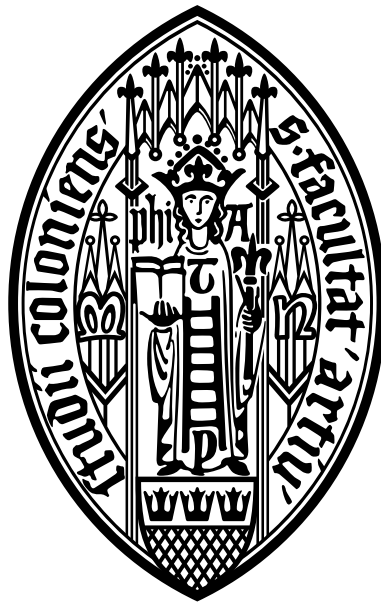


Nuclear Structure and Collectivity of Neutron-Rich Tin Isotopes around ^{132}Sn

Inaugural-Dissertation

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Abstract

The collective properties of the even–even neighbors of the doubly magic nucleus ^{132}Sn were investigated through a series of safe Coulomb-excitation experiments. The aim of this work was to obtain reliable information on the first excited 2^+ states in ^{130}Sn and ^{134}Sn , which provide sensitive experimental insight for the evolution of quadrupole collectivity near the $Z = 50$, $N = 82$ shell closure. Four experimental campaigns were conducted between November 2022 and October 2025 at the HIE-ISOLDE facility at CERN, using post-accelerated radioactive ion beams at 4.4 MeV/u incident on ^{206}Pb and ^{194}Pt targets. De-exciting γ rays were detected with the Miniball spectrometer in coincidence with scattered particles, enabling a clean selection of Coulomb-excited events. From these data, $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values of 0.0486(57) and 0.0449(57) $e^2\text{b}^2$ were extracted for ^{130}Sn and ^{134}Sn , respectively. These results represent the most precise determinations to date and resolve discrepancies between earlier experimental studies and theoretical predictions. In the context of the isotopic chain, the new values confirm the enhanced $B(E2)$ strength at ^{132}Sn and provide consistent evidence for a local maximum of collectivity at the doubly magic nucleus. The spectroscopic quadrupole moments of the 2_1^+ states in ^{130}Sn and ^{134}Sn were found to be compatible with zero within uncertainties, supporting a near-spherical structure of nuclei adjacent to the ^{132}Sn core. The presented results deliver well-constrained experimental benchmarks for large-scale shell-model calculations employing realistic interactions and contribute to a more coherent understanding of quadrupole collectivity and shell evolution in the $Z = 50$ region while also proving the sphericity and double-magic properties of ^{132}Sn .

Zusammenfassung

In einer