of this thesis. Section 2.2 will describe the basic features of the electromagnetic transitions, with a focus on transition probabilities of γ -ray decays and internal conversion electrons. Finally, Section 2.3 will introduce the lifetime measurement techniques, and, specifically, the two Doppler-shift-based methods employed in the data analysis.

2.1 Heavy-ion transfer reactions

Heavy-ion reactions offer an ideal tool to investigate nuclear structure properties and reaction mechanisms. Depending on the impact parameter, the atomic number and mass of the ions, the energy of the beam and the transferred angular momentum, different processes may occur [88, 89]. These collisions are typically classified with the increase of the impact parameter, evolving from fusion reactions, which require a substantial overlap of the two nuclei, through deep inelastic processes, to quasi-elastic reactions, which are associated with the most grazing-type, peripheral collisions. At distances larger than the sum of the radii of the two ions, only Rutherford scattering and Coulomb excitation processes occur. A pictorial illustration of the different collisions is shown in Figure 2.1.

In this framework, heavy-ion transfer reactions are understood as peripheral, quasi-elastic collisions in which only a few nucleons are exchanged between the interacting ions [88, 89, 91–95]. Due to the short interaction times, of the order of 10^{-23} – 10^{-21} s, heavy-ion transfer processes are categorized as direct reactions, in contrast with the longer timescales of 10^{-18} – 10^{-16} s required for the compound nucleus formation in fusion reactions. At reaction energies close to the Coulomb barrier, such processes are typically associated with grazing collisions, where the surfaces of the colliding ions barely come into contact before separating and the ions retain their individuality. Moreover, grazing trajectories are influenced not only by the repulsive Coulomb interaction but also by the attractive nuclear potential, which becomes significant when the distance of closest approach is comparable to the sum of the nuclear radii. As a consequence of this peripheral, surface-dominated mechanism, the differential cross sections of transfer reactions typically exhibit bell-shaped angular distributions with a maximum at the so-called grazing angle, which disappears when the bombarding energy increases and complex intermediate states are formed.

An important feature of heavy-ion transfer processes is the high degree of selectivity on the population of excited levels. Significant enhancement in the cross section is observed for particular matching conditions on Q -value and orbital angular momentum.

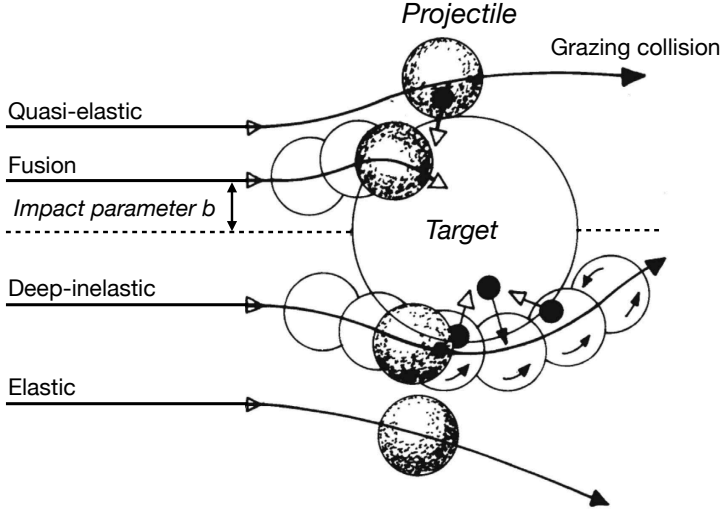


Figure 2.1: Classification of heavy-ion collisions as a function of the impact parameter b . Adapted from Ref. [90].

The reaction Q -value of binary reaction processes $a + A \rightarrow b + B$ is defined as the difference between the initial and final masses of the reacting nuclei or, alternatively, by the difference between the final ($K_b + K_B$) and the initial ($K_a + K_A$) kinetic energy of the recoils. Since in the laboratory frame the target nucleus is typically at rest, i.e. $K_A = 0$, the Q -value can be written as

$$Q = (K_b + K_B) - K_a. \quad (2.1)$$

A further discussion on the reaction Q -value will be found in Section 8.1. Based on a semi-classical approximation, it can be shown that the probability of a transfer channel is maximum when the Q -value is equal to the optimum Q -value (Q_{opt}) [91], defined as

$$Q_{opt} \simeq E_{c.m.} \left(\frac{Z_b Z_B}{Z_a Z_A} - 1 \right), \quad (2.2)$$

where $E_{c.m.}$ is the center-of-mass energy and Z_a, Z_A and Z_b, Z_B are the atomic numbers in the entrance channel and exit channel, respectively. For values different than Q_{opt} , the probability falls rapidly. In the case of neutron transfer, the optimum Q -value is approximately zero, and thus remains constant along an isotopic chain. On the other hand, favorable matching conditions that also account for the transferred angular momentum were derived by Brink [96] for reactions above Coulomb barrier energies:

$$\Delta k = k_0 - \lambda_1/R_a - \lambda_2/R_A \approx 0, \quad (2.3)$$

$$\Delta L = \lambda_2 - \lambda_1 + \frac{1}{2}k_0(R_a - R_A) + \frac{RQ}{\hbar v} - (Z_b Z_B - Z_a Z_A) \frac{e^2}{\hbar v} \approx 0, \quad (2.4)$$

where R is the sum of the radii of the beam R_a and of the target R_A nuclei. The orbital angular momentum and projection on the z -axis for the transferred particle are given by (l, λ) with subscripts 1 and 2 for before and after the transfer, respectively.

In addition, the requirement that the reaction should take place on the reaction plane where the two nuclei meet demands:

$$l_1 + \lambda_1 = \text{even}, \quad (2.5)$$

$$l_2 + \lambda_2 = \text{even}. \quad (2.6)$$

These kinematical conditions depend on the beam energy, the Q -value and the orbital angular momentum of initial and final states, and may lead to strongly enhanced cross sections [91, 96].

Among the heavy-ion transfer collisions, Multi-nucleon transfer reactions (MNT) hold a particular significance. In fact, these reactions can be regarded as a natural bridge from a quasi-elastic regime to a dissipative deep-inelastic process, in which several nucleons and large angular momentum can be transferred between target and beam [95]. These mechanisms are expected to proceed either through sequential single-nucleon exchanges or via more correlated transfers, where the role of pairing and cluster correlations can strongly enhance the probability of transferring multiple nucleons simultaneously. In recent years, multi-nucleon transfer reactions with heavy ions have been used to populate and study moderately neutron-rich nuclei with sizable cross section. This renewed interest was also fostered by the construction of modern large acceptance spectrometer (e.g., PRISMA [97] and VAMOS [98]) with good efficiency and ion-identification capabilities, which, coupled to high efficiency multi-detector array for gamma spectroscopy studies, allow for comprehensive nuclear structure and reaction studies.

Transfer reactions among heavy ions may be described within a quantum-mechanical coupled-channel formalism in the Distorted Wave Born Approximation (DWBA). First-order and second-order approximations are used to describe one- and two-particle transfer, respectively [88]. Alternatively, heavy-ion transfer processes can be treated in a semi-classical approach, which encodes the main features of these reactions [93]. The use of a semiclassical approximation is justified when using heavy ions at low energy ($\approx$ MeV/u), since in this case the De Broglie wave length $\lambda = h/p$ associated with the relative motion is much smaller than the interaction region.

Accurate predictions of transfer cross sections, especially for MNT reactions, have been obtained with the GRAZING model [94]. This code treats the relative motion of the ions on a classical basis and incorporates surface vibrations and transfer in a microscopic way. The relative motion of the system is calculated in a nuclear plus Coulomb effective field, whose parametrization depends on the mass of the interacting nuclei. Collective excitation modes are treated on a macroscopic approximation, while the single-particle states involved in the transfer process are included by using an average single-particle level density. This approach is substantially different with respect to coupled-channel calculations, such as the ones performed by the FRESCO code [99], which are able to describe the exchange of nucleons between specific orbital configurations. However, these require detailed nuclear structure information, such as spectroscopic factors and wave-function overlaps, which are often experimentally unavailable and have to be obtained via Shell Model calculations.

In the context of this thesis, both single- and multi-nucleon, heavy-ion transfer reactions have been employed to populate excited states of even-even Sn isotopes near the valley of stability. GRAZING calculation have been mainly used to estimate both inclusive and differential cross sections of the different reactions. Detailed discussions will be found in the following chapters.

2.2 Gamma-ray spectroscopy

Gamma spectroscopy has become a fundamental technique in nuclear physics, as it provides essential information on the structure of nuclei. By examining γ -ray emission and competing processes, such as internal conversion, one can determine spin and parity of nuclear levels and probe nuclear wave functions through measurements of electromagnetic transition strengths. In the following sections, the basic principles of the γ -ray decay and the internal conversion process will be discussed.

2.2.1 Electromagnetic transitions

After a nuclear reaction or radioactive α and β decays, the final nucleus is often left in an excited state. If the system is bound, the most common de-excitation mode is the γ -ray decay, which consists in the emission of photons.

The typical energy range of γ rays is between 0.1 and 10 MeV, and its dictated by the energy difference of nuclear levels. If we consider a nucleus with mass M that decays from an excited state of energy E_i to a final state E_f via emission of a γ ray, the conservation of energy gives

$$\Delta E = E_f - E_i = E_\gamma + T_R, \quad (2.7)$$

which means that the de-excitation energy is converted into the γ -ray energy and the recoil energy of the nucleus T_R [20]. Using the conservation of momentum, the recoil can be written as $T_R = E_\gamma^2/(2Mc^2)$. Since $\Delta E \ll Mc^2$, the solution for E_γ becomes

$$E_\gamma \cong \Delta E - \frac{(\Delta E)^2}{2Mc^2}. \quad (2.8)$$

Except for light nuclei, the recoil energy correction is generally negligible and the γ -ray energy can be assumed as equal to the level energy difference.

In addition to energy and momentum, electromagnetic transitions also conserve angular momentum and parity. Consider an initial state with spin J_i and parity π_i , and final state with spin J_f and parity π_f . The conservation of angular momentum requires that

$$J_i = L + J_f, \quad (2.9)$$

where L is the angular momentum of the emitted γ ray. Thus, applying the angular momentum coupling rules, L has to satisfy the triangle inequality

$$|J_i - J_f| \leq L \leq |J_i + J_f|. \quad (2.10)$$

The electric or magnetic character of the electromagnetic radiation is determined by the relative parity of initial and final levels. Following the conservation law, the change of parity associated to a radiation field of multipolarity L is determined by

$$\begin{aligned} \Delta\pi(EL) &= (-1)^L, \\ \Delta\pi(ML) &= (-1)^{L+1}, \end{aligned} \quad (2.11)$$

where E and M indicate the electric or magnetic character of the field.

In a quantum mechanical treatment of Maxwell's law, the total transition rate from an excited level i to a level f can be expressed as

$$\lambda_{fi} = \frac{8\pi}{\epsilon_0 \hbar} \frac{L+1}{L((2L+1)!!)^2} \left(\frac{E_\gamma}{\hbar c} \right)^{2L+1} B(\sigma L; J_i \rightarrow J_f), \quad (2.12)$$

Multipolarity (L)	Electric (E)	Magnetic (M)
1	$1.59 \times 10^{15} E_\gamma^3 B(E1)$	$1.76 \times 10^{13} E_\gamma^3 B(M1)$
2	$1.22 \times 10^9 E_\gamma^5 B(E2)$	$1.35 \times 10^7 E_\gamma^5 B(M2)$
3	$5.67 \times 10^2 E_\gamma^7 B(E3)$	$6.28 \times 10^0 E_\gamma^7 B(M3)$
4	$1.69 \times 10^{-4} E_\gamma^9 B(E4)$	$1.87 \times 10^{-6} E_\gamma^9 B(M4)$

Table 2.1: Electromagnetic transition rates for low multipole values L . Transition rates are in s^{-1} , the energies E_γ are in MeV, $B(EL)$ in $e^2 fm^{2L}$ and $B(ML)$ in $\mu_N^2 fm^{2L-2}$. Values taken from Ref. [101].

where L , E_γ and σ are the angular momentum, the energy and the character of the photon, J_i and J_f are the spin of the initial and final states, and $B(\sigma L; J_i \rightarrow J_f)$ is a quantity called *reduced transition probability* [100]. This quantity depends on the matrix elements of the multipole operator $m(\sigma L)$ between the initial and final states, according to

$$B(\sigma L; J_i \rightarrow J_f) = \frac{1}{2J_i + 1} |\langle f || m(\sigma L) || i \rangle|^2. \quad (2.13)$$

A summary of the relation between transition rates and reduced transition probabilities is reported in Table 2.1.

Based on Eq. (2.12), the lifetime τ_i of a state i is defined as

$$\tau_i = \sum_f \frac{1}{\lambda_{fi}}, \quad (2.14)$$

being λ_{fi} the transition rates of all possible final states f . As a consequence, the reduced transition probability, namely the overlap between initial and final wavefunctions, can be obtained via lifetime measurements. In case of multiple γ decays depopulating the initial state, the reduced transition probability can be computed from the partial lifetime of each branch, which is given by

$$\tau_\gamma = \tau \frac{I_{tot}}{I_\gamma}, \quad (2.15)$$

where I_γ is intensity of the γ ray of interest whereas I_{tot} is the sum of the intensities of all possible decay branches, including both γ -ray emission and, if present, internal conversion processes.

For γ rays, it is also possible to have different multiplicities, sharing both the electric and the magnetic characters. In such cases, the *mixing coefficient* δ is introduced as the ratio between the reduced matrix elements of the two multipole operators:

$$\delta \equiv \frac{\langle f || m(\sigma' L') || i \rangle}{\langle f || m(\sigma L) || i \rangle}, \quad (2.16)$$

where the σ and σ' indicate either the electric or magnetic character, while L and L' refers to the multipolarity of the operators. By convention, L' is assumed to be greater than L . Considering Eq. (2.13), δ^2 becomes proportional to the ratio of the reduced transition probabilities of the two excitation modes. As a consequence, the mixing coefficient has to be taken into account when calculating the partial lifetimes of each multipole.

A useful scale for transition rates are the Weisskopf units (W.u.) [102]. One Weisskopf unit is the reduced transition probability of a single nucleon moving between two spherical single-particle orbitals [100], and can be written as

$$\begin{aligned} B_{W.u.}(EL) &= \frac{1}{4\pi} \left(\frac{3}{L+3} \right)^2 1.2^{2L} A^{2L/3} e^2 \text{fm}^{2L}, \\ B_{W.u.}(ML) &= \frac{10}{\pi} \left(\frac{3}{L+2} \right)^2 1.2^{2L-2} A^{(2L-2)/3} \mu_N^2 \text{fm}^{2L-2}. \end{aligned} \quad (2.17)$$

These single-particle estimates allow for a rough evaluation of the number of nucleons involved in the de-excitation, hence in the excited configuration of the nucleus. Probabilities higher than one Weisskopf unit hints at a collective behavior, while a fraction of a Weisskopf unit indicates a hindered transition, which may connect states with different microscopic configurations.

2.2.2 Internal conversion electrons

Internal conversion (IC) is a competing process to γ -ray emission in the de-excitation of atomic nuclei [20, 103, 104]. Instead of releasing the transition energy through a photon, the nucleus transfers this energy directly to one of the bound atomic electrons, due to the interaction with the nuclear electromagnetic field. The electron is subsequently ejected from the nucleus, with a kinetic energy given by

$$T_e = \Delta E - B_e, \quad (2.18)$$

where ΔE is the energy difference between the initial and final levels and B_e is the binding energy of the electron (with positive sign) in its atomic shell. Since electrons belonging to different atomic shells can be involved in the process, the IC electron spectrum typically exhibits several components corresponding to the K , L , M , $\dots$ shells, with the K -shell emission being the most probable. After the conversion process, the vacancy left in an inner electronic shell is rapidly filled by electrons from higher shells, accompanied by the characteristic X-ray emission [20]. The relative probability of internal conversion compared to γ emission is quantified by the *internal conversion coefficient* α , defined as

$$\alpha = \frac{\lambda_{\text{IC}}}{\lambda_{\gamma}}, \quad (2.19)$$

where λ_{IC} and λ_{γ} are the partial transition rates for internal conversion and γ -ray emission, respectively. As the atomic physics aspects of this process are well known, it is possible to theoretically calculate the α coefficients [20, 105]. The contributions from the inner shells can be calculated with good precision, while higher order effects from electrons in outer shells are typically approximated, treating those electrons as free particles.

As introduced in Section 1.2.2, the internal conversion process becomes particularly relevant when considering electric monopole ($E0$) transitions, for which γ -ray emission is strictly forbidden by angular momentum and parity conservation. In such cases, internal conversion (and, at higher energies, electron-positron pair creation) represents the only possible decay mechanism. The monopole operator $m(E0)$, which is obtained from the multipole expansion of the Coulomb interaction between the atomic electrons and the nuclear protons, can be written as

$$m(E0) = \sum_{i=1}^Z e r_i^2, \quad (2.20)$$

where the sum runs over all protons and r_i is the radial coordinate of the i -th proton. The strength of an $E0$ transition is characterized by the square of the dimensionless parameter $\rho(E0)$ (previously defined in Eq. 1. 10):

$$\rho^2(E0) = \left| \frac{\langle f \| m(E0) \| i \rangle}{eR^2} \right|^2,$$

where $R \approx 1.2A^{1/3}$ fm is the conventional nuclear radius, and $\langle f \| m(E0) \| i \rangle$ is the reduced matrix element of the monopole operator between the initial and final nuclear states [52]. The values of $\rho^2(E0)$ are generally $10^{-1} - 10^{-3}$ and so it is standard practice to quote $\rho^2(E0) \times 10^3$. The relationship between the transition rate $\lambda(E0) = 1/\tau(E0)$ and the transition strength $\rho^2(E0)$, is defined by

$$\lambda(E0) = \Omega \rho^2(E0), \quad (2. 21)$$

where Ω is an atomic electron factor, which can be calculated [53]. This term is independent of the nuclear properties, thus, the nuclear structure information is carried by $\rho^2(E0)$.

Experimentally, $\rho^2(E0)$ values are extracted by combining measured internal conversion coefficients with accurate knowledge of level lifetimes [106]. These quantities provide a sensitive probe of changes in nuclear mean-square charge radii between states, and therefore play a central role in the study of shape coexistence, configuration mixing, and collective degrees of freedom in nuclei [12–14, 53, 104].

2.3 Lifetime Measurements

The lifetime of bound excited states is an essential observable for the understanding of the structure of nuclei, providing information about the wavefunction of the states involved in the decay process. In order to cover the large time scale of nuclear lifetimes, namely $10^{-15} - 10^{-6}$ s, a variety of techniques have been developed [101, 107, 108]. As illustrated in Figure 2.2, these techniques can be classified into direct methods, which measure the lifetime τ , and indirect methods, which probe the decay width $\Gamma = \hbar\lambda = \hbar/\tau$ [108].

For the purpose of this work, two direct methods were employed: the Doppler Shift Attenuation Method (DSAM) and the Recoil Distance Doppler Shift (RDDS) method (also known as Recoil Distance Method). Both of them exploit the Doppler shift effect to determine the lifetime of a nuclear excited state, covering ranges of 10 – 1000 fs and 1 – 1000 ps, respectively. In fact, the energy of a γ ray emitted by a moving source (i.e. the nucleus) will be shifted according to the relation

$$E'_\gamma = E_\gamma \frac{\sqrt{1 - \beta^2}}{\beta \cos \theta}, \quad (2. 22)$$

where E_γ is the γ -ray energy in the rest frame, $\beta = v/c$ is the velocity of the nucleus and θ is the relative angle between the direction of motion of the nucleus and the direction of the emitted γ ray. Based on Eq. (2. 22), it is clear that different velocities of the recoiling nucleus will result in different observed γ -ray energies. Thus, by controlling the nucleus' motion, one can measure the lifetime of an excited state.

In the following sections, the basic principles of the DSA and the RDDS techniques will be presented, while specific details on the analysis methods will be discussed in Sections 5.2.1 and 5.2.2.

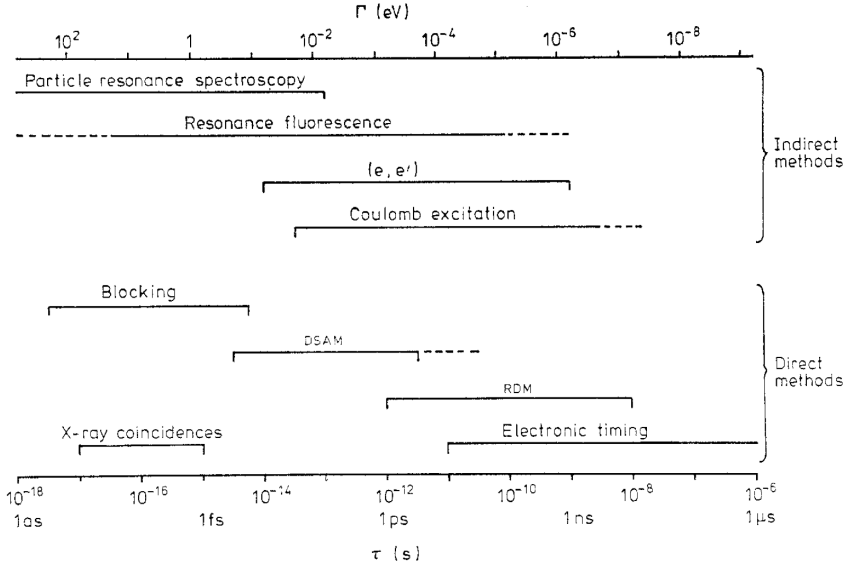


Figure 2.2: Overview of the lifetime measurements methods and their domain of applicability. Reprinted from Ref. [108].

2.3.1 Doppler Shift Attenuation Method (DSAM)

The concept of the DSAM is based on the fact that, in a stopping material, the Doppler shift becomes attenuated due to the slowing down of the recoil nucleus. In a nuclear reaction, the recoiling nucleus is populated in an excited state, and decays by γ -ray emission after a certain amount of time t . If a thick target or a backing layer are used to slow down the nucleus, the corresponding Doppler shift will depend on the instantaneous velocity $\beta(t)$ [107], according to the relation

$$E'_\gamma(t) = E_\gamma \frac{\sqrt{1 - \beta(t)^2}}{\beta(t) \cos \theta} \quad (2.23)$$

which is the time-dependent version of Eq. (2.22). The velocity of the nucleus will decrease continuously from the initial velocity $\beta = \beta_0$ to $\beta = 0$, resulting in a smooth distribution in the γ -ray spectrum. This distribution, which depends on the lifetime of the state, the response function and the angular coverage of the detectors, is denoted as *lineshape* of a γ -ray peak.

In the first implementations of the DSA technique, the lifetime of the state was obtained from the average energy shift $\langle \Delta E_\gamma \rangle$, which can be expressed as

$$\langle \Delta E_\gamma \rangle = F(\tau) E_\gamma \beta_0 \cos \theta, \quad (2.24)$$

where $F(\tau)$ is the attenuation factor and β_0 is the initial velocity ($t = 0$). By comparing the experimental attenuation factor, which is also equal to $F(\tau) = \langle \beta \rangle / \beta_0$, to the theoretical curve calculated from the stopping power theory, one obtains the lifetime of the state [108]. When $F(\tau) \approx 1$ ($F(\tau) \approx 0$), the lifetime is outside the sensitivity range of DSAM, being much smaller (larger) than the stopping time of the nucleus. Despite the strong dependence on the initial velocity and the stopping material, the average stopping time

lies between 100 fs and 1 ps, therefore, the sensitivity range of the DSAM is from tens to thousands of fs.

In more recent applications of this technique, realistic Monte Carlo Simulations are employed to extract the lifetime of a state [109, 110]. The stopping history, the velocity distribution, the feeding of the state and the detector response are taken into account to reconstruct the lineshape of the peak. Thus, the lifetime measurement becomes a minimization problem, in which the simulated γ -ray spectra are compared with the experimental one. A Monte Carlo DSAM implementation for the ROSPHERE array, developed in the context of this thesis, will be exhaustively presented in Section 4.4.

2.3.2 Recoil Distance Doppler Shift (RDDS) Method

The RDDS method allows for lifetime measurements in the 1 – 1000 ps range through the comparison with the time of flight of a nucleus traveling in the vacuum [108]. After leaving the thin, production target at a velocity v , recoil nuclei are then stopped in a thick foil of an high-Z material, placed at a distance d . Those nuclei decaying in flight will emit γ rays whose energies are subject of the Doppler shift of Eq. (2. 22). Conversely, the energy of γ rays emitted from stopped nuclei is unshifted. Therefore, the total γ -ray spectrum shows both the unshifted and the Doppler-shifted components, with intensity I_u and I_s , respectively. Assuming a direct population of a level that decays exponentially with a lifetime τ , the unshifted and shifted intensities are

$$\begin{aligned} I_u(t) &= \int_t^\infty \left(-\frac{dN}{dt'} \right) dt' = N_0 e^{-t/\tau}, \\ I_s(t) &= \int_0^t \left(-\frac{dN}{dt'} \right) dt' = N_0 (1 - e^{-t/\tau}), \end{aligned} \quad (2. 25)$$

where N_0 is the total number of nuclei produced in the excited state and t is the time of flight between target and stopper [10]. Thus, the fraction of the unshifted component can be written as

$$R = \frac{I_u}{N_0} = \frac{I_u}{I_u + I_s} = e^{-t/\tau} = e^{-d/(v\tau)}. \quad (2. 26)$$

In the latter, the time of flight t was expressed in terms of the distance d and component of the velocity v perpendicular to the stopper foil. The lifetime of the state is then obtained from the variation of R as a function of the target-stopper distance, or, alternatively, from the variation of the intensity I_u (or I_s) with respect to a quantity proportional to N_0 . The distance is typically controlled by a device called *plunger*, and for this reason this technique is also referred as the *plunger method*. The basic principles of the plunger method are illustrated in Figure 2.3

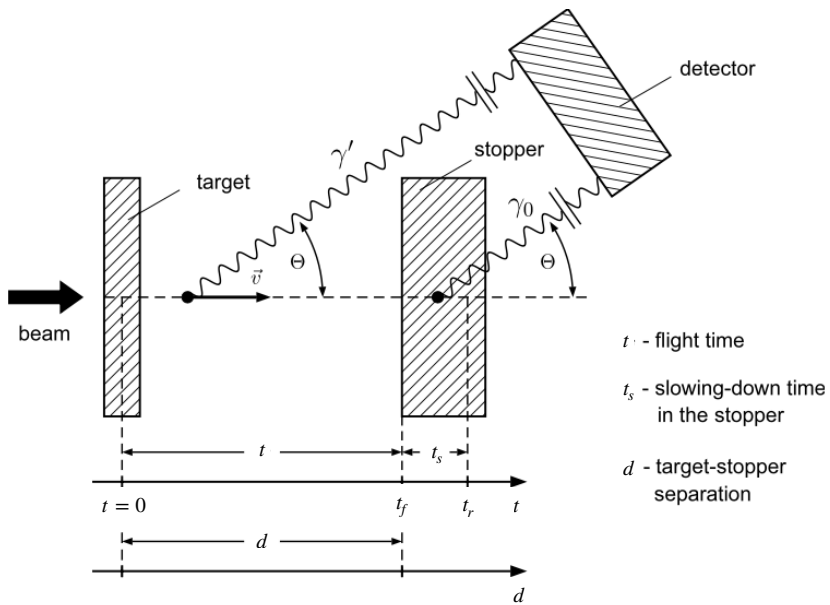


Figure 2.3: Schematic description of the Recoil Distance Doppler Shift method. Adapted from Ref. [111].

Part II

The Experiment at IFIN-HH

The experiments at IFIN-HH

A series of experiments was conducted between 2021 and 2024 at the Tandem facility of the Horia Hulubei National Institute for R&D in Physics and Nuclear Engineering (IFIN-HH), employing the ROSPHERE γ -ray spectrometer coupled to the SORCERER charged particle detector. The Sn isotopes of interest were populated with single- and multi-nucleon transfer reactions at Coulomb barrier energies, using both thick targets and plunger targets for RDDS measurements.

In the following sections, technical aspects of the experiments are presented. Section 3.1 provides an overview of the IFIN-HH facility and its detection systems, namely ROSPHERE, SORCERER and the plunger device, while Section 3.2 describes the reactions and the configurations used in the different experiments.

3.1 Experimental setup

This Section describes the experimental setup that was used in the experimental campaign performed at IFIN-HH between 2021 and 2024. In particular, technical details on the ROSPHERE γ -ray spectrometer, the SORCERER charged particle detector and the plunger device are discussed.

3.1.1 The Tandem facility

The 9 MV FN Pelletron Tandem accelerator is one of the three accelerator facilities of IFIN-HH [112]. It was firstly installed in 1973 and later upgraded in 2006, when the original charging belt was replaced with a Pelletron system. This electrostatic accelerator is capable of providing beams with energies up to 5-10 A MeV, depending on the species.

The majority of the beam species is produced by the Cs Sputter Ion Source, which sputters atoms from the cathode and provides a negative ion beam after their passage through a neutral Cs layer. The negative beam is then extracted and injected into the beam line, where the positive potential of the terminal provides the first acceleration stage. In order to maintain the high voltage on the terminal, the tank of the accelerator is insulated by using a mixture of gases (10 bars of SF₆ and N₂ mixture). The negative ions reaching the terminal region are stripped of electrons using thin carbon foils (5-10 g/cm²), thus changing their charge state from -1 to a positive one. Finally, the positive beam is accelerated away by the same positive terminal potential. Magnetic quadrupole lenses and bipolar magnets are used to focus and analyze the beam. A photograph of the accelerator can be found in Figure 3.1.

During the thick target experiments performed at IFIN-HH, beams of ¹³C, ¹⁶O and ¹⁸O were accelerated at 43, 56 and 59 MeV, respectively. On the other hand, plunger experiments required few additional MeVs of beam energy to compensate the energy

loss in the supporting material of the target (Ta or Au layers of $1\text{--}3\text{ mg/cm}^2$). A fine adjustment of the Tandem voltage was sufficient to obtain the desired beam energy, and the same charge state was maintained.

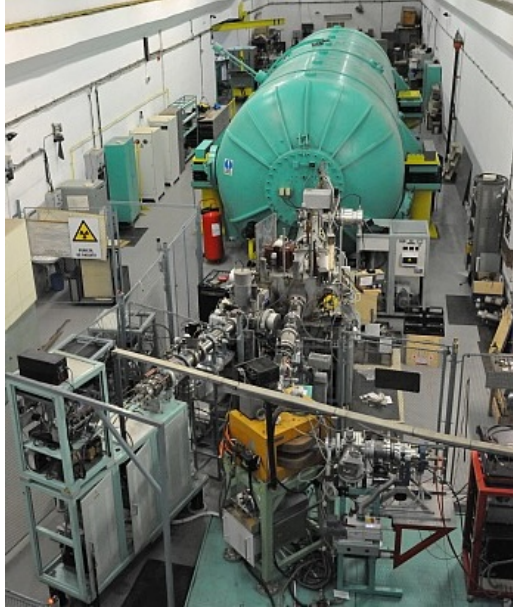


Figure 3.1: The 9 MV FN Pelletron Tandem accelerator installed at IFIN-HH. Reprinted from Ref. [112]

3.1.2 ROSPHERE

The ROmanian array for SPectroscopy in HEavy ion REAction (ROSPHERE) is a Compton suppressed γ -ray spectrometer installed at the 9 MV Tandem accelerator facility of IFIN-HH [113]. From its design phase, the array was built to suit the needs of experiments primarily devoted to measurements of excited nuclear state lifetimes. In particular, the compatibility with a plunger device, the necessity of positioning detectors at various polar angles, and the possibility of mounting both HPGe detectors and large volume scintillators with BGO shields were considered.

The ROSPHERE spectrometer can be used in a mixed configuration, which combines HPGe detectors and $\text{LaBr}_3(\text{Ce})$ scintillators, or in a full HPGe configuration, with up to 25 Compton suppressed detectors. Its support structure has a spherical geometry and is divided into five rings, located at 37° , 70° , 90° , 110° and 143° with respect to the beam axis. Each ring can accommodate up to 5 detectors with their respective anti-Compton shields. An example of configuration is shown in the left picture of Figure 3.2, where 20 HPGe detectors and 5 $\text{LaBr}_3(\text{Ce})$ scintillators were mounted, while a CAD drawing of the entire ROSPHERE structure in a full HPGe configuration is shown on the right side.

Three different type of HPGe detectors are mounted on the array. These devices are commercially available p-type coaxial detectors, produced by ORTEC (type *a*) and CANBERRA (type *b* and *c*). Type *a* and *b* detectors have a relative efficiency of 50-60% and typical resolution (FWHM) of 1.9 keV at 1.33 MeV, while type *c* are larger HPGe detectors with a relative efficiency of about 80% and FWHM ≈ 2.3 keV at 1.33 MeV. In

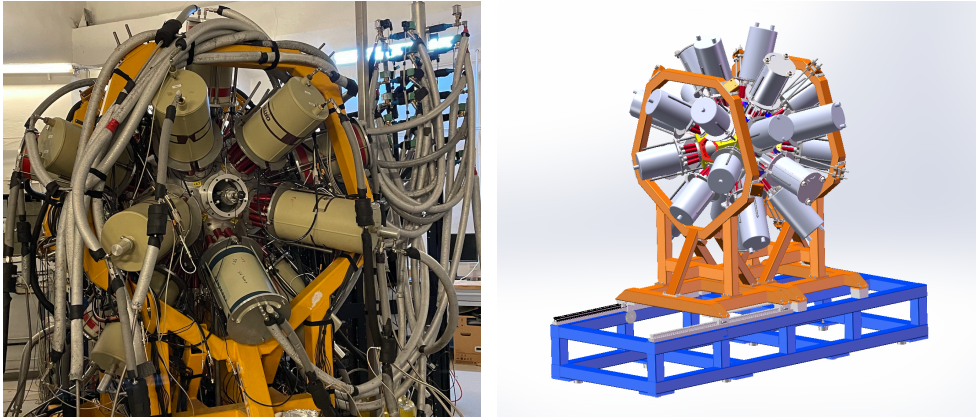


Figure 3.2: (Left) Picture of the ROSPHERE γ -ray spectrometer at IFIN-HH in 2022. (Right) CAD drawing of the ROSPHERE frame, support structure and detectors.

Ring	θ (degrees)	ϕ (degrees)	Distance a/b (mm)
1	37	0, 72, 144, 216, 288	179/210
2	70	36, 108, 180, 252, 324	186/217
3	90	0, 72, 144, 216, 288	176/208
4	110	36, 108, 180, 252, 324	186/217
5	143	0, 72, 144, 216, 288	179/210

Table 3.1: Details on the geometry of the ROSPHERE array in terms of angles and endcap-to-target distances of type a and b detectors. The θ and ϕ angles are expressed in cylindrical coordinates with respect to the beam axis.

Table 3.1, the possible placements of the different HPGe detectors, expressed in terms of (θ, ϕ) coordinates and endcap-to-target distances, are summarized.

The Compton suppression is carried out by using custom designed BGO shield manufactured by Scionix Holland BW. These scintillators consist of 8 optically separated trapezoid shaped BGO crystals, readout by 8 Hamamatsu Photo Multiplier Tubes (PMTs). To shield them from direct γ rays emitted from the target position, conical Tungsten collimators are mounted in front of the crystals. As a consequence, only γ rays which were scattered after the interaction with the HPGe material can be detected. CAD drawings of type a and b detectors with their BGO shield and Liquid Nitrogen dewar are shown in Figure 3.3.

In a full HPGe configuration with 20 type a/b detectors and 5 type c detectors, an absolute efficiency between 2 and 2.5%, and a Peak-to-total (P/T) ratio of about 50% are expected.

The acquisition system consists in a fully digital electronic chain with commercial digitizers modules. After the upgrades of 2020, the previous analog readout system (described in Ref. [113]), was replaced with a simpler chain. CAEN V1725 digitizers modules are employed to read the signals of the default charge sensitive ORTEC and CANBERRA preamplifiers, sampling the signals with a rate of 250 MHz/s. Such mod-

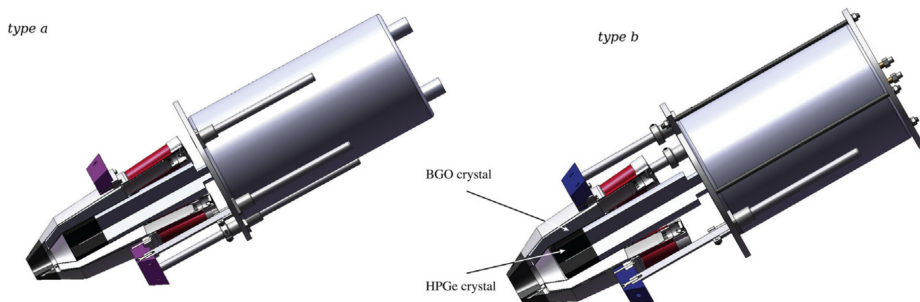


Figure 3.3: CAD drawings of type *a* and *b* Compton suppressed HPGe detectors. Reprinted from Ref. [113].

ules allow for a straightforward synchronization between the boards and the possibility to build up complex acquisition systems. To apply the trapezoidal filter and extract time and energy information, the DPP-PHA firmware is used. As a last step, the acquired information are stored in disks in the Cap'n Proto format [114].

The ROSPHERE array can be coupled to ancillary particle detector systems, such as SORCERER [115], a plastic scintillator with Silicon Photo Multipliers (SiPM) for fast-timing measurements and an annular Si detector. Due to mechanical compatibility reason, the SORCERER array is generally used in combination with the plunger device for lifetime measurements in the 1-100 ps range.

3.1.3 SORCERER

The SORCERER array (SOLar CELls for Reaction Experiments at ROSPHERE) is a charged particle detector array based on commercial photodiodes [115]. A picture of these detectors, mounted in the plunger configuration, can be found in Figure 3.4. It was designed and built to be used in conjunction with the ROSPHERE γ -ray spectrometer, in order to provide measurement of recoiling particles emitted in transfer reactions at Coulomb barrier energies. For this reason, the mounting structure was designed to exploit the features of such reactions in direct kinematics, with recoiling beam-like particles being generally scattered at backward angles in the laboratory frame.

The holding structure of the array is manufactured by performing a laser sintering procedure on a plastic compound, which has the advantage of preventing sealed holes. Thus, the mount can sustain typical reaction chamber pressures of 10^{-6} mbar. Furthermore, the constructed mounts are tailored to the ROSPHERE scattering chamber and the Cologne-Bucharest plunger device.

SORCERER can host six Advanced Photonics PDB-C613-2-ND diodes [116], which are equally spaced and oriented around a central beam hole with a diameter of 5 mm. Each photodiode has an active area of $8.8 \times 9.8 \text{ mm}^2$ which is about 89% of its surface area, with a pairwise electrical isolation. To ensure a rapid exchange of degraded photodiodes, the cells are installed without the usage of adhesive. During the data acquisition, each diode is read out individually and delivered to Canberra 2003BT preamplifiers and mesytec MSCF-16-LN modules for amplification and discrimination. Afterwards, the signals are fed into mesytec MADC-32 and MTDC-32 digitizers for energy and time information, respectively.

In the scattering chamber, the array is placed at a mean distance of 15.7 and 18.9 mm from the target position in the DSAM and RDDS configurations, respectively. Each

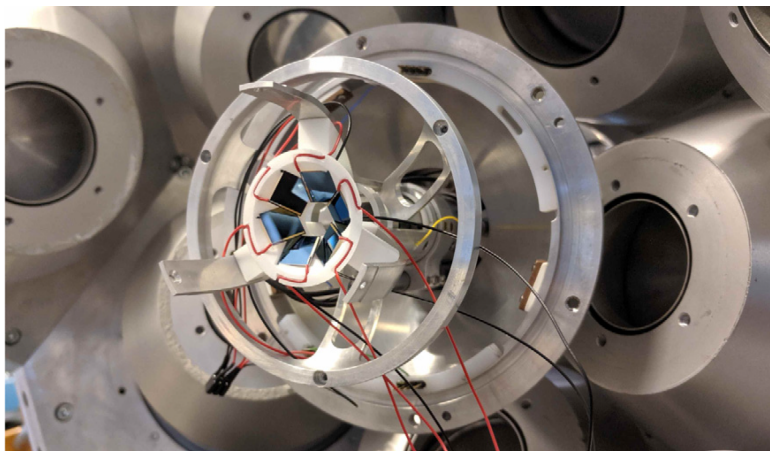


Figure 3.4: The SORCERER Si array mounted in the plunger configuration with the target removed.

photodiode covers an azimuthal angular range of 60° , while the polar angular coverage depends on the configuration. For DSAM experiments, SORCERER covers polar angles between 121.7° and 163.5° , leading to a total subtended solid angle of 15.4%. In the RDDS experiments, the polar angle coverage is reduced to the 130.4° to 165.5° range, and the covered solid angle becomes 11.0%.

3.1.4 The IFIN-HH plunger

In order to perform RDDS lifetime measurements with ROSPHERE, a dedicated plunger device was built on the basis of the plunger developed by the Cologne group [111]. This device is used to obtain the required separation between target and stopper in the micrometer to millimeter range, with sub-micrometer resolution. Its schematics and pictures are shown in Figure 3.5.

It consists of three systems: a mounting system housed in a spherical chamber at the centre of the array that allows the parallel mounting of the target and stopper with micrometer precision; a driving and control system consisting of an off-axis NEXACT N-381 piezo electric actuator with a 25-mm travel range and $0.02\text{-}\mu\text{m}$ resolution, a stack of piezo-ceramics (“piezo-crystal”) for short range corrections and an inductive micrometer probe to measure the travel distance with a $0.1\text{-}\mu\text{m}$ precision; a system of concentric tubes and sliding bearings that transfer the movement of the motor to the target frame. This is the most sensitive part of the plunger, as it has to transfer the off-axis linear motion of the actuator to the target frame while preserving the precision. The target to stopper distance is kept constant by a feed-back system that uses the capacitance method to measure the distance and the piezo-crystal to correct beam induced drifts from the set distance.

3.2 Heavy-ion transfer reactions at the Coulomb barrier

The four experiments of this campaign aimed at studying the population and measure the lifetime of low-lying excited 0^+ states in Sn isotopes in proximity of the valley of stability. Technical details of the different experiments are summarized in Table 3.2. In

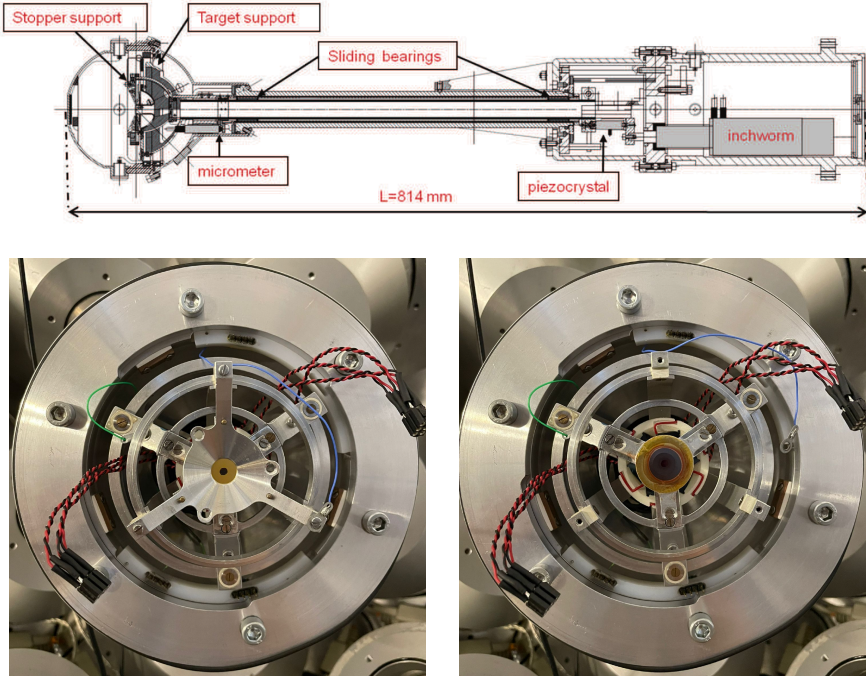


Figure 3.5: (Top) Schematic representation of the plunger device. (Bottom) Photographs of the plunger device installed in the scattering chamber of ROSPHERE, with (left) and without (right) the stopper material. The pictures were taken after the $^{112}\text{Sn}(^{18}\text{O}, ^{16}\text{O})^{114}\text{Sn}$ experiment performed in 2022, as the beam spot is clearly visible.

particular, heavy-ion transfer reactions at Coulomb barrier energies were employed, exchanging from one to four nucleons between beam and target nuclei. These reaction are well-suited to exploit beams from Tandem accelerator facilities, where high-intensity stable beams can be delivered for consecutive weeks.

As mentioned before, all the different reactions were performed at energies close to/slightly below the Coulomb barrier, which was calculated in the c.o.m frame using the Bass radius parametrization for interacting heavy ions [117]:

$$E_B = \frac{1.44Z_A Z_B}{\left[1.12 \left(M_A^{1/3} + M_B^{1/3}\right) - 0.94 \left(M_A^{-1/3} + M_B^{-1/3}\right) + 3\right]}, \quad (3.1)$$

where Z_A , M_A and Z_B , M_B are the atomic number and mass number of the beam A and the target B nuclei, respectively. This choice was done to suppress the fusion-evaporation channel, yet maintaining sizable cross sections for the transfer reaction. Additionally, the coincidence with events in the SORCERER charged particle detector provides a further suppression of the fusion-evaporation channel, since the angular distribution of evaporated particles is approximately isotropic in the laboratory frame for recoil velocity $\beta \ll 1$, while the differential cross section of beam-like particle after the heavy-ion transfer has a maximum at backward angles in the laboratory frame. Therefore, by placing SORCERER at backward angles and using particle- γ coincidences, the observation of transfer reaction products is favored. This is shown in Figure 3.6, where the SORCERER angular ranges in the DSAM and plunger configurations are compared

with the angular distributions in the laboratory frame for the $^{112}\text{Sn}(^{18}\text{O},^{16}\text{O})$ reaction products. The differential cross section was calculated with the FRESKO code [99] for the population of the 0_2^+ state in ^{114}Sn , at about 2 MeV of excitation energy.

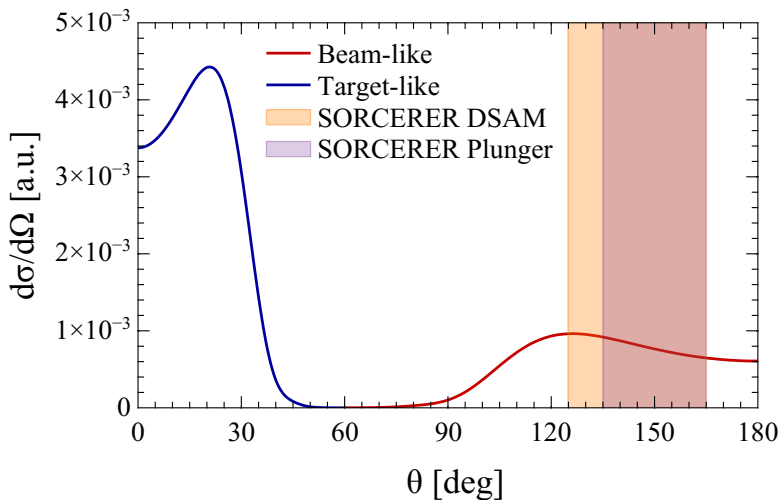


Figure 3.6: Angular distributions of beam-like (red) and target-like (blue) particles in the laboratory frame for the 0_2^+ state of ^{114}Sn , populated in the $^{112}\text{Sn}(^{18}\text{O},^{16}\text{O})$ reaction at 59 MeV. The shaded areas indicate the SORCERER angular coverage for the DSAM (orange) and Plunger (purple) configurations.

For the first two experiments (*Exp. 1* and *2*), the $2n$ transfer reaction was chosen due to its well-known effectiveness in the population of 0^+ states in a multitude of cases [13]. Subsequently, various reaction mechanisms were tested in *Exp. 3*, and the $\pm 1n$ reactions were finally chosen for the RDDS measurements of *Exp. 4*.

3.2.1 Experiment 1

In the first experiment of September 2021, ground state and excited states of ^{114}Sn were populated via the $^{112}\text{Sn}(^{18}\text{O},^{16}\text{O})$ transfer reaction. The ^{18}O beam was accelerated at 59 MeV by the 9 MV Tandem accelerator and impinged on a self-supporting 10 mg/cm^2 -thick ^{112}Sn target. An average beam intensity of 1 pA was delivered for the 7 days of data acquisition. After the transfer reaction, Oxygen-like reaction products were detected by SORCERER, which was mounted in the thick target configuration. The emitted γ -rays were measured by ROSPHERE, with a total of 23 HPGe detectors arranged on the 5 rings.

In the past, the two-neutron transfer (t,p) reaction was studied on a ^{112}Sn target at bombarding energies of 13 MeV (4.33 MeV/u) [74]. Among the identified states, the 0_2^+ and the 0_3^+ states were observed with intensities between 10% and 20 % with respect to the population of the ground state. Thus, by using the similar $(^{18}\text{O},^{16}\text{O})$ heavy-ion reaction, these states are expected to be populated with cross sections of the order of $100 \mu\text{b}$, as compared to the typical cross section of 1 mb for ground state to ground state transfer. For this reaction, the Q -value is 5.85 MeV.

According to (3. 1), the beam energy was set to be about 5% smaller than the Coulomb barrier, in order to suppress the cross section of fusion-evaporation channels. The latter

ID	Beam	Target	Channel	Product	E_{beam} [MeV]	E_B [MeV]	Q [MeV]	Method	Configuration
1	^{18}O	^{112}Sn	$+2n$	^{114}Sn	59	62.0	5.85	DSAM	23 HPGe
2	^{18}O	^{112}Sn	$+2n$	^{114}Sn	63	62.0	5.85	RDDS	20 HPGe
3	^{13}C	^{117}Sn	$-1n/+1n$	$^{116}\text{Sn}/^{118}\text{Sn}$	43	45.6	1.23/4.38	DSAM	18 HPGe
	^{13}C	^{119}Sn	$-1n/+1n$	$^{118}\text{Sn}/^{120}\text{Sn}$	43	45.4	1.69/4.16	DSAM	18 HPGe
	^{16}O	^{116}Cd	$+2p/+ \alpha$	$^{118}\text{Sn}/^{120}\text{Sn}$	56	58.7	-4.82/-2.35	DSAM	18 HPGe
4	^{13}C	^{115}Sn	$-1n/+1n$	$^{114}\text{Sn}/^{116}\text{Sn}$	44.2	45.9	0.63/4.62	RDDS	16 HPGe
	^{13}C	^{117}Sn	$-1n/+1n$	$^{116}\text{Sn}/^{118}\text{Sn}$	44.2	45.6	1.23/4.38	RDDS	16 HPGe
	^{13}C	^{119}Sn	$-1n/+1n$	$^{118}\text{Sn}/^{120}\text{Sn}$	44.2	45.4	1.69/4.16	RDDS	16 HPGe

Table 3.2: Experiments performed at IFIN-HH from 2021 to 2024, using single- and multi-nucleon transfer reactions near Coulomb barrier energies. The different experiments are labeled according to their ID number. The sign of the transferred nucleons refers to the target-like nuclei of interest, namely the Sn isotopes. E_{beam} and E_B are the beam energy and the Coulomb barrier energy in the laboratory frame, while Q is the Q -value of the corresponding reaction channel.

can be estimated using the PACE4 tool of the LISE++ package [118], which predicts a total cross section of 19 mb for the sub-barrier fusion process. The calculated yields of the residual nuclei and evaporated particles for the $^{18}\text{O} + ^{112}\text{Sn}$ fusion-evaporation reaction are shown in Table 3.3. Despite fusion-evaporation channels represent the main source of contamination in the γ -ray energy spectrum, transition from isomeric states of fusion residues could be of particular importance for the plunger measurement of *Exp. 2*. In fact, as it will be discussed in Chapter 5, long-lived states provide an alternative normalization to compare different plunger distances, hence the importance of identifying the evaporation residues.

Z	N	Residue	Channel	Yields [%]	σ [mb]
58	70	^{128}Ce	2n	13.2	2.5
58	69	^{127}Ce	3n	8.3	1.57
57	71	^{128}La	1p1n	12	2.27
57	70	^{127}La	1p2n	44	8.34
56	72	^{128}Ba	2p	0.79	0.149
56	71	^{127}Ba	2p1n	9.03	1.71
56	69	^{125}Ba	2p3n	5.38	1.02
56	68	^{124}Ba	2p4n	0.96	0.182
55	72	^{127}Cs	3p	0.17	0.0322
55	70	^{125}Cs	3p2n	1.72	0.325
55	69	^{124}Cs	3p3n	3.37	0.638
54	70	^{124}Xe	4p2n	0.13	0.0246
54	68	^{122}Xe	4p4n	0.63	0.119
54	67	^{121}Xe	4p5n	0.19	0.036

Table 3.3: PACE4 calculation of the yields and cross section for the $^{18}\text{O} + ^{112}\text{Sn}$ fusion-evaporation reaction at 59 MeV. Only channels with yield greater than 0.1% are reported.

The configuration of this experiment and, in particular, the usage of a 10-mg/cm²-thick target, were chosen to maximize the count rate of HPGe detectors and collect statistics for a detailed γ -ray spectroscopy study in view of a future plunger measurement. Moreover, lifetime information of short-lived states ($\tau \lesssim 3$ ps) could be obtained from the study of the line-shape of γ -ray peaks. Although the thick-target configuration is not ideal for DSA measurements due to the uncertainty on the reaction energy in the target, realistic Monte Carlo simulation are able to account for this effect, allowing for precise lifetime estimations.

3.2.2 Experiment 2

The experiment was performed in October 2022 and the beam time lasted for 20 days. The same $^{112}\text{Sn}(^{18}\text{O}, ^{16}\text{O})$ transfer of *Exp. 1* was used to populate the excited 0^+ states of ^{112}Sn and measure their lifetime with the RDDS method. Due to the poor mechanical

resistance and the low melting point of metallic tin, the plunger target was obtained by depositing 1 mg/cm^2 of ^{112}Sn on a 2.7 mg/cm^2 Au foil. To account for the additional material of the Au backing, the beam energy was raised to 63 MeV, with an average intensity of 0.7 pA.

For this experiment, ROSPHERE comprised 20 HPGe detectors, which were placed on the 37° , 70° , 110° and 143° rings, and SORCERER was operated in the plunger configuration. As the beam time progressed, a significant degradation of the SORCERER cells was observed, resulting in a gradual reduction of the amplitude of the signals as a function of time. Moreover, despite the beam current was kept under 1 pA, beam-like particles scattered by the Au layer incremented the counting rate of the photodiodes and introduced additional noise to the recorded signals.

In order to perform the lifetime measurement, plunger distances had to be chosen according to the reaction kinematics and the expected lifetimes of the states of interest. For the present case, recoil velocities of about $6 \mu\text{m/ps}$ ($\beta \sim 2\%$) and lifetimes in the 1 ps to few tens of ps range were expected. Therefore, the following 4 plunger distances were chosen: 14.5, 22.5, 42.5 and $74.5 \mu\text{m}$, corresponding to time of flights of 2.6, 4.1, 7.8 and 13.6 ps, respectively.

3.2.3 Experiment 3

In the 10 days of beam time of 2023, several reactions were employed to perform a low-spin spectroscopy of $^{116,118,120}\text{Sn}$, aiming at exploring the selectivity in the population of the different states and obtaining lifetime information on short-lived states. In particular, the following reactions were studied:

1. $^{13}\text{C} + ^{117}\text{Sn}$ at 43 MeV, neutron stripping and pickup channels;
2. $^{13}\text{C} + ^{119}\text{Sn}$ at 43 MeV, neutron stripping and pickup channels;
3. $^{16}\text{O} + ^{116}\text{Cd}$ at 56 MeV, two-proton and α transfer channels.

The ^{13}C and ^{16}O beams were delivered with intensities of about 1 pA, and impinged on ^{117}Sn , ^{119}Sn and ^{116}Cd targets with thicknesses of about 14 mg/cm^2 each. The configuration of the ROSPHERE array consisted in 18 HPGe detectors, 5 of which mounted on both the 37° and 143° rings, 1 on both the 70° and 90° rings and 4 on the 110° ring. SORCERER was also used in its DSAM configuration, providing an efficient tagging of the beam-like charged particle.

As it was done in *Exp. 1*, all the reactions were performed at energies just below the Coulomb barrier to suppress fusion-evaporation channels. Moreover, these reactions had favorable Q -values ranging between one and several MeVs, thus ensuring an optimal population of states up to few MeV of excitation energy.

The aforementioned reactions present similar features to (d,p), (t,p), (^3He ,n) and (^6Li ,d), respectively, therefore, they could provide information on the composition of the state wave functions, by comparing the relative cross sections for the populations of different states in different reactions. Moreover, this study was also meant to provide the basis for possible lifetime measurements with the plunger technique. Such an experiment was performed one year later and a detailed description of it is reported in the following Section (3.2.4).

3.2.4 Experiment 4

In this experiment, neutron stripping and pick-up reaction were used to populate excited states in $^{114-120}\text{Sn}$ and perform lifetime measurements with the RDDS technique. The

^{13}C beam was accelerated by the 9 MV Tandem facility at a kinetic energy of 44.2 MeV, with an average beam current of 1 pA. Three different targets were used during the experiment, each of which was irradiated for about 10 days. To improve their mechanical resistance and heat dissipation, tantalum backings were used as supports. The ^{119}Sn and ^{117}Sn targets had thicknesses of 1.0 and 0.8 mg/cm², respectively, both of them being deposited on a 1.1 mg/cm² Ta backing. On the other hand, the 0.8 mg/cm² ^{115}Sn target was deposited on a 1.2 mg/cm² Ta layer, however, an enrichment of only 50.1% was certificated by the producer (as compared to enrichment above 95% for the other cases). Tantalum stoppers of about 10 mg/cm² were produced and stretched on the stopper support. In the target chamber, the target was oriented with the Ta backing on the upstream side, facing the beam. This is the reason why the beam energy was increased at 44.2 MeV with respect to the 43 MeV used in the thick target experiment, given the additional energy losses in the Ta layer.

During the data acquisition with the first target of ^{119}Sn , only 16 HPGe detectors were initially mounted on ROSPHERE, with 5 detectors mounted on the 37°, 90° and 143° rings each. After the target change, two more detectors were added to the 110° ring, where a single HPGe was initially installed. SORCERER was operated in the plunger configuration with 5 working cells. A Si cell was excluded from the acquisition due to a broken preamplifier.

To perform the lifetime measurement, different set of plunger distances were chosen based on the expected lifetimes and the recoil velocities. Depending on the reaction, lifetimes in the range of 1-200 ps were expected, and recoil velocities around 1.3 % were estimated, with minor variations due to the different reaction kinematics. Distances of 13.5, 35.5, 55.5, 75.5, 95.5, 200.5, 300.5 μm were used for the ^{119}Sn target, corresponding to time of flights between 4 and 90 ps. To measure longer lifetimes with the ^{117}Sn , distances of 120, 200, 350, 500, 750 μm were chosen. In this case, the corresponding time of flight range is approximately 30-220 ps. Lastly, the ^{115}Sn target was mainly dedicated to probe short-lived lifetimes in ^{116}Sn that were not covered in the previous case. However, due to a progressive deformation of the target, only the first distance at 7.5 μm could be used.

Preprocessing of IFIN-HH data

In this Chapter, details about the preliminary data processing and analysis, such as detectors calibrations and timing alignments, are reported. Since similar procedures have been applied for all the experiments performed at IFIN-HH (summarized in Table 3.2), the preliminary analysis of *Exp. 4* is here described as a representative example. Deviations between the analyses of the different experiments will be noted and highlighted.

Section 4.1 will describe the data treatment procedures, while Section 4.2 will focus on calibrations, time alignment and performance of the installed detectors. Finally, the implementation of Monte Carlo simulations for the ROSPHERE array is described in Section 4.4.

4.1 Data sorting

The data processing of ROSPHERE-SORCERER experiments involves two different softwares: the custom made `builder_sphere` code and the GASPware package [119]. Raw data files collected with the ROSPHERE and SORCERER DAQ (data acquisition system) in a triggerless mode were written to disk in the Cap'n Proto data format [114], which provides a fast and reliable encoding protocol. Firstly, the `builder_sphere` code was used to read the Cap'n Proto files and perform a preliminary time alignment of all the detectors in the DAQ, using as a reference the timing signal of one of the detectors (generally a BGO or LaBr due to their temporal resolution). Secondly, it performed the event-building procedure by grouping together the time-ordered timestamps which lied in a given coincidence window. At this level, the trigger, the coincidence window and the lookup table (LUT) are set as inputs for the code. In this work, HPGe events were used as triggers and a $1\ \mu\text{s}$ window was chosen. Finally, the program created output files in the GASP format, to be subsequently analyzed via the GSORT program of the GASPware package.

The last step of the data processing consisted in the sorting of GASP files into histograms and multi-dimensional matrices to perform the analysis. This is done after applying energy calibrations, fine-tuned time alignments, coincidence gates and multiplicity filters by means of dedicated keywords in a GSORT setup file. An example of the setup file used for this purpose can be found in Appendix A.

4.2 ROSPHERE data processing

The analysis of the collected datasets, which is meant to provide lifetime measurements using techniques based on the Doppler shift effect, relies on precise energy calibrations of the HPGe detectors, good energy resolution and high P/T ratio. In the following

sections, a description of the calibration, time alignment and Compton suppression procedures of the ROSPHERE detectors is presented.

4.2.1 Source calibrations

The first step of the calibration procedure involves the use of γ -ray sources to convert the amplitude of the collected signals from digitizer channels into energy, thus allowing to sum events observed in the entire array. Sources of ^{56}Co and ^{152}Eu were employed, covering the entire energy range of interest, from 100 keV to 3 MeV.

In *Exp. 1-3*, to correct non-linear behavior of the ADC modules over the range of interest, two different calibration regions were defined, below and above 1.4 MeV. Cubic polynomials have been applied in both of the regions, using ^{152}Eu peaks in the first one and ^{56}Co peaks in the second one. To ensure a smooth merging between the two regions, the 1238-keV peak of ^{56}Co and the 1408-keV peak of ^{152}Eu were shared and used in both of them. Additionally, a 100-keV overlap window was chosen between the two regions, in proximity of the 1.4 MeV endpoint, inside which a random channel was selected for each detector and used as the separation point between the two applied calibration.

On the other hand, a single cubic polynomial was sufficient to calibrate the detectors of *Exp. 4* from 100-keV to 2.5 MeV. In this case, the most intense peaks of both ^{56}Co and ^{152}Eu were used. The results of this procedure are shown in Figure 4.1, where the calibrated energies of all the crystals of *Exp. 4* are presented in three different energy ranges.

In Figure 4.2, the differences between the calibrated energies of the ROSPHERE array and the tabulated energies of the calibration peaks, namely the residuals, are presented. Residuals smaller than 0.1 keV are observed in the range of interest, confirming the positive outcome of the procedure. Similar results were obtained in all the ROSPHERE experiments, with discrepancies generally comprised between ± 0.3 keV.

4.2.2 Energy resolution

The energy resolution of the spectrometer has been determined up to 2.6 MeV, using both of the calibration sources. The FWHM curve as a function of the γ ray energy can be described by the following empirical formula:

$$FWHM = \sqrt{A + BE_\gamma + CE_\gamma^2}, \quad (4.1)$$

where E_γ is the γ -ray energy in keV, and A , B and C are free parameters. Experimental points and the fitting function are shown in Figure 4.3, with the resulting fit parameters being $A = 2.83$, $B = 1.83 \times 10^{-3}$ and $C = 5.46 \times 10^{-7}$.

A total resolution (FWHM) of about 2.4 keV is obtained at 1.3 MeV, which is consistent with the average value of 2.5 keV of the single detectors at 1.3 MeV. Compared to nominal resolutions of the detectors, a 10-20% higher FWHM is obtained, depending on the detector type. In this case, no significant differences have been observed between the different experiments.

4.2.3 Efficiency

Since the sources were not placed in the center of the chamber, only the relative efficiency of ROSPHERE can be determined. Data points from ^{152}Eu and ^{56}Co were normalized

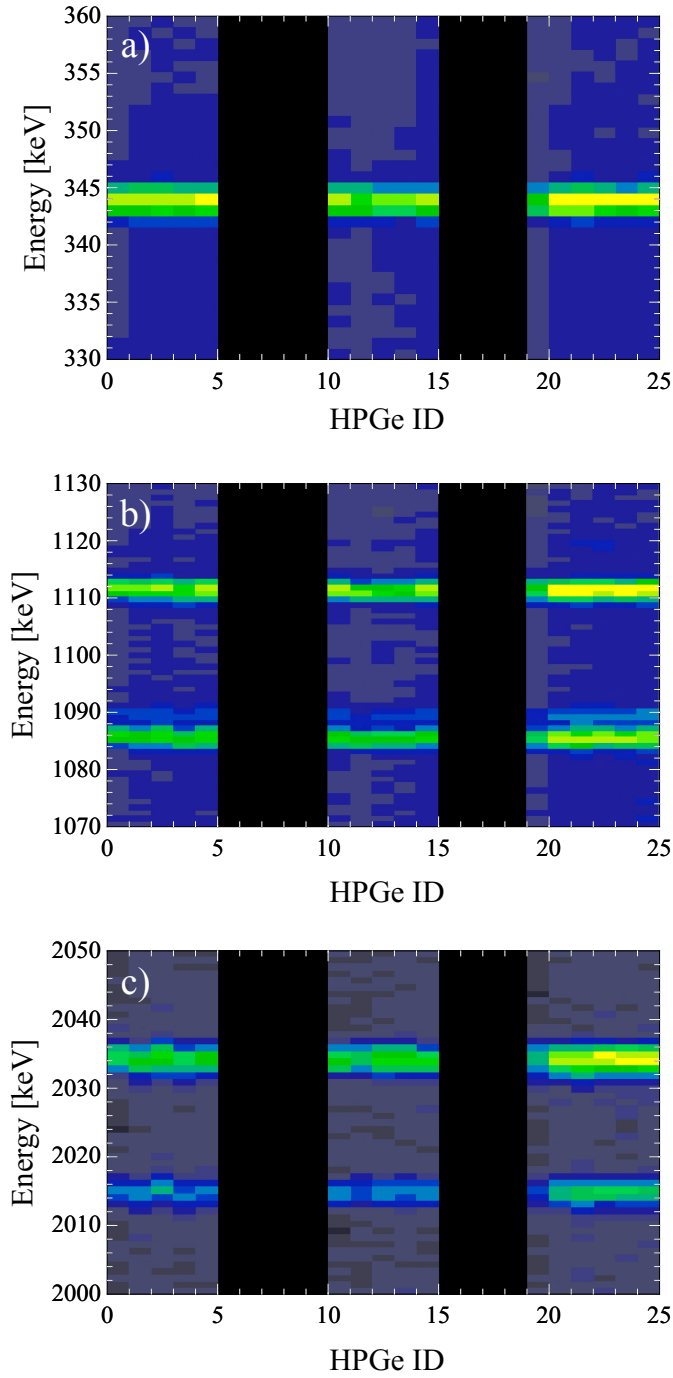


Figure 4.1: HPGe energy versus detector number obtained using a single cubic polynomial for the calibration of the detectors in *Exp. 4*, with 16 working detectors. Panel a) and b) show ^{152}Eu peaks at 344, 1085 and 1112 keV, while Panel c) shows the 2015- and 2035-keV peaks of ^{56}Co .

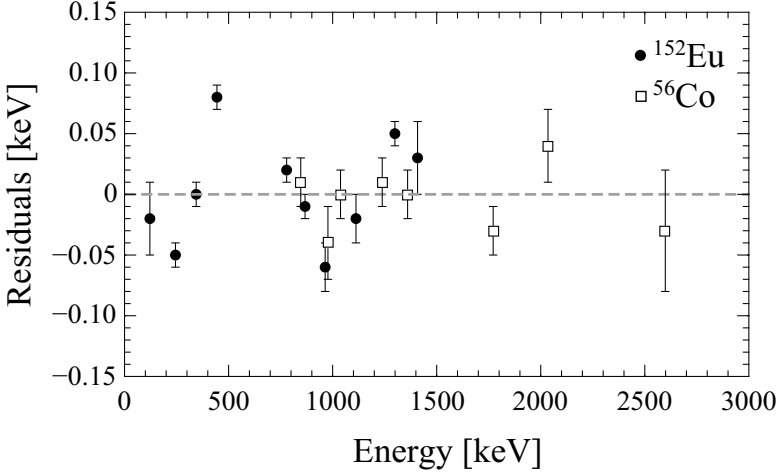


Figure 4.2: Residuals of the measured energy in the ROSPHERE array with respect to tabulated values from the NNDC database [68]. Black circles and white squares indicate ^{152}Eu and ^{56}Co points, respectively. Data collected during *Exp. 4*.

using the ^{152}Eu peak at 867 keV as a reference, and subsequently fitted with the Radware function [120]:

$$\epsilon(E_\gamma) = H \exp \left\{ \left[\left(A + B \ln \left(\frac{E_\gamma}{100} \right) + C \ln \left(\frac{E_\gamma}{100} \right)^2 \right)^{-G} + \left(D + E \ln \left(\frac{E_\gamma}{1000} \right) + F \ln \left(\frac{E_\gamma}{1000} \right)^2 \right)^{-G} \right]^{-1/G} \right\}, \quad (4.2)$$

where E_γ is the γ -ray energy in keV. The coefficients A, B and C describe the efficiency at low energies ($E_\gamma < 100$ keV), D, E and F at high energies ($E_\gamma > 100$ keV), and G is an interaction parameter between the two regions that regulates the sharpness of the turnover region. Figure 4.4 shows the result of the fitting procedure, and the obtained parameters are reported in Table 4.1, with C and G set to the fixed values 0 and 12, respectively.

4.2.4 Stability and in-beam recalibration

During the beam time of all the IFIN-HH experiments, significant gain drifts were observed over time. This behavior is generally correlated with day/night temperature fluctuations, which affect the performance of the digitizer modules. In order to correct it, a run-by-run calibration procedure was applied on each detector. Transitions of backing and stopper materials, such as gold (Au) and tantalum (Ta), and secondary reaction channels were used. In the case of *Exp. 4*, the centroids of the 165- and 301-keV peaks of ^{181}Ta and the 1368- and 2754-keV peaks of ^{24}Mg (produced in the $^{13}\text{C} + ^{12}\text{C}$ fusion-evaporation reaction) were interpolated by a linear polynomial, thus providing additional calibration coefficients for each run. Figure 4.5 shows the spectrum of crystal 24, before (left) and after (right) the in-beam recalibration. This procedure allowed us to

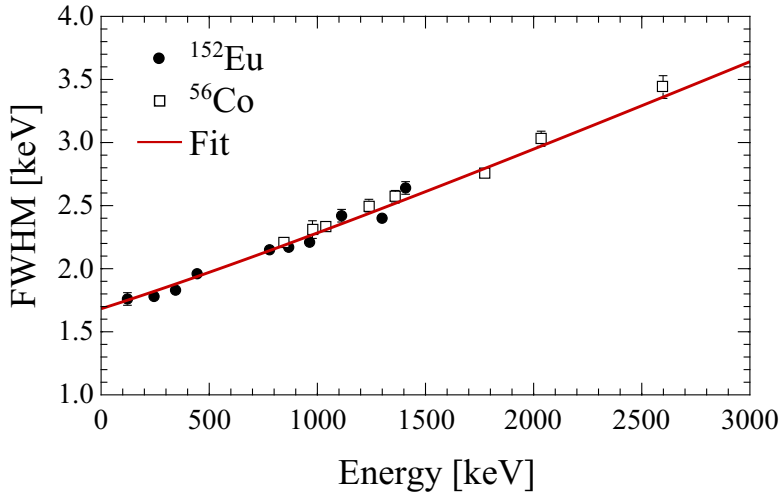


Figure 4.3: Resolution of the ROSPHERE array as a function of the γ -ray energy observed during *Exp. 4*. Black circles and white squares indicate ^{152}Eu and ^{56}Co data points, respectively. The red line indicates the fitting function of Eq. (4. 1).

Parameter	Value
A	1.93(150)
B	-0.37(11)
C	0
D	-0.67(151)
E	0.80(7)
F	-0.10(10)
G	12
H	0.46(70)

Table 4.1: Efficiency parameters obtained from the fit of Eq. (4. 2). Large errors on the vales of some parameters are ascribable to the lack of experimental points below 100 keV.

fully exploit the performances of the array, with a negligible increase in the FWHM with respect to source estimates.

4.2.5 Time coincidence of HPGe detectors

In order to build γ - γ and γ -particle coincidence matrices, a precise time alignment of all the detectors is required. As mentioned in Section 4.1, a first time alignment of the detectors was performed at the level of the event builder and, since the signal of the HPGe detectors was chosen as the START of the trigger window for coincidence events, no further optimization was implemented for this type of detectors.

To select prompt events of HPGe detectors, a time coincidence window is set on events whose time difference ΔT lies within a chosen value. This is shown in Figure

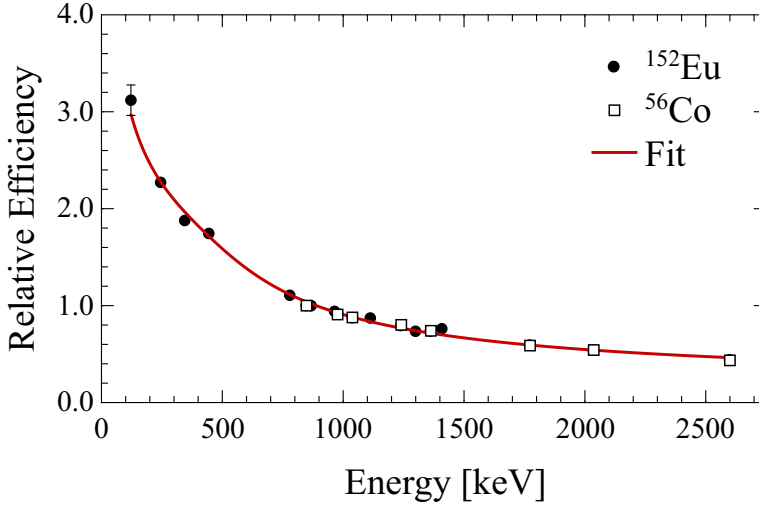


Figure 4.4: Relative efficiency of the ROSPHERE array during *Exp. 4*. Experimental points correspond to peaks of the ^{152}Eu (black) and ^{56}Co (white) sources, which were fitted with the Radware function of Eq. (4.2).

4.6, where the coincidence window (highlighted by a red shaded area) was placed on the main peak appearing at $|\Delta T| = 0$. Due to the poor time resolution of the HPGe detectors, a large width of 400 ns had to be chosen, at the cost of including background contribution that could eventually be subtracted.

4.2.6 Compton suppression

In a typical γ -ray spectrum, a large fraction of the collected data correspond to Compton scattering events, which is the dominant interaction process in HPGe detectors for γ rays between few hundreds of keV and several MeV. Such events are characterized by a partial deposition of the total photon energy, thus polluting the spectrum and reducing the P/T ratio. To reduce the background contribution due to Compton scattering, BGO crystals were installed around each HPGe detector of ROSPHERE. BGOs were operated as veto detectors to discard events with interactions registered in both HPGe and BGO detectors. A time window of about 650 ns was defined to perform the Compton suppression, as shown in the bottom panel of Figure 4.7. Additionally, the time difference spectrum between each BGO and all the HPGe signals is shown in the top panel of Figure 4.7, with the difference being defined as $\Delta T = T_{\text{BGO}} - T_{\text{HPGe}}$ plus an arbitrarily chosen offset. Before setting the time window for the Compton suppression, a fine-tuning of the time alignment of the BGO timing was also performed by manually adjusting the time differences between all the detectors with respect to a reference one.

In figure 4.8, the result of the Compton suppression procedure is highlighted by comparing ^{152}Eu spectra with and without the rejection of Compton scattered events. A reduction of about a factor of 5 can be observed in the background counts between 100 and 1200 keV, corresponding to a significant improvement in the P/T ratio.

Even though some HPGe detectors were not available (only 16 HPGe were installed during *Exp. 4*), all the 25 BGOs were mounted and included in the acquisition. Indeed, BGO detectors were also used as a multiplicity filter, providing additional efficiency to

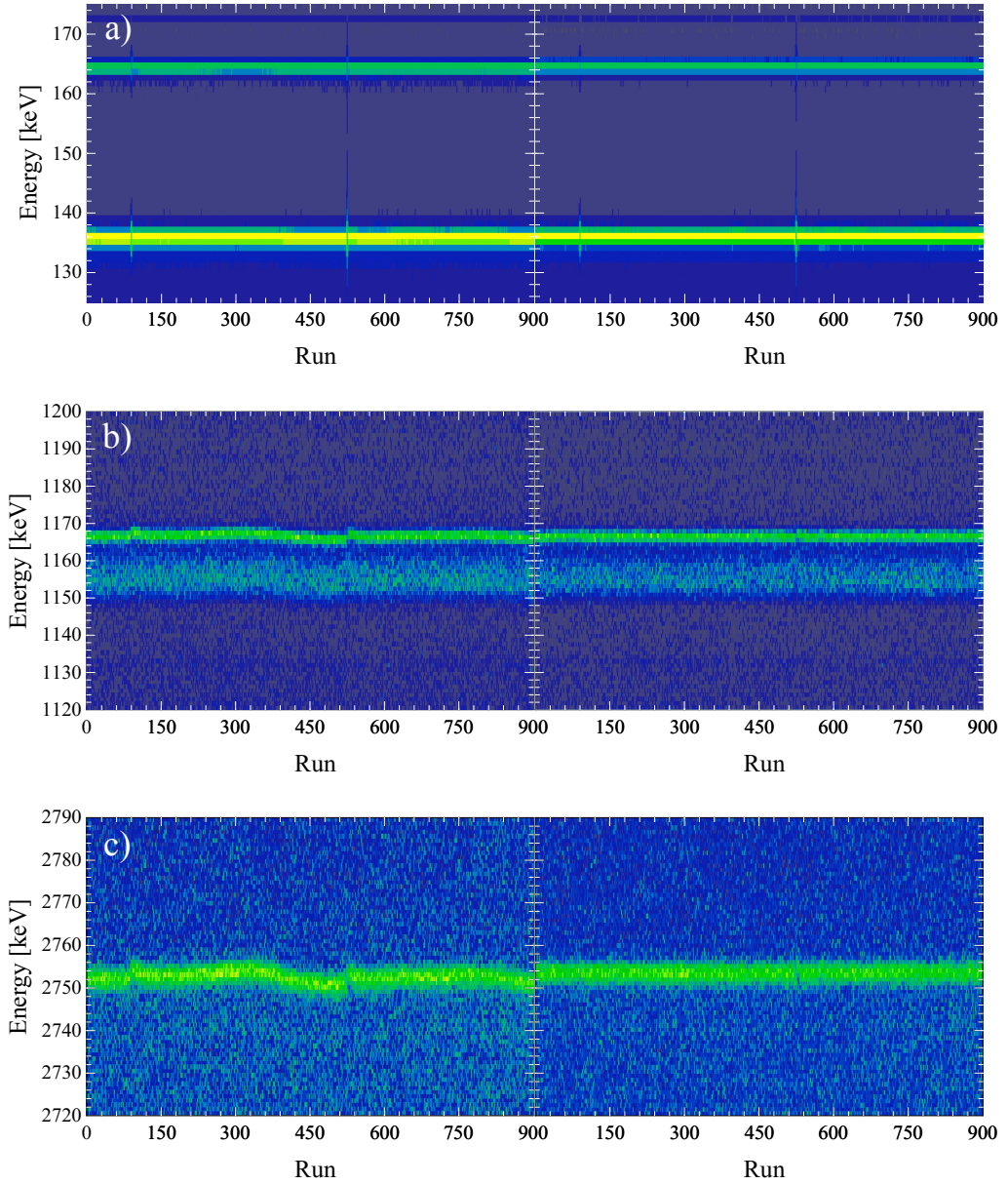


Figure 4.5: Energy spectra of the HPGe crystal 24 before (left) and after (right) the in-beam recalibration procedure, corresponding to data acquired during the run with a plunger distance of 55 μm with respect to the target. Panel a) shows the 131- and 165-keV peaks of ^{181}Ta , while Panel b) and c) show the 1171-keV peak of ^{120}Sn and the 2754-keV peak of ^{24}Mg , respectively. Since the 1171-keV state of ^{120}Sn has a lifetime $\tau = 923(17)$ ss, both the stopped and the shifted component are clearly visible.

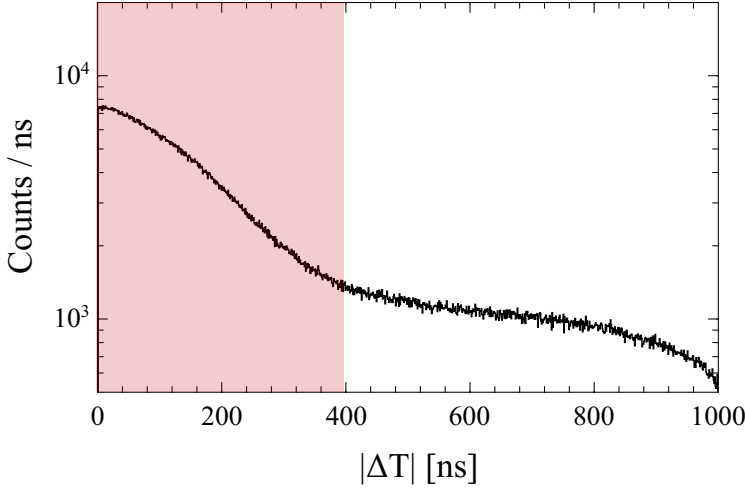


Figure 4.6: Absolute time difference spectrum between subsequent hits of HPGe detectors recorded in the ROSPHERE spectrometer. The red shaded area indicates the time coincidence window used to build γ - γ matrices and determine the γ multiplicity of an event. Data collected during *Exp. 4*.

count the total number of emitted γ rays. To this purpose, the same time window of the Compton suppression was employed, yet used in coincidence mode.

4.3 SORCERER data processing

In the following sections, the pre-analysis procedures applied to the SORCERER array are described.

4.3.1 Gain alignment

Due to its large angular coverage of 125° - 165° , the charged particles detected by SORCERER photodiodes have a wide spread in energy. Therefore, the typical in-beam spectrum of the Si cell is composed of structures ("bumps") that span a large number of channels. For this reason, no α source is used to calibrate the detectors in energy, but only an alignment procedure is performed. As shown in Figure 4.9, the spectra of the different cells are aligned using the centroid channels of the bumps as a reference. During the data acquisition of *Exp. 4*, the cell 4 was disabled due to a malfunctioning of its preamplifier, while cell 2 showed a significantly worse resolution compared to the others.

4.3.2 Time coincidence with SORCERER

Using the combination of γ and particle data, it is possible to improve the P/T ratio of the spectra and remove background from secondary reaction channels (such as fusion-evaporation reactions) and random coincidences. A precise time alignment of the signals of the photodiodes is necessary to select the prompt coincidence peak in the time difference spectrum between HPGe detectors and SORCERER cells. Such a time difference,

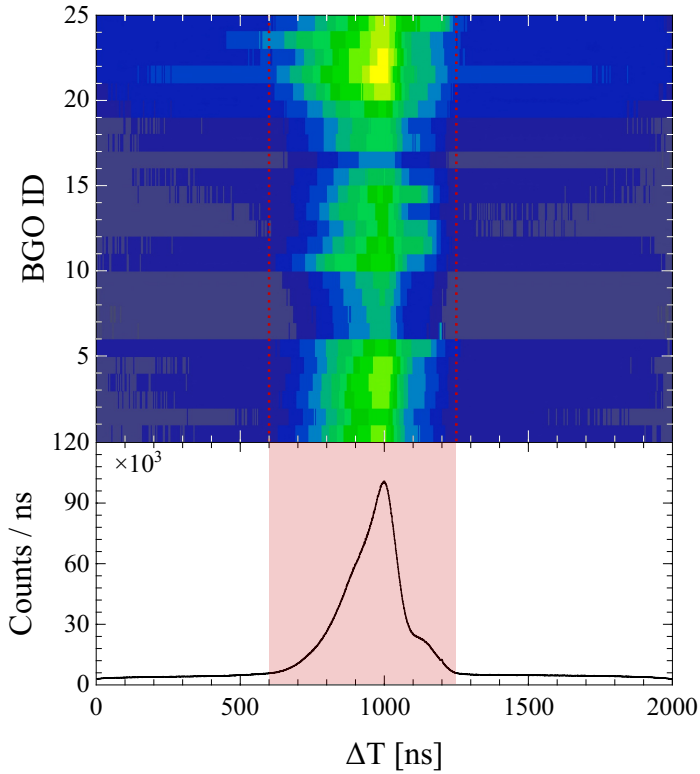


Figure 4.7: Time difference spectra between γ -ray interactions observed in HPGe crystals and BGO shields of ROSPHERE, shown separately for each BGO detector. On the bottom, the total projection on the x axis is shown. The red shaded area indicates the time window used to perform the Compton suppression procedure. Data collected during *Exp. 4*.

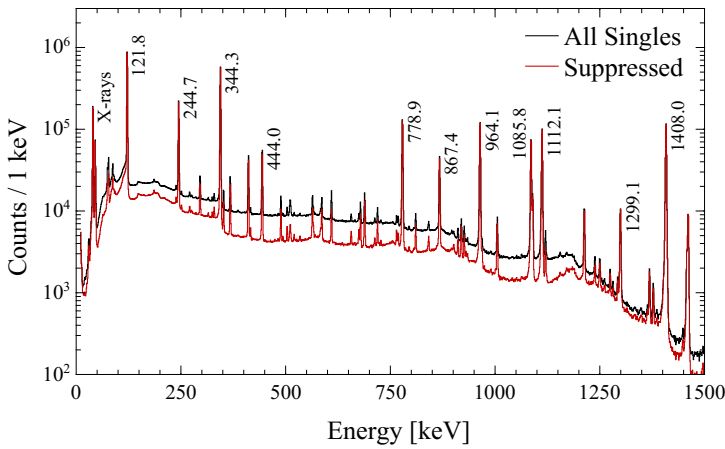


Figure 4.8: Comparison of the ^{152}Eu spectrum observed in ROSPHERE with all singles events (black line) and with the Compton suppression (red line). A reduction of the Compton-induced background by a factor of ~ 5 is obtained with the suppression procedure. Energies of the most intense transitions are labeled in keV units. Data collected during *Exp. 4*.

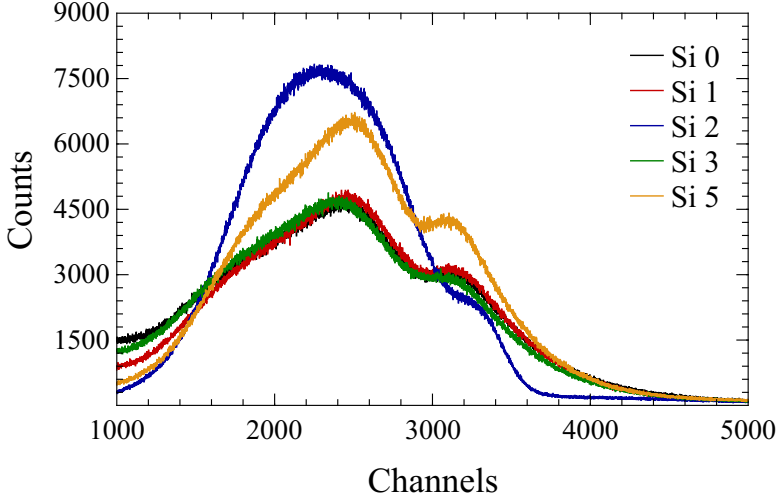


Figure 4.9: Energy spectra of SORCERER detectors after the alignment procedure. The cell 4 was excluded from the data acquisition due to a malfunctioning of its preamplifier. Data collected during *Exp. 4*.

defined as $\Delta T = T_{Si} - T_{HPGe}$ plus an additional offset of 1000 ns, is depicted in Figure 4.10, where the spectrum of each cell (top panel) and the total projection (bottom panel) are shown.

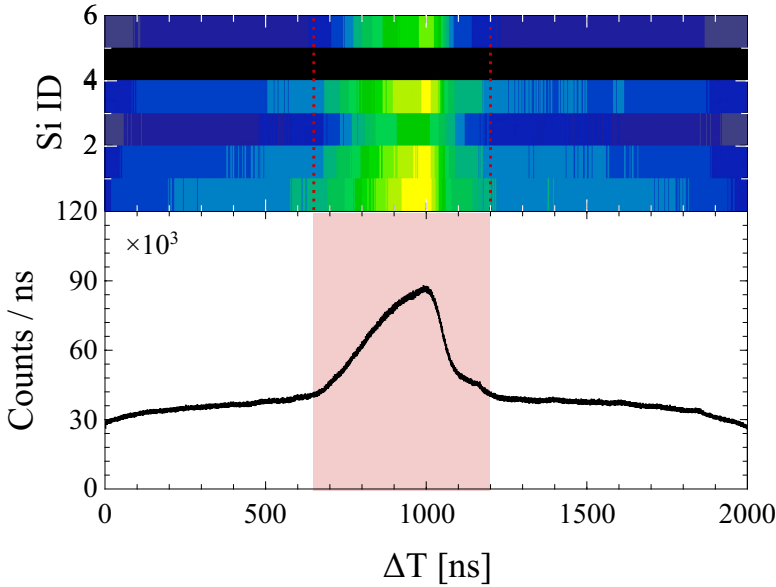


Figure 4.10: Time difference spectra between γ -ray interactions observed in ROSPHERE crystals and particles detected in SORCERER detectors, shown separately for each Si cell. On the bottom, the total projection on the x axis is shown. The red shaded area indicates the time window used to define γ -particle coincidence events. Data collected during *Exp. 4*.

A fine-tuning of the alignment was performed by applying additional offsets to the time differences of the signals. Finally, a time window of 550 ns was set to designate γ -particle coincidence events, ranging from 650 ns to 1200 ns. This window is highlighted as a red shaded area in Figure 4.10.

4.4 Simulations

In order to measure the lifetime of excited states, a newly developed simulation tool for ROSPHERE was implemented and included into the Agata Simulation Code framework [121, 122]. Built on the Geant4 toolkit in its 4.9.6 version [123], the simulation code is able to reconstruct complex geometries of detectors in a virtual three-dimensional space, using Computer Assisted Drawings (CADs). In particular, CAD drawings of ROSPHERE HPGe detectors and BGO shields were included. Once the environment is built, realistic Monte Carlo simulations can be performed to reproduce experimental events, under the assumption that only a single type of reaction takes place in the target. Three main processes are considered in the simulation workflow:

- the reaction mechanism and its kinematics;
- the slowing down of ions in the target material;
- the γ -ray decay of the reaction products.

Each of these steps can be tailored to ensure the best agreement between experiment and simulations. For this purpose, a series of parameters and input files can be included, such as:

- the energy loss tables;
- the thicknesses of target, backing and degrader (if present);
- the interaction profile in the target;
- the angular distribution of the reaction products.

Additionally, background gamma rays and peaks coming from parasitic reactions can be added at the simulation level. The background spectrum is generated by sampling exponential probability distribution or, as an alternative, by considering discrete energies.

The output of the simulations is enclosed in a text file containing events with information on the velocity $\vec{v}$ of the reaction products and all the interaction of the emitted γ rays. Finally, this output is fed to a custom-made algorithm (`readGamma`) that applies smearing on energies, performs the Compton suppression and converts the file into a ROOT TTree [124].

In the following sections, a detailed discussion on the inputs used in this simulation will be given.

4.4.1 Energy losses

As described in Chapter 2, the DSAM strongly relies on the Doppler shift of measured γ -ray energies to determine the lifetime of an excited state, hence, on the knowledge of the velocity of the ion at the moment of the γ -ray emission. For this reason, the accurate knowledge of the stopping powers is of paramount importance. In literature, empirical parametrization of the stopping power of ions can be found [125, 126]. These tables,

tuned on the available experimental data [127], are based on different assumption and might differ in certain energy regions.

In the simulation, Geant4 tables of the stopping power of C in Sn (beam in target) and of Sn in Sn (target-like recoil in target) have been used. As in our thick-target experiments the nuclei are completely stopped in the target, the nuclear stopping power is also considered in addition to the electronic stopping power [128]. In fact, the nuclear component is negligible at high ion velocities, while it becomes significant when the ion decelerates and the probability of interacting with the nuclei of the material is higher. In Figure 4.11, a comparison between experimental data and calculations of the stopping power of C in Sn is shown. A good agreement is found between the tables used by

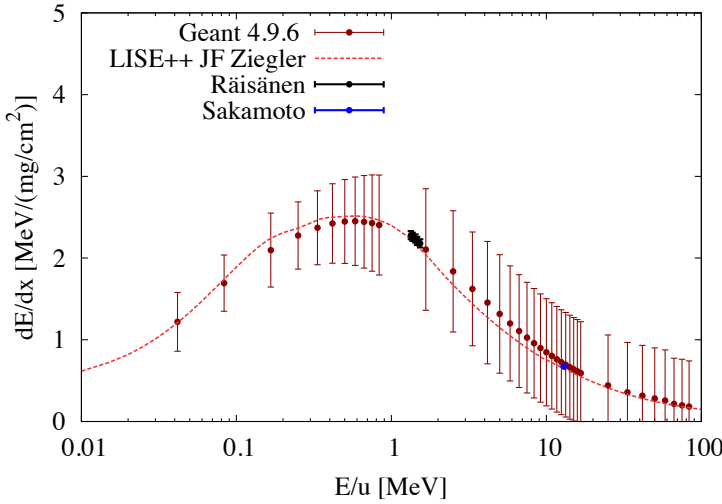


Figure 4.11: Stopping power as a function of energy for C in Sn. Dark red points and bars indicate the energy loss and the straggling calculated by Geant4.9.6 [123], while the red dashed line refers to the energy loss table used in LISE++ [118]. Black and gray points indicate experimental data from Refs. [129, 130].

Geant4 and LISE++. In the range of interest around 3 MeV/u, the straggling calculated by Geant4 dominates over the minor deviations observed between the two curves. Both parametrizations are able to reproduce the experimental data points.

For Sn on Sn, no experimental data is available. Stopping powers provided by Geant4 and LISE++ are reported in Figure 4.12. Despite some discrepancies above 0.3 MeV/u, the two tables are compatible within the straggling of Geant4 in the energy range of about 0.1 MeV/u considered for this experiment.

Since the DSAM techniques critically depends on the knowledge of stopping powers, imprecise tables might lead to systematic uncertainties in the lifetime measurement. For Sn in Sn, deviations up to 15% were observed between LISE++ and Geant4 around 0.2 MeV/u (but much smaller below 0.1 MeV/u). Yet, in the energy range of interest of 0.05-0.10 MeV/u, Geant4 also calculates an energy straggling between 20 to 40%, which significantly exceeds the discrepancies. Thus, one can argue that a possible source of systematic uncertainty is already integrated in the simulation process, without the need of an additional error term to be added after the minimization procedure.

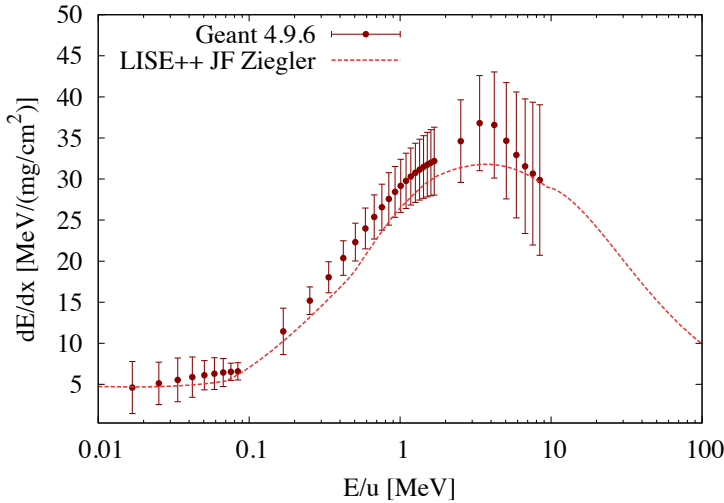


Figure 4.12: Stopping power as a function of energy for Sn in Sn. Red points and bars indicate the energy loss and the straggling calculated by Geant4.9.6 [123], while the blue line refers to the energy loss table used in LISE++ [118]. The region of interest for the experiments of this thesis is from 0.05 to 0.10 MeV/u.

4.4.2 Reaction profile

During the data taking of *Exp. 3*, thick targets of about 14 mg/cm^2 were used. This choice was mostly driven by the need of accumulating statistics for γ -ray spectroscopy, instead of focusing on lifetime measurements. The disadvantage that comes with it is a large uncertainty on the reaction point inside the target, hence on reaction energy and velocity of the final products. To overcome this issue, an interaction profile was calculated for each reaction and used as input for the simulations. At first, the residual beam energy in the target material was evaluated at steps of 0.5 mg/cm^2 . Then, for each of those reaction energies, the total cross section was calculated with the GRAZING code. Lastly, the interaction profile was obtained by correlating the interaction depth with the normalized cross section. The profile obtained for the $^{119}\text{Sn}(^{13}\text{C}, ^{12}\text{C})^{120}\text{Sn}$ reaction is shown in Figure 4.13. Fresco calculations were also tested, but no significant differences were noticed. As expected, the profiles for the other channels remained similar.

Based on Figure 4.13, it is clear that the transfer reaction takes place in the first half of the target, when the beam energy is close to the Coulomb barrier. LISE++ calculations show that, if the reaction happens at the beginning of the target, Sn recoils will have kinetic energies up to 11 MeV. The corresponding range in the Sn target material will be about 2.6 mg/cm^2 and the average stopping time will be of the order of 1 ps. Assuming similar stopping times for reactions at lower beam energies, we find that all the ions will be stopped in the target and states with half-lives longer than 1 ps will mostly emit γ -rays at rest. This hypothesis is corroborated by the absence of lineshapes for γ -rays depopulating long-lived states, i.e. with $T_{1/2} \gtrsim 3 \text{ ps}$ (see Figure 5.6 of Section 5.1).

4.4.3 Angular distributions

In transfer reactions, angular distributions have a strong influence on the velocity distribution of the reaction products. Whenever possible, the usage of measured distributions

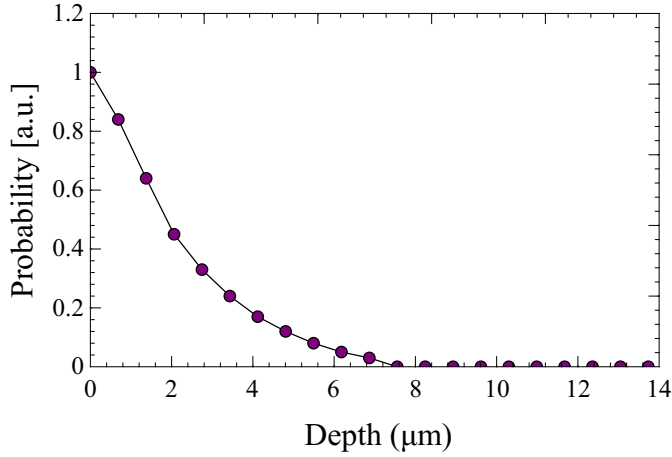


Figure 4.13: Reaction profile as a function of the depth in the target. The probability of the reaction is normalized to 1 at the beginning of the target.

is recommended, as they account for specific features of each reactions. Else, calculated angular distribution are considered.

In the experiments performed at IFIN-HH, no direct measurement of the angular distribution was possible, due to the lack of segmentation of the SORCERER photodiodes. Thus, GRAZING and FRESCO calculations were performed [94, 99]. Although the GRAZING program does not allow to selectively populate single excited states in the reaction products, an excitation energy range can be designated through Wilczynski plots [131], which correlate the Q -value (or the excitation energy) and the scattering angle in the center-of-mass frame. Excitation energies from 1 to 3 MeV were considered.

Between GRAZING and FRESCO calculations, the former were chosen, as they were able to better reproduce the average velocity of Sn recoils observed in the same reactions with the plunger measurement (see discussion in Section 5.2). Indeed, FRESCO distributions led to overestimate the average velocity of ^{116}Sn and ^{118}Sn products in the $-1n$ channel by more than 10%, and thus were discarded. The angular distributions in the center-of-mass frame, which were used as input for the simulations, are shown in Figure 4.14. Among those, small variations are observed. This was expected based on the similar masses and the transfer of a single neutron involved in the reactions.

Angular distribution can also serve the purpose of indirectly including SORCERER in the simulation. In fact, SORCERER is not used to reconstruct the kinematics, but to select an angular range that favors transfer reactions. This can be effectively reproduced by restricting the angular distribution to center-of-mass angles which correspond to the angular coverage of SORCERER in the laboratory frame.

4.4.4 Energy smearing and Compton suppression

A realistic reproduction of the lineshape of γ -ray peaks is crucial for lifetime measurements with the DSAM technique. As a consequence, the response function of the detectors has to be sampled and added in the simulations to avoid the inclusion of systematic errors. Simulated photon energies listed in the output file are then read and analyzed by the `readGamma` executable, which applies a smearing function to reproduce experimental response of ROSPHERE array. Specifically, the smearing is applied by sampling

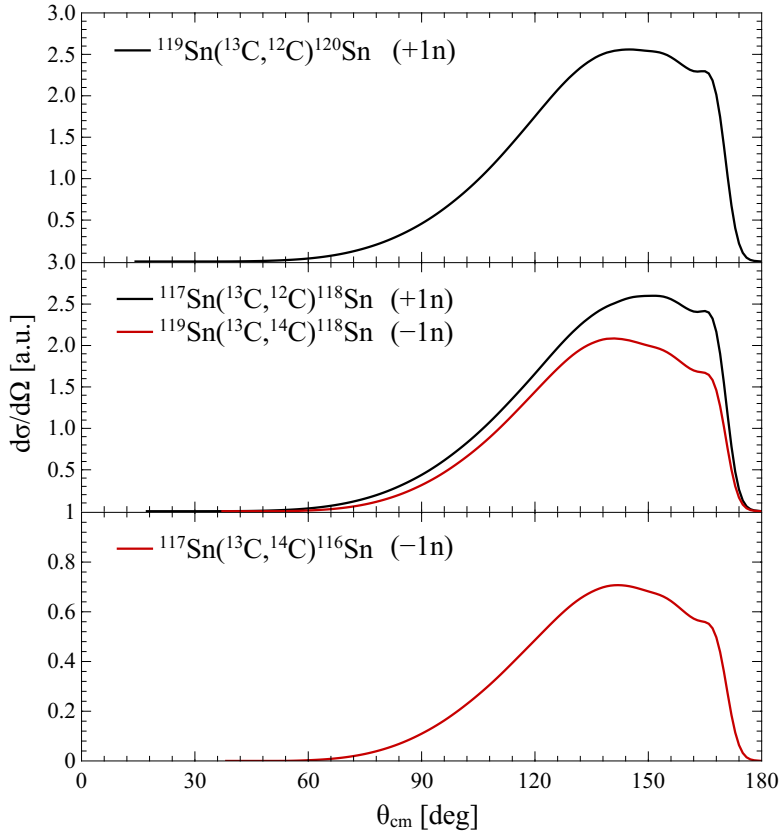


Figure 4.14: Angular distributions in the center-of-mass frame for the $\pm 1n$ reactions of *Exp. 3*. Calculations have been performed with the GRAZING code [94], using Wilczynski plots with $1 < E_x < 3$ MeV. Black and red line indicate $+1n$ and $-1n$ channels, respectively.

a Gaussian distribution with standard deviation parametrized as

$$\sigma = \frac{1}{2.35} \sqrt{A + BE_\gamma + CE_\gamma^2}, \quad (4.3)$$

It is important to note that, since the detectors at backward angles were severely damaged by neutrons, only the forward ring was considered for the lifetime analysis. Thus, a specific set of A,B and C parameters was extracted from the measured in-beam resolution. The experimental data points and the fitted resolution curve are shown in Figure 4.15.

Before the final histograms are produced, the Compton suppression is applied. This is done by excluding simulated events in which a BGO interaction was generated. As shown in Figure 4.16, this procedure effectively reduces the Compton-induced background by a factor of ~ 5 . Similar levels of suppression were observed in the experimental datasets (see Section 4.2).

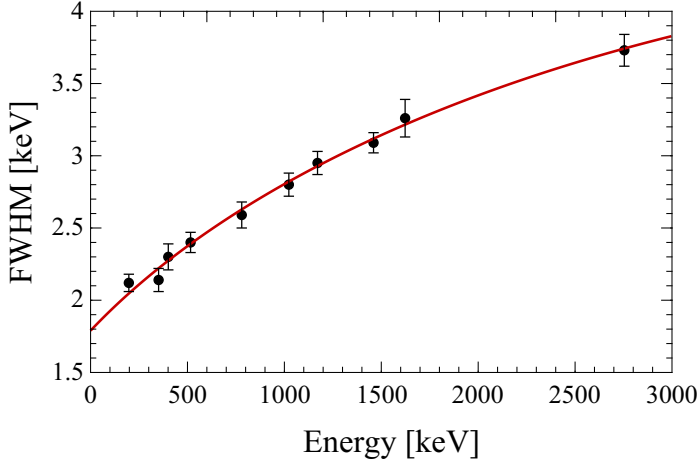


Figure 4.15: In-beam resolution as a function of the γ -ray energy of the ROSPHERE detectors mounted on the 37° ring. Experimental data points (black circles) and the fitted curve (red line) are shown.

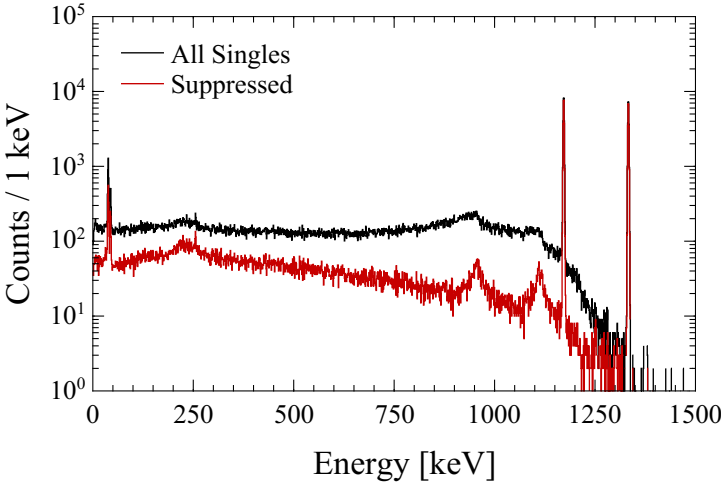


Figure 4.16: Simulated γ -ray spectra of a ^{60}Co source. The full HPGe spectrum is indicated in red, while the Compton-suppressed spectrum is depicted in black.

Analysis of IFIN-HH data

In this chapter, the γ -ray spectroscopy and lifetime analysis of the four datasets is presented. In Section 5.1, the particle- γ and γ - γ coincidence methods are illustrated, and the observed levels schemes are discussed. The lifetime analyses with the DSA and RDDS methods are thoroughly described in Section 5.2.

5.1 Spectroscopy

Using γ -ray spectroscopy techniques, such as γ - γ and particle- γ - γ coincidences, it is possible to study γ -ray cascades, determine the excited levels populated in the different transfer reactions and build level schemes of nuclei of interest.

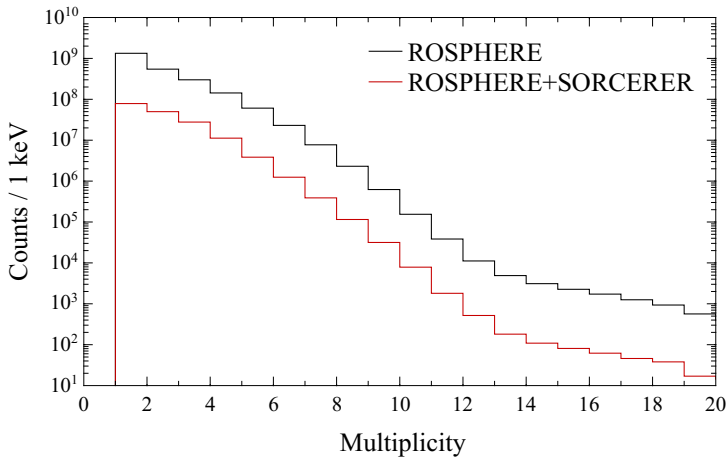


Figure 5.1: Multiplicity of the γ -rays emitted in the $^{13}\text{C}+^{119}\text{Sn}$ reaction of *Exp. 3*, as observed by ROSPHERE considering both HPGe and BGO detectors. The black and red lines indicate ROSPHERE only and ROSPHERE-SORCERER data, respectively.

For this purpose, the thick-target datasets of *Exp. 1* and *Exp. 3* were considered, as they provide cleaner spectra with respect to the plunger datasets. In the latter, the spectra were contaminated by Ta and Au peaks, arising from Coulomb excitation or inelastic scattering induced by the beam impinging onto the target backing and/or the stopper layer. Furthermore, a large fraction of the peaks were split into a shifted and a stopped component, thus increasing the complexity of the analysis.

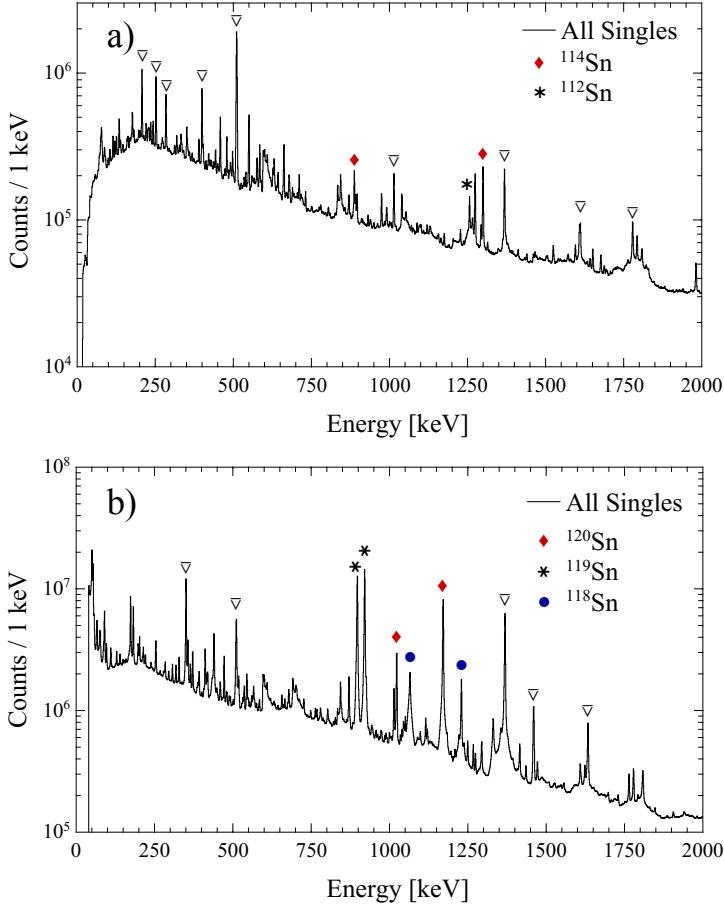


Figure 5.2: a) Singles γ -ray spectrum of the $^{18}\text{O}+^{112}\text{Sn}$ reaction of *Exp. 1*. The most intense transitions of ^{114}Sn are indicated by red diamonds, while transitions from the ^{112}Sn target are indicated with the star symbol. The main transitions from contaminants and secondary reactions (i.e. from ^{25}Mg , ^{127}La , ^{128}Ba and ^{128}Ce) are indicated with white triangles. b) Singles γ -ray spectrum of the $^{13}\text{C}+^{119}\text{Sn}$ reaction of *Exp. 3*. The most intense transitions of ^{120}Sn and ^{118}Sn are indicated by red diamonds and blue circles, respectively. γ -ray peaks resulting from the excitation of the ^{119}Sn target are marked with a star symbol, while the main transitions from contaminants and secondary reactions (i.e. from ^{20}Ne , ^{21}Ne , ^{24}Mg and ^{129}Cs) are indicated with white triangles.

In Figure 5.1, the multiplicity distributions of γ -ray events observed in the $^{13}\text{C}+^{119}\text{Sn}$ reaction are shown. The multiplicity is here defined as the total number of γ rays detected by the whole array. Indeed, besides the Compton rejection, BGO can be used as effective elements for a γ -ray multiplicity filter. A rapid decrease - steeper than exponential - is observed in the ROSPHERE multiplicity distribution, due to both the efficiency of the array, which scales exponentially, and the average number of γ rays emitted by the product nuclei, which is expected to be between 2 and 3 for this kind of transfer reactions. No significant change in the shape of the distribution is observed when the coincidence with SORCERER is required.

Panel a) and b) of Figure 5.2 show the all singles γ -ray spectra in the $^{18}\text{O}+^{112}\text{Sn}$ and $^{13}\text{C}+^{119}\text{Sn}$, respectively, with the $2_1^+ \rightarrow 0_1^+$ and $4_1^+ \rightarrow 0_1^+$ transitions of the Sn products of

interest being marked by red and blue symbols. It can be observed that the peak-to-total ratio of $^{118,120}\text{Sn}$ in the $1n$ transfer spectrum (panel b)) is significantly higher than the one of the $2n$ reaction (panel a)), as expected according to the higher cross section of the $1n$ reaction. Besides neutron transfer, peaks corresponding to target nuclei de-excitations (Coulomb excitation and/or inelastic scattering) were observed in all the experiments. These peaks, belonging to ^{112}Sn (at 1256 keV) and ^{119}Sn (doublets at 897 and 921 keV), are marked with stars symbols in panel a) and b) of Figure 5.2, respectively.

The main sources of background for these experiments come from sub-barrier fusion-evaporation on the target and from fusion evaporation on surrounding materials in the scattering chamber. In the $^{18}\text{O}+^{112}\text{Sn}$ case, the most intense contamination peaks come from ^{25}Mg , which is produced by fusion-evaporation on ^{12}C , and from ^{127}La , ^{128}Ba and ^{128}Ce , which are obtained from fusion-evaporation on the ^{112}Sn target. Instead, for the $^{18}\text{O}+^{112}\text{Sn}$ reaction, peaks from ^{20}Ne , ^{21}Ne , ^{24}Mg and ^{129}Cs were the main source of contamination. These nuclei are produced in fusion-evaporation reactions on ^{12}C (Ne and Mg) and ^{119}Sn (Ce).

Although these background contributions are not negligible, they can be significantly reduced using particle- and $\gamma\text{-}\gamma$ coincidences, which are discussed in the following.

5.1.1 Particle- γ coincidences

As mentioned before, particle- γ coincidences between SORCERER and ROSPHERE can be used to improve the selection on the reaction channel of interest and largely suppress background stemming from electromagnetic decays of fusion-evaporation reaction products. At energies around the Coulomb barrier, beam-like fragments after the transfer process are scattered at backward angles with respect to the beam direction when their mass is much smaller compared to the target. These recoiling fragments are then detected by SORCERER, which is placed at backward angles.

In the analysis, particle- γ coincidences are applied by gating on a time-coincidence window, as described in Section 4.3.2. An example of particle- γ coincidence matrix, built on the $^{13}\text{C} + ^{119}\text{Sn}$ dataset of *Exp. 3*, is shown in Figure 5.3.

Despite the poor resolution due to the lack of angular segmentation of the photodiodes, an additional condition on the energy of the recoils further improves the suppression of fusion-evaporation events, effectively constraining on the two-body scattering kinematics. As shown in the right panel of Figure 5.3, fusion-evaporation and background events are correlated with recoil energies in the entire range (with a small enhancement for $E_{\text{recoil}} < 1000$), while γ -ray decays from Sn isotopes are in coincidence with particle energies between channel 1000 and 3000. Furthermore, the limited mass difference between carbon isotopes resulted in a wide energy spectrum with no clear structure. Therefore, different transfer process were not distinguished and simultaneously selected with the energy gate. Similar considerations apply for all the other reaction performed at IFIN-HH, both with thick and plunger targets, as the kinematics remains analogous.

In Figure 5.4, a comparison between γ -ray singles spectra with and without the coincidence with particles is shown, highlighting the enhancement in the peaks of the transfer channels of interest. In particular, the peak-to-background ratio increases by more than one order of magnitude when the coincidence with particles is required.

As the scattering angles of beam-like recoiling fragments are constrained when SORCERER is considered, the corresponding scattering angle of the target-like recoil is also limited into a smaller range. This means that, in the thick target datasets, Doppler-shifted energy tails of short-lived transitions ($\tau \lesssim 1$ ps) become narrower and better

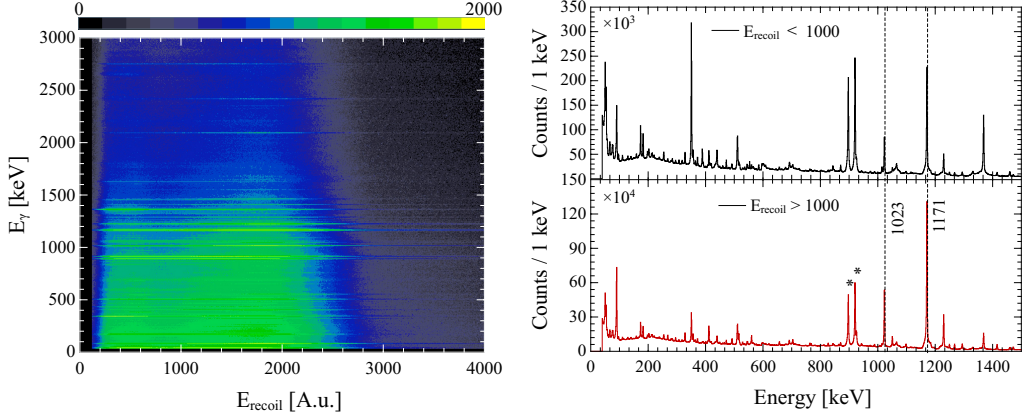


Figure 5.3: (Left) Particle- γ coincidence matrix, correlating the particle energies E_{recoil} detected in SORCERER and the γ -ray energies measured by the ROSPHERE HPGe detectors. The matrix refers to the $^{13}\text{C} + ^{119}\text{Sn}$ reaction of *Exp. 3*. (Right) Projection of the particle- γ matrix, gated on $E_{recoil} < 1000$ (top panel) and on $E_{recoil} > 1000$ (bottom panel). Vertical dashed lines are used as references for the energy of the $2_1^+ \rightarrow 0_1^+$ and $4_1^+ \rightarrow 2_1^+$ transitions in ^{120}Sn at 1171 and 1023 keV, respectively. Star symbols were used to label γ -ray transitions from excited states of the ^{119}Sn target.

defined, due to the smaller variation of both the velocity of the recoil and the relative angle $\theta_{\gamma, recoil}$ between the direction of motion of the recoil and the emission angle of the γ -ray. Such an effect can be seen in the ROSPHERE-SORCERER spectrum of panel a) of Figure 5.4, where the 1256-keV peak of the ^{112}Sn target shows pronounced tails on both sides that were not visible in the ROSPHERE only spectrum.

5.1.2 γ - γ coincidences

In this analysis, information about the population of excited states in the nuclei of interest was extracted by means of γ - γ coincidences. To achieve this, thick target datasets were used, ensuring that the Sn reaction products were completely stopped after the reaction. Nevertheless, γ -ray transition from excited states with lifetimes in the DSAM sensitivity range ($\tau \lesssim 1$ ps) were partially emitted in-flight, as the nuclei slowed down in the target material, and thus appeared at shifted energies due to the Doppler effect. Such γ -ray decays were identified through coincidences between the full ROSPHERE array and the five different rings, using the detectors at 90° to precisely determine the energy of transitions depopulating short-lived states.

In Figure 5.5, an example of a γ - γ matrix is shown. The absence of intense diagonal lines, signature of Compton scattering events, demonstrates the effectiveness of the BGO anti-Compton shields.

In Figure 5.6, the γ -ray spectra in coincidence with the $2_1^+ \rightarrow 0_1^+$ $E2$ transition of ^{120}Sn , ^{118}Sn , ^{116}Sn and ^{114}Sn are shown in panel a), b), c) and d), respectively. To reduce random coincidences and contamination from secondary channels, background regions were defined in proximity of the gating area and subtracted from it. The spectra account for the whole statistics collected during the experiments and coincidences in the full ROSPHERE array are considered. A specific color scheme has been assigned to each set of particles exchanged in the transfer reactions. Based on the spectra of Figure 5.6 and the reconstructed level schemes of Figures 5.7–5.11, it can be observed that in the

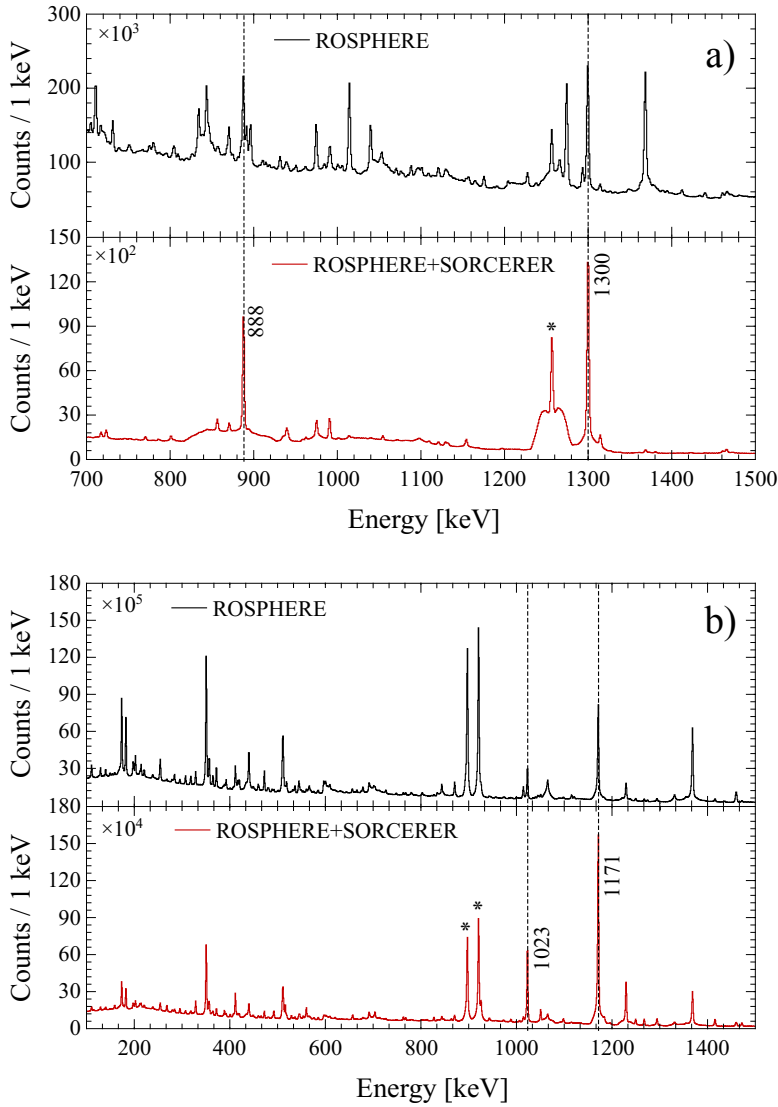


Figure 5.4: Singles γ -ray spectra without (black lines) and with (red lines) energy and time conditions on beam-like ions detected in SORCERER. Panel a) shows the spectra of the $^{18}\text{O} + ^{114}\text{Sn}$ $2n$ transfer reaction of *Exp. 1*, while panel b) refers to the $^{13}\text{C} + ^{119}\text{Sn}$ transfer reaction of *Exp. 3*. Vertical dashed lines are used as references for the energy of the $2_1^+ \rightarrow 0_1^+$ and $4_1^+ \rightarrow 2_1^+$ transitions in ^{114}Sn (1300 and 888 keV) and ^{120}Sn (1171 and 1023 keV), respectively. Star symbols were used to label γ -ray transitions from excited states of the ^{112}Sn and ^{119}Sn targets.

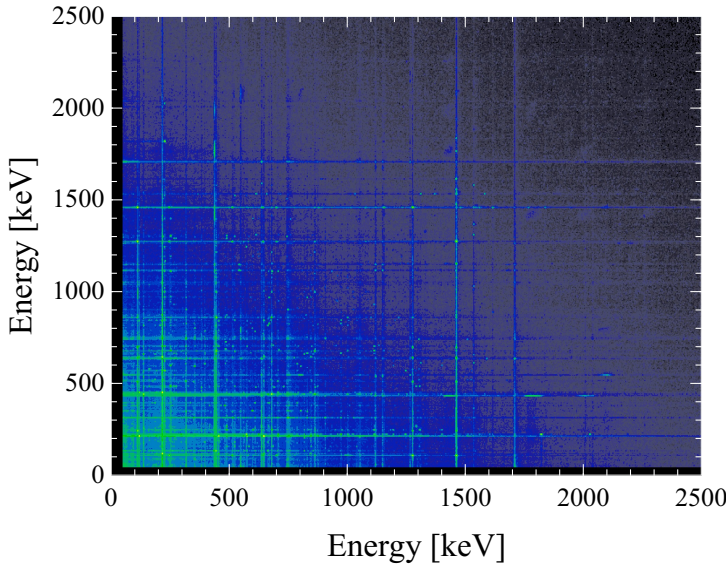


Figure 5.5: γ - γ coincidence matrix of the γ rays produced in the $^{13}\text{C} + ^{119}\text{Sn}$ reaction of *Exp. 3*.

different neutron transfer reaction most of the known states were populated up to ~ 3 MeV of excitation energy, although with different yields depending on the case. On the contrary, the $+1\alpha$ and $+2p$ reactions appeared to be more selective, but the lower cross section and the limited statistics prevented the observation of weak transitions. Multiple Doppler broadened peak shapes of short-lived states can also be observed in all the spectra.

Even without the coincidence with SORCERER, the gate on the γ -ray transition de-exciting the first excited state of Sn isotopes provides an optimal selection of the reaction channel. Only contaminants originating from Coulomb excitation and/or inelastic scattering on the nuclei used for the targets were found in the gated projections. In particular, the observed peaks from ^{119}Sn (doublets at 897 and 921 keV), ^{117}Sn (159 keV) and ^{116}Cd (514, 706 and 867 keV) are indicated with star symbols in Figure 5.6.

At the cost of considering a small fraction of the total statistics (from 20 to 30 %), the coincidence with SORCERER can be required and particle- γ - γ matrices can be built. The latter can be used to lower the background level and confirm the assignment of observed transitions to nuclei produced in the reaction of interest.

5.1.3 Level schemes

In the following, the reconstructed level schemes of $^{114,116,118,120}\text{Sn}$ are discussed. The identification of γ -ray lines to specific levels has been done according to previous literature results (mainly from the NNDC database [68]), particle- γ , γ - γ and particle- γ - γ coincidences. Gates on transitions populating and de-populating excited states (namely gates from "above" and "below") were employed to identify and reconstruct γ -ray cascades. It is important to say that only one new transition (a 621-keV transition in ^{120}Sn) was measured among all the explored transfer channels. In fact, the level schemes of stable Sn isotope are well known, since they can be studied with a multitude of reactions employing stable and high-intensity beams. Moreover, the main goal of the thick target experiments (*Exp. 1* and *Exp. 3*) was to identify the best reaction mechanism for

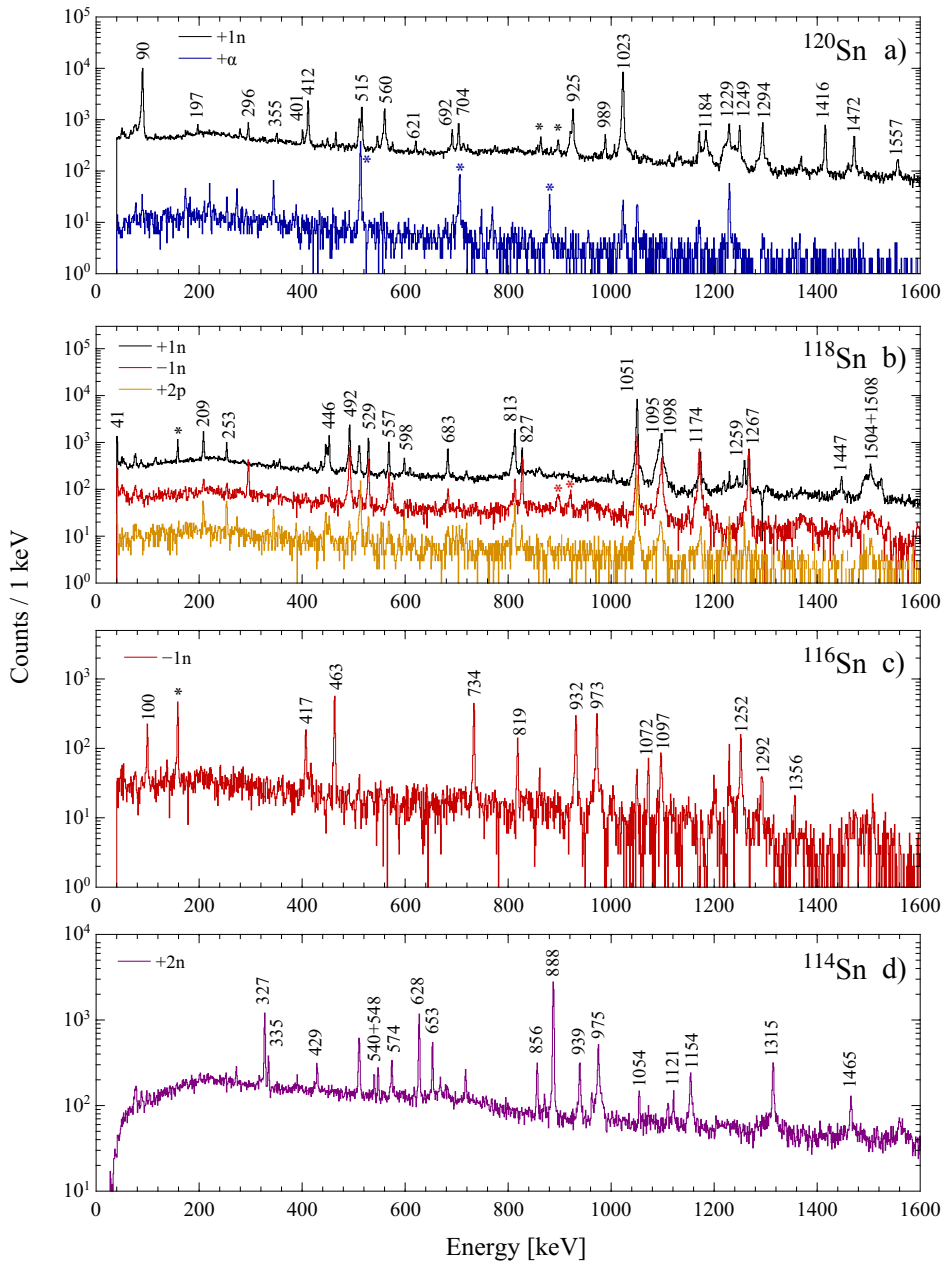


Figure 5.6: Projections of γ - γ matrices obtained in the transfer reactions of *Exp. 1* and *Exp. 3*, gated on the $2_1^+ \rightarrow 0_1^+$ transitions of the Sn isotopes of interest. Panel a) shows the γ -ray spectra gated on the 1171-keV transition of ^{120}Sn , populated with the +1n and + α transfer reactions. The projections of panel b) were gated on the 1230-keV transition of ^{118}Sn , which was populated via +1n, -1n and +2p transfers. Panel c) show the spectrum gated on the 1293-keV γ -ray of ^{116}Sn observed in the -1n transfer reaction, while panel d) depicts the spectrum of ^{114}Sn , gated on the 1300-keV peak, after the pick-up of 2n. The most intense transitions of each nucleus are labeled with their energy in keV, and peaks from the target nuclei are indicated with the star symbol.

the population of excited 0^+ states, rather than expanding the adopted level schemes, thus serving as testing ground for subsequent lifetime measurements with the RDDS technique.

Level scheme of ^{114}Sn

The level scheme of ^{114}Sn , built on the basis of the γ rays detected in the $2n$ transfer reaction of *Exp. 1*, is shown in Figure 5.7. Almost all the known states up to 3.5 MeV of excitation energy were populated, as expected based on the 5.8 MeV Q -value of the reaction. The measured energies and intensities agree with the data available in literature [68]. Spin and parities assignments were recently revised in (p, t) and $(p, p'\gamma)$ experiments [71, 80], and are here adopted accordingly. In particular, two excited states are adopted in the database at 3025 keV, with $J^\pi = 2, 3^+$ and $J^\pi = 0^+$, and another state at 3028 keV is reported with $J^\pi = 2, 3^+$. However, the $J^\pi = 0^+$ assignment for a state at $E_x = 3028(3)$ keV proposed in Ref. [71] and the γ transitions observed by Ref. [80] exclude the previous adoptions. As a consequence, the presence of a $J^\pi = 2^+$ state at 3024.9(2) keV and a $J^\pi = 0^+$ state at 3028.1(2) keV is here assumed. The band structure follows the arrangements proposed in Ref. [75].

Level scheme of ^{116}Sn

In Figure 5.8, the level scheme reconstructed from the $-1n$ transfer reaction dataset of *Exp. 3* is displayed. No discrepancies have been found between the energies and intensities observed in this work and the literature data. In the adopted level scheme [68], two level at 2545.7 and 2790.6 keV are listed with a spin-parity of $J^\pi = (0)^+$. Ref. [132] confirms the 0^+ spin and parity of the former, while suggests a $J^\pi = 4^+$ assignment for the latter. Indeed, the observed angular distribution is reproduced by an $L = 4$ transfer assumption, but disagrees with the $(0)^+$ assignment from the $(n, n'\gamma)$ study [133] and the $J^\pi = 2^+$ deduced from the (d, p) reaction of Ref. [134]. In this work, the assignments from the (p, t) reaction of Ref. [132] are considered. The levels displayed in Figure 5.8 are organized in bands, based on the observed decay scheme and the results from previous works [78]. The "classic" assignment of the 0_2^+ state as band-head of the intruder band is here considered [13, 14].

Level scheme of ^{118}Sn

Figure 5.9 shows levels and γ transitions of ^{118}Sn as observed in the $+1n$ reaction of the present work. In fact, the $-1n$ and $+2p$ reactions seemed to populate the most of the levels seen in the $+1n$ reaction, however, weaker γ ray branches were not observed due to the limited statistics. Energies and intensities of the γ ray transitions are in agreement with the ones from recent β decay and (n, γ) experiments [77, 135]. Discrepancies in the spin and parity assignments have been found between these two works and the NNDC database [68]. Specifically, the excited state at 2878.4 keV was listed as $4, 5, 6^+$ but Refs. [70, 135, 136] suggested a $J^\pi = 5^-$. In addition, a tentative spin of 5 has been proposed for the 3089.3-keV level [77]. Concerning the transition with energy $E_\gamma = 492.3$ keV, conflicting results have been found between the database and the present work. In the NNDC database, such a γ ray connects the 3^- state at 2817.2 keV with the lower lying 3^- at 2324.8 keV. However, in our data the 492.3-keV transition is unambiguously observed in coincidence with both the 40.8- and 1091.5-keV γ rays depopulating the 5^- state at $E_x = 2321.2$ keV, as shown in Figure 5.10. Therefore, the excitation energy of the initial state should be placed at 2813.6 keV, and a $(3)^- \rightarrow 5^-$ γ -ray at 492.3 keV should be

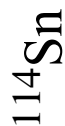


Figure 5.7: Level scheme of ^{114}Sn as observed in the $2n$ transfer reaction of *Exp. 1*.

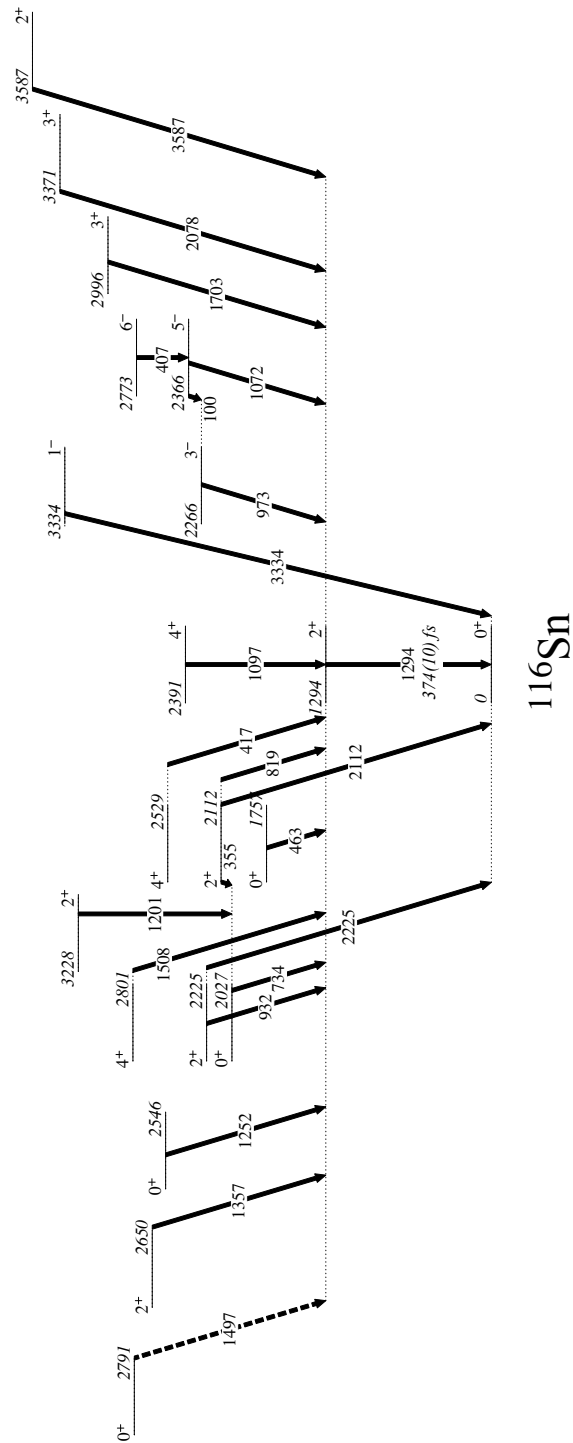


Figure 5.8: Level scheme of ^{116}Sn as observed in the $-1n$ transfer reaction of Exp. 3. Tentative γ rays are indicated with dashed lines.

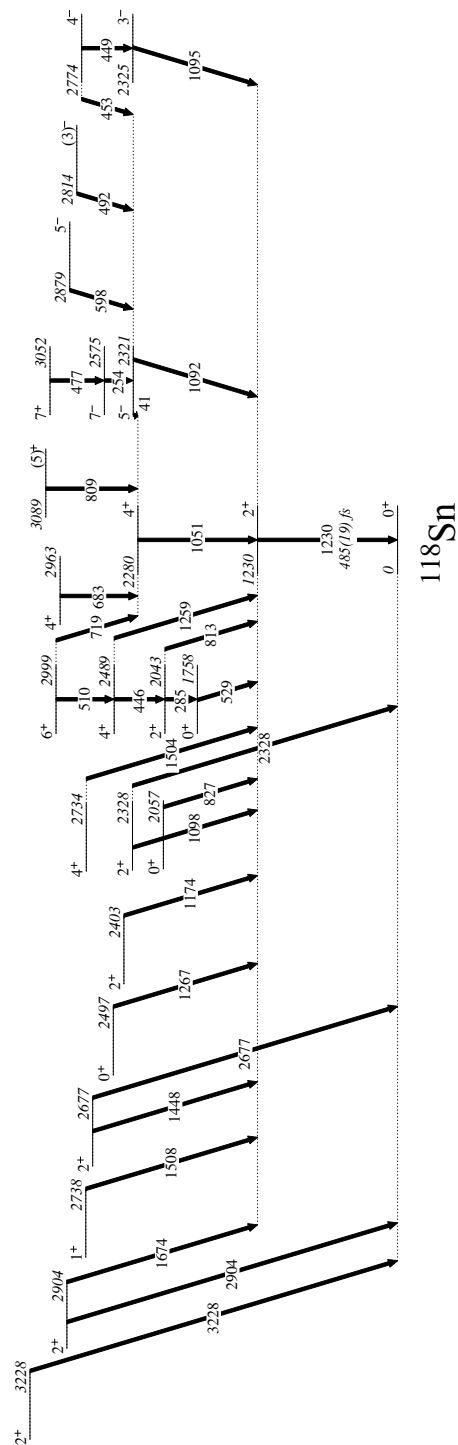


Figure 5.9: Level scheme of ^{118}Sn as observed in the $\pm 1n$ transfer reaction of *Exp. 3*.

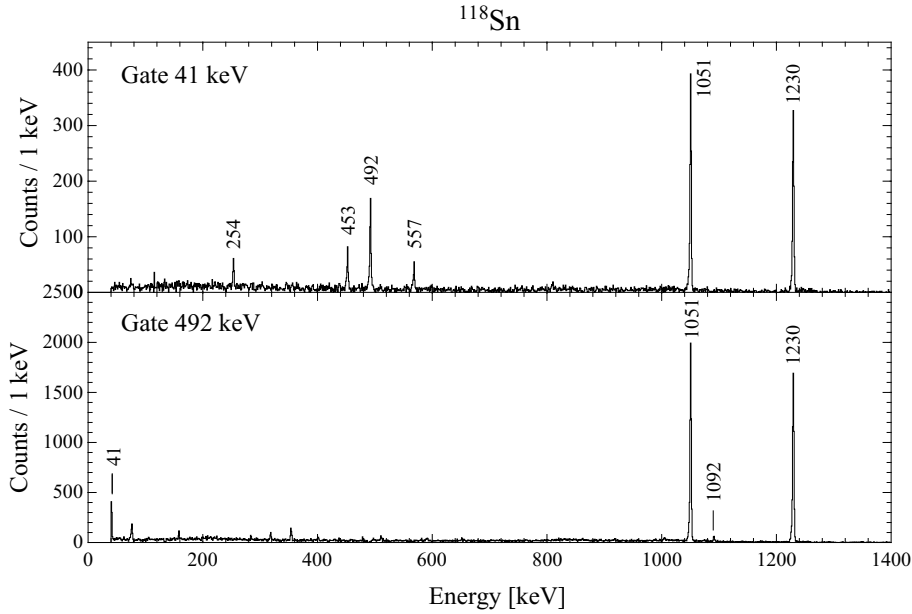


Figure 5.10: Projections of the particle- γ - γ matrix of ^{118}Sn obtained in the $+1n$ transfer reaction of *Exp. 3*. In the top panel, the gate is placed on the $5_1^- \rightarrow 4_1^+$ transition at 41 keV, while in the bottom panel the gate is set on the peak at 492 keV. Since the 41-keV transition is partially cut by the low-energy threshold, only a fraction of the total counts appear in the peak. γ -ray peaks of ^{118}Sn are labeled with their energy in keV.

considered. Yet, the $J^\pi = (3^-)$ assignments of the initial state remains tentative, as in (d, p) and (t, d) experiments a $(4, 5, 6)^-$ level at a similar energy of 2817(3) keV was observed [137, 138].

Level scheme of ^{120}Sn

The level scheme of ^{120}Sn as obtained from the $+1n$ transfer reaction is shown in Figure 5.11. The small amount of counts collected in the $+\alpha$ transfer prevented an adequate comparison between two reactions. As in the previous cases, a good agreement is found between energies and relative intensities observed in the literature and in the present study. Spin and parity assignments were mostly measured with direct reactions and later adopted by the NNDC database [68, 73]. A new γ ray of $E_\gamma = 620.7(2)$ keV was measured in coincidence with the 515.8-keV transition de-exciting the 5^- state at 2800.5 keV and with all the lower lying transitions belonging to the cascade towards the ground state. This coincidence is shown in Figure 5.12, where the projections of the particle- γ - γ matrix, gated on the 516-keV and 621-keV transitions, are presented. Since no known levels have are located at $E_x = 3421.2$ keV, the presence of a new level is here proposed. This state and its γ decay are marked with the red color in the level scheme of Figure 5.11. Levels were arranged in bands based on previous works and observed decays.

Population of 0^+ states

In a quest for the best reaction mechanism for the population of excited 0^+ states, it is useful to study the different excitation patterns in each reaction channel. In Figure 5.13,

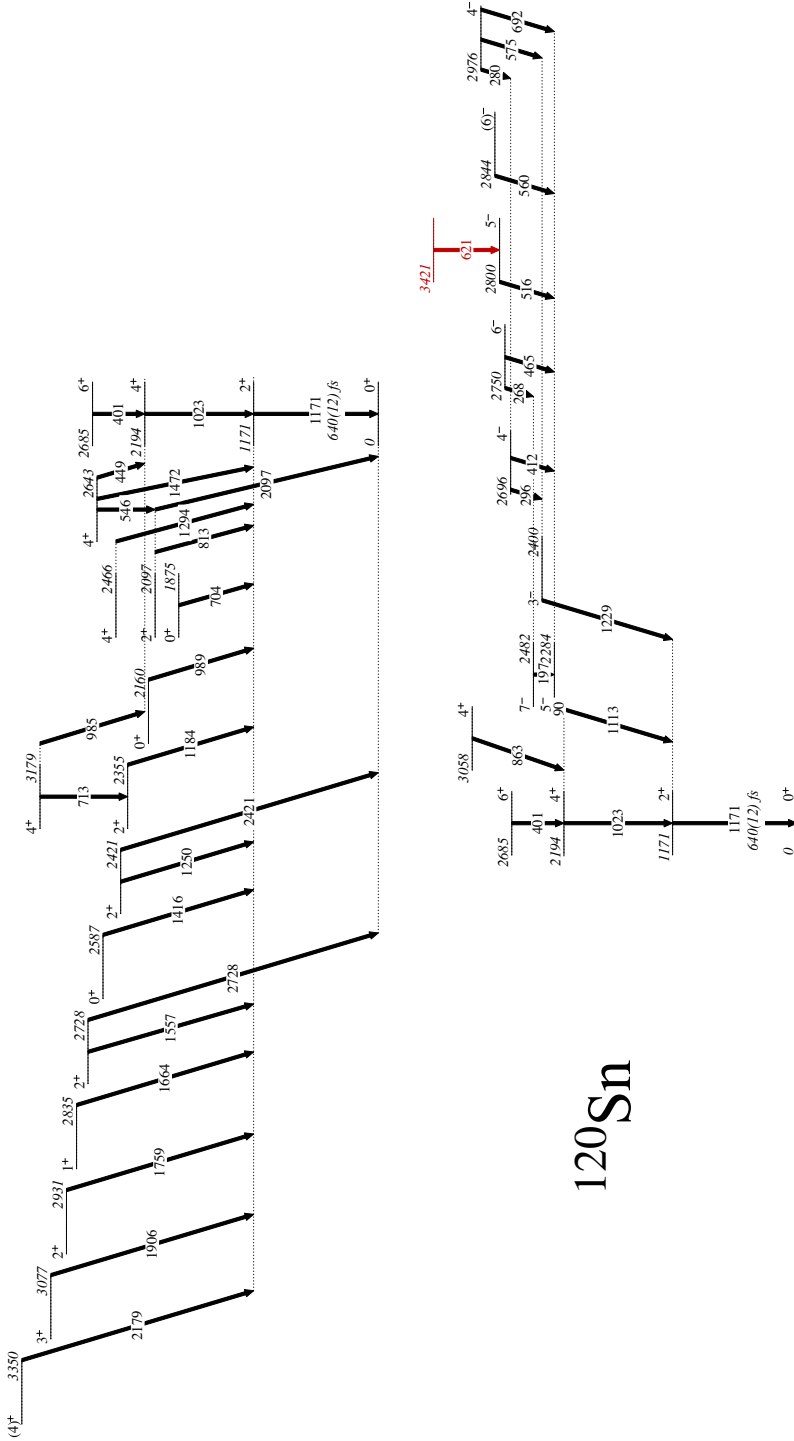


Figure 5.11: Level scheme of ^{120}Sn as observed in the $+1n$ transfer reaction of Exp. 3. The new level and γ ray observed in this work are indicated in red.

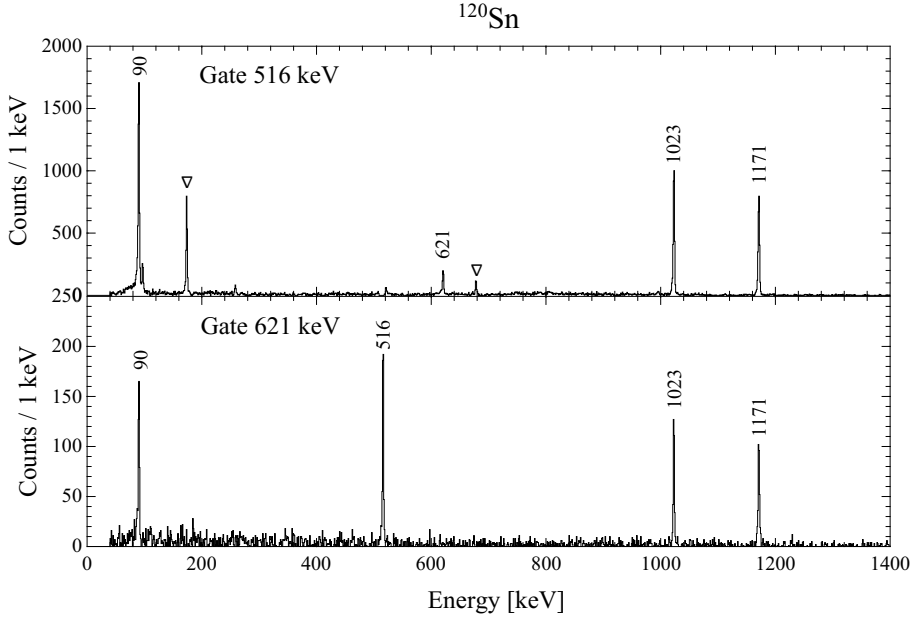


Figure 5.12: Projections of the particle- γ - γ matrix of ^{120}Sn obtained in the $+1n$ transfer reaction of *Exp. 3*. In the top panel, the gate is placed on the $5_2^- \rightarrow 5_1^-$ transition at 516 keV, while in the bottom panel the gate is set on the newly identified transition at 621 keV. γ -ray peaks of ^{120}Sn are labeled with their energy in keV, while contaminant peaks are marked with a white triangle.

the measured intensities of the γ rays depopulating 0_2^+ , 0_3^+ and 0_4^+ in $^{114,116,118,120}\text{Sn}$ are shown. In order to compare them, intensities have been corrected for the detection efficiency and normalized to the one of the $0_2^+ \rightarrow 2_1^+$ transition, which is set to 100%. By doing so, these values become good indicators of the population yields of the excited 0^+ states, accounting for both direct and indirect excitations. In this case, the contributions given by internal conversion transitions, whose intensities typically sit at $I_e \lesssim 10^{-3} I_\gamma$ for Sn isotopes [87, 104], are neglected.

From the panels of Figure 5.13, one observes that from ^{114}Sn to ^{120}Sn , the population of 0_2^+ and 0_3^+ states remains similar, with variations up to $\pm 30\%$ depending on the transferred nucleons. The only exception is the α reaction for ^{120}Sn , which shows a weaker population of 0_3^+ . However, due to the very limited statistics, the intensities of the α -transfer channel have to be considered as tentative. On the other hand, the relative yields of the 0_4^+ state change dramatically based on the reaction channel. While being hindered in ^{114}Sn , its population becomes favored in the $\pm 1n$ and $+1n$ channels of ^{118}Sn and ^{120}Sn , respectively. This is evident for the $-1n$ transfer of ^{118}Sn , where the normalized yield of the 0_4^+ state is more than 3 times the one of 0_2^+ state. Moreover, no correlation with the Q -values of the reactions is observed. A summary of the relative intensities and errors is also reported in Table 5.1.

It is important to note that, although precious information for this work, these yields cannot be properly converted into cross sections due to the lack of a precise external normalization. Nevertheless, these values can be used to formulate hypothesis concerning the character of these excitations in terms of main wavefunction components: such a discussion will be carried out in Chapter 9.

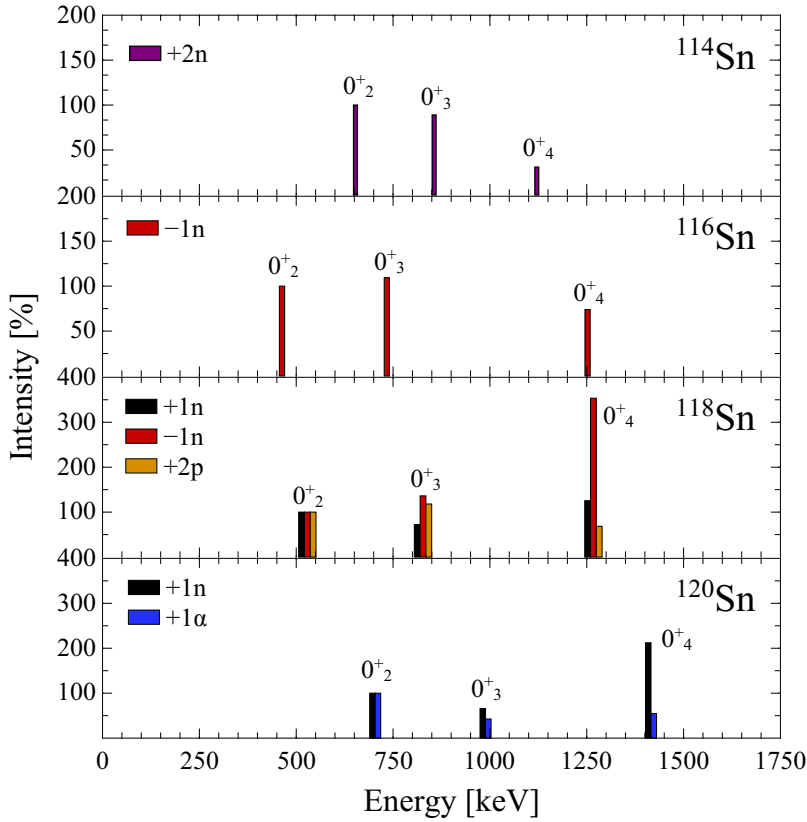


Figure 5.13: Relative intensities of the $0_{2,3,4}^+ \rightarrow 2_1^+$ transitions of $^{114,116,118,120}\text{Sn}$ observed in coincidence with the corresponding $2_1^+ \rightarrow 0_1^+$ transitions. The intensities are corrected for the efficiency and normalized to the intensity of the γ ray depopulating the 0_2^+ states.

Nucleus	Channel	$I(0_2^+)[\%]$	$I(0_3^+)[\%]$	$I(0_4^+)[\%]$
^{114}Sn	$+2n$	100(6)	89(12)	31(17)
^{116}Sn	$-1n$	100(4)	112(7)	78(9)
^{118}Sn	$+1n$	100(7)	75(8)	132(9)
^{118}Sn	$-1n$	100(9)	127(10)	338(25)
^{118}Sn	$+2p$	100(36)	122(55)	73(29)
^{120}Sn	$+1n$	100(7)	71(10)	222(14)
^{120}Sn	$+1\alpha$	100(67)	44(13)	58(29)

Table 5.1: Summary of the relative intensities of the $0_{2,3,4}^+ \rightarrow 2_1^+$ transitions of $^{114,116,118,120}\text{Sn}$ observed in coincidence with the corresponding $2_1^+ \rightarrow 0_1^+$ transitions. The intensities and the associated errors are corrected for the efficiency and normalized to the intensity of the γ ray depopulating the 0_2^+ states.

5.2 Lifetime measurements

This Section will present the analysis and results of the lifetime measurements performed on the IFIN-HH datasets. Section 5.2.1 will focus on the lifetime measurements of excited states with the DSA technique, while Section 5.2.2 will describe the results obtained on the RDDS datasets with the Decay Curve Method (DCM). A summary of the final results will be presented in Section 5.2.3.

5.2.1 DSAM measurements

As described in Chapter 2, the DSAM is a Doppler-based technique that exploits the lineshape of a γ -ray peak to determine the lifetime of short-lived nuclear excited state (i.e. with $\tau \lesssim 1$ ps). In practical terms, line shapes of γ transitions are simulated for a set of possible lifetimes of the initial level, and then compared with the experimental spectra to find the best agreement.

For the experimental data, γ - γ coincidence matrices were employed. The gate was placed on the axis with events recorded by all the ROSPHERE detectors, while only coincidence events detected in the 37° ring were considered. This choice was motivated by the fact that, during *Exp. 3*, the HPGe detectors mounted on the backward ring suffered neutron damage, leading to a deterioration in energy resolution and the appearance of low-energy tails on the γ -ray peaks. Detectors on rings at 70° , 90° and 110° were also discarded due to their lower sensitivity to lineshape effects and the reduced available statistics (few detectors mounted on each of them). An example of γ -ray spectra for different ROSPHERE rings is shown in Figure 5.14, where the prominent shifted components of the 1664-keV transition of ^{120}Sn are depicted. This γ -ray transition, which depopulates the 1^+ state at 2835 keV with half-life $T_{1/2} = 130_{-30}^{+60}$ ps [68], will also be used as benchmark for the lifetime estimation with the DSAM method.

The simulated spectra were produced using the novel software implementation presented in Section 4.4. These simulation accounted for the reaction mechanism, the angular distributions of the reaction products and the energy losses in the target material to generate realistic events. After applying energy smearing and Compton suppression, histograms were obtained.

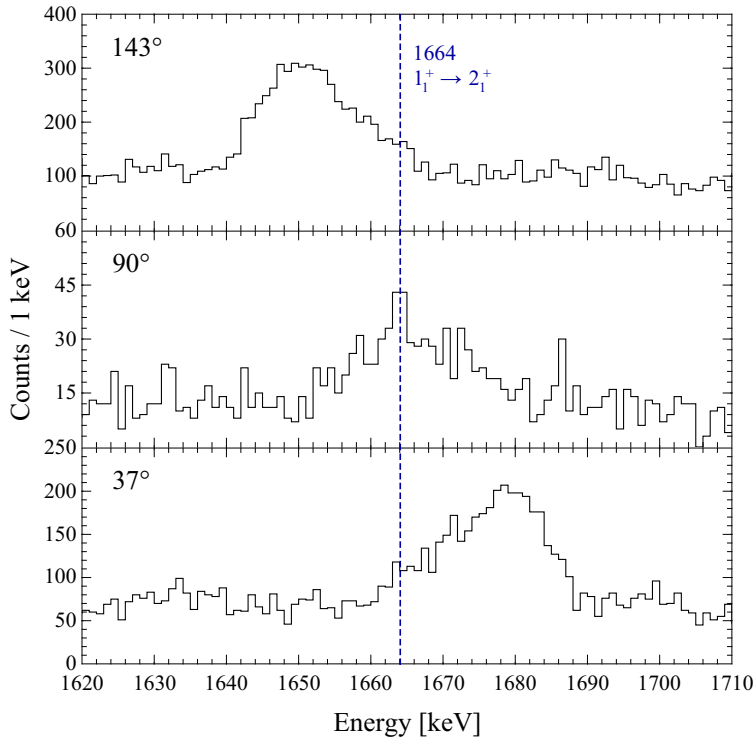


Figure 5.14: Partial γ -ray spectra recorded in the 37° , 90° and 143° rings of ROSPHERE, gated on the $2_1^+ \rightarrow 0_1^+$, 1171-keV transition of ^{120}Sn . The focus is on the 1664-keV $1_1^+ \rightarrow 2_1^+$ transition, which appear as mostly shifted due to the $T_{1/2} = 130^{+60}_{-30}$ fs half-life of the 1_1^+ state. The coincidence gate was set on the full ROSPHERE array. Data collected during Exp. 3.

In order to quantify the agreement between simulated and experimental spectra, the χ^2 test was used. According to the Neyman statistic [139], the χ^2 can be defined as:

$$\chi^2 = \sum_i \left(\frac{y_{exp}(i) - y_{sim}(i)}{\sigma_{exp}(i)} \right)^2, \quad (5.1)$$

where i enumerates the bins in the region considered for the comparison (region of sensitivity), $y_{exp}(i)$ is the number of counts in the i -th bin of the experimental histogram, $y_{sim}(i)$ the number of counts in the i -th bin of the simulated histogram, and $\sigma_{exp}(i)$ is the experimental uncertainty on $y_{exp}(i)$. In a raw experimental histogram, where the number of counts in each bin is governed by the Poisson statistics, the error is given by $\sqrt{y_{exp}}$. Instead, when the background is subtracted, the error has to be propagated accordingly. This is the case of our experimental spectrum: two background regions were set on the sides of the gating windows and the background was subtracted from the coincidence spectrum. A normalization was also applied to take into account the different widths of the gate and background windows. Finally, the uncertainty becomes $\sigma_{exp} = \sqrt{y_{exp} + y_{bkg}}$, being y_{bkg} the counts in the normalized background spectrum.

Using the Neyman's χ^2 definition, the statistical error of the simulation is assumed to be negligible, which is justified here because the number of simulated events greatly

exceeds that of the experimental data. In particular, 2×10^6 events were generated for each lifetime simulation.

For each transition of interest, the χ^2 was calculated in a specific region of interest. This region, which identifies portions of the spectrum which are sensitive to lineshape effects, can be defined by simulating a very short decay ($T_{1/2} = 1$ fs) and a very long one ($T_{1/2} = 5$ -10 ps). The corresponding γ rays will then be characterized by only in-flight or stopped components, thus providing the boundaries of the sensitivity region. This is shown, as an example, in Figure 5.15, where the 1664-keV γ ray was simulated for half-lives of 1, 500, 1000 and 5000 fs. In this case, the region of interest was set between 1658 and 1691 keV. Once the region of interest is chosen, simulated counts have to be

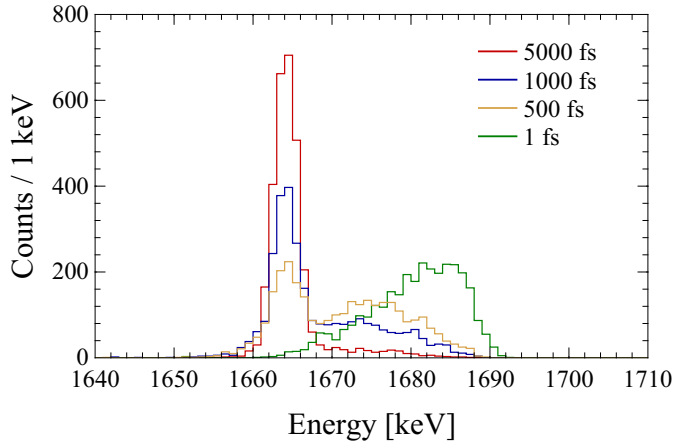


Figure 5.15: Simulated γ -ray spectra of the 1664-keV transition of ^{120}Sn for the 37° ring of RO-SPHERE, assuming different half-lives of the depopulated state.

normalized to match the experimental ones. As the minimization procedure requires the same number of counts in both the experimental and the simulated spectrum, the latter has to be scaled accordingly. Finally, the χ^2 curve is obtained by scanning the half-life of the state in a range of values. Steps of 5 and 10 fs were chosen for expected half-lives below and above 1 ps, respectively. To avoid point-like fluctuations, the results were obtained by fitting the χ^2 with 3^{rd} or 4^{th} order polynomials and finding their minimum $\chi^2_{\min}$. Similarly, the one standard deviation confidence intervals, provided by the $\chi^2_{\min} + 1$ level, were determined from the fitted curves.

In the following, results for the 2835- and 1171-keV levels of ^{120}Sn and for the 2545-keV level of ^{116}Sn will be presented. The first two results can be considered as benchmarks for this new implementation of the simulation code for ROSPHERE, as their lifetime was previously measured [68].

The 1_1^+ state of ^{120}Sn

The 2835-keV 1^+ state was found to be a good candidate for the validation of the simulation procedure, given the large statistics collected in its 1664-keV γ -ray peak in the $+1n$ transfer dataset of *Exp. 3*, and the absence of feeding from higher-lying states. The χ^2 curve resulting from the minimization procedure is shown in the left panel of Figure 5.16. Data points corresponding to different simulated lifetimes have been fitted with a 3^{rd} -order polynomial, allowing for a better estimation of the minimum χ^2 value. The

latter is found for a simulated half-life of $T_{1/2} = 174^{+14}_{-13}$ fs, with error bars corresponding to the $\chi^2_{min} + 1$ level. A comparison between the experimental spectrum and the best simulation is illustrated in the right panel of Figure 5.16, demonstrating the good agreement achieved. We note that the γ -ray spectrum utilized here is a projection of a γ - γ matrix, gated on the 1171-keV transition depopulating the lower lying 2^+_1 state.

The half-life obtained in the present work is fully compatible with the adopted value of 130^{+60}_{-30} fs, which was measured with the DSA method in $(n,n'\gamma)$ [68]. For this reason, this can be considered a first proof of the performance of this new implementation of a simulation package for ROSPHERE.

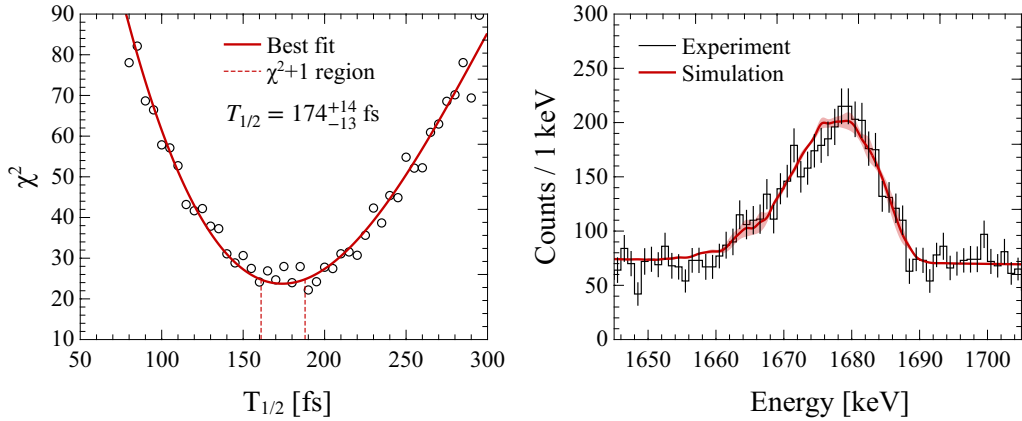


Figure 5.16: (Left) χ^2 curve as a function of the half-life of the 1^+ state at 2835 keV. Solid and dashed red lines indicate the cubic polynomial fit and the $\chi^2_{min} + 1$ range extracted from it, respectively. (Right) Comparison between experimental (black line) and simulated (red line) spectra for the $1^+_1 \rightarrow 2^+_1$ transition of ^{120}Sn at 1664 keV. Red shaded bands correspond to the $\chi^2_{min} + 1$ confidence range.

The 2^+_1 state of ^{120}Sn

Since most of the excited states of ^{120}Sn decay towards the first excited 2^+ state at 1171 keV, its lifetime estimation requires a precise control on the feeding from higher-lying states. This is achieved by using the experimental γ -ray spectrum gated on the 1664-keV transition of the 2835-keV state. In this way, it is possible to restrict the simulation to the $1^+ \rightarrow 2^+ \rightarrow 0^+$ cascade, assuming a direct population of the 1^+ state.

The lifetime of the 2835-keV state becomes an input parameter, and both the literature value and that of our measurement (see Section 5.2.1) were tested. The left panel of Figure 5.17 shows the χ^2 curve obtained using the literature value of 130^{+60}_{-30} fs for the half-life of the 2835-keV level. Based on this, a half-life $T_{1/2} = 644^{+60}_{-53}$ fs is obtained. The simulation corresponding to this value is shown in the right panel of Figure 5.17, and similar lineshapes between simulated and experimental spectra are observed. Instead, if our result of $T_{1/2}(1^+_1) = 174^{+14}_{-13}$ fs is used as input, the minimization procedure yields a half-life of $T_{1/2} = 581^{+59}_{-54}$ fs for the 2^+_1 state. Although a slightly lower half-life was determined in the second case, our results are well aligned with the adopted value of 640 ± 12 fs, which is the weighted average of several measurements [68].

It is important to note that the error bars of the 2835-keV level half-life are not included in the error estimation. Therefore, slightly larger statistical errors have to be

expected. Yet, the measured values are intended to be a benchmark of a complex analysis which required the inclusion of feeding. The agreement with previous experimental results provides further evidence of the goodness of this lineshape simulation procedure and opens to the possibility to measure new lifetimes for the first time. Such a case will be discussed in the following.

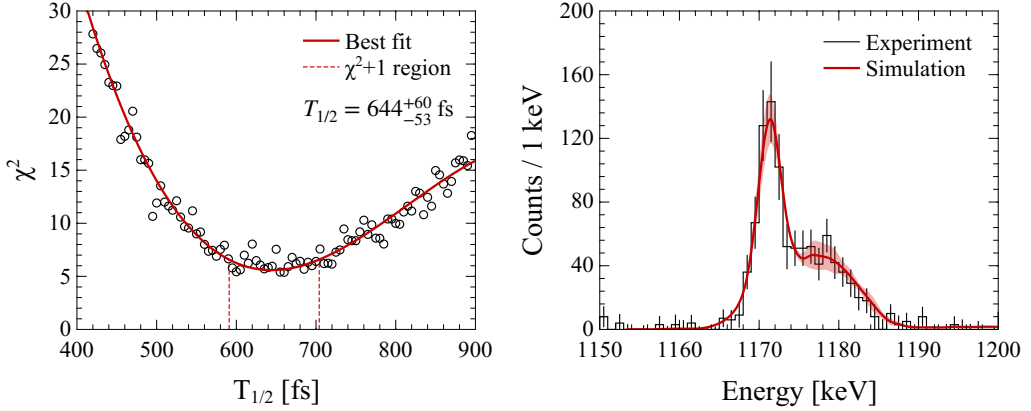


Figure 5.17: (Left) χ^2 curve as a function of the half-life of the 2_1^+ state at 1171 keV. Solid and dashed red lines indicate the cubic polynomial fit and the $\chi^2_{min} + 1$ range extracted from it, respectively. (Right) Comparison between experimental (black line) and simulated (red line) spectra for the $2_1^+ \rightarrow 0_1^+$ transition of ^{120}Sn at 1171 keV. Red shaded bands correspond to the $\chi^2_{min} + 1$ confidence range.

The 0_4^+ state of ^{116}Sn

In the dataset of *Exp. 3*, the 0_4^+ state of ^{116}Sn was populated in the $^{117}\text{Sn}(^{13}\text{C}, ^{14}\text{C})$ reaction, and a noticeable lineshape was observed in γ -ray peak corresponding to its de-excitation towards the 2_1^+ state. Thanks to that, the lifetime of this state could be estimated with the DSA method.

The χ^2 curve obtained by comparing simulated spectra with the experimental spectrum, resulting after gating on the $2_1^+ \rightarrow 0_1^+$ transition, is shown in the top-left panel of Figure 5.18. By fitting this curve with a 4th-order polynomial, a clear minimum is determined, yielding a half-life of 1172^{+98}_{-88} fs. The corresponding simulated spectrum is shown as a red line in the top-right panel of Figure 5.18, and a good agreement with the experimental spectrum is observed.

Since a few bins contained less than 10 counts, the same procedure was repeated after a rebinning of the spectrum to 2 keV per channel. As it can be observed in the bottom panels of Figure 5.18, no significant differences were found in the two cases. In fact, a half-life of 1156^{+97}_{-86} fs was determined, which confirms the result obtained with the unbinned spectrum.

5.2.2 RDDS measurements

In this analysis, the RDDS technique has been applied to measure the lifetime of several excited 0^+ states in $^{114,116,118,120}\text{Sn}$. Results were obtained with the Decay Curve Method (DCM), which is better suited than Differential Decay Curve Method (DDCM) for analyses with a small number of distances and low-statistics datasets [111].

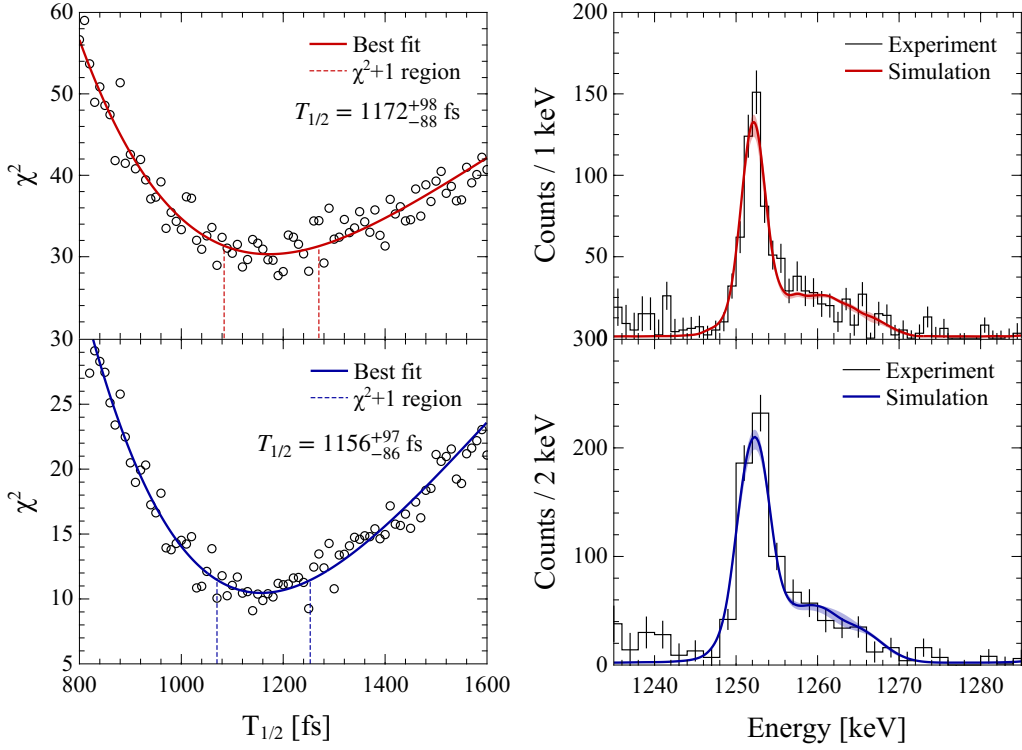


Figure 5.18: (Left) χ^2 curve as a function of the half-life of the 0_4^+ state at 2546 keV. Solid and dashed red lines indicate the cubic polynomial fit and the $\chi^2_{min} + 1$ range extracted from it, respectively. (Right) Comparison between experimental (black line) and simulated (red line) spectra for the $0_4^+ \rightarrow 2_1^+$ transition of ^{116}Sn at 1252 keV. Red shaded bands correspond to the $\chi^2_{min} + 1$ confidence range.

As introduced in Chapter 2, if the lifetime of a state is longer than ≈ 1 ps, it can be extracted using the plunger device from the intensity of the unshifted (I_u) and shifted (I_s) components of a γ -ray peak that depopulates the state of interest. If the state is populated directly, the fraction of residual nuclei in the initial excited state, after a time t , is given by:

$$R(t) = \frac{I_u(t)}{I_s(t) + I_u(t)} = e^{-\frac{t}{\tau}} = e^{-\frac{d \ln 2}{v \cos(\theta) T_{1/2}}}, \quad (5.2)$$

where, I_u and I_s are the intensity of the stopped and shifted components, respectively, τ is the lifetime of the state, d is the target-stopper distance, v is the velocity of the emitted ion and θ is the scattering angle. Knowing the average velocity and scattering angle of the emitting ion, the unknown parameter $T_{1/2}$ can be determined by fitting the experimental ratio R , sampled at different plunger distances, with an exponential decay curve. In a general case, when higher-lying levels are involved in the γ cascade, additional feeding terms have to be accounted for using the Bateman equations [140]. Such a procedure is described in Refs. [111, 141].

In this work, a slightly modified version of this procedure was applied. Indeed, the lifetime analysis is performed by using projections of γ - γ coincidence matrices, gated on the unshifted component of the $2_1^+ \rightarrow 0_1^+$ transition of the Sn products of interest.

Reaction	Nucleus	E_x [keV]	E_γ [keV]	$T_{1/2}$ [ns]
$^{13}\text{C} + ^{119}\text{Sn}$	^{120}Sn	2284	90	5.55 ± 0.03
$^{13}\text{C} + ^{117}\text{Sn}$	^{118}Sn	2575	254	230 ± 10
$^{18}\text{O} + ^{112}\text{Sn}$	^{128}Ce	207	207	0.30 ± 0.03

Table 5.2: γ -ray transitions used for the normalization of the stopped intensities at different plunger distances. Energies and lifetimes are the adopted ones [68].

This choice not only improves the quality of the spectra, but also eliminates most of the shifted components of long-lived states, which appear as broadly shaped due to the large dispersion in the scattering angle and velocity of the ion. As a result, only the unshifted component of the peak of interest can be fitted, and thus, the $(I_u + I_s)$ normalization is replaced with a different term, which is required to be proportional to the number of nuclei produced in the excited state of interest. The normalization of the stopped intensities is performed by measuring the intensity of transitions from isomeric states ($T_{1/2} \gtrsim 1$ ns). This ensures that the γ rays are always emitted at rest in the stopper, regardless of the target-stopper distance. Thus, the area of this peak is proportional to the total number of ions produced in the reaction. Table 5.2 lists the transitions used for the normalization in the different reactions. When possible, transitions from the reaction channels of interest were preferred, as in the case of the $^{13}\text{C} + ^{117,119}\text{Sn}$ transfer reactions. For the $2n$ transfer case, no isomeric transition was observed with sufficient statistics, therefore, γ rays from fusion-evaporation channels were considered. A similar procedure was successfully applied in Ref. [141].

Before going into the details of the lifetime analysis, it is important to briefly describe the spectra used to extract the areas of the peaks of interest. To exploit all the available statistics without introducing systematic errors, asymmetric γ - γ matrices were considered. On one axis, all the rings of ROSPHERE are used, while on the other one just the outer rings at 37° and 143° are included. The gate is then placed on the total ROSPHERE projection, and the analysis is performed on the projection from the 37° and 143° rings. In this way, one avoids the in-flight component of γ rays that are emitted at relative recoil- γ angles of about 90° , hence with recorded energies that overlap with the unshifted component. Detectors at 70° , 90° and 110° – mostly affected by this issue – are then ruled out. In addition, the coincidence with SORCERER had to be excluded because of the huge cut in statistics (70%), making the lifetime estimation unfeasible.

In the following, the velocity determination and the lifetime analysis are discussed in detail.

Velocity estimation

As discussed previously, the velocity of the reaction products plays a key role in the lifetime analysis. Generally, the velocity of the emitting ion can be estimated either from the kinematical reconstruction provided by particle detectors (if possible), or from the shape of the γ -ray energy distribution. In fact, the energy difference between the shifted and the unshifted component directly depends on the Doppler shift formula. Therefore, by measuring the centroid of the shifted component for each detector ring, one obtains independent measurements of the ion velocity. Such a procedure can be applied if shifted components have peaked (Gaussian-like) distributions which allow for precise fits. Conversely, broad, asymmetric structure prevent a clear evaluation of the centroid, and thus

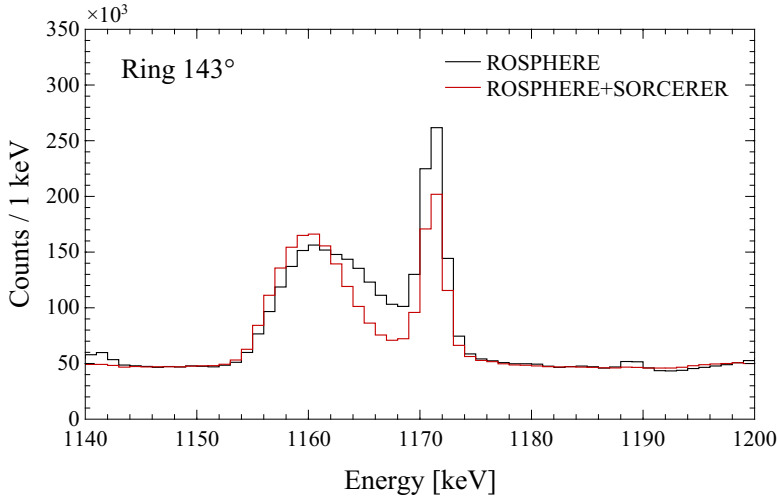


Figure 5.19: Comparison between normalized γ -ray spectra of ^{120}Sn obtained with ROSPHERE only (black line) and with the SORCERER coincidence (red line). Only the backward ring of ROSPHERE, positioned at 143° with respect to the beam direction, is considered.

require a specific treatment.

Figure 5.19 shows the 1171-keV γ -ray of ^{120}Sn observed in the backward ring of ROSPHERE, without and with the coincidence with SORCERER. The Doppler shift of the peak can be directly estimated only when the coincidence with SORCERER is required, as the velocity and the emission angle of the ions are constrained into a smaller range. However, since SORCERER was not employed for the lifetime analysis due to the 70% reduction in statistics that introduced, an alternative method is needed.

In order to measure the velocity of ions without particle coincidences, a Doppler-correction optimization method was applied. At first, a set of average Doppler-correction coefficients was generated for each of the (β, θ_{ion}) point allowed by the kinematics of the reaction. Then, a series of Doppler-corrected histograms was generated, each corresponding to a different average velocity $\langle \beta \rangle$. Two different corrections are compared in Figure 5.20. It clearly appears that a better resolution is obtained for $\langle \beta \rangle = 1.25\%$. Once applied, the quality of the Doppler correction was evaluated with a Figure of Merit (FoM), defined as

$$FoM = \sum_i \sigma_i, \quad (5.3)$$

where σ_i is the standard deviation of the Gaussian fit of the i -th peak. Only peaks associated to short-lived states, i.e. with $T_{1/2} \lesssim 1$ ps, were considered. In the ^{120}Sn case, peaks at 1184, 1230, 1250, 1294 and 1664 keV were fitted, as indicated in Figure 5.20. Finally, the FoM was plotted as a function of the average beta $\langle \beta \rangle$, and fitted with a cubic polynomial, as shown in the left column of Figure 5.21. The minimum of the best fit determined the average beta used for the lifetime analysis. In addition, the average emission angle is extracted from the univocal correlation between β and θ_{ion} provided by the kinematic line. The latter were calculated with the LISE++ suite [118], taking into account the energy losses in the target material.

Unfortunately, this procedure did not provide an immediate way to estimate the uncertainty on the optimal beta. Thus, the uncertainty was extracted using a bootstrap

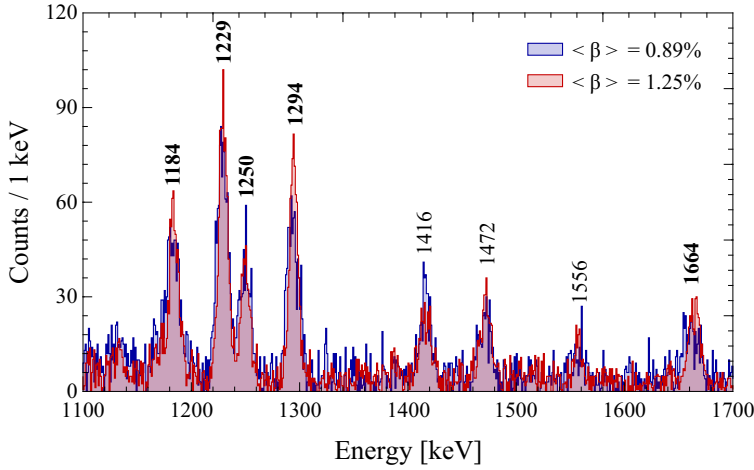


Figure 5.20: Doppler-corrected γ -ray spectra, gated on the shifted component of the $2_1^+ \rightarrow 0_1^+$ transition of ^{120}Sn at 1171 keV. Blue and red colors indicate spectra with Doppler corrections associated to average velocities of 0.89% and 1.25%, respectively. Bold labels indicate the peaks considered for the optimization procedure (see text for details).

Recoil	Channel	β [%]
^{120}Sn	+1n	1.25 ± 0.01
^{118}Sn	-1n	1.28 ± 0.03
^{118}Sn	+1n	1.29 ± 0.01
^{116}Sn	-1n	1.37 ± 0.02
^{114}Sn	+2n	1.93 ± 0.01

Table 5.3: Velocities of the Sn recoils in the different transfer reactions obtained from the optimization procedure of the Doppler correction.

technique [142]. This consists in resampling the experimental data point of the FoM as a function of β , shown in the left column of Figure 5.21, with a Monte Carlo randomization. At each iteration, a new data set was generated by sampling each point 100000 times from a Gaussian distribution, with centroid and width equal to the experimental value and its error, respectively. Then, cubic polynomial fits provided the minimum beta of each simulated set. The final beta distributions are presented in the right column of Figure 5.21, and their standard deviation is used as uncertainty on the velocity. The results obtained for each reaction are summarized in Table 5.3.

Results for ^{114}Sn

The 0_2^+ state at 1953 keV In Figure 5.22, the projected spectra of ^{114}Sn for the 14.5, 22.5, 42.5 and 74.5 μm plunger distances, obtained after gating on the stopped component of the 1300-keV $2_1^+ \rightarrow 0_1^+$ peak, are shown. Despite the good statistics collected in the $0^+ \rightarrow 2_1^+$, γ -ray transition at 653 keV, background fluctuations introduced uncertainties up to 30% in the peak area estimation. Normalized areas were extracted for each distance

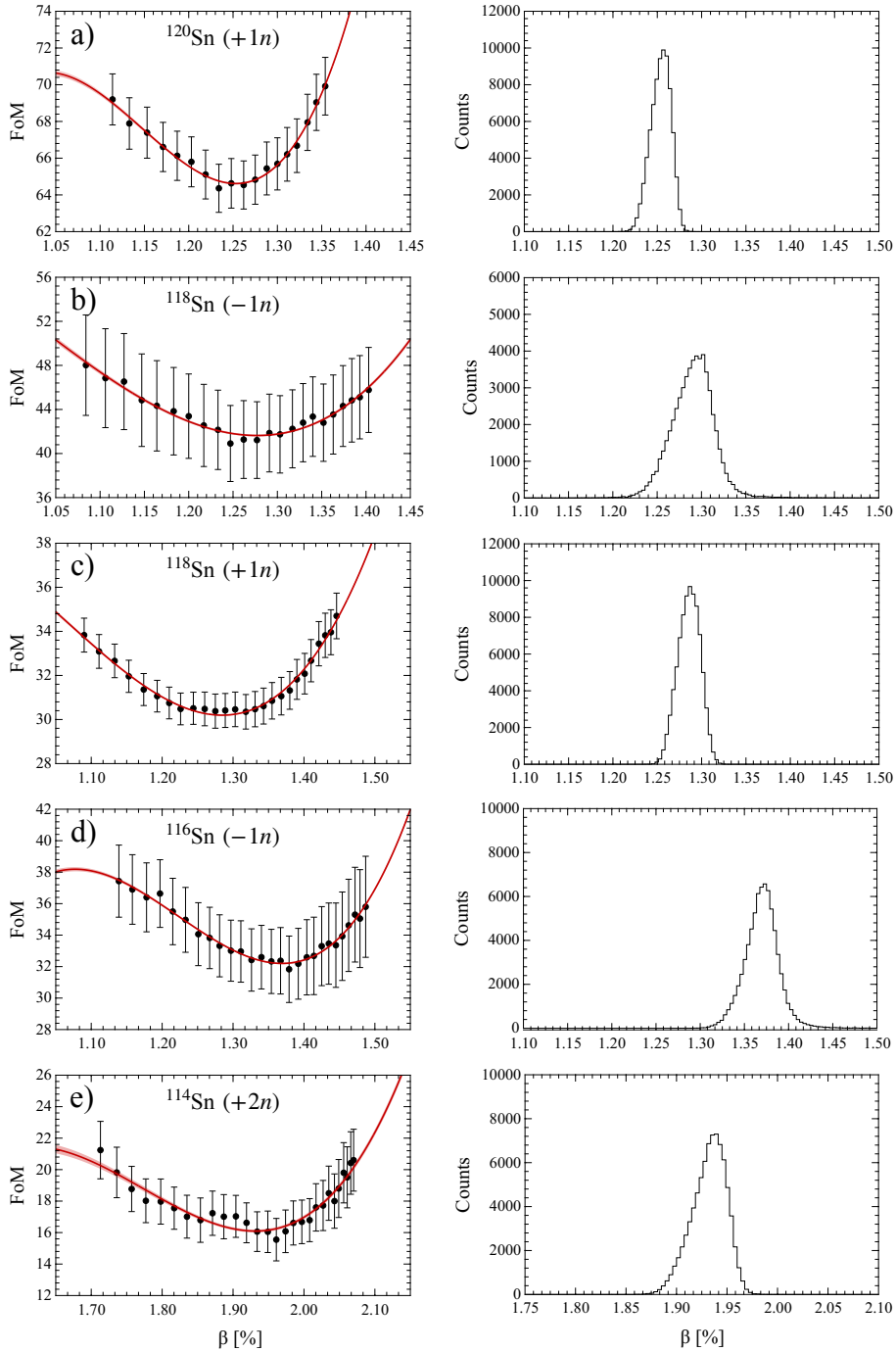


Figure 5.21: Optimization of the average velocity $\langle \beta \rangle$ of the recoil Sn ions. Each row corresponds to Sn products produced by different reaction channels. (Left) FoM as a function of the average velocity used for the Doppler correction. Experimental points and errors are indicated in black, while the cubic polynomial fit is in red. (Right) Velocity distributions as obtained with the bootstrap method [142] (see text for details).

and fitted with an exponential decay curve, as shown in the top-left panel of Figure 5.23. Indeed, no feeding component was observed in both the thick-target and the plunger datasets. From the DCM fit, a half-life of 11.1 ± 5.6 ps is obtained.

The obtained half-life result is affected by a significant error due to the large uncertainties in the experimental data points. This suggests that the symmetric error estimate obtained with the conventional likelihood minimization procedure [143], adopted here, may be too simplistic. Therefore, uncertainties were estimated using a Monte-Carlo bootstrap resampling technique [142]. Each experimental data point, i.e. the normalized areas for each plunger distance, was resampled 100000 times from a Gaussian distribution defined by its measured value and experimental uncertainty. For every resampled dataset, the exponential decay curve was refitted, and a corresponding half-life value extracted. The resulting ensemble of $T_{1/2}$ values forms a skewed distribution with a sizable right tail, as shown in the top-right panel of Figure 5.23. The final result of $T_{1/2} = 11.0^{+9.3}_{-4.0}$ ps is quoted in terms of the median and the central 68% interval, using the 16%-84% percentiles as lower and upper boundaries, respectively. We note that the median of the bootstrap distribution and the value from the DCM fit are well aligned. This bootstrap procedure is applied consistently to all the subsequent half-life estimations.

The half-life result for the 0_2^+ state can also be compared with the value of 6.5 ± 2.3 ps deduced by Bäcklin et al. [87] from Coulomb excitation, which is about 40% smaller. However, considering the large error bars of our measurement, the two results remain compatible.

The 0_3^+ state at 2156 keV Similarly to the 0_2^+ case, the high background level prevented a precise estimation of the area of the 856-keV peak, depopulating the 0_3^+ of ^{114}Sn , at different plunger distances. This is indeed reflected in the large error bars observed in the normalized area, as reported in the bottom panel of Figure 5.23.

No feeding transitions were observed in this work. The 1405-keV transition, observed in coincidence with the 856-keV transition [68, 144], de-excites the 3561-keV levels, which was not populated in the employed $2n$ transfer reaction.

Through the DCM analysis, a half-life of 13.0 ± 5.9 ps is extracted from the exponential decay fit. As in the 0_2^+ case, a more reliable estimate of the uncertainty is obtained from the bootstrap distribution of the half-life, shown in the bottom-right panel of Figure 5.23. Since the distribution is significantly skewed, the asymmetric uncertainties are determined from the central 68% interval, and the half-life result for the 0_3^+ state becomes $T_{1/2} = 13.0^{+9.4}_{-4.2}$ ps. A good agreement is found with the lower limit of $T_{1/2} > 7.6$ ps proposed by Bäcklin et al. [87].

Results for ^{116}Sn

The 0_2^+ state at 1757 keV In Figure 5.24, the projections of the γ - γ for the $^{13}\text{C} + ^{117}\text{Sn}$ reaction, gated on the stopped component of the 1294-keV $2_1^+ \rightarrow 0_1^+$ transition of ^{116}Sn , are shown as a function of the plunger distance. Among the five distances, only the first 3 were employed to estimate the lifetime of the excited 0_2^+ state of ^{116}Sn , which primarily appeared at shifted energies at 500 and 750 μm .

No significant feeding contributions to the 0_2^+ state were expected or measured in our spectroscopic analysis. The only feeding that could eventually affect our measurement comes from the 271-keV, $E0$ transition, depopulating the 0_3^+ state. Its intensity is currently unknown, however, based on the systematics in the Sn chain ^{118}Sn [14, 87, 104,

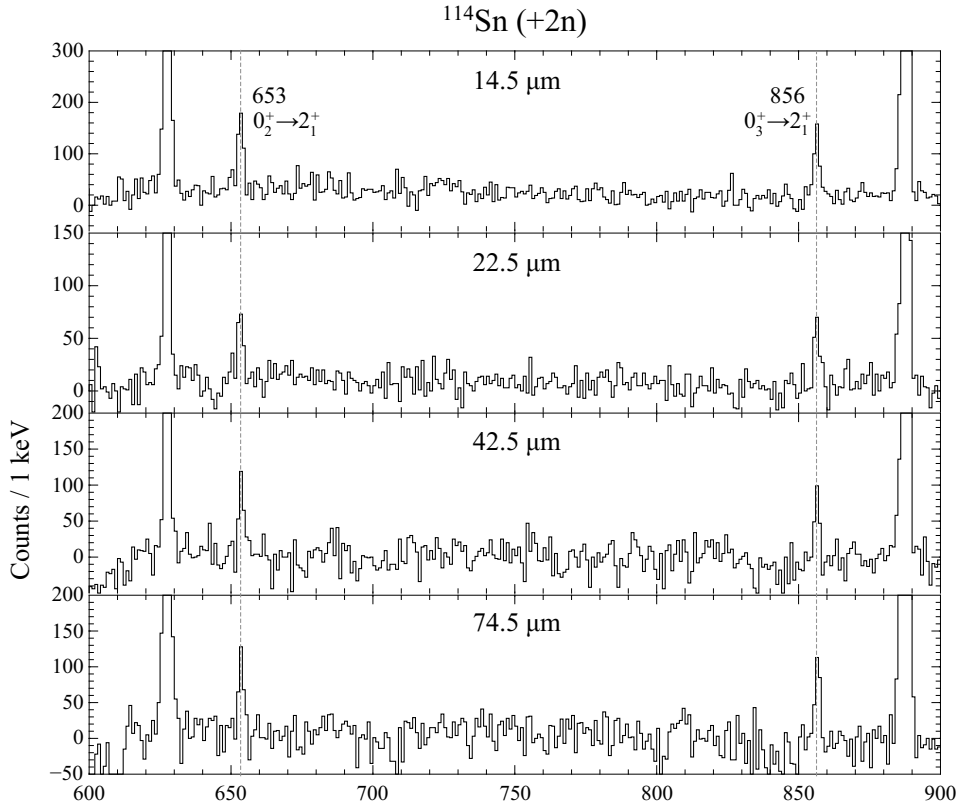


Figure 5.22: Partial γ -ray spectra of ^{114}Sn for different plunger distances. The spectra are projections of the γ - γ matrices produced in the $^{18}\text{O} + ^{112}\text{Sn}$ reaction, gated on the stopped component of the $2_1^+ \rightarrow 0_1^+$ transition at 1300 keV. Grey dashed lines indicate the stopped components of the transitions at 653 and 856 keV, depopulating the 0_2^+ and 0_3^+ states, respectively.

[135], one could expect values of the order of a few percent. As such, the error introduced will be well within the statistical error bars of our measurement.

The top-left panel of Figure 5.25 shows the results of the DCM analysis. Despite the limited number of data points, the fitting procedure provided error bars of about 25%. On the other hand, the bootstrap resampling method, used to improve the uncertainty estimation, provides an asymmetric distribution of $T_{1/2}$. Using the median and the central 68% interval of this distribution, shown in the top-right panel of Figure 5.25, a half-life of $T_{1/2} = 51.1^{+18.8}_{-11.7}$ ps is obtained for the 0_2^+ state. We note that the bootstrap median and the fitted half-life from the DCM analysis are again consistent, indicating that the bootstrap procedure only affects the uncertainty determination.

This result is in good agreement with the previous Coulomb excitation measurement of 44 ± 7 from Ref. [87], as well as the fast-timing results of 45 ± 10 and 51 ± 12 reported in Refs. [145] and [78], respectively.

The 0_3^+ state at 2027 keV After gating on the stopped component of the 1294-keV transition, depopulating the first excited state in ^{116}Sn , the $0_3^+ \rightarrow 2_1^+$ transition at 734 keV can be clearly identified. At all the target-stopper distances, the stopped component emerges

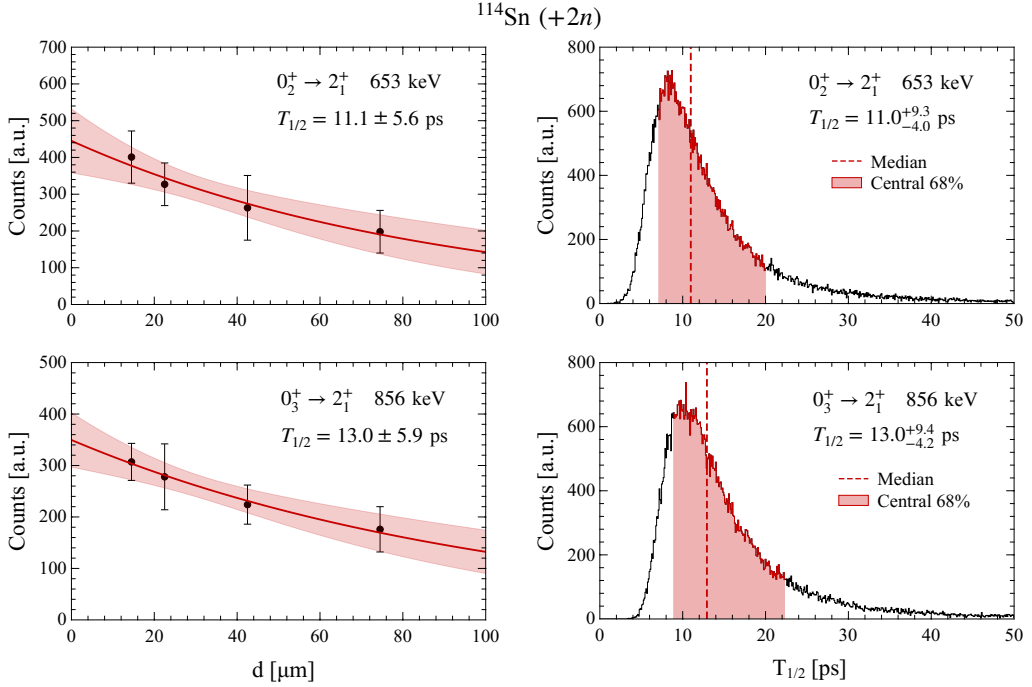


Figure 5.23: (Left) Normalized area as a function of the plunger distance for the stopped component of the $0_2^+ \rightarrow 2_1^+$ (top) and $0_3^+ \rightarrow 2_1^+$ (bottom) transitions of ^{114}Sn , populated in the $+2n$ transfer reaction. Half-lives are extracted with the DCM method using an exponential decay fit (red lines). Red shaded areas correspond to the 68% error bands of the fit. (Right) Monte Carlo distributions of the half-lives based on the normalized intensities of the $0_2^+ \rightarrow 2_1^+$ (top) and $0_3^+ \rightarrow 2_1^+$ (bottom) transitions of ^{114}Sn (see text for details). Red dashed lines indicate the median of the distributions, while red shaded areas correspond to the central 68% regions.

from the flat background and almost no hints of shifted structures appear on the sides.

In the thick-target dataset, the 1201-keV γ -ray transition, depopulating the 2^+ state at 3228 keV, was measured in coincidence with the 734-keV transition. A direct feeding of $25 \pm 6\%$ was determined. However, the 1201-keV peak appeared as partially shifted when detected in the forward and backward angles, thus pointing to lifetimes in the ps range. Furthermore, by looking at the summed statistics of all the plunger distances, hints of shifted components appear around 1201 keV.

In the fitting procedure of the DCM analysis, the presence of a feeder was taken into account, using a conservative upper limit of 10 ps for the lifetime of the 3228-keV state. Yet, its influence on the fit turned out to be negligible, as the time-of-flight range for this measurement was comprised between 30 and 200 ps. The observed stopped intensities and the decay fit are depicted in the bottom-left panel of Figure 5.25. A half-life of 184 ± 50 ps is extracted for the 0_3^+ state via the DCM fit. After applying the bootstrap procedure, asymmetric uncertainties are obtained from the central 68% interval of the skewed half-life distribution (see the bottom-right panel of Figure 5.25), resulting in $T_{1/2} = 184^{+66}_{-40}$ ps. Although slightly larger, this result agrees with both the adopted value of 160 ± 20 ps [68, 145] and the recent measurement of 164 ± 10 ps [78], both determined with fast-timing techniques.

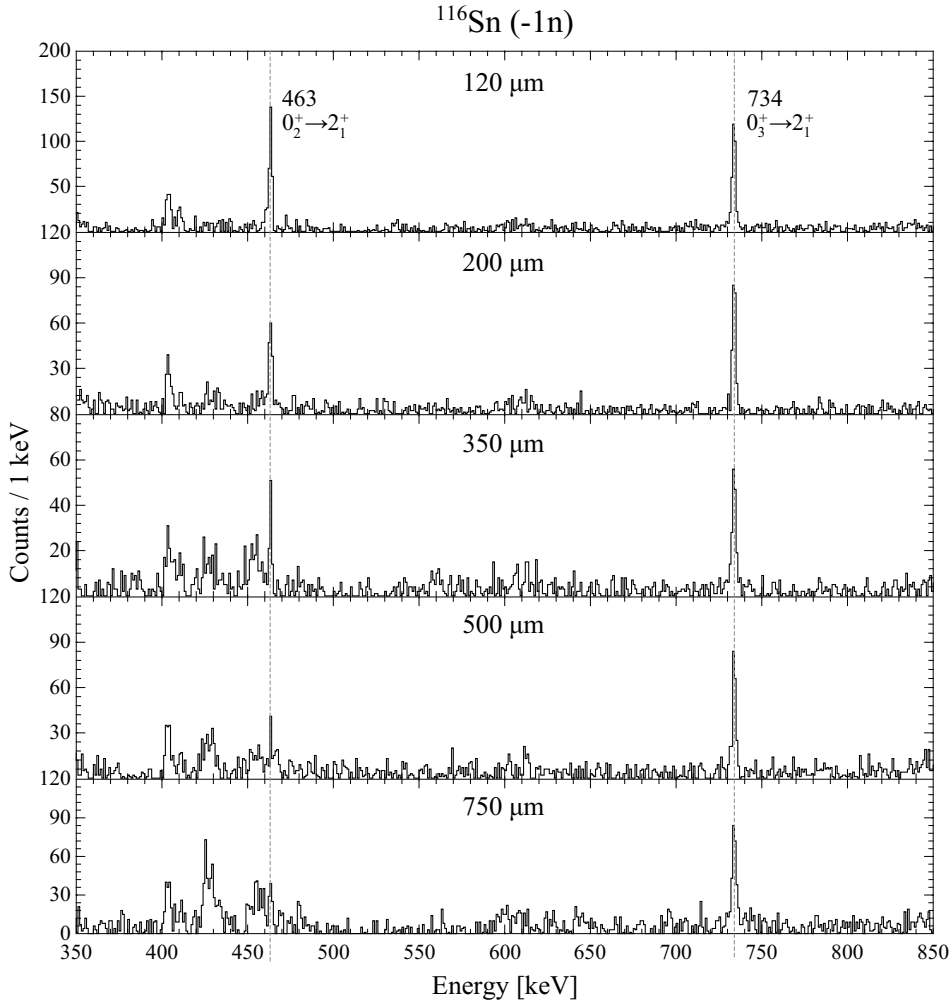


Figure 5.24: Partial γ -ray spectra of ^{116}Sn for different plunger distances. The spectra are projections of the γ - γ matrices produced in the $^{13}\text{C} + ^{117}\text{Sn}$ reaction, gated on the stopped component of the $2_1^+ \rightarrow 0_1^+$ transition at 1294 keV. Grey dashed lines indicate the stopped components of the transitions at 463 and 734, depopulating the 0_2^+ and 0_3^+ states, respectively.

Results for ^{118}Sn

The 0_2^+ state at 1758 keV Figure 5.26 shows the plunger γ -ray spectra of ^{118}Sn as observed in the $-1n$ channel of the $^{13}\text{C} + ^{119}\text{Sn}$ reaction, after gating on the stopped component of the 1230-keV $2_1^+ \rightarrow 0_1^+$ transition. In particular, the transition at 529 keV, depopulating the 0_2^+ state at 1758 keV, is well separated and identifiable up to the 95.5 μm target-stopper distance, after which it becomes hidden below the background.

In the top-left panel of Figure 5.27, the normalized counts for each distance and the DCM fit are shown. Being the intensities of the feeding transitions from the 2_2^+ and 0_3^+ transition at the level of 0.8% and 2.5%, respectively [77, 135], and the direct population cross section similar to that of the 0_3^+ state, their contribution was considered to be negligible.

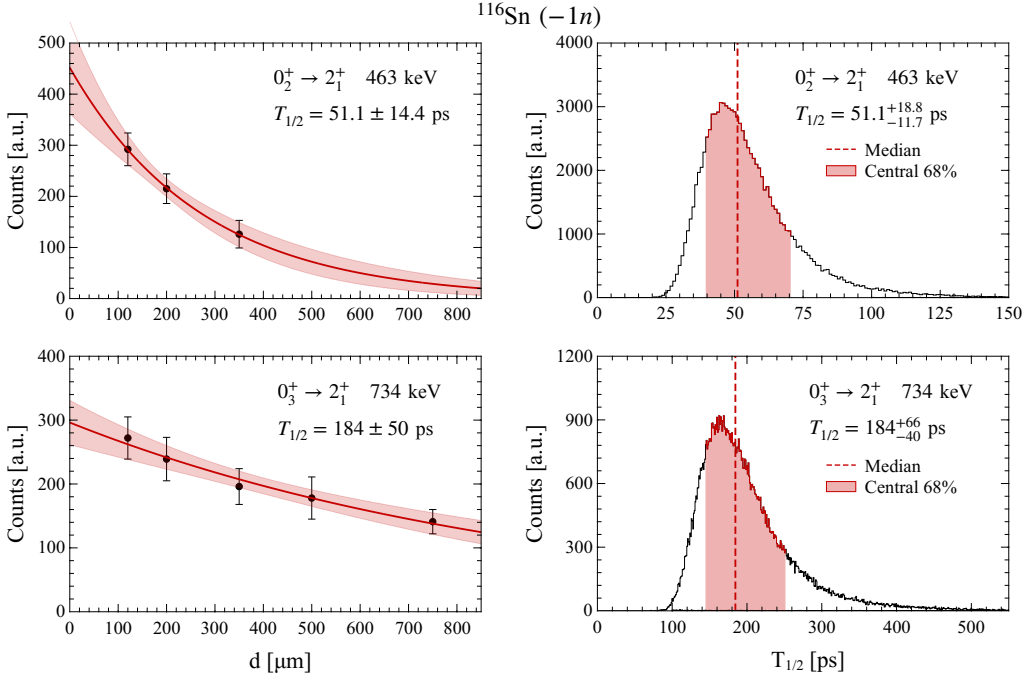


Figure 5.25: (Left) Normalized area as a function of the plunger distance for the stopped component of the $0_2^+ \rightarrow 2_1^+$ (top) and $0_3^+ \rightarrow 2_1^+$ (bottom) transitions of ^{116}Sn , populated in the $-1n$ transfer reaction. Half-lives are extracted with the DCM method using an exponential decay fit (red lines). Red shaded areas correspond to the 68% error bands of the fits. (Right) Monte Carlo distributions of the half-life based on the normalized intensities of the $0_2^+ \rightarrow 2_1^+$ (top) and $0_3^+ \rightarrow 2_1^+$ (bottom) transitions of ^{116}Sn (see text for details). Red dashed lines indicate the median of the distributions, while red shaded areas correspond to the central 68% regions.

The fitting procedure provides a half-life of $T_{1/2} = 18.2 \pm 2.2$ ps. In this case, being the relative error just above 10%, a minor deviation from the Gaussian shape is observed in the bootstrap distribution of the half-life, which is shown in the top-right panel of Figure 5.27. The outcome of the bootstrap procedure, expressed as median and central 68% interval, is $T_{1/2} = 18.2^{+2.4}_{-2.0}$ ps. Such a value is slightly lower than the 21 ± 3 ps half-life obtained by Bäcklin et al. with Coulomb excitation [87], but still in good agreement with it.

The 0_3^+ state at 2057 keV For the lifetime measurement of the 0_3^+ state of ^{118}Sn , both of the $^{119}\text{Sn}(^{13}\text{C}, ^{14}\text{C})^{118}\text{Sn}$ ($-1n$ channel) and the $^{117}\text{Sn}(^{13}\text{C}, ^{12}\text{C})^{118}\text{Sn}$ ($+1n$ channel) reaction datasets were exploited. In the first case, 7 plunger distances, ranging from 13.5 μm to 300.5 μm , could be used. The spectra in Figure 5.26 show the good statistics collected in the stopped component of the peaks, as well as a low background level on its sides. Due to the presence of the $2_2^+ \rightarrow 2_1^+$ transition at 813 keV, which depopulates the 2043-keV level with half-life of 2.9 ± 0.4 ps, shifted components appear on the left side of the $0_3^+ \rightarrow 2_1^+$ peak.

The result obtained in the $-1n$ channel is depicted in the middle-left panel of 5.27. Based on what we observed in the thick-target dataset and on the γ -ray branches measured in literature [68, 135], no feeding contribution was found nor included. Minor

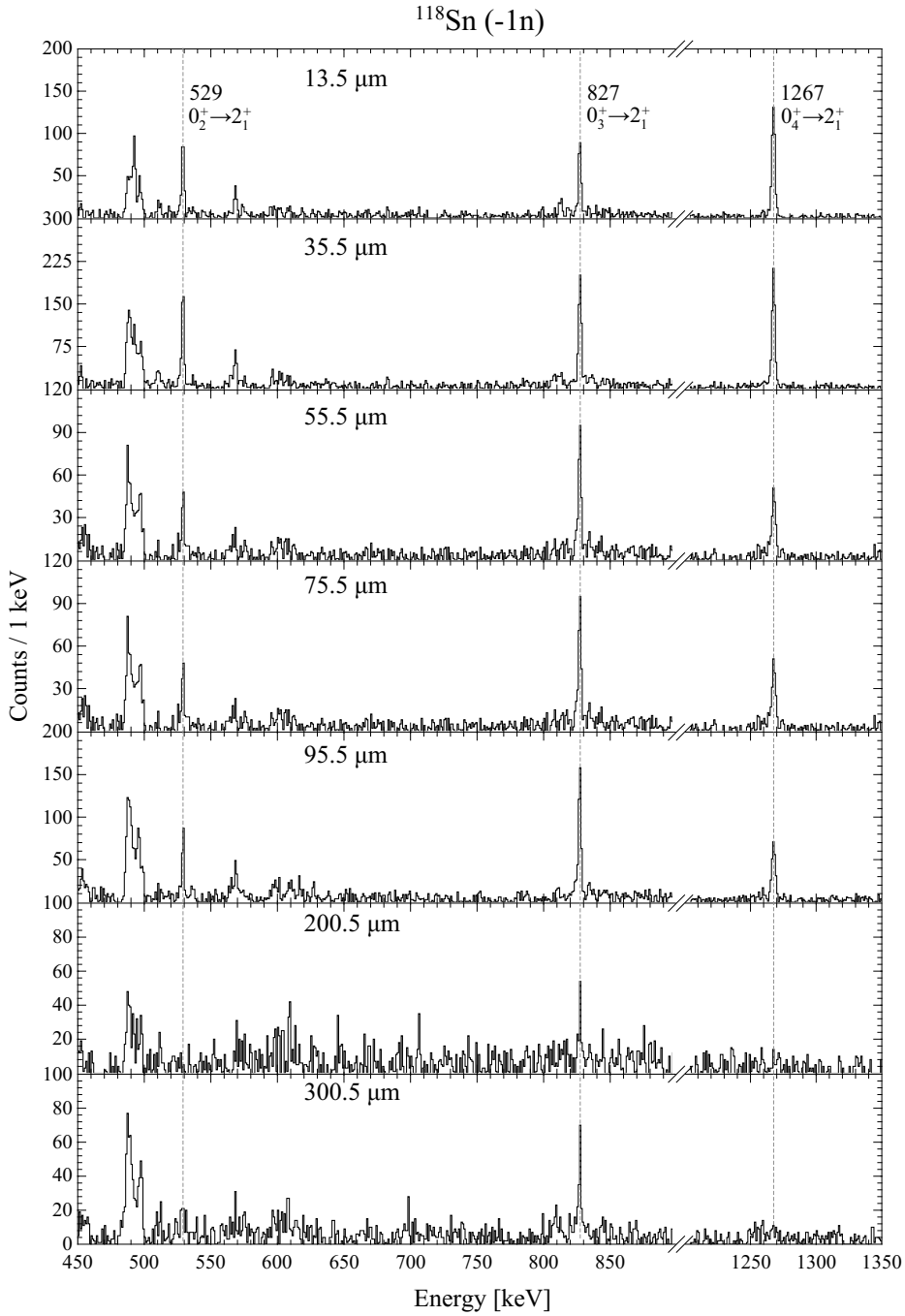


Figure 5.26: Partial γ -ray spectra of ^{118}Sn for different plunger distances. The spectra are projections of the γ - γ matrices produced in the $^{13}\text{C} + ^{119}\text{Sn}$ reaction, gated on the stopped component of the $2_1^+ \rightarrow 0_1^+$ transition at 1230 keV. Grey dashed lines indicate the stopped components of the transitions at 529, 827 and 1267 keV, depopulating the 0_2^+ , 0_3^+ and 0_4^+ states, respectively.

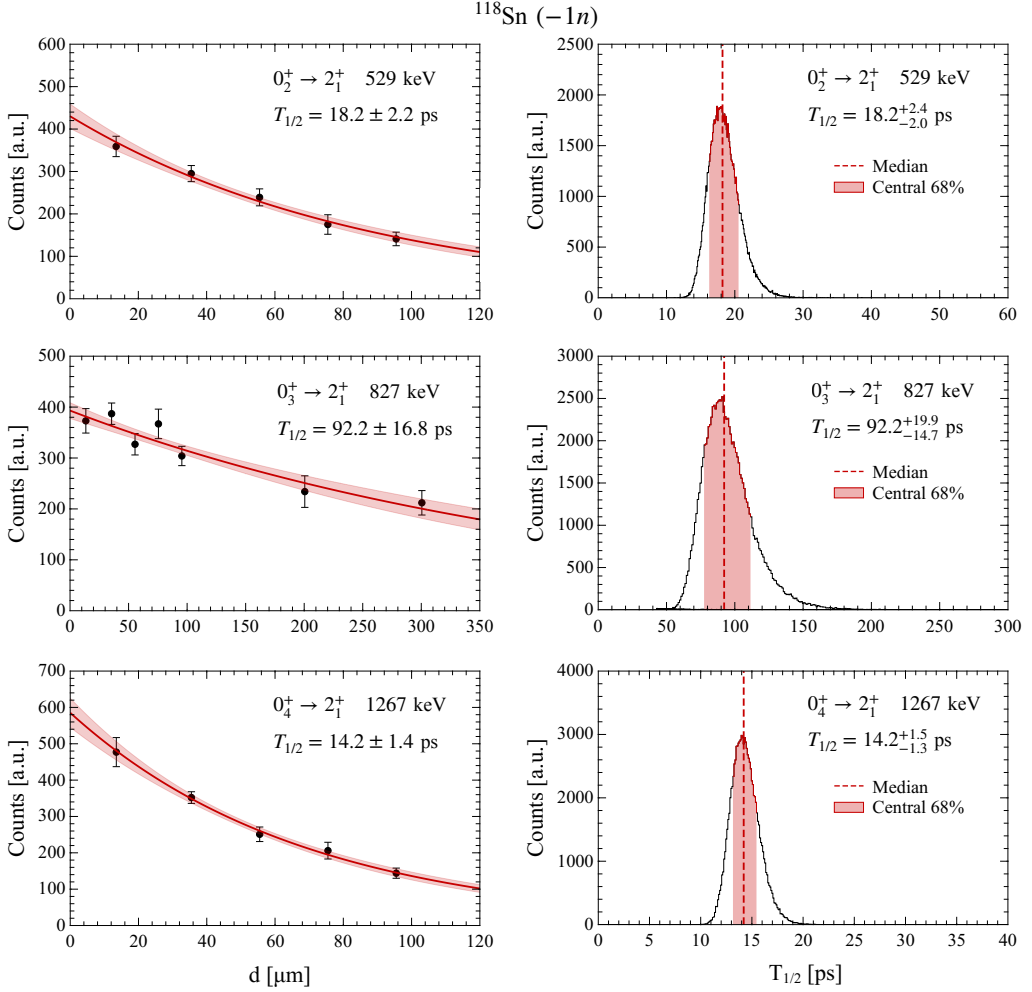


Figure 5.27: (Left) Normalized area as a function of the plunger distance for the stopped component of the $0_2^+ \rightarrow 2_1^+$ (top), $0_3^+ \rightarrow 2_1^+$ (middle) and $0_4^+ \rightarrow 2_1^+$ (bottom) transitions of ^{118}Sn , populated in the $-1n$ transfer reaction. Half-lives are extracted with the DCM method using an exponential decay fit (red lines). Red shaded areas correspond to the 68% error bands of the fits. (Right) Monte Carlo distributions of the half-life based on the normalized intensities of the $0_2^+ \rightarrow 2_1^+$ (top), $0_3^+ \rightarrow 2_1^+$ (middle) and $0_4^+ \rightarrow 2_1^+$ (bottom) transitions of ^{118}Sn (see text for details). Red dashed lines indicate the median of the distributions, while red shaded areas correspond to the central 68% regions.

fluctuations can be observed at small distances, where the experimental points are close to each other and the measurement was less sensitive. Nevertheless, the deviations from the fit are compatible with statistical fluctuations. From the DCM analysis of this dataset, a half-life of 92.2 ± 16.8 ps is obtained. By performing the bootstrap resampling procedure, a slightly skewed half-life distribution is produced, as shown in the middle-right panel of Figure 5.27. Although the uncertainties become asymmetric, the bootstrap median remains consistent with the DCM fit. The half-life result is then quoted as $T_{1/2} = 92.2^{+19.9}_{-14.7}$ ps.

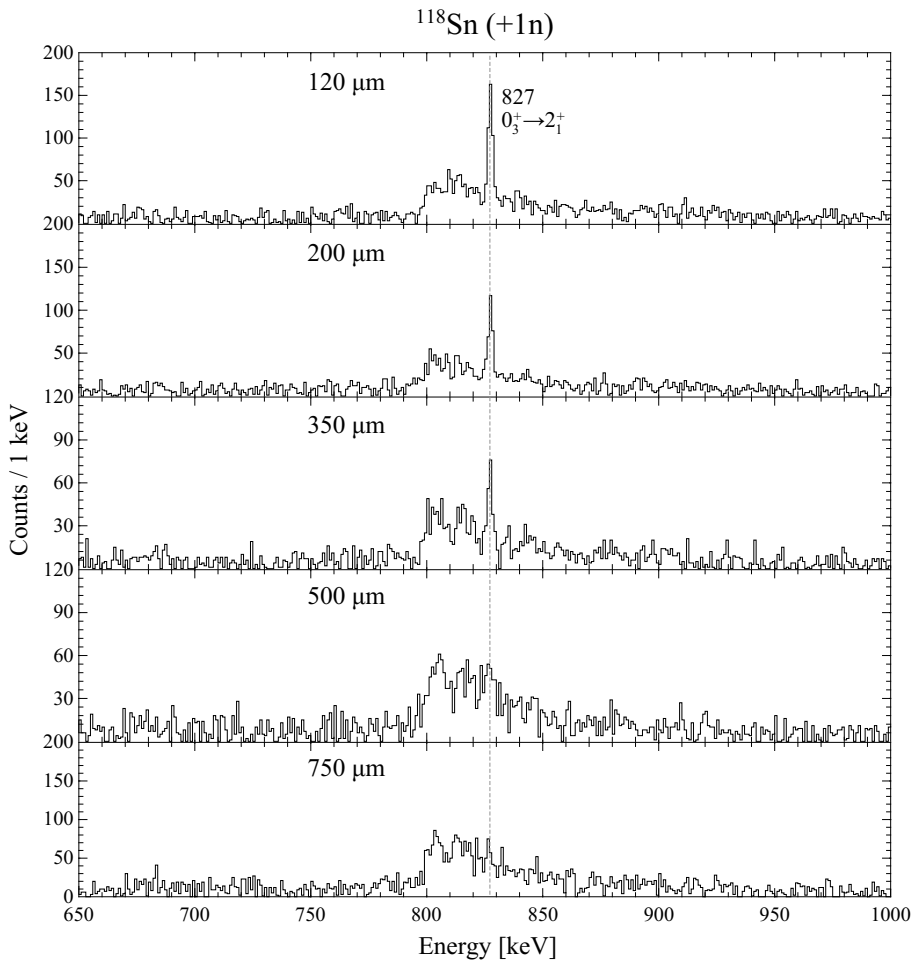


Figure 5.28: Partial γ -ray spectra of ^{118}Sn for different plunger distances. The spectra are projections of the γ - γ matrices produced in the $^{13}\text{C} + ^{117}\text{Sn}$ reaction, gated on the stopped component of the $2_1^+ \rightarrow 0_1^+$ transition at 1230 keV. The gray dashed line indicates the stopped components of the 827-keV transition depopulating the 0_3^+ state.

For the analysis of the 0_3^+ state in ^{118}Sn , target-stopper distances of 120, 200, 350, 500 and 750 μm were also available in the $+1n$ dataset. The corresponding γ -ray spectra, gated on the stopped component of the lower-lying, 1230-keV transition, are shown in Figure 5.28. Based on the measurement from the $-1n$ reaction channel, about 70% of the γ -rays should be emitted in-flight at 500 μm (time-of-flight of ~ 150 ps). Indeed, no clear stopped peak is observed in the 500- and 750- μm spectra, so only the first 3 distances were considered. In addition, shifted components of the 813-keV peak are clearly visible and extend up to a few keV before the 827-keV peak. Still, the area of the peak can be extracted by carefully including all the components.

The normalized stopped areas and the exponential fit are shown in the left panel of Figure 5.29. The half-life result from the DCM method is $T_{1/2} = 81.9 \pm 47.7$ ps. The lifetime uncertainty was also evaluated using the bootstrap procedure described previously, which in this case is particularly significant due to the limited number of data points and

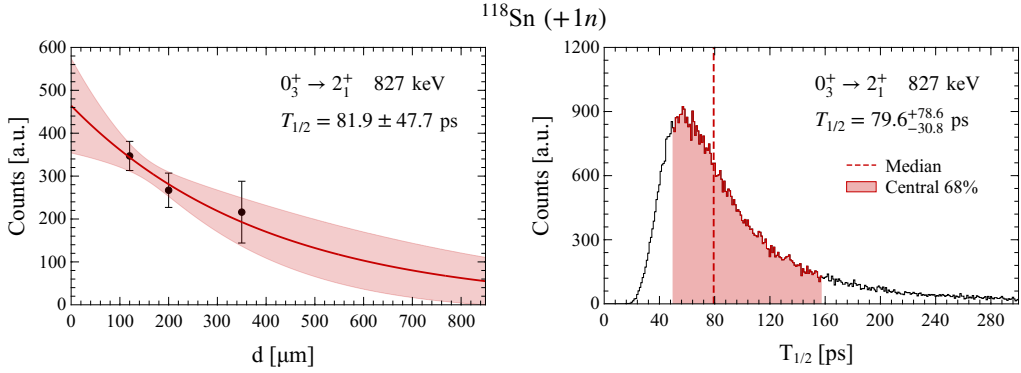


Figure 5.29: (Left) Normalized area as a function of the plunger distance for the stopped component of the $0_3^+ \rightarrow 2_1^+$ transition of ^{118}Sn , populated in the $+1n$ transfer reaction. The half-life is extracted with the DCM method using an exponential decay fit (red lines). Red shaded areas correspond to the 68% error bands of the fit. (Right) Monte Carlo distributions of the half-life based on the normalized intensities of the $0_3^+ \rightarrow 2_1^+$ (top) transition of ^{118}Sn (see text for details). The red dashed line indicates the median of the distribution, while the red shaded area corresponds to the central 68% region.

the large errors on the last one. The bootstrap distribution of the half-life is shown in the right panel of Figure 5.29. Using the median and central 68% interval, the result becomes $T_{1/2} = 79.6^{+78.6}_{-30.8}$ ps. The bootstrap median and the DCM fit slightly differ, likely due to the large experimental uncertainties of this dataset.

Although results from the $-1n$ and $+1n$ datasets are compatible within uncertainties, the $+1n$ dataset shows significantly larger error bars. Therefore, we decide to quote the result of the $-1n$ dataset, $T_{1/2} = 92.2^{+19.9}_{-14.7}$ ps, as the half-life of the 2057-keV 0_3^+ state of ^{118}Sn .

The 0_4^+ state at 2497 keV As discussed in Section 5.13, the cross section for the 0_4^+ state in the $-1n$ transfer reactions was by far the largest with respect to the other excited 0^+ states. This is reflected in the statistics collected in the stopped component of the γ -ray peak at 1267 keV, of the order of a few hundred counts per distance, as shown in Figure 5.26. Specifically, the evolution of the stopped intensity of the $0_4^+ \rightarrow 2_1^+$ decay appears similar to that of the $0_2^+ \rightarrow 2_1^+$ transition, with the peak disappearing after the $95.5 \mu\text{m}$ distance was reached. Thus, the first 5 distances were considered for its lifetime estimation. Moreover, γ - γ and particle- γ - γ coincidences from the thick-target dataset of *Exp. 3* showed no other coincidence than the one with the γ -ray de-exiting the first 2^+ excited state. Indeed, the feeding observed from Ref. [135] comes from higher-lying states with $E_x > 3300$ keV (and with γ -ray branches to the state of interest generally smaller than 20%), which were not populated in the $^{119}\text{Sn}(^{13}\text{C}, ^{14}\text{C})^{118}\text{Sn}$ reaction of this work.

In the bottom-left panel of Figure 5.27, DCM data points and fit are shown. A value of 14.2 ± 1.4 ps is obtained for the half-life of the 0_4^+ state at 2497 keV, which was here measured for the first time. Similarly, when using the bootstrap method to improve the uncertainty estimation (see the bottom-right panel of Figure 5.27), a half-life $T_{1/2} = 14.2^{+1.5}_{-1.3}$ ps is determined from the median and central 68% interval of the distribution. This demonstrates that the DCM analysis and the bootstrap procedure tend to converge

when the relative errors are less than or equal to 10%.

Results for ^{120}Sn

The 0_2^+ state at 1875 keV As it can be seen in Figure 5.30, the data collected in the $0_2^+ \rightarrow 2_1^+$ γ -ray peak of ^{120}Sn is around 100 events per distance. Particular care was dedicated to the background fitting. Indeed, the 704-keV peak lies on a background generated by the shifted components of the 692-keV peak of ^{120}Sn and the 701-keV peak of ^{129}Ba (fusion-evaporation product), which increases as a function of the distance. Since the area of the stopped component rapidly decreased, only the first 3 distances at 13.5, 35.5 and 55.5 μm were considered.

In the top-left panel of Figure 5.31, the result of the fitting procedure is shown, along with 68% error bands. Since no feeding contributions were observed in the thick target experiment and the branching ratio of the $0_3^+ \rightarrow 0_2^+$ $E0$ transition is negligible [81, 87, 104], an exponential decay fit was considered. The DCM half-life for this state is $T_{1/2} = 8.2 \pm 2.0$ ps.

An improved estimation of the half-life uncertainty can be obtained with the bootstrap method. According to the median and the central 68% interval of the half-life distribution, shown in the top-right panel of Figure 5.31, the half-life of the 0_2^+ state of ^{120}Sn is $T_{1/2} = 8.2_{-1.7}^{+2.6}$ ps. Good agreement is found between this result and the previous Coulomb excitation measurement of Bäcklin et al. [87], which deduced a half-life of 7.4 ± 1.0 ps.

The 0_3^+ state at 2160 keV In the 989-keV transition depopulating the excited 0_3^+ state, few hundreds of counts were observed for each plunger distance. As shown in Figure 5.30, its stopped component remained clearly visible and isolated up to the 95.5 μm distance, while background fluctuations are more evident for the 200 and 300 μm case.

The normalized counts of the stopped components are shown in the middle panel of Figure 5.30, as a function of the target-stopper distance. Although a decay trend can be observed from the first 5 distances, the lifetime of the state is heavily influenced by the 200.5- and 300.5- μm points.

From the spectroscopic analysis, feeding transitions were not observed, therefore they are excluded from the fitting procedure. A half-life of 40.4 ± 7.2 ps is obtained from the DCM analysis. As shown in the middle-right panel of Figure 5.31, by applying the bootstrap resampling method, a slightly skewed distribution is found, and, as a consequence, a better estimation of the uncertainty involves the use of the central 68% interval. From this procedure, a half-life of $40_{-6.2}^{+7.9}$ ps is determined for the 0_3^+ state of ^{120}Sn . This result can be compared with the recent measurement of Wu et al. [81]. Using a fast-timing technique after a neutron capture reaction, they measured a half-life $T_{1/2} = 34.7 \pm 4.9$, which is compatible with our result.

The 0_4^+ state at 2587 keV Among the excited 0^+ states populated in the $^{119}\text{Sn}(^{13}\text{C}, ^{12}\text{C})^{120}\text{Sn}$, the 0_4^+ had the highest cross section. The good statistics and the cleanliness of the spectra after the γ coincidence make this state an optimal candidate for a lifetime measurement. Distances up to 95.5 μm were considered, as the stopped component becomes no longer visible from the 200.5 μm spectra (see Figure 5.30).

No coincidence with transitions depopulating higher lying states was measured in our data nor in other experiments from literature. As a consequence, an exponential decay curve was fitted and a half-life of 14.8 ± 1.1 ps was determined for the first time.

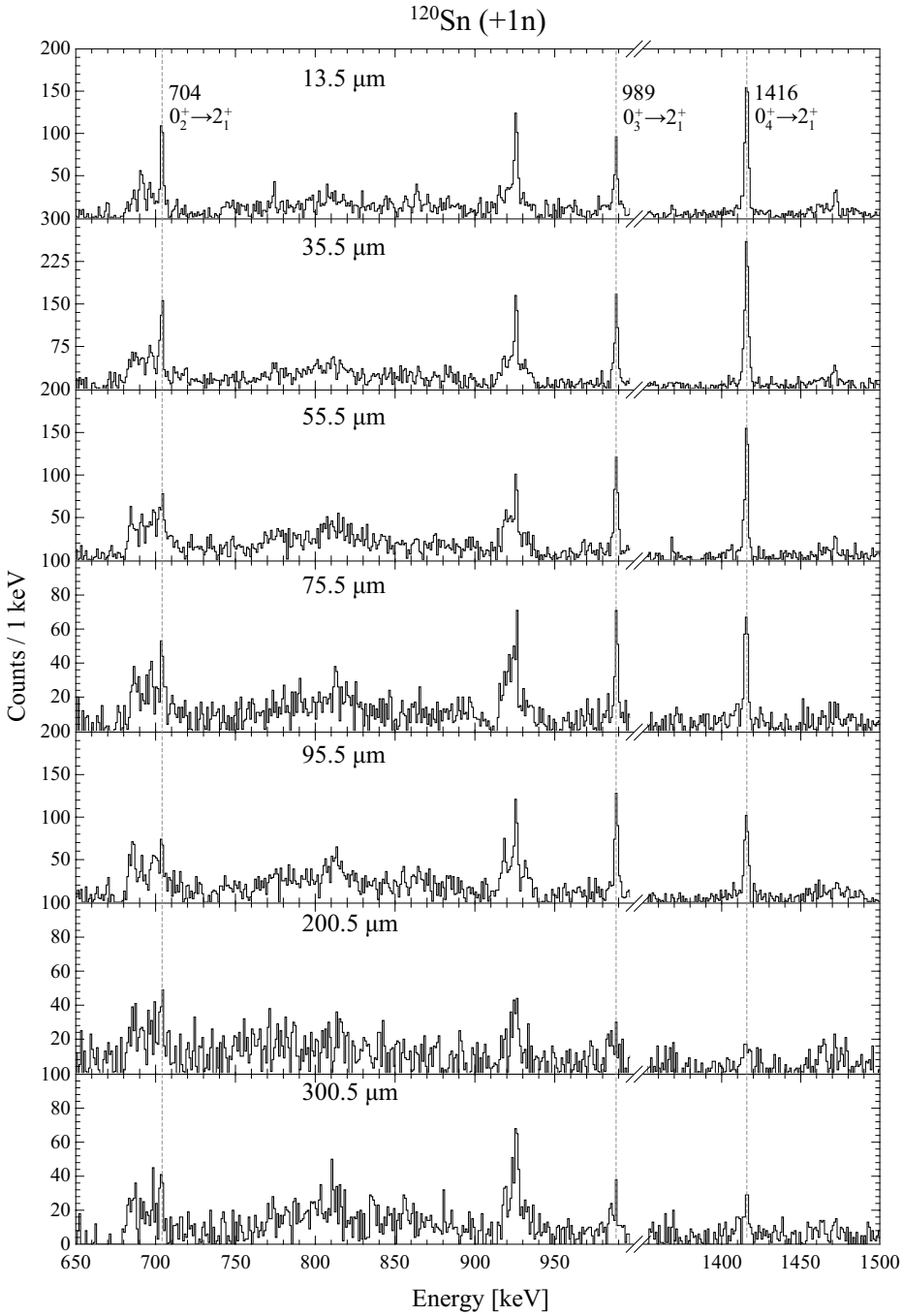


Figure 5.30: Partial γ -ray spectra of ^{120}Sn for different plunger distances. The spectra are projections of the γ - γ matrices produced in the $^{13}\text{C} + ^{119}\text{Sn}$ reaction, gated on the stopped component of the $2_1^+ \rightarrow 0_1^+$ transition at 1171 keV. Grey dashed lines indicate the stopped components of the transitions at 704, 989 and 1416 keV, depopulating the 0_2^+ , 0_3^+ and 0_4^+ states, respectively.

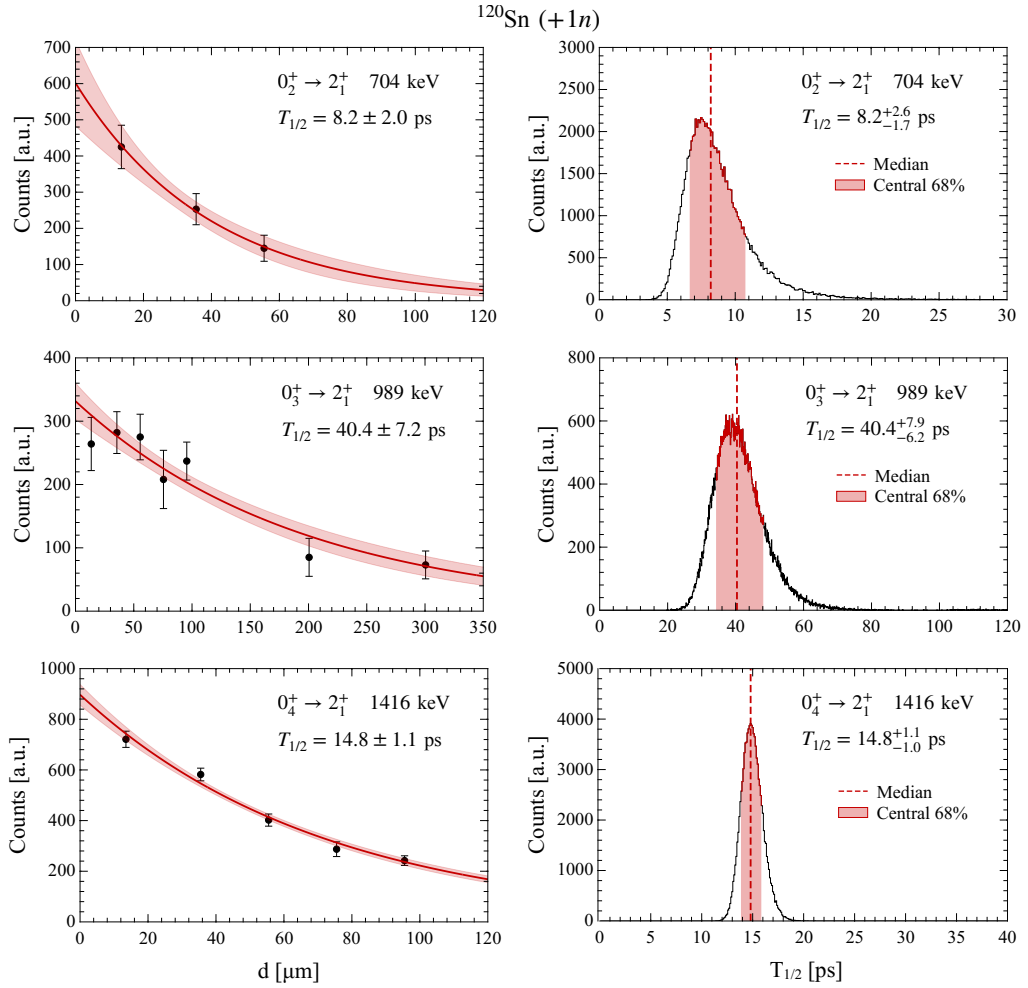


Figure 5.31: (Left) Normalized area as a function of the plunger distance for the stopped component of the $0_2^+ \rightarrow 2_1^+$ (top), $0_3^+ \rightarrow 2_1^+$ (middle) and $0_4^+ \rightarrow 2_1^+$ (bottom) transitions of ^{120}Sn , populated in the $+1n$ transfer reaction. Half-lives are extracted with the DCM method by an exponential decay fit (red lines). Red shaded areas correspond to the 68% error bands of the fits. (Right) Monte Carlo distributions of the half-life based on the normalized intensities of the $0_2^+ \rightarrow 2_1^+$ (top), $0_3^+ \rightarrow 2_1^+$ (middle) and $0_4^+ \rightarrow 2_1^+$ (bottom) transitions of ^{120}Sn (see text for details). Red dashed lines indicate the median of the distributions, while red shaded areas correspond to the central 68% regions.

The experimental points and the DCM fit are shown in the bottom-left panel of Figure 5.31. Although the relative uncertainties obtained from the DCM fit are well below 10%, the bootstrap procedure was applied for consistency with the other analyses. As shown in the bottom-right panel of Figure 5.31, the half-life distribution is almost Gaussian, with a minor tail on the right side. Using the median and central 68% interval of this distribution, the final result for the half-life of the 0_4^+ state of ^{120}Sn becomes $14.8^{+1.1}_{-1.0}$ ps.

5.2.3 Summary of the measured lifetimes

In this work, the half-lives of several excited states of $^{114,116,118,120}\text{Sn}$ were measured using the DSA and RDDS methods. In addition, the bootstrap resampling technique was applied to improve the error estimates for the RDDS analysis. Table 5.4 summarizes the obtained half-life results, comparing them with the values available from the literature.

Nucleus	E_i [MeV]	I_i^π	$T_{1/2}^{lit}$ [ps]	$T_{1/2}^{exp}$ [ps]
^{114}Sn	1.953	0_2^+	$6.5(23)^a$	$11.0^{+9.3}_{-4.0}$
	2.156	0_3^+	$> 7.6^a$	$13.0^{+9.4}_{-4.2}$
^{116}Sn	1.757	0_2^+	$44(7)^a, 45(10)^b, 51(12)^c$	$51.1^{+18.8}_{-11.7}$
	2.027	0_3^+	$160(20)^b, 164(10)^c$	184^{+66}_{-40}
	2.546	0_4^+		$1.17^{+0.10}_{-0.09}$
^{118}Sn	1.758	0_2^+	$21(3)^a$	$18.2^{+2.4}_{-2.0}$
	2.057	0_3^+	$> 0.7^d, < 200^a$	$92.2^{+19.9}_{-14.7}$
	2.497	0_4^+		$14.2^{+1.5}_{-1.3}$
^{120}Sn	1.875	0_2^+	$7.4(10)^a$	$8.2^{+2.6}_{-1.7}$
	2.160	0_3^+	$34.7(49)^d$	$40.4^{+7.9}_{-6.2}$
	2.587	0_4^+	$> 0.34^f$	$14.8^{+1.1}_{-1.0}$

^a From Coulomb excitation [87]. ^d From DSAM [146].
^b From fast-timing [145]. ^e From fast-timing [81].
^c From fast-timing [78]. ^f From DSAM [147].

Table 5.4: Half-lives of excited states of $^{114-120}\text{Sn}$ measured in the present work with the DSA and RDDS methods, in comparison with adopted values from the NNDC dataset [68, 87, 145–147] and from recent unevaluated datasets [78, 81].

For ^{114}Sn , only the half-life of the 0_2^+ state and a lower limit for the 0_3^+ state were previously measured in Coulomb excitation by Bäcklin et al. [87]. These values agree with those obtained in the present direct measurement. We note that the $T_{1/2}(0_2^+) = 6.5 \pm 2.3$ ps of Ref. [87] is about 40% smaller than our measurement, but still compatible within uncertainties.

In the case of ^{116}Sn , Coulomb excitation [87] and fast-timing measurements [78, 145] determined the half-life of the 0_2^+ state, with the latter two also determining that of the 0_3^+ state. Our results for the 0_2^+ and 0_3^+ states are fully consistent with previous measurements. In addition, the half-life of the 0_4^+ state is measured here for the first time with the DSA method.

While the half-life of the 0_3^+ state of ^{118}Sn was previously deduced via Coulomb excitation [87], only a lower limit was determined using DSAM in a $(n, n'\gamma)$ experiment [146], and an upper limit was obtained using fast-timing techniques in a (d, p) experiment [87]. Good agreement is found between the half-life of the 0_2^+ state measured in the present work and the adopted value, and the half-life of the 0_3^+ state, previously unknown, falls between the adopted lower and upper limits. No information was available in the literature for the half-life of the 0_4^+ state, for which we provide a first estimate.

As for $^{114,116,118}\text{Sn}$, the half-life of the 0_2^+ state of ^{120}Sn was previously deduced from Coulomb excitation [87], while that of the 0_3^+ state was recently obtained in a (n, γ) experiment using fast-timing techniques. Our half-life results for these two states are in good agreement with the previous ones. Additionally, the half-life of the 0_4^+ state, determined here for the first time, is fully compatible with the lower limit estimated by Demidov et al. [147] using the DSA method in an (n, γ) experiment.

Part III

The Experiments at LNL

The experiment at LNL

The experiment discussed in this Chapter was at Legnaro National Laboratories (LNL) in July 2022. The experimental setup involved the large acceptance magnetic spectrometer PRISMA, coupled to the Advanced GAMMA Tracking Array (AGATA). Both thick-target and plunger-target configurations were employed to study Sn and Cd isotopes of interest, which were populated via a multi-nucleon transfer reaction.

Section 6.1 describes in detail the experimental setup, focusing on the Tandem accelerator, the AGATA array, the PRISMA spectrometer and the plunger device, while Section 6.2 illustrates the features of the $^{32}\text{S}+^{110}\text{Cd}$ multi-nucleon transfer reaction and the different setup configurations used during the data acquisition.

6.1 The Experimental setup

6.1.1 The Tandem facility

The Tandem-XTU is the largest electrostatic accelerator installed at Legnaro National Laboratories. Thanks to its maximum operating voltage of 15 MV, this accelerator is able to provide stable beams from ^1H to ^{238}U with energies up to 5-10 A-MeV. The basic functioning principles of this tandem accelerator are analogous to those discussed in Section 3.1.1 for the case of the IFIN-HH 9 MV Tandem, and thus a detailed description will be omitted here. A schematic view of the Tandem facility is shown in Figure 6.1,

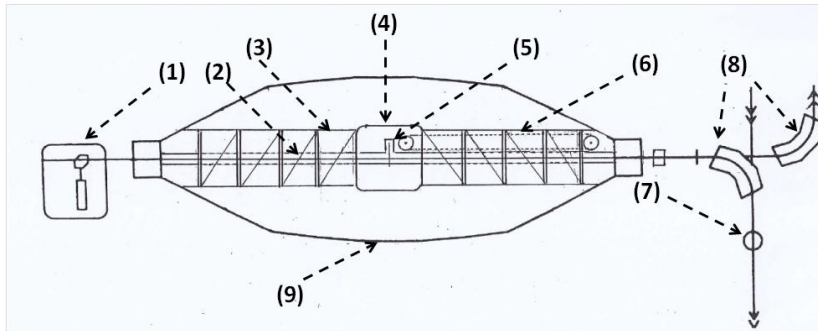


Figure 6.1: Schematic drawing of the Tandem accelerator installed at LNL: (1) the ion source platform, (2) the accelerating pipe, (3) the column, (4) the high voltage terminal, (5) the ion "stripping" station, (6) the laddertron charging belt, (7) the diagnostic station, (8) the bending magnets, (9) the tank filled in SF₆ at 7 atm pressure. Reprinted from Ref. [112].

where its main elements are highlighted. In particular, a ^{137}Cs ion source (1) is used

as an electron gun to negatively charge the ions, which are subsequently attracted and accelerated through the pipe (2) by the high-voltage terminal (4). The latter, which is charged with a laddertron belt (6), is encapsulated in a safety tank (9) filled with SF₆ to prevent electrical discharges. A ¹²C stripping foil (5) is used to remove electrons and positively ionize the species, which are then exposed to the second accelerating stage and finally delivered to the experimental area.

In the present experiment, the ³²S beam was accelerated at energies between 133 and 170 MeV, depending on the configuration.

6.1.2 The PRISMA spectrometer

PRISMA is a large acceptance magnetic spectrometer designed to identify reaction products of heavy-ion reactions at energies between 5 and 20 MeV/A [97, 148, 149]. The primary features of this spectrometer are a very large solid angle (80 msr) and momentum acceptance ($\pm 10\%$), good mass resolution (1/300) via Time-of-Flight (ToF) measurement, energy resolution up to 1/1000, and capability of rotation around the target in a wide angular range from -30° to 130° . A summary of these characteristics is reported in Table 6.1, while a schematic representation of PRISMA is shown in Figure 6.2. The main components of the spectrometer are:

- a Micro-Channel Plate entrance detector;
- a Multi-Wire Parallel Plate Avalanche Counter (MWPPAC) focal plane detector;
- a Ionization Chamber (IC);
- a dipole magnet;
- a quadrupole magnet;

The combined information of entrance and focal plane positions, energy and ToF provided by these detector leads to the reconstruction of the trajectories, the determination of the velocity and the identification of the atomic number Z and mass number A of the incoming ions. Such procedures will be discussed in Chapter 7.

Characteristics of PRISMA			
Solid angle	$\Delta\Omega$	80	msr
Azimuthal acceptance	$\Delta\theta_{lab}$	± 6	deg
Polar acceptance	$\Delta\varphi_{lab}$	± 11	deg
Energy acceptance	ΔK	± 20	%
Momentum acceptance	Δp	± 10	%
E resolution	$\Delta E/E$	1/1000	
Z resolution	$\Delta Z/Z$	1/60	
A resolution	$\Delta A/A$	1/300	

Table 6.1: Main characteristics of the PRISMA spectrometer.

In the following, the different elements of the PRISMA spectrometer will be described in detail.

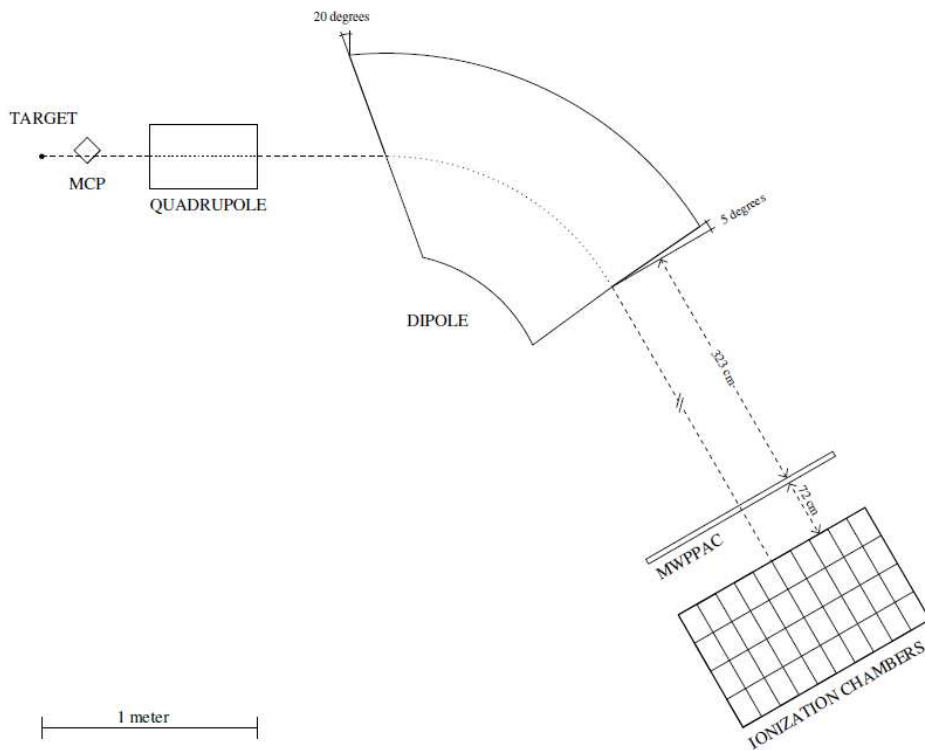


Figure 6.2: Schematic view of the PRISMA magnetic spectrometer. Reprinted from Ref. [150].

The MCP detector

The entrance detector of PRISMA is a Micro-Channel Plate detector [149], and it is placed at 25 cm from the target. The active area of detection is $8 \times 10 \text{ cm}^2$, covering the whole solid angle of PRISMA. The detector, which is rotated by an angle of 135° with respect to the optical axis of the spectrometer, provides a time signal and two position measurements along the x and y directions, namely along and transversal to the dispersion plane, respectively. When passing through the $20 \mu\text{g}/\text{cm}^2$ carbon foil, the incoming ions generate secondary electrons, which are accelerated towards the anode by an electric field. During the migration towards the anode, a magnetic field is employed to drive the electrons into helicoidal trajectories, limiting their transverse motion. In addition to the signal amplification, the position-sensitive anode provides the position information based on the difference in the arrival time of the signal of its two orthogonal delay lines, one for each axis. For position calibration purposes, a cross-shaped mask with reference points is placed before the carbon foil.

The intrinsic time resolution of this device is approximately 350 ps, while the position resolution is $\sim 1 \text{ mm}$ in each direction. Furthermore, the detection efficiency for heavy ions at energies of a few MeV/A is close to 100%. A scheme and a photograph of this device are shown in Figure 6.3.

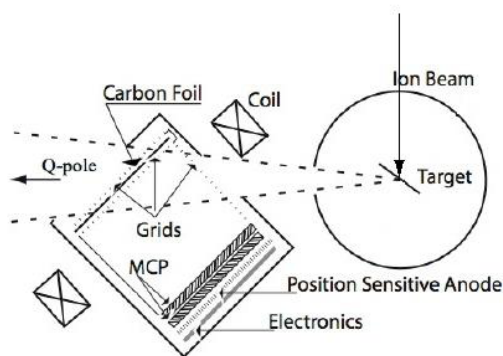


Figure 6.3: Schematic view (left) and picture (right) of the MCP detector of PRISMA. Reprinted from Ref. [149].

The magnetic elements

Two optical elements — a quadrupole magnet and a dipole magnet — are placed after the MCP entrance detector. Their function is to firstly focus the reaction products onto the dispersion plane and secondly bend their trajectories towards the focal plane detectors.

The quadrupole singlet is mounted at a distance of 50 cm from the target. Its entrance window has a diameter of about 320 cm and a width of about 42 cm. Due to the presence of fringing fields which extend over the geometrical dimension of the magnet itself, effective geometrical parameters are utilized in the analysis to optimize the trajectory reconstruction. Its magnetic field is set to focus the ions along the y vertical axis while dispersing them along the x horizontal axis.

The dipole magnet is placed 60 cm downstream of the quadrupole exit window, bending the trajectories of the incoming particles according to their magnetic rigidity. The bending angle of the magnet is fixed at 60° , with a curvature radius of 1.2 m. Since the maximum value of the magnetic field is 1 T, the maximum magnetic rigidity $B\rho$ is 1.2 Tm.

The MWPPAC detector

The focal plane detector, encountered by the incoming particles after about 3.2 m from the exit window of the dipole, is an array of Multi-Wire Parallel Plate Avalanche Counters [148]. It measures 100 cm along the x horizontal axis and 13 cm along the y vertical axis. Furthermore, the detector is segmented in 10 independent sections, or pads, each with an active area of $10 \times 13 \text{ cm}^2$. Each section has a three-electrode structure: a central cathode and two orthogonal wire planes. The cathode is polarized at high voltage (500–600 V) and provide the timing information, while the two anodes are at ground potential and provide the x and y position information. By combining the timing information of the MCP and the MWPPAC, it is possible to reconstruct the Time-of-Flight (TOF), hence the velocity, of the incoming ions. The whole detector has a total of 1000 wires on horizontal plane, 100 for each section with a step of 1 mm. Therefore, the position resolution in the x direction is 1 mm. Conversely, the wires on the vertical plane are not segmented and cover the entire width of the detector. Since these wires are connected in groups of two, the position resolution is 2 mm. A photograph and a three dimensional schematic

view of the MWPPAC detector of PRISMA is depicted in Fig. 6.4.

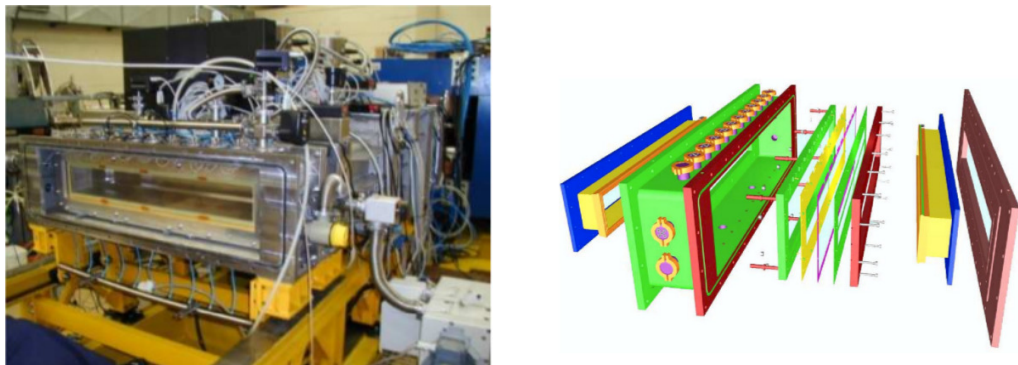


Figure 6.4: Picture (left) and schematic drawing (right) of the MWPPAC detector of PRISMA. Reprinted from Ref. [148].

The IC detector

At the end of the spectrometer, 60 cm downstream of the MWPPAC detector, a segmented Ionization Chamber (IC) is mounted [97]. The purpose of this array is to measure the total and the partial energy loss of the ions in the gas, which allows for the determination of the atomic number Z of the ions via the $\Delta E - E$ method (see discussion in Chapter 7). Its total sensitive volume of $110 \times 20 \times 120 \text{ cm}^3$ is separated from the vacuum by a $1.5\text{-}\mu\text{m}$ -thick mylar window, and can be filled with either high-purity methane (CH_4) or carbon tetrafluoride (CF_4), kept at pressures between 15 and 100 mbar. These gases provide a good balance between stopping power and counting rate capabilities thanks to their high electron and ion drift velocities. As shown in the two panels of Figure 6.5, the cathode and anode electrodes of the IC are composed of 40 pads (4 in the longitudinal z direction and 10 in the transversal x direction), operating as independent sections. The first (or the first two sections) can be used to determine the energy loss ΔE , whereas the sum of the energy losses in all the four sections is taken as the total energy E . With this method, the Z identification is typically obtained with a resolution of $\Delta Z/Z \approx 1/60$.

In order to eliminate events from ions with trajectories exiting the chamber (shown in the right panel of Figure 6.5), additional sections are placed on both sides of the IC, acting as veto.

6.1.3 The AGATA array

The Advanced GAMMA-ray Tracking Array (AGATA) [151, 152] is a European project devoted to the design and development of a state-of-the-art 4π γ -ray spectrometer. Similarly to its counterpart GRETINA [153] in the United States, AGATA is based on the innovative concept of γ -ray tracking and pulse-shape analysis (PSA) in electrically segmented, high-purity germanium crystals. The tracking technique consists in the reconstruction of the 3D track of each γ -ray detected by the spectrometer, according to the energy deposition and the location of the interaction points inside the array. The position resolution of a few mm required by the tracking algorithm can only be achieved with the electrical segmentation of the HPGe crystals, digital electronics and pulse-shape analysis

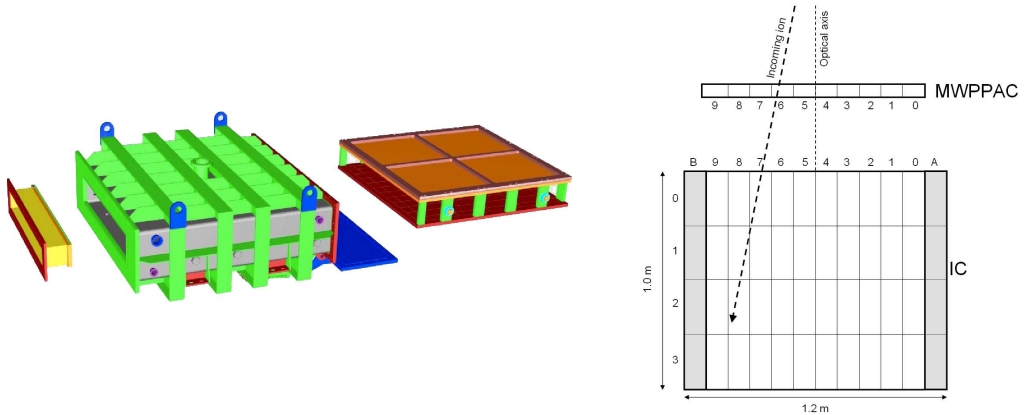


Figure 6.5: (Left) Three-dimensional model of the PRISMA Ionization Chamber. (Right) Schematic top view of the Ionization Chamber. The side sections used to reject bent trajectories are indicated with a gray shade. Reprinted from Ref. [150].

methods [151, 154]. With this technological advancement, new-generation γ -ray spectrometers grant unprecedented performances in terms of efficiency, position resolution and Doppler-correction capabilities compared to conventional, Compton-suppressed γ -ray spectrometers that have been commonly used since the 1980s (e.g., GASP [155], EUROBALL [156], GAMMASPHERE [157], GALILEO [158], ROSPHERE [113]).

The AGATA array is conceived as an itinerant spectrometer to exploit both intense stable beams and radioactive ion beams provided by the different European laboratories involved in the collaboration. After a first campaign at LNL (Legnaro, Italy) started in 2009 [159], AGATA was moved to GSI (Darmstadt, Germany) [160] and afterwards to GANIL (Caen, France) [161]. Since March 2022, AGATA has been operational at LNL, where it is expected to remain for the next few years.

In the current experimental campaign at LNL, AGATA can be coupled to either charged particle detectors, such as SPIDER [162] and TRACE [163], or to the magnetic spectrometer PRISMA [97]. A photograph of the AGATA-PRISMA configuration [152], which was the one used in the present experiment, is shown in Figure 6.6. Specific details about this configuration will be given in Section 6.1.4.

The Crystals

The AGATA detectors are segmented, encapsulated coaxial n-type high-purity germanium crystals. As shown in Figure 6.7, the crystals have a rear diameter of 8 cm and a length of 9 cm. To optimize the solid angle coverage, the front face of the crystals is tapered into three different hexagonal geometries, identified with letters A (red), B (green) and C (blue). These crystals are arranged into triplets, called AGATA Triple Clusters (ATCs), contained in a common cryostat. In the foreseen final configuration, a total of 180 detectors, arranged in 60 ATCs, will be mounted on the support structure (honey-comb), covering more than 80% of the solid angle.

Each crystal is encapsulated into a hermetically sealed aluminum canister with a 0.8-mm-thick wall. The external contact of the crystals is electrically divided into 36 segments, each providing an independent output (see left panel of Figure 6.7). Additionally, a core contact is used to apply the positive bias voltage and collect two full-energy

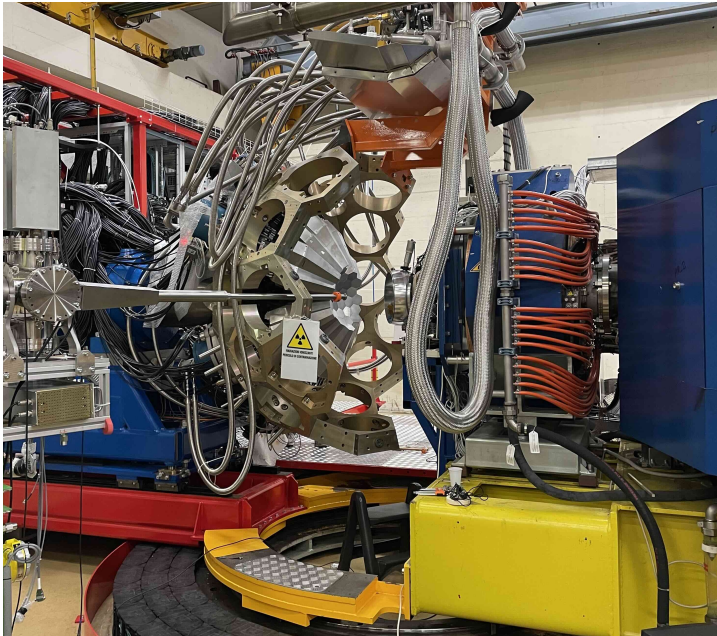


Figure 6.6: Photo of the AGATA-PRISMA configuration taken in July 2022 before the experiment discussed in this thesis.

signals, which correspond to a 5 MeV (high gain) and a 30 MeV (low gain) range.

The Front-End Electronics

In order to ensure the maximum performances and minimal noise, custom-made preamplifiers have been developed by the AGATA collaboration. Each channel from the segment and core contacts, which collect holes and electrons, respectively, is equipped with charge-sensitive resistive feed-back preamplifiers [151]. These are divided in a cold part and a warm part. The cold part consists of a low-noise Field Effect Transistor (FET) placed close to the electrodes of the detector, operating at cryogenic temperatures. Conversely, the warm part operates at room temperatures at about 15 cm from the first part, and comprises a low-noise amplifier, a pole-zero stage, a differential output buffer and fast-reset circuitry. An additional de-saturation circuit is employed to detect saturated signals caused by high-energy or pile-up events. This circuit preserves good energy resolutions (~ 2 keV at 1.33 MeV) up to rates of about 50 kHz.

The 38 output signals from the preamplifiers of each crystals are then digitized at 100 MHz by a 14-bit Fast Analog-to-Digital Converters (FADC). The sampled signals are subsequently transmitted to the pre-processing electronic boards. The energy of the signal is determined using either a digital trapezoidal filter or the Moving Window Deconvolution (MWD) method [151], both of which account for the response function of the preamplifier. The timing information, primarily used for triggering, is derived from the high-gain core signal via a leading-edge algorithm.

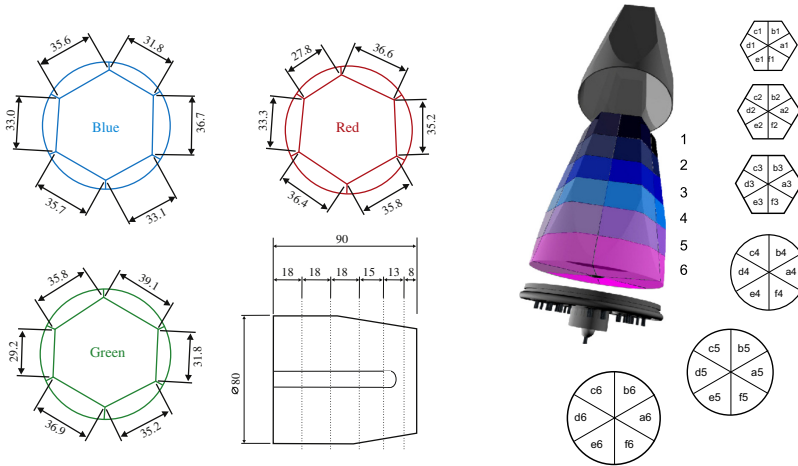


Figure 6.7: Schematic representation of the AGATA crystals, indicating the dimensions in mm of the three crystal types (A - red, B - green, C - blue) and the 36-fold segmentation. Reprinted from Ref. [151].

6.1.4 The AGATA-PRISMA configuration

During the first part of the experimental campaign at LNL, the AGATA-PRISMA configuration was used. This setup is particularly suited to perform γ -ray spectroscopy studies with multi-nucleon transfer reactions, combining the γ -ray detection capabilities of AGATA with the ion identification provided by PRISMA. Here, AGATA and PRISMA are mounted facing each other on a common platform that can rotate from 20° to 110° with respect to the beam direction (see Figure 6.8 and Ref. [152]). In this way, the angle of PRISMA can be adjusted to match the grazing angle of the reaction, maximizing the detection efficiency. The honeycomb structure of AGATA has two additional degrees of freedom. Thanks to its central shaft, it is able to rotate and translate around its symmetry axis. The distance between the target and AGATA detectors can thus be adjusted from the nominal position of 23.5 cm, to the close-up position of 18 cm, depending on the requirements of each experiment.

The scattering chamber was designed to meet the mechanical requirements of this configuration, such as the possibility of a common rotation of the two spectrometers. It comprises a set of movable shells and a fixed shell. The movable shells are aluminum hemispheres with a radius of 17 cm and a thickness of 2 mm, which cover angles between 7° and 88° . The fixed body of the chamber consists of the scattering chamber and the expansion chamber, which houses the beam dump and the electronics for ancillary devices.

During the present experiment, 33 detectors (11 ATCs) were installed in the AGATA honeycomb, of which 30 were operational. This configuration is shown in Figure 6.9, where the detectors are labeled according to the conventional numbering scheme [151].

Data Acquisition and Trigger

The Data Acquisition system (DAQ) processes the signals of the segments from the read-out electronics and stores them on disk. In this configuration, the data flows of AGATA and PRISMA follow two parallel routes, which are later merged based on the timestamp

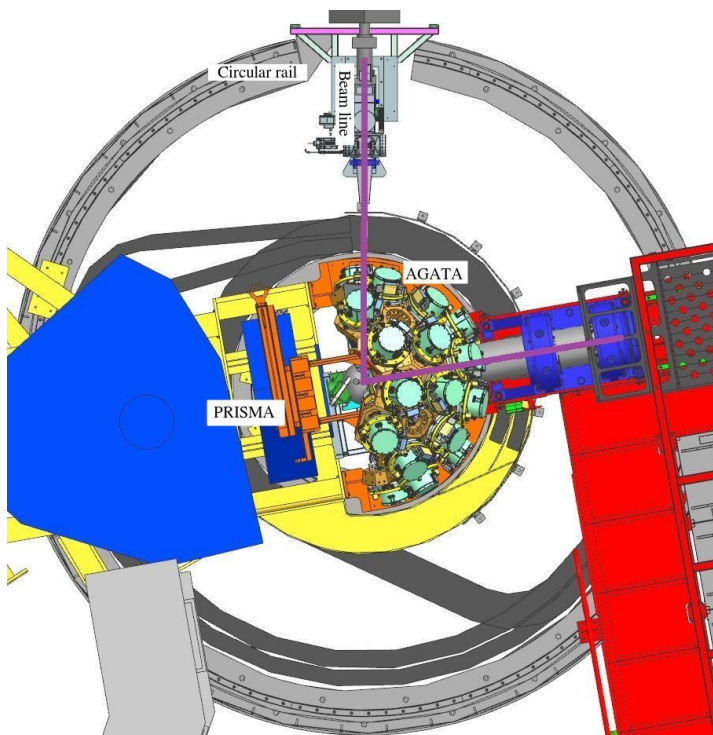


Figure 6.8: CAD drawing of AGATA coupled to PRISMA. Reprinted from Ref. [152].

of the events.

The synchronization of the AGATA signals is managed by the Global Trigger and Synchronization system (GTS) [164]. When the core signal of an AGATA crystal is above a given threshold, it generates a local trigger, enabling the acquisition of the 36 segments. For each local trigger, a request is sent to the GTS, which synchronizes and validates the events. On the other hand, the readout signals of PRISMA are collected by VME modules. The merging is performed via dedicated modules, the AGATA VME adapters (AGAVA), which allow complementary detectors to participate in the trigger decisions and obtain correlated AGATA-PRISMA events.

In the present experiment, the global trigger of the DAQ was provided by the OR of all cathode signals of the MWPPAC of PRISMA. Thus, AGATA events were stored only if a reaction product reached the focal plane of PRISMA. An additional trigger condition was applied during the final part of the experiment. The trigger was set as an OR between PRISMA MWPPAC events and AGATA events with multiplicity of at least 3.

6.1.5 Plunger

During the second part of the experiment, the GALILEO plunger device was mounted in the chamber [158]. This device was developed in collaboration with the Cologne group (IKP), sharing similar features with the ROSPHERE plunger described in Section 3.1.4.

As shown in Figure 6.10, a compact design was chosen, with the motor lying below the target and the stopper foils. In this configuration, a good γ -ray transparency for the AGATA detectors, which are placed at backward angles, can be achieved. Conical sup-

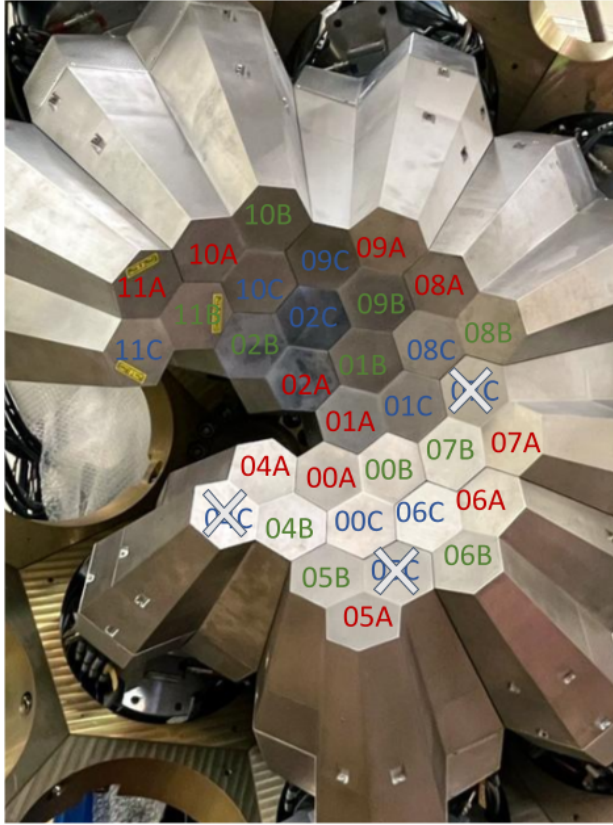


Figure 6.9: Photo of the AGATA array in the configuration of the present experiment. The array was composed of 33 detectors, each marked with its ID. The crystals excluded from the analysis are marked with white crosses.

ports hold and stretch the target and stopper foils, with their tension being manually adjusted using three screws. Maintaining the parallelism between the foils is particularly important, and it is achieved through an iterative procedure by bringing the foil to electrical contact and correcting it using three screws. While the stopper foil is placed at a fixed position, the target support is placed on the sliding rails of the piezo motor, and can be moved with sub-micrometric precision over a range of distances from a few μm up to 15 mm. During the experiments, the distance fluctuations caused by the beam irradiating the target and stopper foils are corrected using a capacitance motor.

The right panel of Figure 6.10 shows a photo of the plunger with the cadmium target used in the experiment of this thesis.

6.2 The multi-nucleon transfer experiment

This experiment was performed in July 2022, with the aim of studying excited states of Sn and Cd isotopes around mass $A=110$. These nuclei were populated with a multi-nucleon transfer reaction, induced by ^{32}S beam impinging on ^{110}Cd targets. Beam-like reaction products were detected by PRISMA, which was placed at 67° with respect to the beam direction, whereas the emitted γ rays were measured by the AGATA spectrometer,

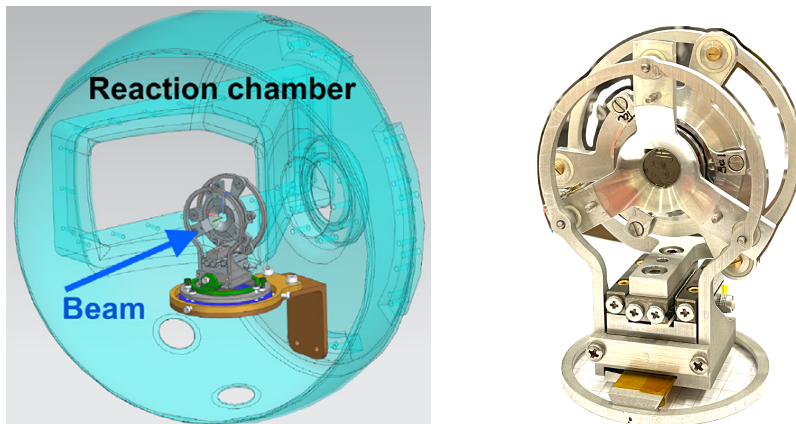


Figure 6.10: (Left) CAD drawing of the plunger device installed in the AGATA scattering chamber [152]. (Right) Picture of the plunger device taken before the installation in the scattering chamber. The cadmium target used in the present experiment was mounted.

centered at 123° . Two main target configurations were employed during the data acquisition, here denoted as the thick-target and the plunger configuration. In the first case, a 1.05-mg/cm^2 -thick Cd target, deposited on a Ta layer of 2.0 mg/cm^2 , was mounted in the scattering chamber at an angle of -43° with respect to the beam direction (see left panel of Figure 6.11). The TANDEM facility accelerated a ^{32}S beam to 164 MeV, which, after accounting energy loss in the Ta layer, resulted in a mid-target reaction energy of 144 MeV. According to Eq. (3.1), this energy is approximately 15% above the Coulomb Barrier. In this configuration, the target-like recoils are scattered nearly parallel to the rotated target, losing energy until they stop within it, effectively creating a “thick” target. In the plunger case, the 1.1-mg/cm^2 -thick ^{110}Cd target, glued onto a 2.1-mg/cm^2 -thick Ta backing, was tilted at $+52^\circ$ with respect to the beam line. Furthermore, a thick Ta foil was used as stopper material. In order to have a similar reaction energy, the beam energy was increased to 170 MeV, accounting for the slightly thicker target materials and the larger target angle. A schematic view of the two configurations of the experimental setup are shown in Figure 6.11.

According to GRAZING calculations [94], the $^{32}\text{S} + ^{108}\text{Cd}$ multi-nucleon transfer reaction is expected to primarily populate $^{106-110}\text{Cd}$, $^{109-111}\text{In}$ and $^{110-112}\text{Sn}$ isotopes, with cross sections from tenths to hundreds of mb. Relevant details about the reaction channels of main interest are summarized in Table 6.2, such as the calculated cross sections and the reaction Q -values. According to the optimal excitation energy window estimated with LISE++ [118], low-lying levels up to a few MeV are expected to be populated. In the past, similar $2p$ and $2n$ transfer reactions were successfully employed to populate excited states in Sn and Cd isotopes. In particular, $^{110-116}\text{Cd}(^3\text{He},n)$ reaction studies [69] observed a large enhancement in the cross section of excited 0^+ states in Sn isotopes. These states, built on intruder configurations, are of particular interest for the present work.

As shown in the top panel of Figure 6.12, differential cross sections calculated with the GRAZING code for the ^{30}Si beam-like (PL) and the ^{112}Sn target-like (TL) products are peaked at approximately 67° and 47° , respectively. Similar maximum angles are found for the angular distributions of the $(-2p)$ and $(-2n)$ channels (see Table 6.2). This motivated the placement of PRISMA at 67° , which maximizes the detection efficiency

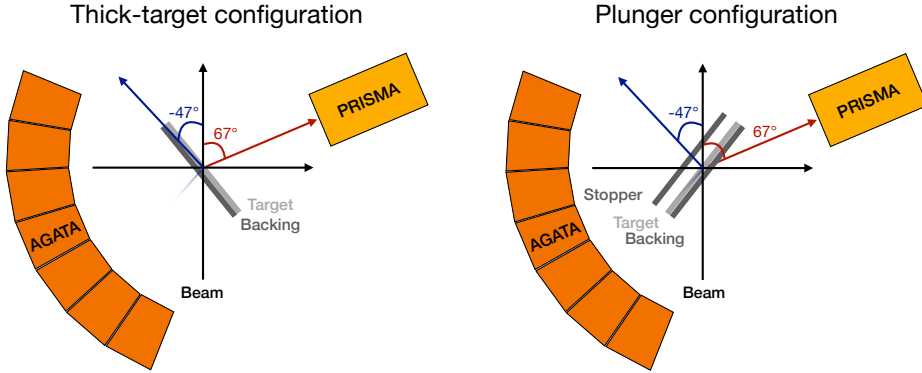


Figure 6.11: Thick- (left) and plunger-target (right) configurations used in the experiment, with the target rotated at -43° and 52° , respectively. PRISMA was placed at the grazing angle of 67° , with AGATA facing it. Angles and directions corresponding to the maximum of the differential cross section of the beam-like and target-like products are indicated in red and blue, respectively (see text for details).

Nucleus	σ_{tot} [mb]	θ (PL)	θ (TL)	Q [MeV]	E_X [MeV]
$^{112}\text{Sn}(-2p)$	1.5	67°	47°	-3.3	6.6
$^{110}\text{Sn}(-2p, 2n)$	0.3	67°	47°	-6.4	3.4
$^{108}\text{Cd}(+2n)$	10	67°	50°	+2.8	2.8

Table 6.2: Summary of the main properties of the $^{32}\text{S}+^{110}\text{Cd}$ reaction at 144 MeV for the transfer channels of interest. The channels are indicated with respect to the beam-like particles detected by PRISMA. σ_{tot} is the inclusive cross section calculated with the GRAZING code, whereas θ (PL) and θ (TL) are the angles at the peak of the differential cross sections of the beam-like (PL) and target-like (TL) products, respectively. Q indicates the Q -value of the reaction, while E_X indicates the excitation energy available for the total system.

of beam-like reaction products. Moreover, the chosen reaction energy of 144 MeV ensures that the kinetic energies of the beam-like target products match the acceptance of PRISMA, which requires a few MeV/A. The kinematic lines calculated with the LISE++ suite [118] for the $^{110}\text{Cd}(^{32}\text{S}, ^{30}\text{Si})^{112}\text{Sn}$ reaction are shown in the bottom panel of Figure 6.12.

The experiment lasted about 200 hours of effective beam on target, of which 120 hours were analyzed in the present thesis. About 30 hours were dedicated to the thick-target configuration, with an average beam current of 9 enA. A total of 93 hours of data corresponded to the plunger configuration, however, with only 3 enA of beam intensity, resulting in comparable statistics for the two configurations.

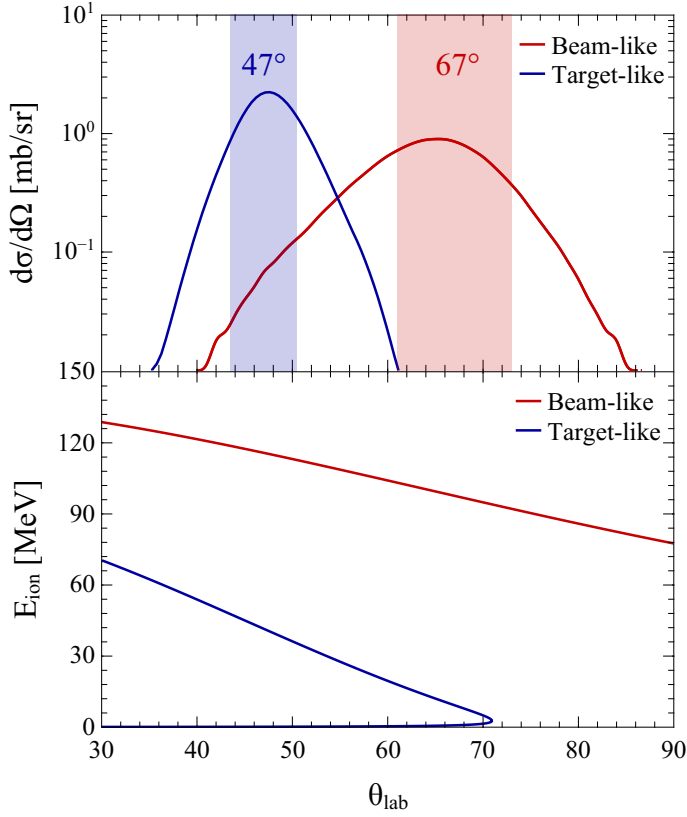


Figure 6.12: (Top) Differential cross sections for the $^{110}\text{Cd}(^{32}\text{S}, ^{30}\text{Si})^{112}\text{Sn}$ reaction (-2p channel), calculated with the GRAZING code [94]. The distribution of the ^{30}Si beam-like product is in red, while that of the ^{112}Sn target-like product is in blue. The angles corresponding to the maximum of the distributions are labeled accordingly. (Bottom) Kinematic lines of ^{30}Si (red) and ^{112}Sn for reactions at mid-target, calculated with LISE++ [118].

Preprocessing of LNL data

This chapter describes the preliminary analysis procedures of the AGATA-PRISMA data. In Section 7.1, the data sorting procedure and the analysis software are briefly illustrated. In Section 7.2 the processing of PRISMA data is discussed, focusing on the calibrations, the reconstruction of the trajectories of ions and the ion identification. Finally, the AGATA data sorting procedures and the performance of the array are addressed in Section 7.3.

7.1 Data sorting

The sorting of AGATA-PRISMA data involves several stages, which are handled by a series of software codes. As discussed in Section 6.1.4, PRISMA and AGATA data are treated in a separate way and later merged based on the timestamps of the events.

The PRISMA data analysis is performed via the `prismafilters` executable, which exploits the tools of the `prisma_library` codebase [165]. Calibrations, trajectory reconstruction and ion identification are performed at this level, and the analyzed data are then stored in ROOT TTree files [124].

In parallel, the `femul` code of the `agapro` package [166, 167] is used to process the AGATA data. This operation is often called *data replay*, and consists of a series of processes managed by dedicated modules (called *actors*). These will be described in detail in Section 7.3. At the end of the replay process, ROOT files with sorted events are produced.

Depending on the analysis workflow, the merging can be done by the `femul` event builder or at a later stage by matching the synchronized timestamps. Final histograms and matrices are then produced by the `agataselector` code [168], which reconstructs the reaction kinematics, evaluates the energy losses and applies the Doppler correction to the γ -ray spectra.

7.2 Preprocessing of PRISMA data

As introduced in Section 6.1.2, the PRISMA spectrometer is composed of several detectors. Alignments, calibrations and detailed optimizations of the detectors are required to accurately identify the reaction products in terms of velocity β , atomic number Z and mass A . This involves the following steps:

- Gates and thresholds are set for the different detectors of PRISMA. The MCP and the MWPPAC are calibrated using reference values, while the sections of the IC are aligned using a charge injector.

- The event-by-event identification of the atomic number Z of the ions is achieved with the $E - \Delta E$ method, using a combination of signals from the IC.
- The time of flight (ToF) of the ions between the MCP and the MWPPAC is aligned among the different sections of the focal plane detector.
- The trajectory of the ions is reconstructed based on the interaction position in the MCP and the MWPPAC, the magnetic fields, and the ToF. From the reconstructed length of the trajectories, the velocity β and the A/q ratio are obtained.
- The charge state q of the ions is determined based on the magnetic rigidity and the total energy loss in the IC. The A/q is then calibrated and the mass spectrum is finally obtained.

In the following sections, each step will be described in detail. Since data were acquired with two distinct configurations of the experimental setup, the PRISMA data processing had to be performed twice. The thick-target configuration is here taken as a representative example and any deviations in the analysis procedures will be highlighted.

7.2.1 MCP Calibration

The MCP entrance detector of PRISMA, previously described in Section 6.1.2, measures the (x,y) interaction point and provides the stop time signal for the ToF measurement. In order to accurately reconstruct the trajectories of the incoming ions, the interaction position has to be converted from ADC channels to millimeters. The left panel of Figure 7.1 shows the two-dimensional spectrum of the x and y raw positions as recorded by the ADC. A polygonal gate was applied to the matrix to reject spurious signals. The reference points of the cross-shaped metallic mask, marked with green crosses, are clearly visible.

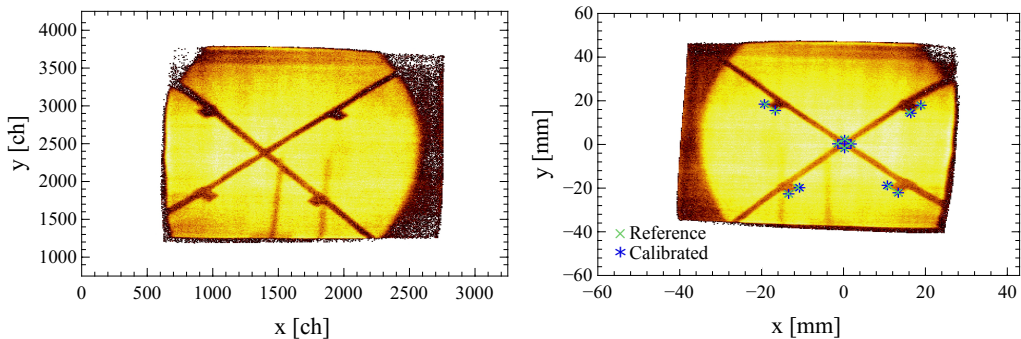


Figure 7.1: Raw (left) and calibrated (right) x and y positions of the MCP detector of PRISMA. Reference points from the shadow of the metallic mask and calibrated points are indicated with green crosses and blue asterisks, respectively.

The calibration process consists of in three steps, involving multi-parametric transformations to take into account possible deformations of the MCP surface. Firstly, a

linear transformation is applied to stretch the raw coordinates:

$$\begin{pmatrix} x' \\ y' \end{pmatrix} = \begin{pmatrix} A_0 & A_1 \\ B_0 & B_1 \end{pmatrix} \begin{pmatrix} x_{raw} \\ y_{raw} \end{pmatrix}, \quad (7.1)$$

where the diagonal elements $A_0 = B_1 = 1$. Then, the transformed coordinates are calibrated using quadratic and linear polynomials for the x' and y' coordinates, respectively.

$$\begin{aligned} x'' &= a_2 x'^2 + a_1 x' + a_0, \\ y'' &= b_1 y' + b_0. \end{aligned} \quad (7.2)$$

Finally, a rotation of an angle θ is implemented to account for the orientation of the MCP with respect to the optical axis:

$$\begin{pmatrix} x_f \\ y_f \end{pmatrix} = \begin{pmatrix} \cos \theta & \sin \theta \\ -\cos \theta & \cos \theta \end{pmatrix} \begin{pmatrix} x'' \\ y'' \end{pmatrix}, \quad (7.3)$$

Despite the seven independent parameters involved in the procedure, the problem is well-constrained as a total of eleven reference points were used. The values of the parameters are obtained from a least-squared fit of the distance between the reference and the calibrated points. The calibrated MCP is shown in the right panel of Figure 7.1, demonstrating the good outcome of the procedure.

7.2.2 MWPPAC calibration

As introduced in Section 6.1.2, the MWPPAC detector measures the time and position of the reaching the focal plane of PRISMA. For a single event, each of the 10 sections of the MWPPAC provides three signals, namely the right x_{right} , left x_{left} , and cathode $x_{cathode}$ positions on the x axis. In addition, two common signals for the up y_{up} and down y_{down} positions are recorded. Since the y information has only a second order effect on the trajectory reconstruction, it is used to check the beam centering and is therefore excluded from the analysis. The cathode signal is also sent to a TAC module, acting as the start signal for the ToF of the ions from the MCP to the MWPPAC.

Before the calibration of the x coordinate, a threshold is set on the signals from each section, and a polygonal gate is applied on the matrix correlating $x_{right} + x_{left}$ and $x_{cathode}$ to suppress background contributions. The x_{FP} position on the horizontal plane is obtained from the difference between the right and left signal

$$x_{FP} = x_{right} - x_{left}, \quad (7.4)$$

As an example, the left panel of Figure 7.2 shows a spectrum resulting from the difference between the right and the left signals of a single section of the PPAC. Red dashed lines indicate reference values used to perform the linear calibration of the focal plane position, corresponding to the initial, central and final points of the section. As the two wires in the middle of each section are short-circuited, a spike and a subsequent drop in counts are produced in the central channels, providing the central reference point for the calibration.

If one of the right and left signals is missing, the cathode information can be used to reconstruct the position from the relation $2x_{cathode} = x_{right} + x_{left}$. The focal plane

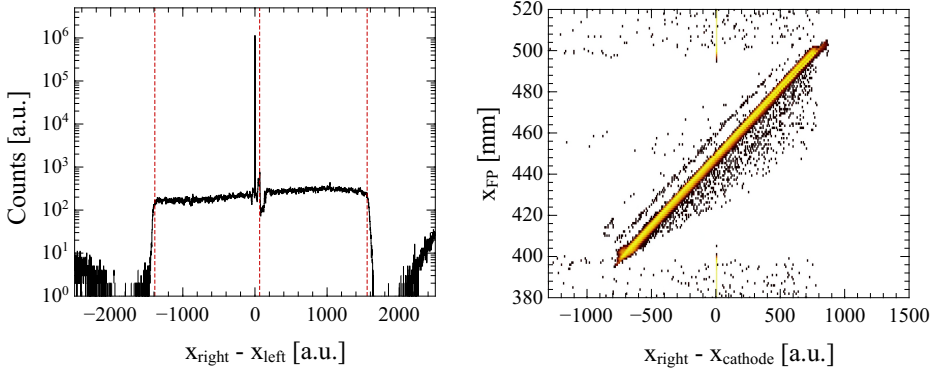


Figure 7.2: (Left) Difference between right and left signals for section 4 of the MWPPAC detector. Red dashed lines indicate the reference values used for the calibration procedure. (Right) Matrix correlating the difference between the right and cathode signal with the calibrated x focal plane position of section 4, which is used to recover the position information when the left signal is missing (see text for details).

position is therefore given by

$$\begin{aligned} x_{FP} &= x_{right} - x_{cathode}, \\ x_{FP} &= x_{cathode} - x_{left}. \end{aligned} \quad (7.5)$$

A linear calibration is then required for each section to convert these differences into position in mm. An example of the correlation between $x_{right} - x_{cathode}$ and x_{FP} is shown in the right panel of Figure 7.2.

The focal plane position of the 10 sections of the MWPPAC, after the linear calibration in mm, is shown in Figure 7.3. The peak structures correspond to the two main charge state of ^{32}S ions (quasi-elastic channel), which dominate over the contributions from the transfer channels.

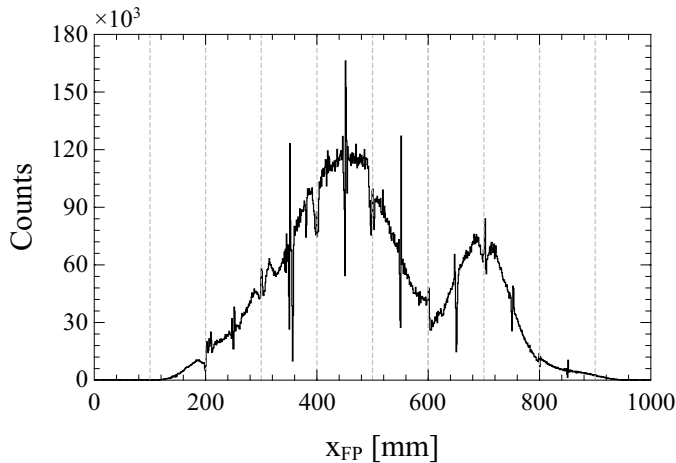


Figure 7.3: Calibrated x position of the ions detected by the MWPPAC focal plane detector. Grey dashed lines indicate the separation point between different sections.

7.2.3 TOF calibration

The ToF of the detected ions is derived from the difference between the timing signal of the MWPPAC and the MCP. The timing signals from the sections of the cathode act as a start, while the delayed time signal of the MCP is the common stop. In this way, one avoids the acquisition of events with ions that pass through the MCP but do not reach the focal plane. Independent TAC modules are then used to process the 10 signals from the cathode, which need to be calibrated from channels (arbitrary units) to nanoseconds. Considering the stability of the electronic modules, a set of coefficients is generally provided by the local team of LNL.

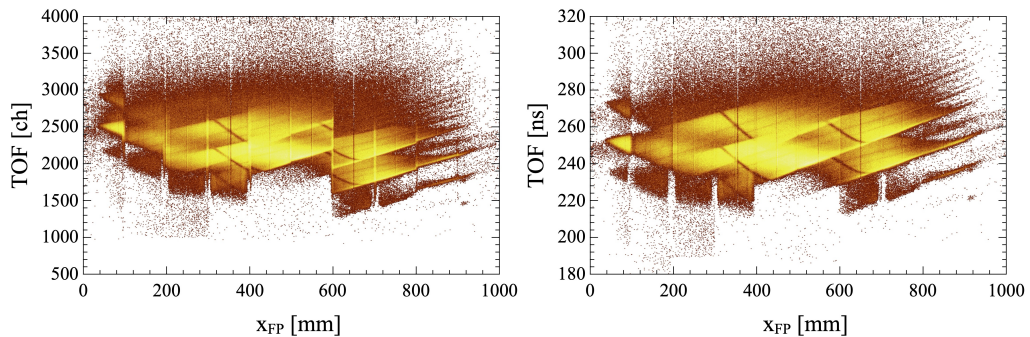


Figure 7.4: Raw (left) and calibrated (right) ToF as a function of the x focal plane position.

The ToFs of the different sections are not aligned after the calibration, therefore they have to be adjusted with individual offsets, as shown in Figure 7.4. In this way, the whole MWPPAC can be used as a single detector to determine the velocity β and the A/q ratio of the ions.

Once the alignment procedure is performed, a common ToF offset has to be set to retrieve the correct time-of-flight values. A first estimation of this value can be obtained based on the approximate path length of 6 m between the MCP and the MWPPAC, and the expected velocity $\beta \sim 8\%$ of the ions after the energy losses in the target material.

In a second stage, a fine tuning of the ToF alignment is achieved from the A/q versus x_{FP} correlations and, more precisely, from the Doppler-corrected γ -ray spectra observed in coincidence. In fact, the final ToF value results from the optimization of the centroid and FWHM of the Doppler-corrected γ -ray spectra, which depends on the ToF, hence the velocity, of the detected ions.

7.2.4 Z identification

At the end of their trajectories in the PRISMA spectrometer, the ions are stopped in the ionization chamber, which is used to perform the Z identification of the ions. As discussed in Section 6.1.2, the IC consists of a total 40 central pads and 8 side pads, each providing an independent signal. At first, lower and upper thresholds are set on each signal. In particular, setting lower thresholds for the side pads is crucial, as they act as a veto to reject trajectories of ions exiting the detector. Then, the signals of all pads are aligned with respect to each other, matching the gain of the different sections. A proper conversion from ADC units to MeV is not possible, since the calibration coefficients are determined from a linear interpolation of signals injected in the pads with a pulser.

Thanks to its segmentation, the ionization chamber can be used as an E- ΔE telescope. The loss of energy of heavy charged particles in the matter follows the Bethe-Block formula [169]:

$$\frac{dE}{dx} \approx -Z^2 \left(\frac{e^2}{4\pi\epsilon_0} \right) \frac{4\pi z N_A \rho}{m_e c^2 \beta^2 A} \left[\ln \left(\frac{2m_e c^2 \beta^2}{I} \right) - \ln(1 - \beta^2) - \beta^2 \right], \quad (7.6)$$

where Z is the atomic number of incident particle, while z , A and ρ are atomic number, the mass number and the density of the medium. $\beta = v/c$ is the velocity of the ion, N_A is the Avogadro number, m_e is the electron mass and I is the ionization potential. In the case of non-relativistic particles ($\beta \ll 1$), the dominant terms of the previous expression are given by

$$\frac{dE}{dx} \propto \frac{Z^2}{E}. \quad (7.7)$$

Thus, the integrated energy loss ΔE of a particle over a finite path length is proportional to the ratio Z^2 , and inversely proportional to the its kinetic energy E . From the IC signals, the ΔE is determined as the energy loss in the first two longitudinal sections, whereas the total energy E is obtained as the sum of the energy losses in all the four sections. In addition, the pressure of the gas is optimized in order to completely stop the reaction products and have a Bragg peak typically centered in the third section.

By studying the $E - \Delta E$ correlation matrix, the Z identification can be performed, as shown in Figure 7.5. Different branches of hyperbola correspond to different elements, ranging from S ($Z = 16$) down to Mg ($Z = 12$), which are selected using polygonal gates (shown in Figure 7.5 as red lines).

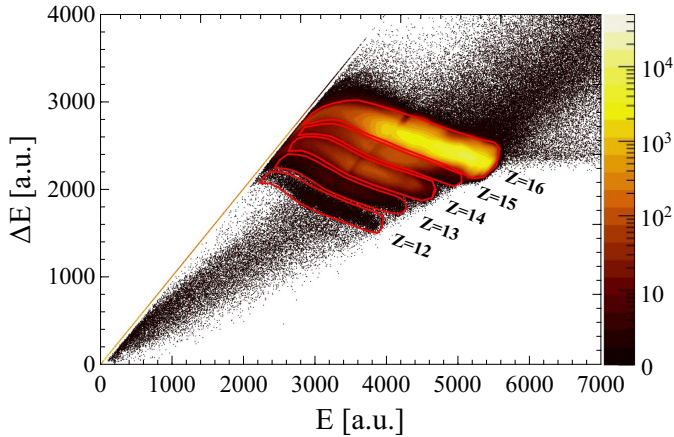


Figure 7.5: $E - \Delta E$ matrix of ions detected in the IC of PRISMA. ΔE is the energy observed in the first two longitudinal sections, while E is the total energy released in the IC. Structures corresponding to elements from S ($Z=16$) to Mg ($Z=12$) are selected using the polygonal gates shown in red.

7.2.5 Trajectory reconstruction

A pivotal step in the analysis routine of PRISMA is the reconstruction of the trajectories, which leads to the isotopic identification of the incoming ions. Combining time and position information provided by the MCP and the MWPPAC with the knowledge of

the magnetic field of the optical elements, it is possible to reconstruct the trajectory of the ions and determine their velocity on an event-by-event basis.

The PRISMA algorithm performs the kinematic reconstruction with a few assumptions, simplifying the calculations and compensating the lack of a precise mapping of the magnetic fields. The ions are assumed to be emitted from the center of the target, the motion of the ions is considered planar in the horizontal plane of PRISMA, and the magnetic elements are supposed to be ideal.

In a quadrupole magnet, the magnetic field is defined as

$$\vec{B} = -\vec{\nabla}U(x, y), \quad (7.8)$$

where U is given by

$$U(x, y) = \frac{B_{max}}{R_q}xy. \quad (7.9)$$

In the latter, B_{max} is the maximum intensity of the magnetic field and R_q is the inner radius of the quadrupole magnet. Ions passing through the magnet are then subject to the Lorentz force, which can be expressed as

$$\vec{F} = q\vec{v} \times \vec{B}. \quad (7.10)$$

Since the ions are assumed to be moving only in the z direction, the effect of the magnet is to focus on the horizontal direction x ($F_x > 0$), while defocusing on the vertical direction y ($F_y < 0$). Using the coordinates of the MCP interaction and the initial position at the target center, the analysis algorithm is able to solve these equations and estimate the hyperbolic motion induced by the quadrupole.

After exiting the quadrupole, the ions enter in the dipole magnet, which is characterized by a uniform magnetic field oriented in the y vertical direction. Here the Lorentz force induces the ions to follow a uniform circular motion on the horizontal plane with radius ρ . By applying the equivalence between the Lorentz force and the centripetal force mv^2/ρ , one obtains the magnetic rigidity $B\rho$ as

$$B\rho = \frac{p}{q} = \frac{A}{q}v, \quad (7.11)$$

where $p = mv$ is the momentum, q is the charge and A is the mass of the ion. Since the trajectory from the magnetic dipole to the MWPPAC is supposed to be a straight line, the arrival position on the focal plane completely depends on the exit point from the dipole, hence on ρ . An iterative procedure is applied by the algorithm to determine the optimal trajectory. From the MCP position and an initial ρ value of 120 cm, it minimizes the difference between the calculated and the measured position on the focal plane, changing the ρ value at each step. The convergence is reached when the difference is below 1 mm. As an output, the algorithm provides the total length L of the trajectories from the MCP to the MWPPAC.

At this stage, the velocity of the ions can be calculated as

$$\beta = \frac{L}{ToF \cdot c}. \quad (7.12)$$

Thus, combining Eq. (7.11) and (7.12), one obtains the expression for the A/q ratio:

$$\frac{A}{q} = \frac{B\rho}{L}ToF. \quad (7.13)$$

In the present derivation, relativistic effects have been neglected since the velocity β of the ions entering PRISMA is of the order of 8%.

7.2.6 Q identification

Since the atomic number Z and the A/q ratio are provided by the $E-\Delta E$ method and the trajectory reconstruction procedure, respectively, the determination of the charge states would directly lead to the mass identification of the incoming ions. This can be achieved by combining information from the MWPPAC and the IC, as described in the following. According to Eq. (7.11), the curvature radius of the ions depends on their charge state q , therefore, the same isotope follows different trajectories based on its charge state. In a non-relativistic regime, by multiplying both sides of Eq. (7.11) by the velocity v , one obtains the expression

$$B\rho v = \frac{mv^2}{q} = 2\frac{E}{q}, \quad (7.14)$$

where $E = \frac{1}{2}mv^2$ is the kinetic energy of the ion. Since B is constant and $\beta = v/c$, this leads to

$$E \propto q\rho\beta. \quad (7.15)$$

Experimentally, the total energy E of the ion is measured by the IC, while ρ and β are determined from the trajectory reconstruction. Therefore, different charge states can be separated by correlating the total energy E and the product $\rho\beta$. In the analysis, this identification is performed by gating on the atomic number Z of the ions, drastically reducing the background contributions. Bi-dimensional gates are then applied to select events corresponding to different charge states, as shown in the correlation matrix of Figure 7.6, where the 11^+ , 12^+ and 13^+ charge states of the $Z = 14$ Si ions are clearly identified. Since the absolute value in MeV of the kinetic energy is not measured, the charge state of each band structure is initially assigned based on the theoretical distributions estimated with the REACTION code [170]. These assignments can be verified in a second stage from the mass number identification.

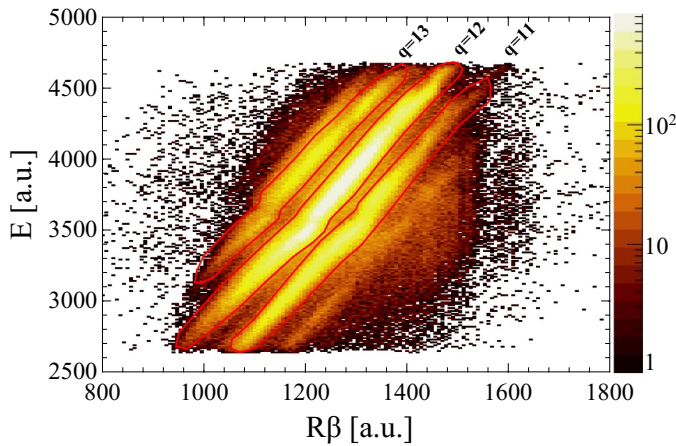


Figure 7.6: Matrix correlating the total energy E observed in the IC with the product of the radius ρ and the velocity β , gated on $Z = 14$. Structures corresponding to different charged states are selected using polygonal gates (red lines), which are labeled accordingly.

7.2.7 Mass identification

The last step in the PRISMA data processing involves the event-by-event identification of the mass number A of the reaction products. Ideally, the mass spectrum could be determined by simply multiplying the reconstructed A/q ratio for each charge q , summing over contributions from the different charge states. However, this is done after a series of optimizations and calibrations that noticeably improve the quality of the mass spectra.

A significant reduction in the mass resolution is caused by the presence of optical aberrations, which originate from higher-order magnetic fields and fringe fields. Such effects are not considered in the trajectory reconstruction routine, which treats the optical element and the motion of the ions in a first order approximation, and have to be corrected in a second stage.

As a first step, the nominal values of the quadrupole length and the distance target-quadrupole are slightly modified (within 10%) to take into account the presence of fringe fields and higher-order magnetic fields. An iterative procedure was applied to determine the optimal effective parameters that maximized the A/q resolution.

The optimization procedure, however, is not sufficient to completely eliminate aberration effects. As shown in the left panel of Figure 7.7, drifts are still observed in the A/q ratio as a function of the x position on the MCP. Similar trends are observed as a function of the y position on the MCP and x_{FP} focal plane position. Starting from the A/q versus

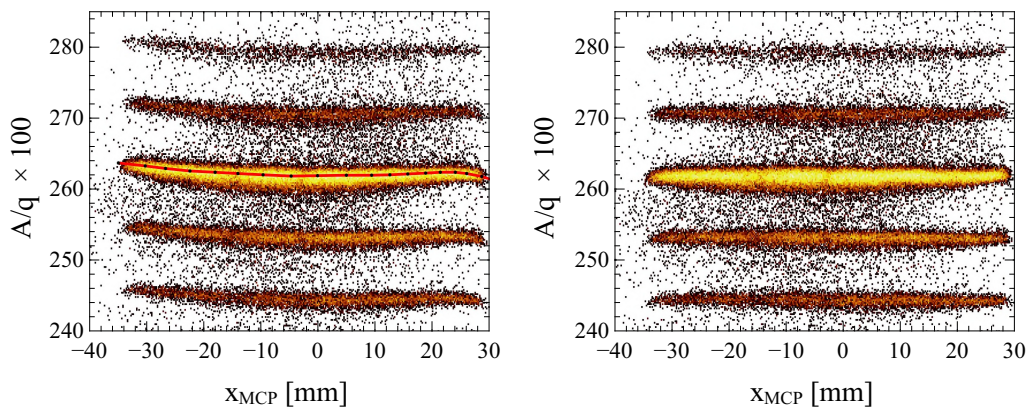


Figure 7.7: A/q spectrum of the $Z = 14$ ions, gated on $q = 12$, as a function of the x MCP position before (left) and (after) the aberration correction procedure. In the left panel, the spline function used to correct the A/q drifts is shown with a red line.

x MCP matrix, the aberration correction is applied by fitting the most intense structure of A/q ratio, gated on a charge state q , with spline functions (or alternatively high-order polynomial). The result of this correction is shown in the right panel of Figure 7.7. The same procedure was then repeated, in order, on the A/q versus y MCP matrix and on A/q versus x focal plane matrix.

Once the aberrations are corrected, the final mass spectra can be obtained after calibrating the A/q ratio. This is done to take into account nonlinearities in the A/q reconstruction, which result in discrepancies between the measured values and the calculated value. The latter is estimated from an assumed mass distribution, which can be verified by looking at the γ rays in coincidence with a specific A/q ratio. After the linear calibration, the mass A is obtained by multiplying the A/q ratio for an effective charge q_{eff} , which is defined as

$$q_{eff} = \frac{1}{N} \sum_{i=1}^N \frac{A_i}{(A/q)_i^{cal}}, \quad (7.16)$$

where N is the number of peaks in the A/q spectrum, $(A/q)_i^{cal}$ is the centroid of the i -th peak and A_i is the expected mass number. As an example, Figure 7.8 shows the calibrated A/q ratio, as a function of the x focal plane position, gated on $Z = 14$.

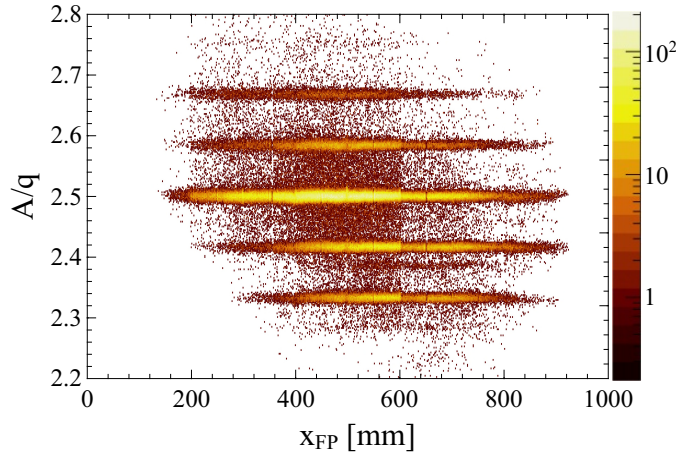


Figure 7.8: A/q of the $Z = 14$ ions as a function of the x focal plane position, as obtained after the aberration corrections and the A/q calibration procedures.

The final mass distribution is obtained by summing over contributions from each charge state:

$$A = \sum_j q_j^{eff} \left(\frac{A}{q} \right)_j. \quad (7.17)$$

In Figure 7.9, the mass spectra obtained for each Z are shown. An excellent separation between mass peaks is observed, corresponding to an average mass resolution of $\Delta A/A = 1/230$.

In the analysis, the mass number of each ion is typically assigned by rounding the calibrated mass to the closest integer. If needed, polygonal gates can also be set in the mass versus x focal plane matrix to improve the mass selection.

7.3 AGATA data processing

The data processing of AGATA consists of two main parts, called the *Local Level processing* and *Global Level processing*. At the local level, where each crystal is treated individually, the data are handled in a series of stages (or actors in the AGATA language), namely the *Preprocessing*, *Pulse-Shape Analysis (PSA)*, and *Post-PSA*. At the global level, the array is considered as a whole, and data streams from the crystals and complementary detectors, such as PRISMA, are combined. At this stage, the *Event Builder* and the *Tracking* routines are performed, and a final output is provided in the ROOT TTree format [124]. In Figure 7.10, all the steps of the AGATA data processing are summarized.

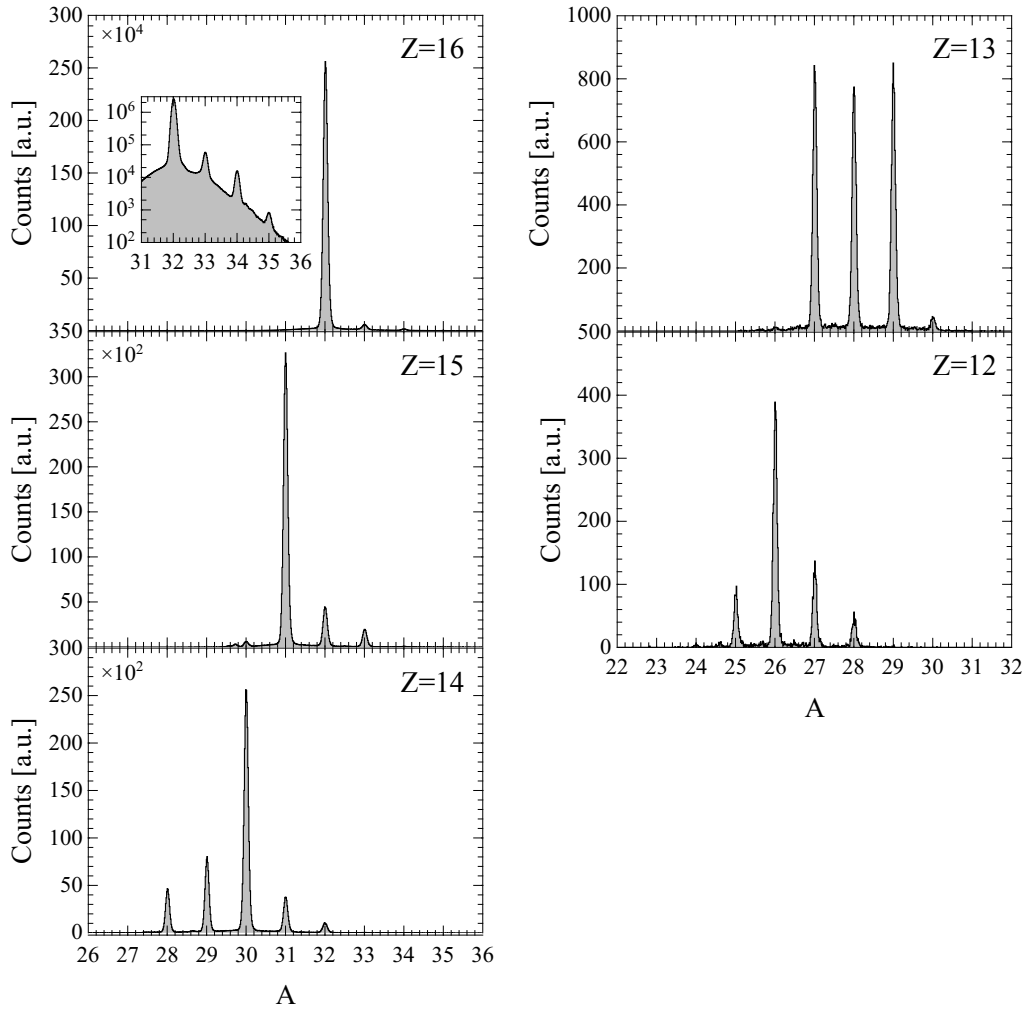


Figure 7.9: Mass spectra, gated on Z , of the ions detected in PRISMA. In the inset of the top-left panel, the $Z=16$ mass distribution is shown in logarithmic scale to better display the neutron-transfer channels.

7.3.1 Preprocessing

The Preprocessing stage is devoted to the preparation of the data for the PSA algorithm. The first calibration, the cross-talk corrections and the time alignment of the segments are performed.

Initial calibration

The raw signals of segments and cores are first calibrated with a ^{60}Co source, which is needed for an adequate cross-talk correction. At this level, the offset of the linear calibration coefficient is set to zero, and only events with a segment multiplicity equal to 1 are considered. This is done to rule out cross talk effects from the calibration procedure. In Figure 7.11, raw and calibrated segment spectra of crystal 02A are shown.

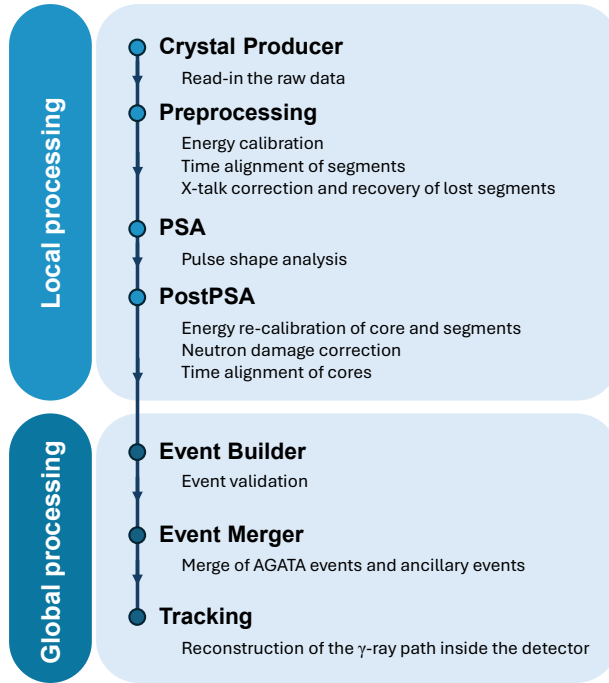


Figure 7.10: Schematic view of the processing stages of the AGATA data.

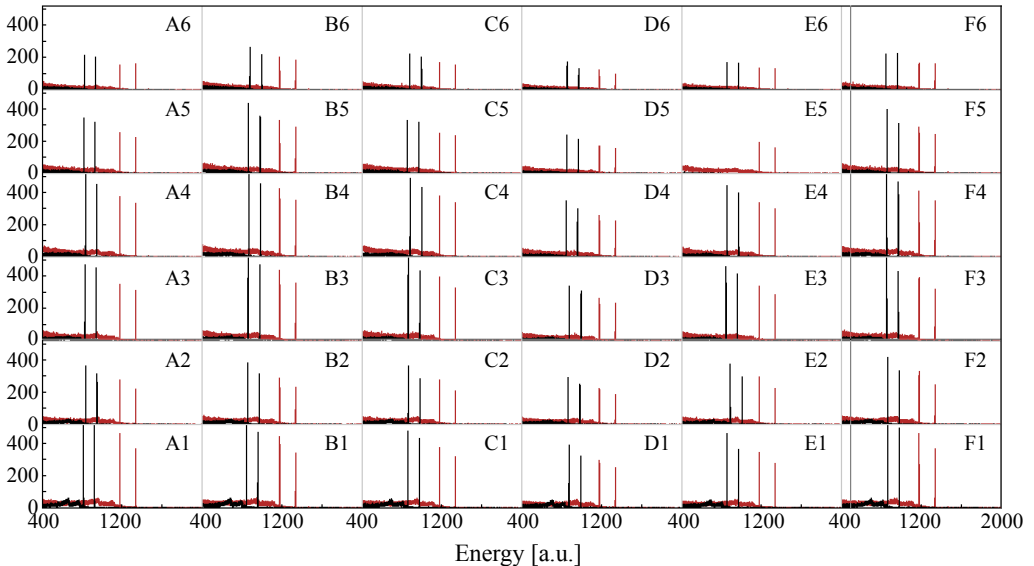


Figure 7.11: ^{60}Co spectra before (black) and after (red) the initial calibration of crystal 02A. The raw spectrum of crystal E5 is not visible due to its lower gain.

Cross-talk correction

Cross-talk is a phenomenon observed in any electrically segmented detector, and its effects on the signals have to be corrected to properly reconstruct the energy of incident

γ -rays. Due to the electrical couplings between the segments, the collection of a signal in one segment modifies the signals in the neighboring ones. One of the consequences is a dependence of the segment energies as a function of the multiplicity, as shown in Figure 7.12.

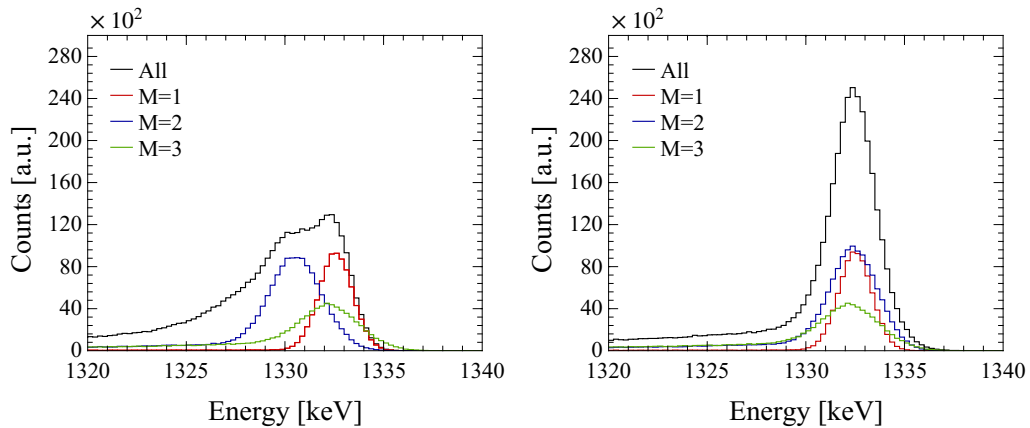


Figure 7.12: Partial γ -ray spectra of ^{60}Co before (left) and after (right) the cross-talk corrections. Spectra corresponding to different segment multiplicities are indicated with colored lines, while the spectrum associated to the sum of all multiplicities is shown in black.

In the Preprocessing routine, the cross-talk correction is implemented according to the linear cross-talk theory of Bruyneel et al. [171]. Using ^{60}Co data, the `xTalkSort` executable of the [167] package sorts the events according to their multiplicity and calculates the corresponding energy shift with respect to the nominal energy of the cobalt peaks. A set of coefficients is then extracted and applied to the data. An example of cross-talk corrected spectra can be found in the right panel of Figure 7.12.

The cross-talk correction can also be used to recover the signal of broken or missing segments. In case of only one faulty segment, the signal is recovered by subtracting the sum of the energies of all the other segments from the core energy. If more than one segment is malfunctioning, this procedure cannot be applied.

Time Alignment

As a last step of the Preprocessing stage, a time alignment of the segment signals is performed. This is crucial to maximize the performance of the PSA algorithm, which assumes that all traces start at the same time. Therefore, mismatches in time could lead to an unreliable reconstruction of the interaction point in the crystal, characterized by clusterization patterns. To prevent such effects, the difference between the CFD zero-crossing time of the segments and the core trigger timestamp is minimized by applying time shifts to each segment. The result of this procedure is shown in Figure 7.13, where the time spectra of all segments are aligned such that their mean values correspond to 500 ns.

Additionally, the core signals of all the detectors are aligned to the same reference value.

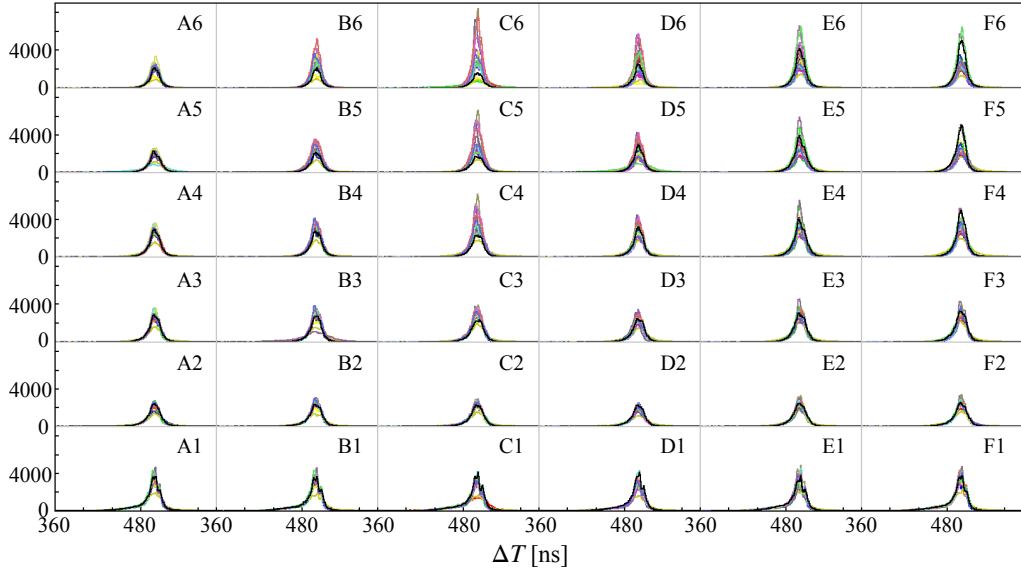


Figure 7.13: Time difference between segment and core signal after the alignment procedure for all the detectors used in the present experiment.

7.3.2 PSA

The PSA actor is in charge of extracting the position of each interaction of the γ rays in the HPGe crystals, based on the signal traces observed in the different segments. A γ ray interacting in a germanium medium creates electron-hole pairs which are collected by the electrodes. In the charge-collection process, a net signal is generated in the hit segment, whereas a transient signal is induced on neighboring segments. Due to the detector geometry, high-voltage, impurities and mobilities of the charge carries, the induced signals exhibit a clear spatial dependence. This phenomenon can thus be exploited to reconstruct the interaction position within a segment of a crystal.

By comparing the signal traces with the position-dependent pulses of a reference library, the PSA is able to determine the interaction position with a FWHM of 4 mm [172]. The library of traces used in AGATA is the AGATA Data Library (ADL) [173], which is generated via Geant4 simulations and contains pulses simulated within the crystal volume in a spatial grid of 2 mm. The Adaptive Grid Search algorithm handles the χ^2 minimization procedure [172], providing a good balance between computational performance and accuracy.

7.3.3 PostPSA

At the Post-PSA stage, additional refinements can be applied on a local level to improve the performances of the detectors. In particular, the neutron-damage correction, the final linear recalibration and the global time alignment are executed.

Neutron damage correction

Exposure to fast neutrons is known to deteriorate the energy resolution of n -type HPGe detectors, including the AGATA crystals. Neutron interactions displace atoms in the

germanium lattice, creating defects that act as trapping centers for charge carriers, particularly holes. These traps are characterized by very long release times on the order of milliseconds. As a consequence, the charge collection process is incomplete, and the measured spectrum exhibits an exponential tail on the low-energy side of the photopeak.

In AGATA, the segmentation of the detectors combined with the PSA are used to mitigate the effects of neutron-induced damage. The standard correction procedure employed by the AGATA collaboration follows the method proposed by Bruyneel *et al.* [174]. In this approach, the measured energy of each channel i (whether core or segment) is recalibrated according to its sensitivity to electrons and holes, denoted as $s_{e,h}$, and two parameters $\lambda_{e,h}$ that quantify the extent of neutron damage:

$$E' = \frac{E}{1 + s_e/\lambda_e + s_h/\lambda_h}. \quad (7.18)$$

The parameters $\lambda_{e,h}$ are extracted using calibration data from the 1332.5 keV transition of a ^{60}Co source. The `agapro` executable `SortPsaHits` performs a grid search over possible parameter values, optimizing them by maximizing the photopeak integral while simultaneously minimizing the resolution (FWHM) of the peak. The parameters yielding the best compromise between peak intensity and resolution are then fixed and applied consistently throughout the analysis.

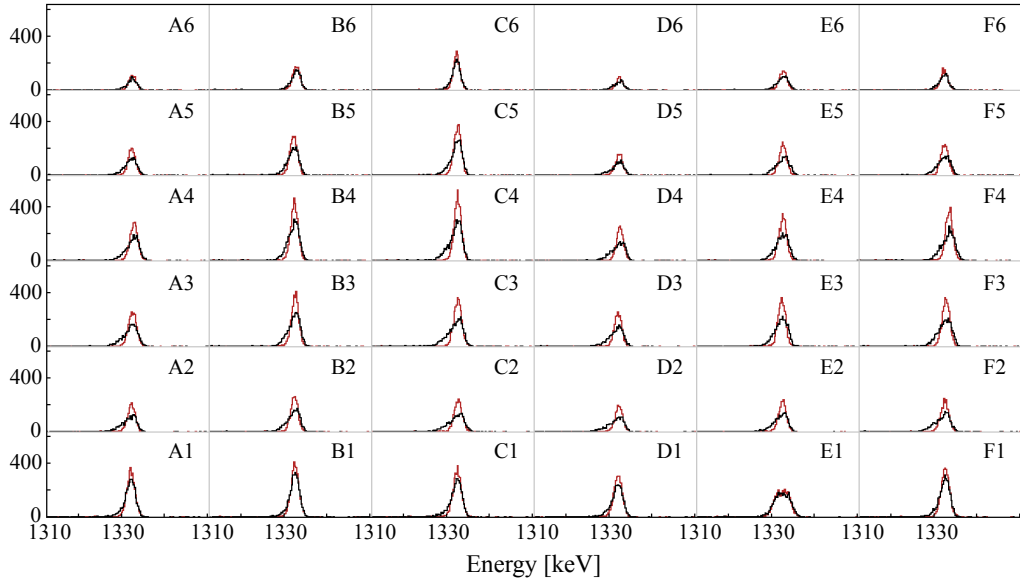


Figure 7.14: Partial γ -ray spectra of ^{60}Co before (black) and after (red) the neutron-damage correction of crystal 11A.

In the present experiment, neutron damage effects were clearly visible only in the 05B, 08A, 08B, 08C, 11A and 11B crystals. In fact, the experiment was performed at the beginning of the LNL campaign of AGATA, when most of the crystal were in optimal condition. As an example, the spectra of detector 11A before and after the neutron damage correction are shown in Figure 7.14.

Final calibration

After the neutron damage correction, a final recalibration is performed using data from a ^{152}Eu source. A linear behavior is observed in most of the crystals, therefore, first-order polynomials are used. The good quality of the calibrated core spectra will be later discussed in Section 7.3.6, in comparison with that obtained from the tracking.

At this stage, one can renormalize the energy of the sum of the segments to that of the corresponding core. This function minimizes the impact of faulty segments on the performance of a single core, improving the peak-to-total ratio. Moreover, it reduces the tails on the left-side of peaks caused by the neutron damage. For this reason, this option was applied on this dataset.

Global time alignment

The final time alignment of the crystal cores is performed at the Post-PSA level. Since the Event Builder assumes that all the interacting γ -rays are detected simultaneously by the array, this step reduces random coincidences and improves the tracking performance. The time differences between each pair of crystals are calculated with the `SolveTT.py` executable [167], and the resulting time shifts are subsequently applied to optimize the alignment.

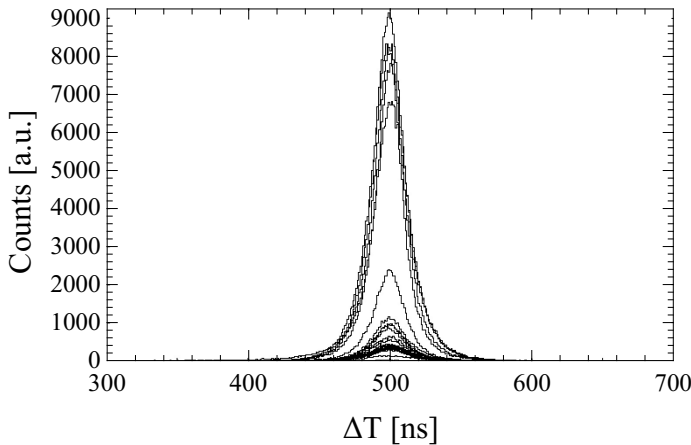


Figure 7.15: Time differences between crystal 02A and all the other detectors present in the acquisition, as obtained after the global time alignment procedure.

7.3.4 Event Builder and Merger

In the Event Builder routine, the synchronized signals of the AGATA crystals are merged with the ancillary data, such as PRISMA, by matching their recorded timestamps. A large time window of $1\ \mu\text{s}$ is used to build coincidence events. At a later stage in the analysis, a narrower gate on the PRISMA-AGATA time-difference spectrum is applied (see Section 8.2).

7.3.5 Tracking

As introduced in Section 6.1.3, new-generation γ spectrometers, such as AGATA and GRETINA, are able to identify the sequence of interactions associated with each incident

photon. This task is performed using *tracking algorithms*, which reconstruct the photon trajectory through the detector material, providing the total deposited energy and incoming direction. The tracking procedure accounts for the various mechanisms through which γ rays interact with germanium and energy ranges in which these processes dominate: the Photoelectric effect (below ~ 100 keV), Compton scattering (~ 100 keV–10 MeV), and pair production (above ~ 10 MeV).

Several tracking methods have been developed, which can be broadly categorized into forward tracking, backward tracking, Bayesian approaches, and more recently, machine-learning-based algorithms [175]. Both forward and backward tracking techniques rely on the Klein–Nishina formalism of Compton scattering, which represents the dominant interaction mechanism in the energy range most relevant to AGATA. The fundamental assumption is that, except at very low γ -ray energies, Compton scattering tends to predominantly occur in the forward direction. As a consequence, consecutive Compton interactions are expected to be located close to each other in the (θ, ϕ) angular plane.

In the Orsay Forward Tracking (OFT) code [176], the one used by the AGATA Tracking actor, the reconstruction starts with the identification of the first interaction point. Interaction points are then grouped into spatial clusters, and, within each cluster, the most probable interaction sequence is determined by minimizing a figure of merit (FOM). The latter evaluates the consistency of the measured energy and photon scattering angles with the Compton formula.

The AGATA tracking routine uses three main parameters [172]:

- σ_θ (**SigmaTheta**): the average spatial resolution of the interaction-point positions,
- P_{track} (**MinProbTrack**): the minimum probability for a γ ray to undergo a Compton interaction,
- P_{sig} (**MinProbSig**): the minimum probability for a γ ray to penetrate a given depth and be absorbed via the photoelectric effect.

In the most recent implementation of the OFT code, the parameter P_{sig} has been fixed to a constant value. Values of $\sigma_\theta = 0.8$ and $P_{\text{track}} = 0.05$ were used in the present analysis.

7.3.6 Performances

At the end of the preliminary processing of AGATA data, the performance of the array is evaluated in terms of calibration accuracy, energy resolution and absolute efficiency.

As discussed in Sections 7.3.1 and 7.3.3, the calibration of segments and cores is performed in two steps: an initial gain matching with ^{60}Co data and, after the cross-talk and the neutron-damage corrections, a second calibration with ^{152}Eu data. The final calibration residuals, i.e. the differences between tabulated and calibrated energies, obtained for the cores energies of each functioning detector are shown in Figure 7.16. It is clear that the residuals are mostly concentrated in the ± 0.2 keV range, with only three detectors showing discrepancies larger than 0.4 keV. The good quality of the calibration is confirmed by the tiny discrepancies observed in the tracked energies, which are consistently smaller than 0.1 keV.

Figure 7.17 shows the energy resolution of the AGATA array as a function of the γ -ray energy. In particular, outlier data points correspond to detectors 00A and 02A, whose core signal showed a significantly worse resolution with respect to the others. At 1.3 MeV, the FWHM is 2.5 keV for the average over the cores, while a value of 2.8 keV is

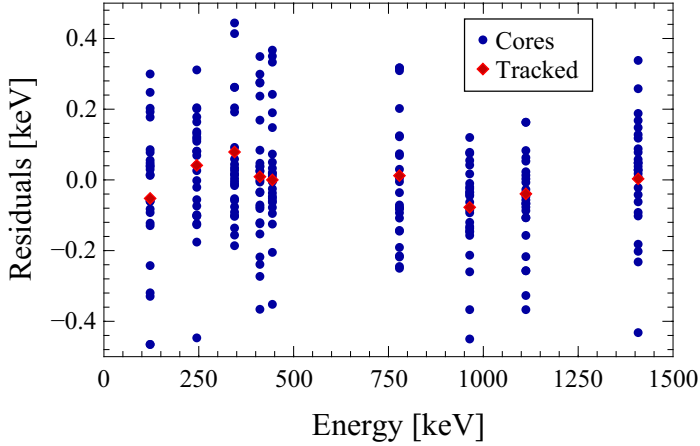


Figure 7.16: Residuals of the peak energy measured in the different core (blue circles) and tracked (red diamond) spectra, using ^{152}Eu source data.

obtained with tracking. The tracked resolution curve has been fitted with the empirical expression of (4. 1), and the result is displayed in Figure 7.17 with a red line.

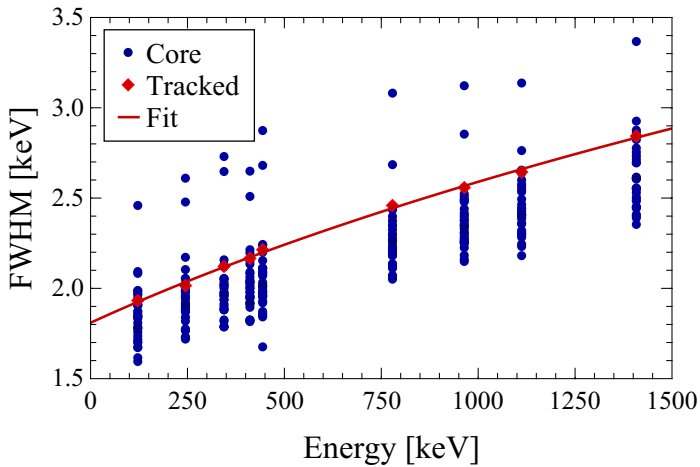


Figure 7.17: Resolution (FWHM) of AGATA as a function of the γ -ray energy with data from a ^{152}Eu source. The single core and the tracked resolutions are indicated with blue circles and red diamonds, respectively. The fit of the tracked curve is indicated with a red line.

Finally, the absolute efficiency of the array was estimated using an activity-calibrated ^{152}Eu source, placed in the center of the chamber. As in the whole experiment, AGATA was in the nominal position, and no absorbers were mounted.

Based on the activity A of the source at the time of the experiment, the acquisition time Δt , the branching ratio $BR(E_\gamma)$ and the observed integral $I(E_\gamma)$ of a γ -ray transition, the absolute efficiency can be calculated as

$$\varepsilon(E_\gamma) = \frac{I(E_\gamma)}{BR(E_\gamma) \cdot A \cdot \Delta t}. \quad (7.19)$$

Since the average count-rate was approximately 700 Hz for each AGATA crystal, the dead-time effects were found to be negligible. In Figure 7.18, the measured absolute efficiency for the total core and the tracked spectra are shown. Data points were also fitted with Radware functions [120] (see Eq. (4.2)), which are superimposed to the experimental data points. At the reference energy of 1.3 MeV, absolute efficiencies of about 2.9% and 4.0% are obtained for core and tracked spectra, respectively. These values are in good agreement with the simulated efficiency estimated in Ref. [152] for a similar configuration of the AGATA array.

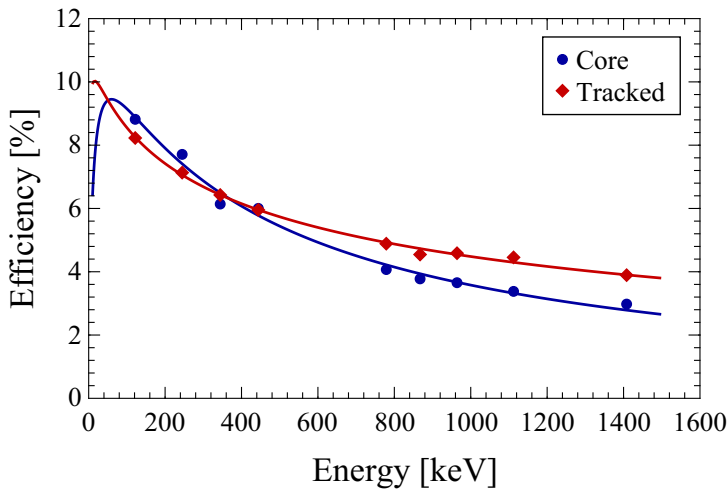


Figure 7.18: Absolute efficiency of AGATA in the nominal position as a function of the γ -ray energy with data from a ^{152}Eu source. Experimental points and fitted Radware functions [120] from the total core and tracked spectra are indicated in blue and red, respectively. The statistical errors are not visible as they are smaller than the symbol sizes.

In this Chapter, the data analysis and the results of the AGATA-PRISMA experiment are presented. Section 8.1 focuses on the study of the $^{32}\text{S} + ^{110}\text{Cd}$ multi-nucleon transfer reaction, describing the procedure through which differential and inclusive cross sections have been determined. Section 8.2 illustrates the spectroscopic study of the target-like reaction products, which was performed via particle- γ coincidence techniques. Finally, Section 8.3 addresses the possibility of estimating upper limits and lifetimes for excited states of reaction products populated in the current dataset.

8.1 Reaction study

This Section focuses on the study of the $^{32}\text{S} + ^{110}\text{Cd}$ reaction mechanism. After the introduction of physical quantities, such as the Q -value, the experimental procedure to determine the cross section of elastic and transfer channels is described.

8.1.1 Reaction Q -value

Consider a binary reaction process $a + A \rightarrow b + B$, where the beam nucleus a impinges on the target nucleus A , at rest, and leads to the production of the beam-like nucleus b and the target-like nucleus B . If the total kinetic energy before and after the collision are $K_i = K_a$ and $K_f = K_b + K_B$, the total kinetic energy loss ($TKEL$) of the reaction can be defined as

$$TKEL = K_i - K_f = K_a - (K_b + K_B), \quad (8.1)$$

where K_b and K_B are the kinetic energies of the beam-like and of the target-like reaction products, respectively. This quantity is the opposite of the reaction Q -value, namely

$$Q = -TKEL. \quad (8.2)$$

Since the total kinetic energy is conserved in a binary reaction process, one obtains

$$m_a c^2 + m_A c^2 + K_a = m_b c^2 + m_B c^2 + K_b + K_B + E_b^* + E_B^*, \quad (8.3)$$

being E_b^* and E_B^* the excitation energies of the two reaction products. Therefore, Eq. (8.1) can be written as

$$TKEL = K_i - K_f = (m_b + m_B - m_a - m_A)c^2 + (E_b^* + E_B^*). \quad (8.4)$$

When $E_b^* = E_B^* = 0$, the equation reduces to the opposite of the ground-state Q value (Q_{gg}), which corresponds to the mass difference of the nuclei involved in the reaction. In

the case of elastic scattering, both Q_{gg} and the energies E_b^*, E_B^* are zero, and the $TKEL$ distribution is centered around 0 MeV.

In the analysis, the $TKEL$ is reconstructed based on experimentally measured quantities. By applying the momentum conservation law

$$\vec{p}_a = \vec{p}_b + \vec{p}_B, \quad (8.5)$$

and recalling that $p^2 = 2mK$, it is possible to rewrite Eq. (8.1) as

$$TKEL = \frac{m_B - m_a}{m_B} K_a + \frac{2\sqrt{m_a m_b K_a K_b}}{m_B} \cos \theta_b - \frac{m_b + m_B}{m_B} K_b. \quad (8.6)$$

The beam-like kinetic energy K_b and scattering angle θ_b are directly measured by PRISMA, while K_a is the chosen beam energy. It is important to note that the measured energies have to be corrected for the energy losses in the target material, which are not negligible for targets thicker than a few hundred micrometers.

Figure 8.1 shows the total $TKEL$ spectrum of ^{32}S , integrated over the full angular acceptance of PRISMA (black curve). In the thick-target configuration, ^{32}S ions scattered on both the ^{110}Cd target and the Ta backing were simultaneously detected, giving rise to two distinct structures. The peak at negative $TKEL$ values corresponds to scattering on Ta, which is associated with longer times of flight, while the peak centered near 0 MeV is attributed to scattering on Cd. The finite thicknesses of the target and backing layers (1 and 2 mg/cm², respectively) introduce a non-negligible uncertainty in the $TKEL$ reconstruction. Although energy-loss corrections are applied event-by-event based on the ion identification, the uncertainties in the interaction depth, the energy straggling, and the intrinsic resolution of the PRISMA setup lead to an energy resolution of approximately 7 MeV.

For the purposes of the present analysis, only the Cd peak is relevant, as it provides an estimate of the intensity of the elastic channel. The procedure to extract its cross section is discussed in the following section.

8.1.2 Cross section of the elastic channel

The investigation of elastic scattering is the first step in the analysis of cross sections and angular distributions of nuclear reactions. At forward angles, where the effects of the nuclear interaction are negligible, the elastic cross section is assumed to be described by the Rutherford formula. Therefore, it can be used as a normalization reference for other reaction channels. In practice, once the intensity of the elastic channel is determined, it provides a conversion factor between the observed counts and mb, allowing experimental data and theoretical calculations to be compared on an absolute scale.

After identifying the beam-like ^{32}S ions in PRISMA through their Z and A, elastic scattering events can be selected using the method described in Ref. [177], by comparing $TKEL$ spectra with and without the requirement of γ coincidences. These spectra, shown in Figure 8.1, correspond to the total (black line) and to the inelastic (red line) contributions. The latter is normalized to match the inelastic tail of the total spectrum (in the region above 10 MeV), thus compensating for the AGATA efficiency. By subtracting the normalized inelastic contribution from the total spectrum, one obtains the elastic component (blue line).

In order to determine the elastic cross section, the experimental counts in the Cd peak of the elastic $TKEL$ spectrum are first measured for each angle θ_{lab} covered by PRISMA, in steps of 1 degree. Events associated with scattering on the Ta backing were discarded

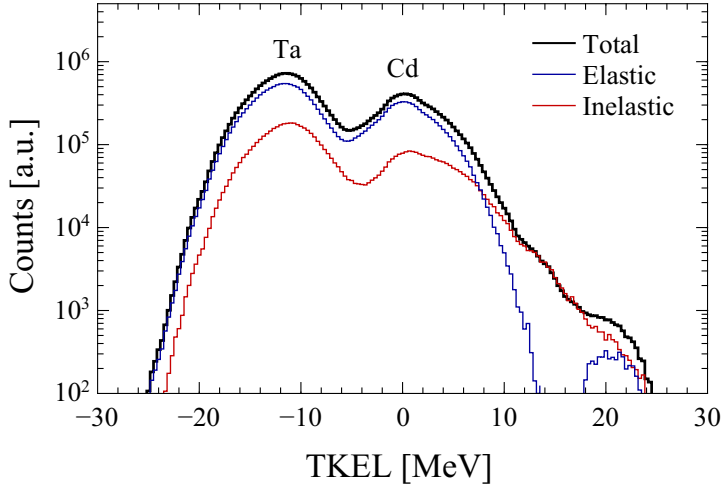


Figure 8.1: *TKEL* spectra for ^{32}S ions. The two main peaks correspond to the scattering on the Ta backing (left) and on the Cd target (right), as labeled in the figure. The total distribution measured by PRISMA is in black, while elastic and inelastic contributions are in red and blue, respectively. The elastic spectrum is obtained by subtracting from the total spectrum the inelastic distribution, for which the coincidence with γ rays in AGATA is required (see text in Section 8.1.2 for details).

by gating on $TKEL > -5$ MeV. Subsequently, the measured counts are corrected for the solid angle $\Delta\Omega$ corresponding to each θ_{lab} angle. This can be calculated using

$$\Delta\Omega = \int_{\varphi_i}^{\varphi_f} d\varphi \int_{\theta_i}^{\theta_f} \sin\theta d\theta = \Delta\varphi [\cos(\theta_i) - \cos(\theta_f)], \quad (8.7)$$

where $\theta_{i,f} = \theta_{lab} \mp 0.5^\circ$, and $\varphi_{i,f}$ are the boundaries of the φ range associated to θ_{lab} . At this stage, the conversion factor C can be obtained by normalizing the experimental data in counts/sr to the elastic cross section, setting their ratio to one at an angle where the Rutherford scattering dominates. This is illustrated in Figure 8.2, where experimental points, normalized at 67° , are compared to the elastic angular distribution calculated with the GRAZING code [94], both divided by the Rutherford differential cross section. Based on this procedure, the resulting conversion factor is

$$C = 1 \text{ mb} = 3.67 \times 10^5 \text{ counts}. \quad (8.8)$$

In Figure 8.2, a good agreement is observed between experimental points and the theoretical calculations. The experimental errors, which are smaller than the symbol size, only account for statistical errors. Systematic uncertainties have not been included in the present analysis, but could be evaluated in future work through a detailed study of the response function of the PRISMA spectrometer [178].

It must be noted that GRAZING calculations were performed by assuming a reaction energy of 142 MeV. Despite being slightly lower than the expected mid-target energy of 144 MeV, this energy provides an improved agreement with the experimental data points, and could be explained by an uncertainty of a few degrees in the target rotation. This would result in higher energy losses of the beam in the first half of the target, thus reducing the average reaction energy. This assumption will be adopted and maintained throughout the analysis.

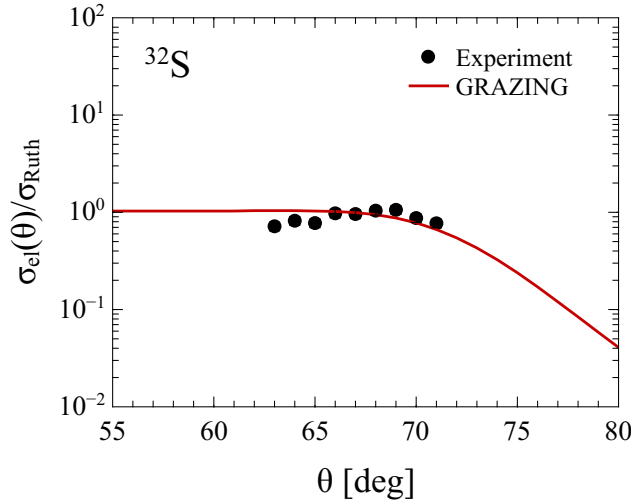


Figure 8.2: Ratio of the elastic cross section σ_{el} and the Rutherford cross section for $^{32}\text{S} + ^{110}\text{Cd}$ reaction. Experimental data points (black circles) are compared with theoretical calculations (red line) performed with the GRAZING code [94]. At this stage, the response function of the PRISMA spectrometer [178], mostly affecting the edge of the angular distribution, is not included yet.

8.1.3 Inclusive cross sections

Inclusive cross sections represent the total probability of producing a nucleus in a given reaction process. In the present work, inclusive differential and total cross sections have been studied for the identified transfer channels.

In Figure 8.3, the angular distributions of the main transfer channels are shown. Experimental data points were first obtained by integrating angle by angle the $TKEL$ measured by PRISMA. Then, the differential cross section has been calculated by normalizing the counts with the previously determined conversion factor C . These results are compared with theoretical calculations performed with the GRAZING code [94], including both low-lying and high-lying excitations. As illustrated in the central panels of Figure 8.3, GRAZING distributions are able to reproduce the single nucleon transfer with a reasonable agreement. Instead, when multiple nucleons are involved, larger discrepancies are observed. Such deviations can be attributed to the involvement of pair transfer and evaporation processes, which cannot be accounted for within the GRAZING model.

From the inclusive angular distribution, it is possible to derive the total cross sections measured in the angular acceptance of PRISMA. These can be obtained through the following integration:

$$\sigma = \sum_i \left(\frac{\Delta\sigma}{\Delta\Omega} \right)_i \times \Delta\Omega_i = \sum_i \left(\frac{\Delta\sigma}{\Delta\Omega} \right)_i \times 2\pi \times \Delta \cos \theta_{lab,i}, \quad (8.9)$$

where $\left(\frac{\Delta\sigma}{\Delta\Omega} \right)_i$ is the angular distribution at the laboratory angle $\theta_{lab,i}$ in steps of 1 degree. The term $\Delta \cos \theta_{lab,i}$ provides the integration of the solid angle $\Delta\Omega_i$ over $\theta_{lab,i}$, whereas the 2π factor results from the integration over the entire φ_{lab} range, for which these reactions are symmetric.

Such an integration procedure has been applied to both experimental and GRAZING

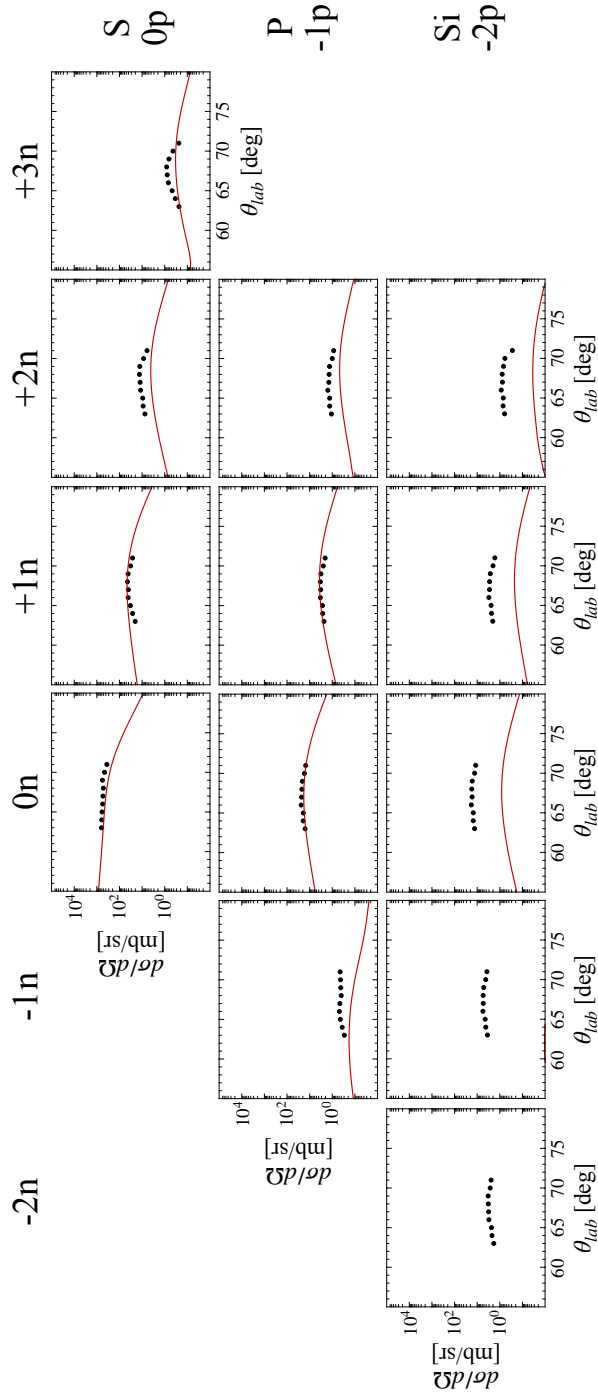


Figure 8.3: Angular distributions of the main reaction channels observed in the $^{32}\text{S} + ^{110}\text{Cd}$ reaction. Experimental data points are indicated by black symbols, while the theoretical calculations performed with the GRAZING code are displayed with red lines. Statistical errors bars are not visible, as they are smaller than the symbol size.

angular distributions. A comparison between the two is shown in Figure 8.4, including all the identified transfer channels. As previously observed for the angular distributions, GRAZING estimates (black lines) are in good agreement with the experimental values in single-nucleon transfer cases. The discrepancies increase gradually along the proton stripping side, where theoretical calculations underestimate the total cross section of channels involving several nucleons by orders of magnitude. Down to Si isotopes, the centroid of GRAZING distributions matches the experimental one, while Al and Mg distributions are substantially shifted, due to the increased relevance of neutron evaporation. Similar deviations have been observed in several recent works [95, 177, 179–181], and are mainly due to the limited capabilities of the GRAZING code, which does not include complex reaction processes in its treatment, such as the transfer of correlated pairs of nucleons.

8.2 Spectroscopy

In a multi-nucleon transfer experiment, particle- γ coincidences are powerful tools for data analysis. The event-by-event ion identification provided by PRISMA allows for an accurate selection of the reaction channel of interest, leading to coincidence γ -ray spectra with low background levels and negligible contamination from other channels.

As mentioned in 7.3.4, data from AGATA and PRISMA are initially merged, at the Event Builder stage, in a $1\mu\text{s}$ window. Prompt coincidence events from the multi-nucleon transfer reaction are then selected by applying an additional time condition. In Figure 8.5, the difference between AGATA and PRISMA timestamps is shown. To reduce the background from uncorrelated particle- γ events, a 140-ns time gate is applied on the narrow coincidence peak, which is highlighted with a red shaded area.

After selecting meaningful events, particle- γ matrices can be constructed. In the following Sections, the analysis methods used to perform the spectroscopic study are illustrated. The reconstructed level schemes for the main reaction channels are finally discussed in 8.2.3.

8.2.1 Doppler correction

In the experimental configuration of the present experiment, PRISMA detects and identifies beam-like reaction products. The latter, if excited, will emit γ -rays while in motion. As a consequence, the energy of the detected photons will be shifted according to the Doppler effect (see Eq. (2.22)). Similar considerations apply to the target-like partner of the reaction, which could emit γ rays either stopped in the target or in motion, if its kinetic energy is sufficiently large.

Since the PRISMA spectrometer measures the velocity v_b and the direction of the scattered beam-like ions b and AGATA measures the photon emission angle, it is possible to perform an event-by-event correction of the Doppler-shifted γ rays for both the beam-like and target-like products. The procedure is straightforward for the beam-like nucleus, as all the relevant quantities are directly measured. Conversely, the Doppler correction for the target-like nucleus requires the reconstruction of the reaction kinematics.

As introduced in Section 8.1.1, under the assumption of binary reaction process

$$a + A \rightarrow b + B, \quad (8.10)$$

it is possible to determine the magnitude v_B and the direction θ_B of the target-like binary partners B . A schematic view of this process is depicted in Figure 8.6, where a denotes

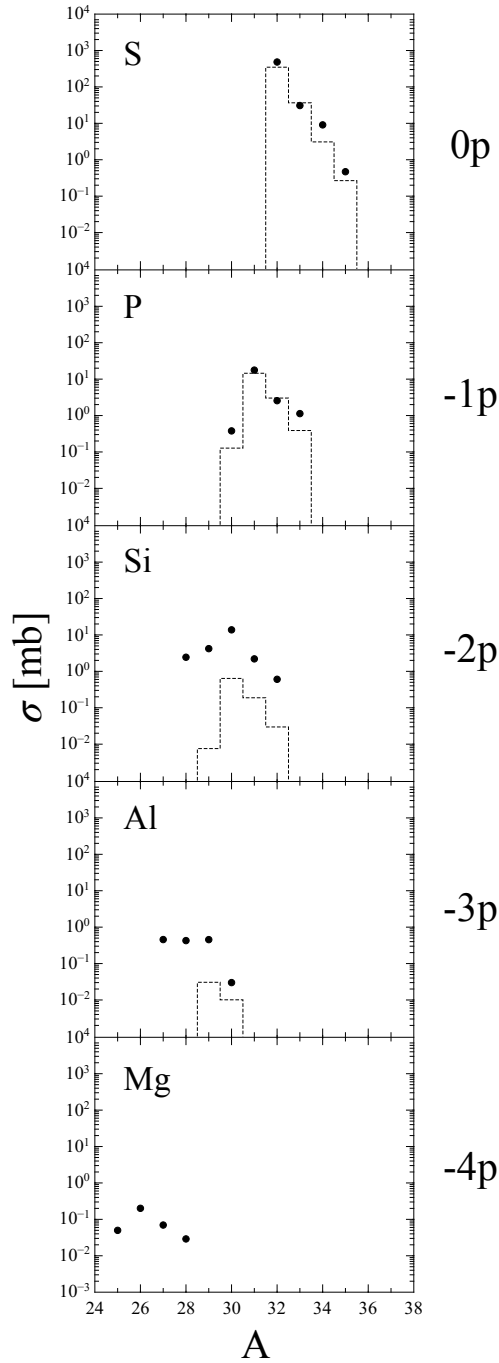


Figure 8.4: Inclusive cross sections for all the identified nuclei produced in the $^{32}\text{S}+^{110}\text{Cd}$ experiment, integrated over the angular acceptance of PRISMA (63° - 71°). Experimental data points are indicated with black symbols, while theoretical calculations performed with the GRAZING code [94] are displayed with dashed black lines. In the case of Mg, GRAZING cross sections are smaller than 10^{-3} mb.

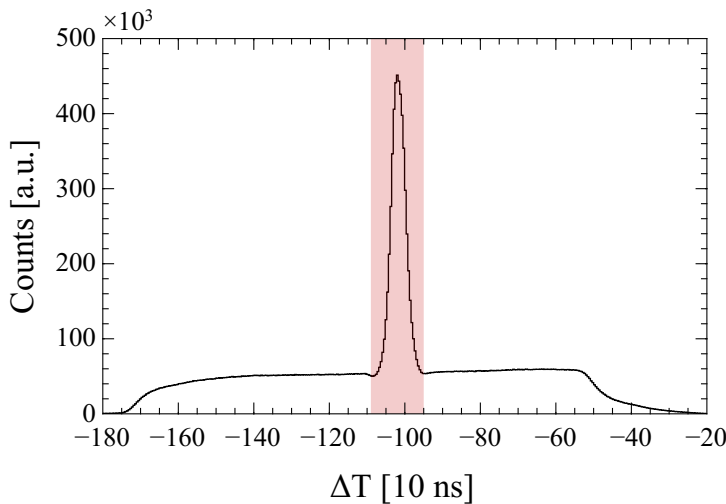


Figure 8.5: Time difference between AGATA and PRISMA timestamps in units of 10 ns. The red shaded area indicates the 140-ns gate used to select coincidence events.

the beam impinging on the target A , and b and B are the beam-like and target-like reaction products, respectively.

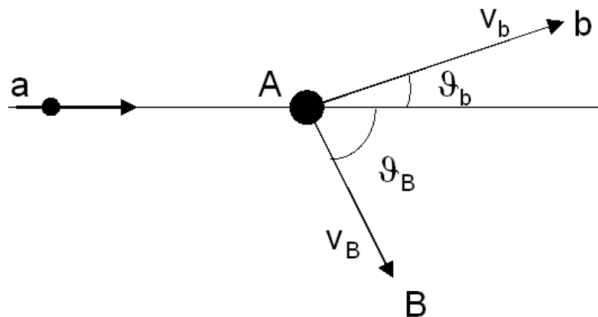


Figure 8.6: Schematic illustration of a two-body reaction collision.

By applying the momentum conservation law in the laboratory frame, projected on the beam axis x and the perpendicular axis y , one finds:

$$m_a v_a = m_b v_b \cos \theta_b + m_B v_B \cos \theta_B, \quad (8.11)$$

$$m_b v_b \sin \theta_b = m_B v_B \sin \theta_B. \quad (8.12)$$

The solution of this system of equations gives the relations

$$\theta_B = \arctan \left(\sqrt{2m_b K_b} \frac{\sin \theta_b}{\frac{\sqrt{2m_a K_a}}{\sqrt{2m_b K_b}} - \cos(\theta_b)} \right), \quad (8.13)$$

$$v_B = \sqrt{2m_b K_b} \frac{m_b \sin \theta_b}{m_B \sin \theta_B}, \quad (8.14)$$

where K_a and K_b are the kinetic energies of the beam-like and target-like nuclei, respectively.

This simple derivation provides the basic ingredients for an event-by-event Doppler-shift correction of the γ rays emitted by the target-like reaction product, even though they are not directly measured by PRISMA. Such a procedure was particularly relevant for the present work, specifically in the analysis of the plunger-target dataset. On the contrary, γ rays of the target-like nuclei of interest were mostly emitted at rest in the thick-target configuration, therefore, no Doppler correction was necessary.

As an example, Figure 8.7 shows the γ -ray spectrum from the plunger dataset in coincidence with ^{30}Si ions identified in PRISMA. Panel a) shows the γ -ray spectrum as detected by the AGATA array, without any correction, while the spectra obtained after the Doppler correction for ^{112}Sn and ^{30}Si are shown in Panel b) and c), respectively.

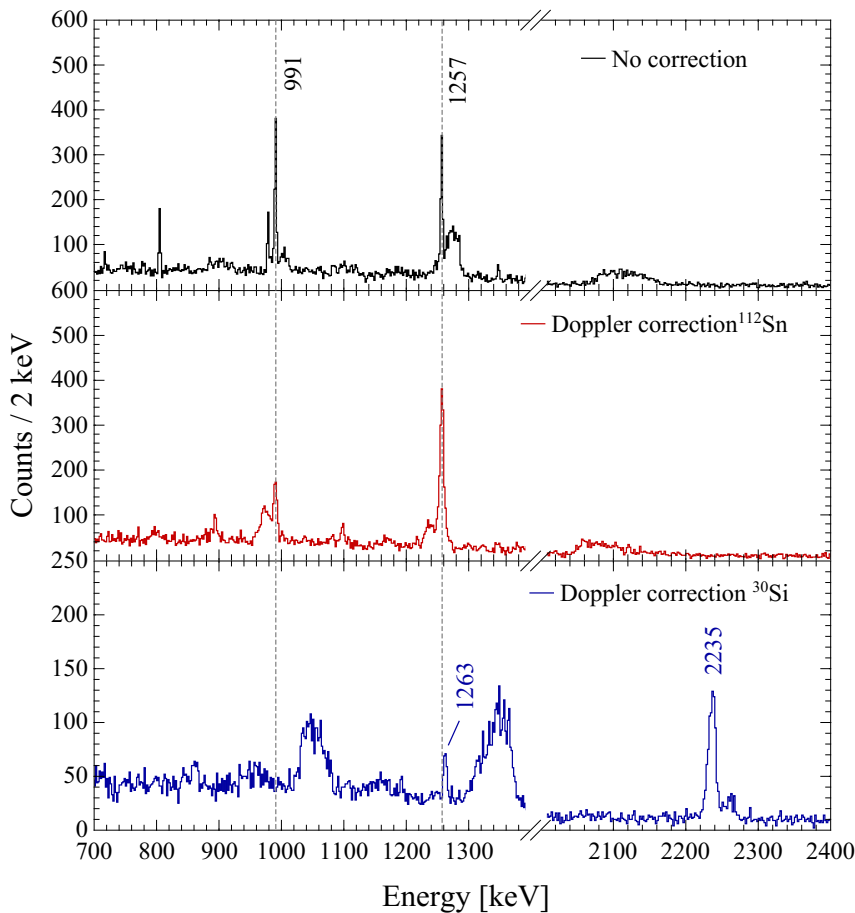


Figure 8.7: Comparison between γ ray spectra observed in AGATA in coincidence with ^{30}Si nuclei identified in PRISMA, with data collected in the plunger configuration. The spectra without any Doppler correction is shown in panel a), while panel b) and c) display spectra corrected for the ^{112}Sn target-like binary partner and for the ^{30}Si beam-like nucleus, respectively. Vertical dashed lines are used as reference for the 1257-keV $2_1^+ \rightarrow 0_1^+$ and the 991-keV $4_1^+ \rightarrow 2_1^+$ transitions of ^{112}Sn . In panel c), identified γ -ray peaks of ^{30}Si are labeled with their energies in blue.

The $2_1^+ \rightarrow 0_1^+$ transition at 1257 keV and the $4_1^+ \rightarrow 2_1^+$ transition at 991 keV of ^{112}Sn are clearly visible in both panel a) and b). The former highlights the stopped components of the two peaks, corresponding to γ rays emitted at rest in the plunger, whereas the latter brings evidence for the shifted component emitted in flight. In panel c), peaks belonging to the ^{30}Si nucleus (1263 and 2235 keV) emerge from the background. Conversely, ^{112}Sn peaks appear as broad structures at shifted energies.

Thanks to the detection capabilities of the AGATA-PRISMA setup, which measure the interaction position of ions and γ rays with a resolution of $\sim 0.2^\circ$ and $\sim 1^\circ$, respectively, a limited worsening in the energy resolution is observed after the Doppler correction. At 1.3 MeV, the FWHM of the beam-like corrected spectra is approximately 5 keV. Due to the uncertainty introduced by the energy losses in the kinematic reconstruction, a slightly worse resolution of about 7 keV is found in the spectra corrected for the target-like nuclei.

8.2.2 TKEL selection

As introduced in Section 8.1.1, the reconstruction of the kinematics yields the *TKEL* of the reaction, which expresses the total excitation energy of the reaction products, both the target-like and the beam-like.

Although the experimental resolution of the TKEL spectra is of several MeV, it can be used, to some degree, to control the population of low-lying and higher-lying states in the nuclei of interest. This is done by correlating the *TKEL* spectra obtained with PRISMA with the γ rays observed in coincidence with AGATA.

Figure 8.8 shows, as an example, *TKEL*-gated projections of the γ -ray spectra of ^{110}Cd . The γ -ray spectrum in panel a) is in coincidence with the peak at negative TKEL energies, namely that related to the ^{32}S scattering on Ta (see discussion in Section 8.1.1). This is evident from the γ detected in AGATA, as several transitions between low-lying states of Ta are observed. Conversely, the projection of panel b), resulting from a narrow gate is set on Cd peak around 0 MeV, is dominated by the 658-keV $2_1^+ \rightarrow 0_1^+$ transition of ^{110}Cd . Peaks corresponding to the de-excitation of the 2_1^+ and 4_1^+ states at 818 and 885 keV, respectively, are also clearly visible. Transitions from higher lying states are only visible when gating on the high-energy region of the *TKEL* spectrum, shown in panel c), such as the 937-keV and the 997-keV γ rays depopulating the 6_1^+ and the 5_1^- states, respectively. In addition, one can notice a contamination from ^{112}Cd and ^{114}Cd , whose $2_1^+ \rightarrow 0_1^+$ peaks are marked with triangular symbols in panels b) and c). Based on the measured counts, an isotopic purity of $\sim 97\%$ was estimated for the ^{110}Cd target.

In the case of proton transfer channels, the γ -ray spectra include not only transitions from the target-like binary partner of the reaction, but also those from nuclei produced after evaporation. With the *TKEL* selection, these contributions can be effectively separated, as evaporation is a dissipative phenomenon that occurs when a large amount of energy is exchanged between beam and target [95].

8.2.3 Level schemes of Cd and Sn isotopes

Using the particle- γ coincidence techniques discussed in the previous sections, it is possible to perform a detailed study of the γ -ray spectra and precisely identify transitions of target-like nuclei of interest. In particular, level schemes of even-even Cd and Sn isotopes were built according to the γ -rays observed in both the thick-target and the plunger-target datasets. While the former provides cleaner spectra, the latter proved to be well-suited for the observation of transitions from short-lived states ($\tau \lesssim 1$ ps), oth-

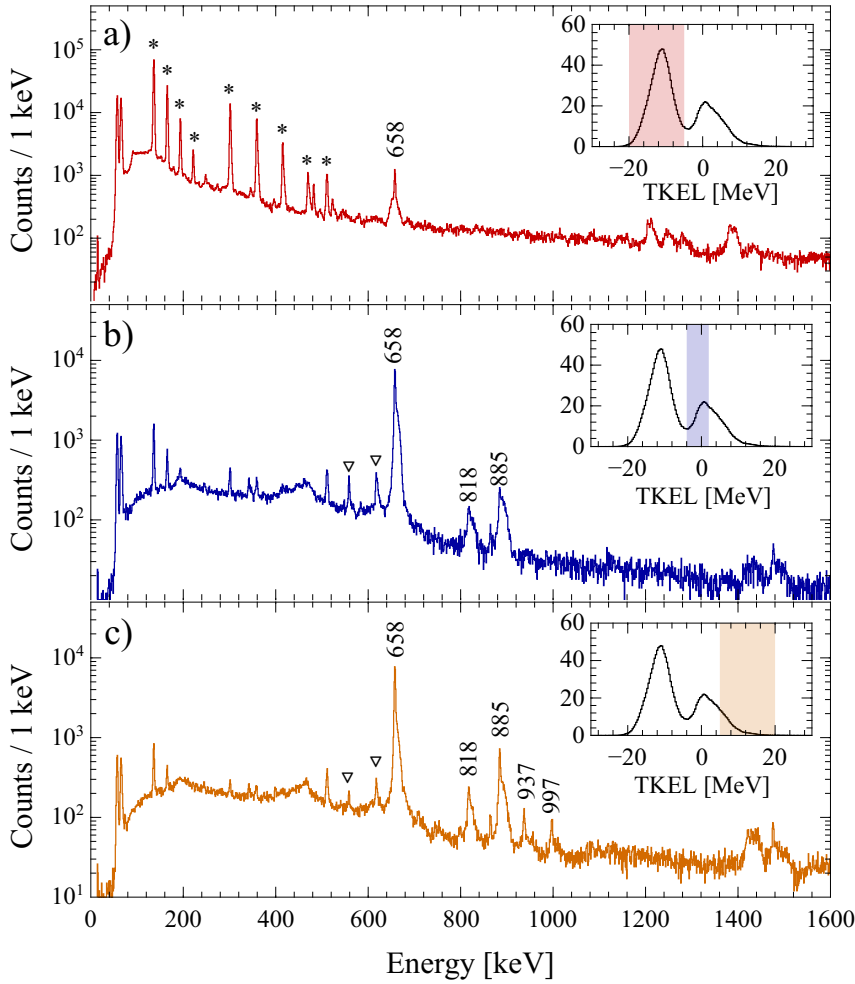


Figure 8.8: Partial γ -ray spectra observed in coincidence with ^{32}S nuclei identified in PRISMA, gated on the reconstructed $TKEL$ of the reaction. In the insets, the regions selected by the different gates on the $TKEL$ are shown (see text for details). Identified γ -ray transitions of ^{110}Cd are labeled with their energy in keV, while transitions from the $^{112,114}\text{Cd}$ contaminants and the Ta backing are indicated with triangles and asterisks, respectively. Due to the geometry of the setup in the thick-target, a right tail appears on ^{110}Cd peaks, corresponding to events with ions exiting the target and emitting γ rays in flight.

erwise hidden in the thick-target configuration. Although all the identified transitions were previously known, the present results supply valuable information on the population of excited states associated with multi-nucleon transfer reactions above Coulomb barrier energies.

Level scheme of ^{108}Cd

In the present experiment, ^{108}Cd nuclei were produced in the $2n$ pick-up channel. As illustrated in Figure 8.9, most of the yield is condensed in the 2^+ and 4^+ states, with the $4^+ \rightarrow 2_1^+$ transition being observed with about 30% of the $2_1^+ \rightarrow 0_1^+$ intensity.

The level scheme presented in Figure 8.10 was built on the basis of the observed γ rays from data acquired in either the thick- or plunger-target configurations. Only low-spin levels up $J = 6$ and excitation energies up to 2.8 MeV were populated. Although (p,t) reactions proved to be effective in the population of excited 0^+ states in similar Sn isotopes (see discussion in Chapter 5 and Refs. [70–73, 132, 182]), such excitations were not observed in the present experiment. This could be ascribed to the different nature of this heavy-ion transfer reaction and to the specific angular momentum coupling of the nucleons exchanged between ^{32}S and ^{110}Cd . Additionally, the energy acceptance of the PRISMA spectrometer required the reaction energy to be about 15% above the Coulomb barrier, as compared to the near-Coulomb-barrier energies of the heavy-ion transfer reactions described in Chapter 5).

In the Level scheme of Figure 8.10, level energies and spin assignments are the adopted ones [68]. Levels are gathered into bands according to the band structures proposed in Refs. [183, 184].

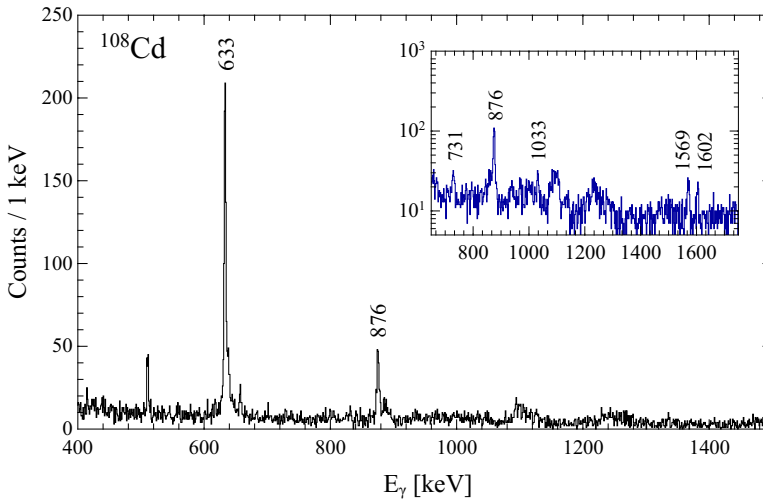


Figure 8.9: Partial γ -ray spectrum observed in AGATA in coincidence with ^{34}S nuclei identified in PRISMA, with data collected in the thick-target configuration. In the inset, the plunger spectrum with Doppler correction for ^{108}Cd nuclei is shown, providing evidence of transitions from short-lived states ($\tau \lesssim 1$ ps). Transitions from ^{108}Cd are labeled accordingly.

Level scheme of ^{110}Cd

Excited states in ^{110}Cd were populated through inelastic scattering on the ^{110}Cd target. A partial γ -ray spectrum obtained in coincidence with identified ^{32}S ions in PRISMA is illustrated in Figure 8.11, where the most intense transitions of ^{110}Cd are labeled accordingly. The reconstructed level scheme, shown in Figure 8.12, is characterized by low-lying and low-spin excitations, with $2 < J < 6$ and $E_x < 2.5$ MeV. Energies and

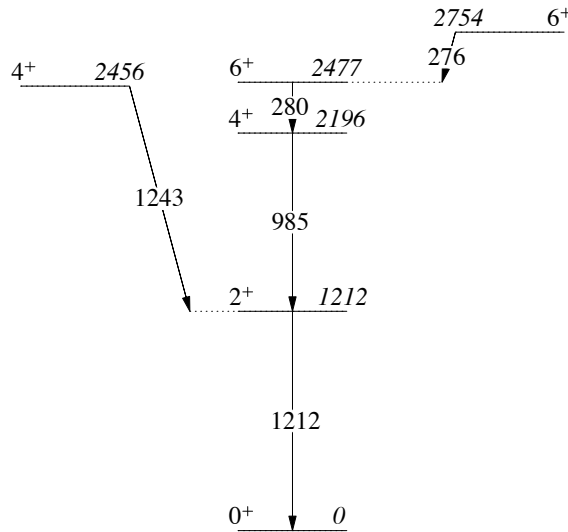


Figure 8.14: Level scheme of ^{110}Sn as observed in the present multi-nucleon transfer experiment.

^{111}Sn in the spectrum. In fact, neutron evaporation phenomena become relevant with increasing numbers of transferred nucleons, leading to the population of different nuclei with respect to the binary partner of the reaction.

By applying a gate on the low-side of the $TKEL$ spectra, a small peak at 934 keV appears in the projected spectrum, as shown in Figure 8.16. This corresponds to the γ -ray transition connecting the first excited 0_2^+ state with the 2_1^+ state. Despite this additional selection, no other transition from excited 0^+ state was observed. Such a result is in stark contrast with the previous $2p$ transfer experiment performed with $(^3\text{He},n)$ reactions, where an enhanced cross section towards the 0_2^+ of ^{112}Sn was observed [69]. Possible explanations of this behavior could be found in the different nature of the heavy-ion transfer reaction mechanism with respect to the $(^3\text{He},n)$ process, and in the available excitation energy window of 6.6 MeV (see Table 6.2), which could be too large for the direct population of 0^+ states. Furthermore, based on the reconstructed level scheme shown in Figure 8.17, it is evident that the reaction preferentially populated yrast states up to the $(9)^-$ level at 3693 keV. On the contrary, low-lying, off-yrast states are indirectly populated according to the decay patterns of the γ cascades, which, however, tend to avoid the 0^+ states of interest.

As in the previous cases, the energies and spin-parities in the levels scheme are the adopted ones [68], as no discrepancies were found between those and the present experimental data.

Level scheme of ^{114}Sn

In this work, ^{114}Sn nuclei were populated in both the experimental studies performed at IFIN-HH (see Chapter 5) and the one performed at LNL. Nevertheless, significant differences were observed in the populated levels and the corresponding γ -ray transitions.

Figure 8.18 shows the thick-target spectrum of γ rays observed by AGATA, gated on ^{28}Si ions measured by PRISMA. As in the ^{112}Sn case, yrast states of ^{114}Sn up to the 9^- at 3511 keV were found to be strongly populated in this reaction. Additionally, transitions

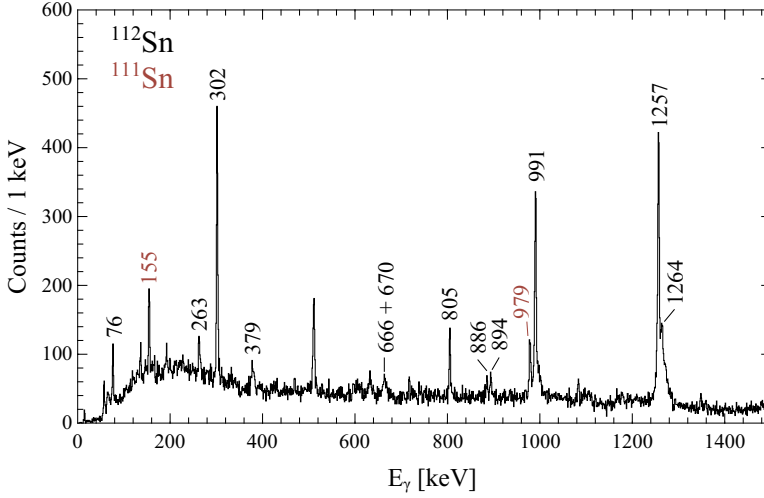


Figure 8.15: Partial γ -ray spectrum observed in AGATA in coincidence with ^{30}Si nuclei identified in PRISMA, with data collected in the thick-target configuration. Transitions from ^{112}Sn and ^{111}Sn are labeled in black and red, respectively.

from ^{113}Sn , such as the 661-keV γ decay from the isomeric $11/2^+$ state, give evidence of one-neutron evaporation processes.

The reconstructed level scheme for ^{114}Sn as observed in the $^{32}\text{S} + ^{110}\text{Cd}$ multi-nucleon transfer reaction can be found in Figure 8.19. While in the IFIN-HH experiment almost all the states up to 3.5 MeV were populated, as shown in Figure 5.7 of Section 5.1, here a reduced number of levels is observed. Furthermore, no transitions of de-exciting 0^+ states were observed, similarly to the other reaction channels.

8.3 Feasibility of lifetime measurements

In the present analysis, the possibility of performing lifetime measurements with the available datasets was investigated. The analysis is ongoing, however, its complex nature and a series of additional reasons, listed below, prevented its completion within the time frame of this PhD project.

First, the limited population of the states of interest and lower beam current of 3 nA that was available for the plunger dataset led to longer acquisition time in the same configuration. Because of this, data corresponding to only a single plunger distance have sufficient statistics to be analyzed in detail. In this condition, one could estimate lifetimes with large error bars or upper limits, if the distance calibration of the plunger is found to be reliable. ^{110}Cd targets are notoriously difficult to stretch due to their poor mechanical properties, therefore, benchmarks with states of known lifetimes have to be used to check the quality of the distance calibration. Considering the nominal plunger distance of $50\text{ }\mu\text{m}$ and target-like velocities in the $7\text{--}8\text{ }\mu\text{m}/\text{ps}$ range, this procedure requires states with lifetimes of a few picoseconds and no feeding from above. Unfortunately, these are stringent constraints for the present experiment, given the population pattern and the short-lived nature ($\tau \lesssim 1\text{ ps}$) of higher lying states.

A possible solution could be provided by the use of realistic simulations with the AGATA simulation framework [122], which are able to reconstruct decay patterns and

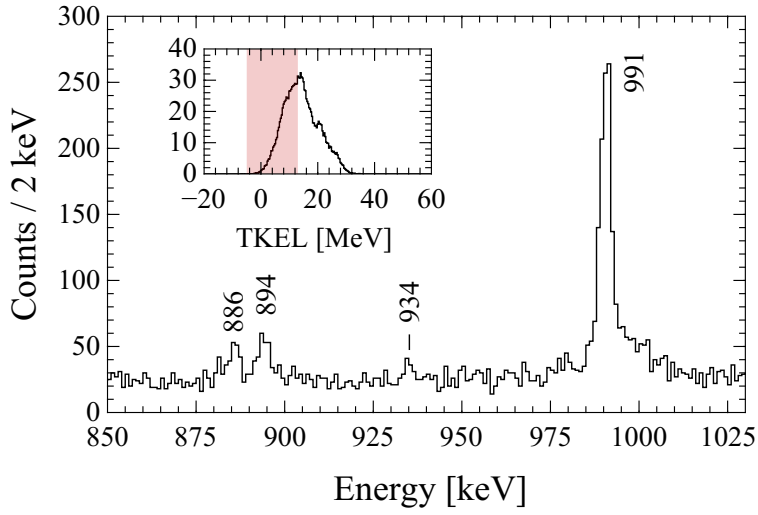


Figure 8.16: Partial γ -ray spectrum gated on the reconstructed $TKEL$, in coincidence with ^{30}Si nuclei identified in PRISMA. The gate set on the $TKEL$ spectrum is displayed with a red shaded area on the spectrum in the inset. Transitions from ^{112}Sn are labeled accordingly. Data collected in the thick-target configuration.

effectively account for feeding from higher lying states. This remains a potential follow-up of the present research.

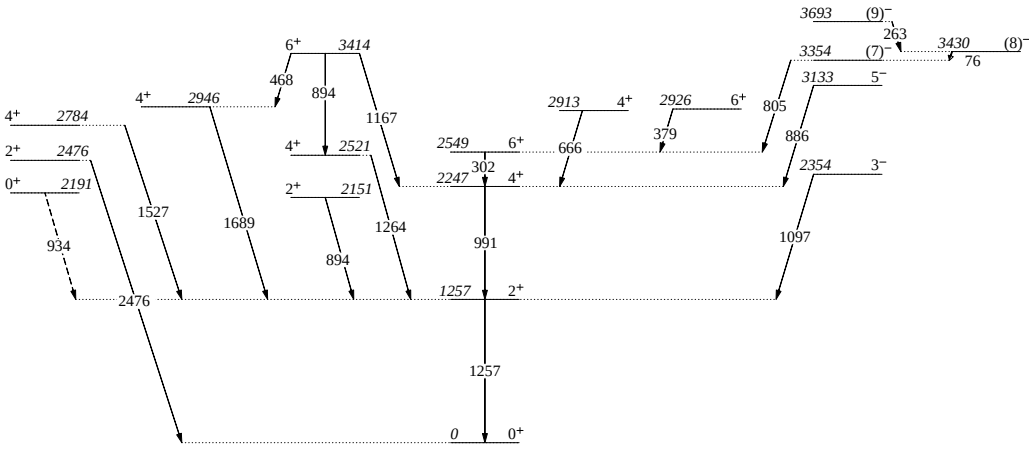


Figure 8.17: Level scheme of ^{112}Sn as observed in the present multi-nucleon transfer experiment.

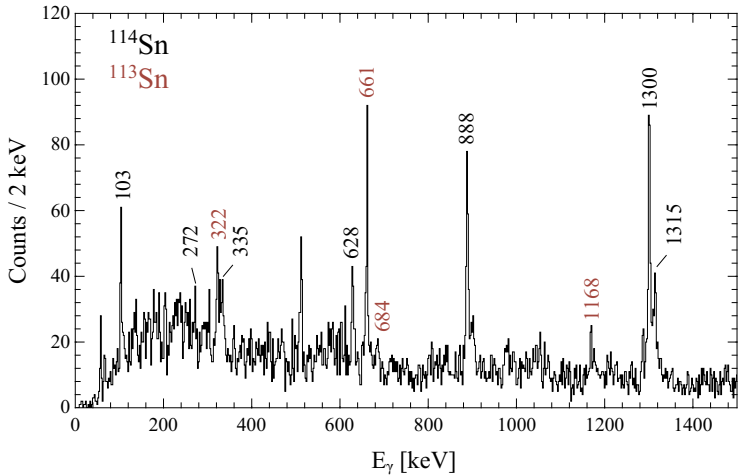


Figure 8.18: Partial γ -ray spectrum observed in AGATA in coincidence with ^{28}Si nuclei identified in PRISMA, with data collected in the thick-target configuration. Transitions from ^{114}Sn and ^{113}Sn are labeled in black and red, respectively.

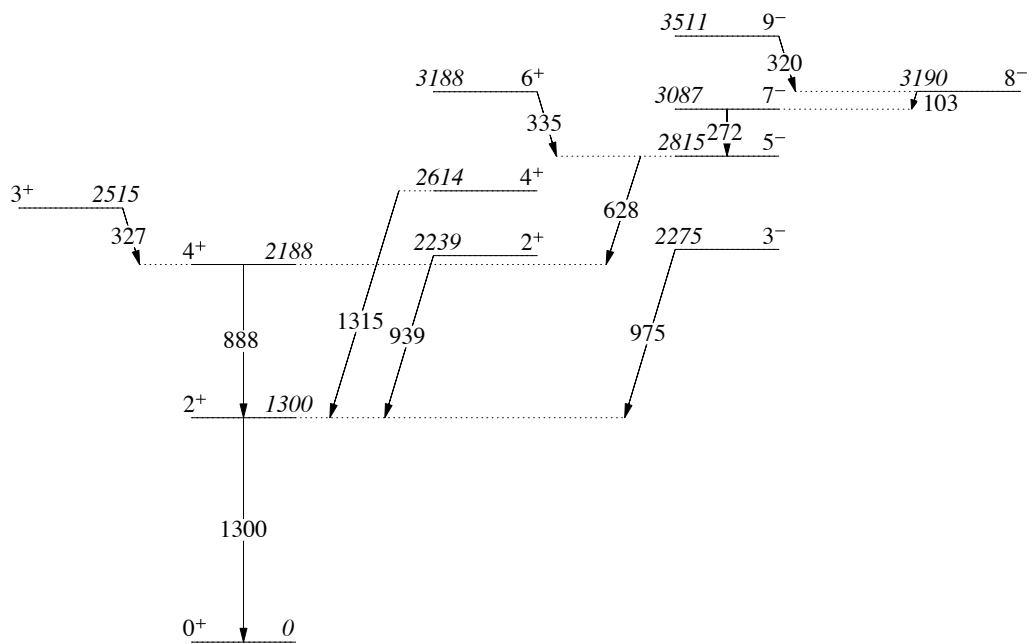


Figure 8.19: Level scheme of ^{114}Sn as observed in the present multi-nucleon transfer experiment.

Part IV

Discussion and Conclusion

Results and discussion

In this work, the $B(E2)$ transition probabilities from low-lying excited 0^+ states towards the 2_1^+ states have been deduced from lifetime measurements in even-even Sn isotopes at the valley of stability. The results obtained for $^{114-120}\text{Sn}$, together with the available literature data, are summarized in Table 9.1 and illustrated in Figure 9.1.

In the following discussion, the experimental results are compared with theoretical predictions from IBM [80], mean-field HFB [85, 186], and MCSM [27, 83] calculations. As outlined in Section 1.3, both mean-field and MCSM approaches predict the presence of a prolate-deformed secondary minimum in the potential energy surface of even-even $^{110-120}\text{Sn}$ isotopes, separated by a sizable barrier from the main spherical-oblate minimum. This opens the possibility of deformed 0^+ states residing in the prolate minimum, characterized by hindered transitions towards states of spherical-oblate nature. In the Sn mass region, the most advanced MCSM calculations are still preliminary, while they proved to be effective in describing shape coexistence phenomena, for instance, in the lighter mass region of Ni [32, 33] and Zr [187, 188] isotopes. In particular, for Sn nuclei, potential energy surfaces and wave-function components have been calculated only up to the 0_3^+ states, as shown in Figure 9.2, and $B(E2)$ transition probabilities are not yet available.

A direct comparison of our experimental results with literature data and up-to-date theoretical calculations is presented for each of the studied $0_{2,3,4}^+$ states of $^{114-120}\text{Sn}$ in the sections below.

9.1 Results for $^{110,112}\text{Sn}$

For the case of $^{110,112}\text{Sn}$, our measurements did not provide additional information on transition probabilities, as the 0^+ states of interest were not effectively populated by the employed multi-nucleon transfer reaction. From the available data in the literature, the 0_2^+ appears as the band-head of the intruder configuration, based on energy systematics and two-proton transfer cross sections. In addition to the 0_2^+ state, the 2^+ and 4^+ members of this positive-parity intruder band were also recently identified in ^{112}Sn by Spieker et al. [80] in terms of their inter-band transitions, which emerged as the most collective $E2$ transitions in the nucleus. However, the information concerning the lifetime of the excited 0^+ states in $^{110,112}\text{Sn}$ is lacking, with only a single upper limit reported for the lifetime of the 0_2^+ of ^{112}Sn . This points to a reduced transition probability of $B(E2; 0_2^+ \rightarrow 2_1^+) \leq 9.2 \text{ W.u.}$, which is slightly smaller than the observed 10-20 W.u. values for $A = 114 - 120$ tin isotopes. Yet, the existing data are not sufficient to draw firm conclusions about the presence of coexisting deformed configurations in $^{110,112}\text{Sn}$, calling for further experimental investigation.

Nucleus	E_i [keV]	I_i^π	$T_{1/2}^{exp}$ [ps]	E_f [keV]	I_f^π	E_γ [keV]	BR [%]	$E/M\lambda$	$B(E2)^{exp}$ [W.u.]	$B(E2)^{lit}$ [W.u.]
^{114}Sn	1953	0_2^+	$11.0^{+9.3}_{-4.0}$	1300	2_1^+	653	99.7	[E2]	13^{+8}_{-6}	$22(8)^a$
	2156	0_3^+	$13.0^{+9.4}_{-4.2}$	1300	2_1^+	856	100	[E2]	$2.9^{+1.4}_{-1.2}$	$< 4.9^a$
^{116}Sn	1757	0_2^+	$51.1^{+18.8}_{-11.7}$	1294	2_1^+	463	99.7	[E2]	15^{+5}_{-4}	$18(3)^a, 18(4)^b, 15(4)^c$
	2027	0_3^+	184^{+66}_{-40}	1294	2_1^+	734	100	[E2]	$0.43^{+0.12}_{-0.11}$	$0.49(6)^b, 0.48(3)^c$
	2546	0_4^+	$1.17^{+0.10}_{-0.09}$	2112	2_2^+	434	2.9	[E2]	27^{+2}_{-2}	
				1294	2_1^+	1252	97.1	[E2]	$4.5^{+0.4}_{-0.4}$	
^{118}Sn	1758	0_2^+	$18.2^{+2.4}_{-2.0}$	1230	2_1^+	529	100	E2	22^{+3}_{-3}	$19(3)^a$
	2057	0_3^+	$92.2^{+19.9}_{-14.7}$	1230	2_1^+	827	97.6	E2	$0.45^{+0.09}_{-0.08}$	$> 0.2^a, < 59^d$
	2497	0_4^+	$14.2^{+1.5}_{-1.3}$	1230	2_1^+	1267	100	E2	$0.35^{+0.04}_{-0.03}$	
^{120}Sn	1875	0_2^+	$8.2^{+2.6}_{-1.7}$	1171	2_1^+	704	100	[E2]	$11.3^{+3.0}_{-2.7}$	$12.6(17)^a$
	2160	0_3^+	$40.4^{+7.9}_{-6.2}$	1171	2_1^+	989	100	[E2]	$0.42^{+0.08}_{-0.07}$	$0.50(7)^e$
	2587	0_4^+	$14.8^{+1.1}_{-1.0}$	1171	2_1^+	1416	100	[E2]	$0.19^{+0.01}_{-0.01}$	$< 8.3^f$

^a From Coulomb excitation [87].^b From fast-timing [145].^c From fast-timing [78].^d From DSAM [146].^e From fast-timing [81].^f From DSAM [147].

Table 9.1: Summary of the spectroscopic information on excited 0^+ states in $^{114-120}\text{Sn}$, as obtained in the present work. E_i , E_f and I_i , I_f are the initial and final excitation energies and spin-parities, respectively. $T_{1/2}^{exp}$ is the half-life measured in the present work. E_γ is the energy of a γ -ray transition, BR is the corresponding branching ratio with respect to the all the possible decay channels, including internal conversion intensities when available [14, 68, 135], and $E/M\lambda$ is the multipolarity. $B(E2)^{exp}$ and $B(E2)^{lit}$ are the transition probabilities from the present work and from the data available in the literature, respectively.

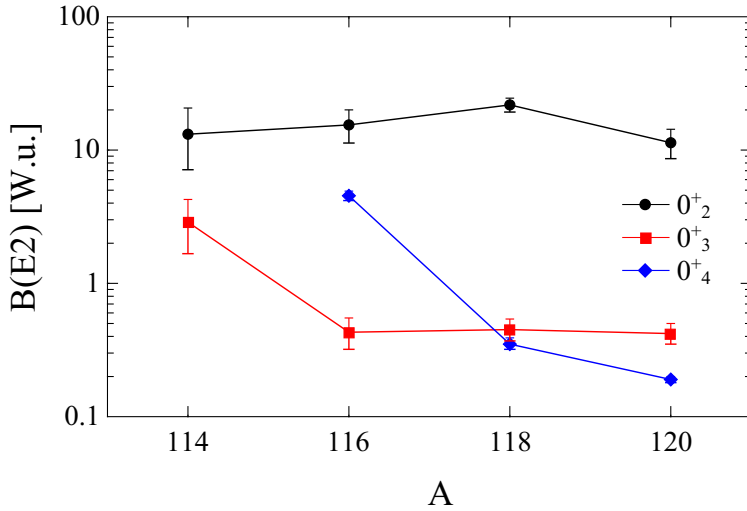


Figure 9.1: $B(E2; 0_{2,3,4}^+ \rightarrow 2_1^+)$ transition probabilities of $^{114-120}\text{Sn}$ as obtained in the present work.

9.2 Results for ^{114}Sn

From our measurements, we determined $B(E2)$ transition probabilities of 13_{-6}^{+8} and $2.9_{-1.2}^{+1.4}$ W.u. for the γ decay of the 0_2^+ and 0_3^+ states towards the 2_1^+ state in ^{114}Sn , respectively. These results, which are in agreement with the previous values of 22(8) and < 4.9 W.u. by Bäcklin et al. from Coulomb Excitation [87], are shown in Figure 9.1 in comparison with the ones obtained for heavier tin isotopes. The large $B(E2; 0_2^+ \rightarrow 2_1^+)$ indicates a significant overlap between the 0_2^+ and the 2_1^+ wavefunctions. While the latter and the ground state are traditionally considered of spherical nature, experimental observations of Allmond et al. [189] and calculations of MCSM [83] point to modest but different deformations. The electric quadrupole moment $Q(2_1^+)$ measured by Allmond et al. is positive and decreases from ^{112}Sn to ^{118}Sn , becoming negative for $^{122,124}\text{Sn}$, thus indicating prolate to oblate evolution along the isotopic chain [189]. Conversely, the MCSM points to 0_1^+ and 2_1^+ states with a spherical-oblate character, with their main components residing in the primary minima of the potential energy surfaces [14, 83]. This is shown in Figure 9.2, where the main wavefunction components of the ground states are superimposed to the potential energy surfaces of $^{110-120}\text{Sn}$. Based on the observed rotational bands [75] and reaction data [69], the 0_2^+ was originally interpreted as the band-head of the intruder configuration. Instead, the results of Spieker et al. [80] opened the possibility of the 0_3^+ as the intruder band-head, although a large level of mixing between the 0_2^+ and 0_3^+ states is expected according to the IBM calculations they performed. Such calculations predict $B(E2)$ values which are in agreement with our measurements, especially concerning the $B(E2; 0_3^+ \rightarrow 2_1^+)$, which they estimate to be 2 W.u. compared to the $2.9_{-2.1}^{+0.9}$ W.u. we obtained. In the IBM picture, the 0_2^+ and the 0_3^+ states are associated to the normal and the intruder configurations, respectively. Similarly, MCSM [14, 83] predicts similar spherical-oblate components for the lower-lying 0^+ states up to the 0_3^+ , hinting at a certain level of mixing between these states. It is also important to note that the transition probability of $B(E2; 0_3^+ \rightarrow 2_1^+)$ is almost an order of magnitude higher than values observed in $^{116-120}\text{Sn}$ isotopes, as shown in Figure 9.1,

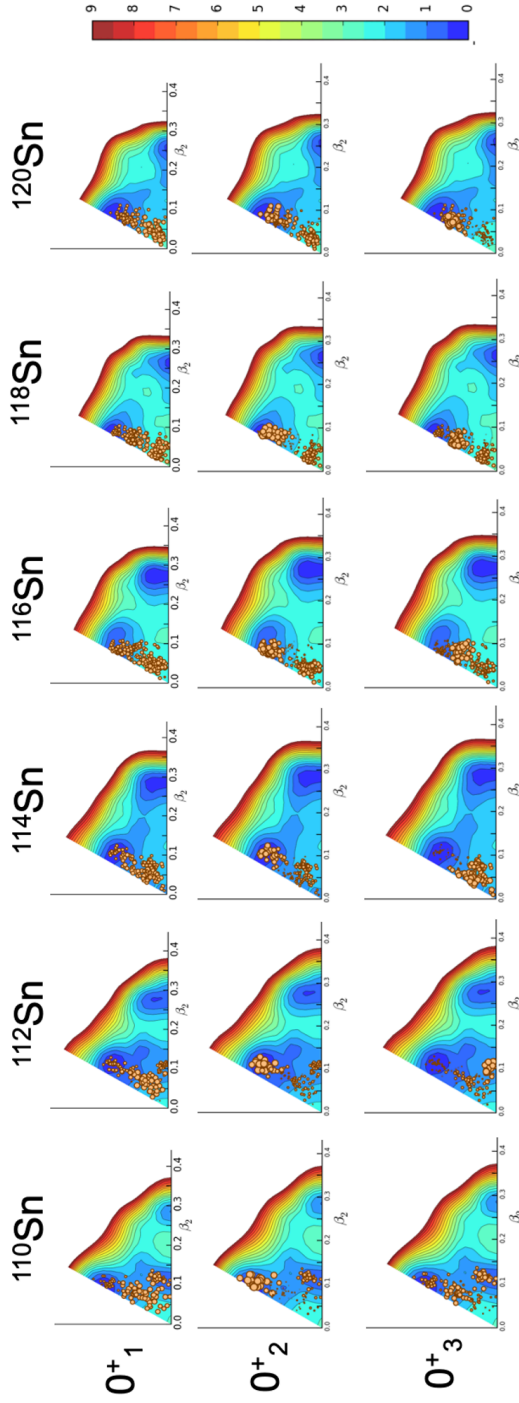


Figure 9.2: MCSM calculations of potential energy surfaces for $A = 110 - 120$ Sn, obtained with the constrained Hartree-Fock-Bogoliubov method, using the A3DA-m interaction of Togashi et al. [83]. The main components of the wavefunctions of the 0_1^+ , 0_2^+ and 0_3^+ states are shown in the different rows.

whereas the $B(E2; 0_2^+ \rightarrow 2_1^+)$ falls within 10-20 W.u. for the isotopes in this region. This drastic change could be evidence of a change in the microscopic structure of the 0_3^+ state along the isotopic chain.

9.3 Results for ^{116}Sn

In ^{116}Sn , our results agree with the previous values of the $B(E2)$ of the 0_2^+ and 0_3^+ obtained by Bäcklin et al. with Coulomb excitation [87], Julin et al. in a fast-timing (p,p') study [145] and Petrache et al. in a fast-timing (n, γ) investigation [78]. From the observed values of several tens of W.u. for the $B(E2; 2_2^+ \rightarrow 0_{2,3}^+)$ and the $B(E2; 4_{1,2}^+ \rightarrow 2_2^+)$, Petrache et al. concluded that a strong mixing occurs between the normal and intruder configurations, in agreement with IBM calculations. Similar conclusions arise from the strong overlap in the main wavefunction components in the $0_{1,2,3}^+$ states calculated by MCSM, as shown in Figure 9.2. A recent study of Jalili et al. [186] that combined the Bohr collective model, DFT calculations, and the Two-State Mixing Model (TSMM), also supports this strong mixing picture. According to their analysis, mixing between β and γ bands in ^{116}Sn should lead to the coexistence of oblate and prolate deformations.

Additional insights in ^{116}Sn are provided by $E0$ transitions. The large $10^3 \times \rho^2(E0; 0_3^+ \rightarrow 0_2^+) = 100(20)$ value observed by Kantele et al. [82] provides further indication, in a two-state-mixing assumption, of either strong admixture between the states or a large difference in their deformation parameter β . Based on the mixing amplitudes $\alpha^2 = 0.69$ and $\beta^2 = 0.31$ calculated by Pore et al. [79], a difference $\Delta\beta \geq 0.22$ between the deformations of the two states was estimated. From the analysis of this β -decay experiment [79], the 0_3^+ was suggested as the bandhead of the intruder configuration, as a consequence of the observation of an enhanced $B(E2; 2_2^+ \rightarrow 0_3^+)$. Although this assumption is in agreement with the IBM predictions of Ref. [78], the strong interaction between normal and intruder configurations prevents a firm assignment.

The $B(E2; 0_4^+ \rightarrow 2_1^+) = 4.5_{-0.4}^{+0.4}$ W.u. and $B(E2; 0_4^+ \rightarrow 2_2^+) = 27_{-2}^{+2}$ W.u., determined for the first time in the present work, support the interpretation of a strong overlap between the normal and intruder configurations. In addition, the 0_4^+ in ^{116}Sn is associated to a normal configuration in the IBM calculations, which is consistent with our results. In this case, microscopic MCSM calculations are not available yet.

9.4 Results for ^{118}Sn

In ^{118}Sn , the 0_2^+ band head decays toward the near-spherical 2_1^+ state with a rather fast transition. The $B(E2; 0_2^+ \rightarrow 2_1^+)$ of 22_{-3}^{+3} W.u. measured in this work is in good agreement with the previous measurement of Bäcklin et al., and it is the largest observed among the 0_2^+ states in $^{114-120}\text{Sn}$.

The reduced transition probability of the $E2$ transition depopulating the 0_3^+ was previously unknown. Our result of $0.45_{-0.08}^{+0.09}$ W.u. perfectly fits within the $0.2 - 59$ W.u. adopted range [68, 87, 146], and it is comparable to the values obtained in the neighboring ^{116}Sn and ^{120}Sn even-even isotopes (see Figure 9.1). The newly measured half-life allows for the determination of the $E0$ transition strength as well. Based on the indirect measurement of a 2.5(2)% branching ratio for the $E0$ $0_3^+ \rightarrow 0_2^+$ transition relative to the $E2$ $0_3^+ \rightarrow 2_1^+$ transition, determined in the (n, γ) experiment of Ref. [135], we obtain $10^3 \times \rho^2(E0) = 82_{-15}^{+16}$, which is comparable with the measured values of 86(16) and 120(50) milliunits of ^{116}Sn and ^{120}Sn [14, 81, 104], respectively. Therefore, a significant degree of mixing could be expected between the 2057-keV, 0_3^+ state and the 2p-2h,

0_2^+ state. The similar character of these two states was also suggested by previous $E0$ measurements [190].

A significant reduction of about one order of magnitude is observed in the transition probabilities of the $0_4^+ \rightarrow 2_1^+$ γ decay moving from ^{116}Sn to ^{118}Sn , which could hint at a change in the structure of this state along the isotopic chain. The $B(E2)$ value of $0.35_{-0.03}^{+0.04}$ W.u. suggests a smaller admixture with the 2_1^+ state, which is thought to reside in the main, near-spherical minimum of the PES. According to the classification of shape coexistence cases proposed by Leoni et al. [14], the hindrance factor of 14, derived from the $B(E2; 0_4^+ \rightarrow 2_1^+)$, suggests the presence of a sizable barrier in the potential energy surfaces between the two states, which may reside in minima corresponding to different deformations. Specifically, the main components of the wavefunction of the 0_4^+ state could be hosted by the prolate-deformed secondary minimum of the potential energy surface, which is well defined in ^{118}Sn (see Figures 1.9 and 9.2). In recent years, a hindrance factor of about 24 observed in ^{66}Ni for the $0_4^+ \rightarrow 2_1^+$ led to the identification of the 0_4^+ state, in the MCSM framework, as a shape-isomer-like structure with an axial prolate deformation [32]. A similar scenario could occur in ^{118}Sn , however, a comparison with updated microscopic calculations is needed to firmly assess the nature of this state.

9.5 Results for ^{120}Sn

The ^{120}Sn case appears well aligned to that of ^{118}Sn , as shown in Figure 9.1. The $B(E2; 0_2^+ \rightarrow 2_1^+)$ observed in this work and in that of Bäcklin are just above 10 W.u., as in the other even-even Sn isotopes of this region. In the same way, our measured $B(E2; 0_3^+ \rightarrow 2_1^+)$ value of $0.42_{-0.07}^{+0.08}$ W.u. is in good agreement with the value of 0.50(7) W.u. determined in the (n, γ) experiment of Wu et al. [81]. In their work, Wu et al. also argue that, assuming a complete admixture of spherical and deformed unperturbed states, a minimum $\Delta\beta = 0.23(2)$ is required to reproduce the $10^3 \times \rho^2(E0; 0_3^+ \rightarrow 0_2^+)$ value of 120(50) observed in their work.

As in the ^{118}Sn , the 0_4^+ of ^{120}Sn decays with a hindered transition towards the 2_1^+ state. Our measured $B(E2) = 0.19_{-0.01}^{+0.01}$ W.u. corresponds to a hindrance factor $HF = 26$, which is the highest value observed so far among the transitions from excited 0^+ states in Sn isotopes. Such a result suggests that the 0_4^+ state could reside in the prolate-deformed secondary minimum of the PES, shown in Figure 9.2, with a moderately high barrier separating it from the main ground-state minimum.

9.6 Systematics and overview

The systematic trends of the $B(E2; 0_{2,3,4}^+ \rightarrow 2_1^+)$ — in particular the dramatic drops in the $B(E2; 0_3^+ \rightarrow 2_1^+)$ and $B(E2; 0_4^+ \rightarrow 2_1^+)$ at $A = 116$ and $A = 118$, respectively (see Figure 9.1) — could be linked to the evolution of the neutron single-particle energies. In this context, the low-lying single-particle states of odd Sn isotopes reflect, in good approximation, the main components of the corresponding single-particle orbitals [80]. As shown in Figure 9.3, the energies of the $5/2^+$ and $7/2^+$ states increase rapidly from ^{113}Sn to ^{115}Sn , whereas the $3/2^+$ and $11/2^-$ states decrease to low excitation energy starting from ^{115}Sn . This evolution in the neutron single-particle structure likely alters the wavefunction composition of the excited 0^+ states, accounting for the observed drops in the $B(E2; 0_{3,4}^+ \rightarrow 2_1^+)$. Furthermore, the $h_{11/2}$ orbital is particularly relevant for the Type II shell evolution mechanism, as described in Section 1.2.1. The lowering of the neutron

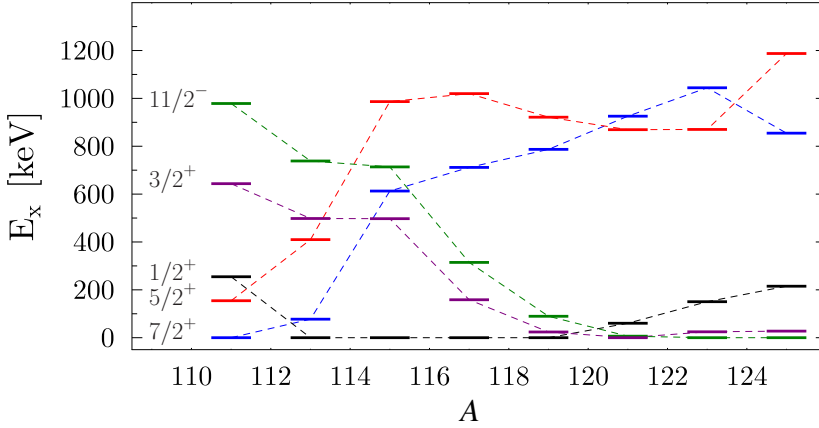


Figure 9.3: Systematics of the low-lying single-particle states in the odd Sn isotopes. The data have been compiled from Ref. [68]. The experimentally observed states might in good approximation reflect the corresponding single-particle levels, i.e., $1g_{7/2}$ (blue), $2d_{5/2}$ (red), $3s_{1/2}$ (black), $2d_{3/2}$ (purple), and $1h_{11/2}$ (green). Reprinted from Ref. [80].

$h_{11/2}$ orbital could facilitate its occupation and, in a MCSM picture, trigger the action of the tensor force. This would result in a spontaneous symmetry breaking, leading to the appearance of prolate deformations [14, 26]. Because of their hindered decay towards the 2_1^+ state, the 0_4^+ states in ^{118}Sn and ^{120}Sn might be good study cases to verify this hypothesis through updated MCSM calculations.

To summarize, the present results support previous measurements and provide additional information in terms of $E2$ and $\rho^2(E0)$ transition probabilities, completing the picture between ^{114}Sn and ^{120}Sn . The complex scenario of coexisting structures call for state-of-the-art microscopic calculations to interpret the rich experimental data that are now available.

Conclusions and perspectives

The present thesis is dedicated to a systematic investigation of 0^+ excitations (up to 0_4^+) in the context of shape coexistence studies in Sn isotopes with mass number $A = 110 - 120$. For this purpose, spectroscopy and lifetime measurements of low-lying excited states were performed in dedicated experimental campaigns. At IFIN-HH, four experiments (denoted as *Exp. 1-4*) were conducted between 2021 and 2024 with the ROSPHERE-SORCERER setup, while at LNL, an AGATA-PRISMA experiment was carried out in 2022. In this thesis, the full analysis of the collected data is presented.

In the experiment at IFIN-HH, transfer reactions induced by ^{13}C and $^{16,18}\text{O}$ beams on Sn and Cd targets were employed at energies close to the Coulomb barrier ($\sim 3 \text{ MeV/A}$). Using particle- γ , γ - γ and particle- γ - γ coincidence techniques, level schemes were established and the population of excited 0^+ states observed in the thick-target datasets of *Exp. 1* and *3* was investigated, with the aim of optimizing intensity yields for subsequent plunger measurements. In addition, the thick-target dataset of *Exp. 3* was used to perform lifetime measurements in the picosecond range with the Doppler Shift Attenuation Method (DSAM). For this purpose, a new code implementation for ROSPHERE simulations was developed as an extension of the AGATA simulation framework and validated by successfully reproducing the known lifetimes of transitions in ^{120}Sn . This led to the determination of the lifetime of the 0_4^+ state in ^{116}Sn .

Since one- and two-neutron transfer processes proved effective in populating the states of interest in $^{114-120}\text{Sn}$, they were used in the plunger measurements of *Exp. 2* and *4*. The analysis with the Decay Curve Method provided lifetimes for the $0_{2,3}^+$ states in $^{114,116}\text{Sn}$ and the $0_{2,3,4}^+$ states in $^{118,120}\text{Sn}$.

At LNL, the $^{32}\text{S} + ^{110}\text{Cd}$ multi-nucleon transfer reaction was studied at a beam energy about 15% above the Coulomb barrier. Owing to the identification of beam-like reaction products by the PRISMA spectrometer, inclusive angular distributions and total cross sections were extracted and compared with GRAZING semiclassical calculations. Particle- γ coincidences, combined with *TKEL* gating, made it possible to perform a spectroscopic study of the reaction channels. The resulting γ -ray spectra showed strong population of yrast states in the $(2p, xn)$ transfer channels of interest, whereas the population of excited 0^+ states was hindered. This suggests that the angular-momentum coupling of transferred nucleons between the beam and target, together with the chosen reaction energy, could have been suboptimal for maximizing 0^+ state population. Further analysis of this dataset, in particular with realistic Monte Carlo simulations, could provide additional information on excited-state lifetimes, especially in neighboring systems like ^{108}Cd .

From the measured lifetimes, $B(E2)$ and, in one case, $\rho(E0)$ transition probabilities from $0_{2,3,4}^+$ states in even-even $^{114-120}\text{Sn}$ isotopes were derived. The present results provide a significant extension of the data available in literature. The $B(E2)$ values for transitions depopulating the 0_2^+ state of $^{114-120}\text{Sn}$ and the 0_3^+ state of $^{116,120}\text{Sn}$ were previously known, and our results are consistent with those. Conversely, the transition probabilities from the 0_3^+ of $^{114,118}\text{Sn}$ and the 0_4^+ of $^{116-120}\text{Sn}$ were determined here for the first time. Clear systematic trends have been observed in the $B(E2)$ values, hinting at an evolution of the microscopic configuration of the 0_3^+ and 0_4^+ along the isotopic chain. According to the present result and literature data, the 0_2^+ and 0_3^+ are expected to be strongly mixed, in agreement with predictions of IBM calculations for ^{114}Sn and ^{116}Sn . Similar conclusions can be drawn from preliminary MCSM calculations, which include states up to the 0_3^+ . In $^{116-120}\text{Sn}$, this hypothesis is also supported by the $\rho^2(E0; 0_3^+ \rightarrow 0_2^+)$ values lying close 100 m.u., of which the one of ^{118}Sn was determined in this work.

Of particular interest are the newly measured $B(E2; 0_4^+ \rightarrow 2_1^+)$ values. While being fast in ^{116}Sn , the transition becomes significantly hindered in ^{118}Sn and, specifically, in ^{120}Sn . Such a behavior could indicate substantial differences in the wavefunction of the 0_4^+ states with respect to that of the 2_1^+ states, which are expected to reside in the main spherical-oblate minimum of the potential energy surfaces of these nuclei. As a consequence, these states emerge as plausible candidates for being located in the well-separated, prolate-deformed secondary minimum of the potential energy surfaces of ^{118}Sn and ^{120}Sn .

In conclusion, this thesis work provides a significant amount of new spectroscopic information on 0^+ excitations in $^{110-120}\text{Sn}$ nuclei, which, together with the available literature, can serve as a remarkable testing ground for state-of-the-art shell model and beyond-mean-field calculations. A deeper insight into the nature of shape coexistence in this pivotal mass region is therefore expected to emerge in the near future.

The GSORT code

The GSORT program is the data-sorting code of the GASPware package, which was employed to analyze the experiments performed at IFIN-HH. The data processing was based on calling specific subroutines as keywords in a setup file, to perform calibrations, apply gates and produce matrices. In the following, an example of the setup file used for the analysis of *Exp. 4* (see Table 3.2) is shown.

A.1 The GSORT setup file

format GASP

```
*****
*** HEADER *****
*****
* Parameters (suppressedDetector - Ge)
* 0      1      2      3      4      5
* Energy  Time(ns) Time(ps) Energy BGO  Time BGO(ns)  Time BGO(ps)
*
* Parameters (Detector - La, Si)
* 0      1      2
* Energy  Time(ns)  Time(ps)
*
* G = HPGe
* L = LaBr/LaCe
* S = Solar Cells
* T = Pulser
* F0 = Run Timestamp in seconds

header      F      1  65536 PLUS 5  65536  65536  65536  65536  65536
detector    G  25      6  65536  65536  65536  65536  65536  65536
detector    S   6      3  65536  65536  65536
detector    T  25      3  65536  65536  65536

detector    A  25 PLUS  6  65536  65536  65536  65536  65536  65536

pairdef Cal/Ring2Ring.pairs

*****
*** CALIBRATIONS *****
*****

kill S      4

*** Global *****
recal      G0 Cal/EnSource      NORUN  ZERO  0.000  1.000  10  16384  0  25
```

```

recal          S0 Cal/SiAlignment  NORUN  ZERO  0.000  1.000  0  16384  0  6

*** Run-by-Run *****
recal          G0 Cal/EnBeam      RUN    ZERO  0.000  1.000  10 16384  0  25

copy G A

*****
*** GLOBAL GATES AND FOLD *****
*****

gate G1   900   1400 IN  1 25
gate G4   600   1250 OUT 1 20

*** Multiplicity filter *****
dmult F1 2
      G1 900 1400
      A4 600 1250

*****
*** HPGE ONLY MATRICES *****
*****

store

sort2d G0 G1 GeE-GeT      Res 8192 8192 Step 128 128
sort2d Gi G0 GeE-i       Res 25 8192 Step 1 128
sort2d Gi G1 GeT-i       Res 25 8192 Step 1 128
sort2d G0 F1 Ge-M        Res 8192 32 Step 128 2

sort2d_symm G0 Ge-symm  Res 8192 Step 128

sort3d_pair G0 G0 P0 Ge-R3D Res 8192 8192 64 Step 128 128 4

*****
*** Si-HPGE MATRICES *****
*****

*** Si matrices *****
sort2d S0 S1 SiE-SiT      Res 8192 8192 Step 128 128
sort2d Si S0 SiE-i       Res 8192 8192 Step 128 128
sort2d Si S1 SiT-i       Res 8192 8192 Step 128 128

*** Gates *****
gate S0   10   16000 IN  1 6
gate S1   650   1200 IN  1 6

*** Coincidence matrices *****
sort2d S0 G0 SiE-GeE      Res 8192 8192 Step 128 128

sort2d_symm G0 Ge-symm-SiGate Res 8192 Step 128

sort3d_pair G0 G0 P0 Ge-R3D-SiGate Res 8192 8192 64 Step 128 128 4

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List of Publications

As of 1st December 2025

Refereed publications

- **G. Corbari**, et al., Physics Letter B 855, 138851 (2024)

Publications in conference proceedings

- S. Bottoni, ..., **G. Corbari**, ..., et al., Nuclear Physics A 1062, 123138 (2025)
- **G. Corbari**, et al., Acta Physica Polonica B, Proc. Supp., 18, 2A27 (2025)
- A. Giaz, ..., **G. Corbari**, ..., et al., Acta Physica Polonica B, Proc. Supp., 18, 2A34 (2025)
- F. Galtarossa, ..., **G. Corbari**, ..., et al., European Physics Journal Web Conf. 324, 00029 (2025)
- G. Ciconali, ..., **G. Corbari**, ..., et al., Acta Physica Polonica B, Proc. Supp., 17, 3A5 (2024)
- **G. Corbari**, et al., Acta Physica Polonica B, Proc. Supp., 17, 3A1 (2024)
- **G. Corbari**, et al., Il Nuovo Cimento 46 C, 90 (2023)
- **G. Corbari**, et al., Acta Physica Polonica B, Proc. Supp., 16, 4A33 (2023)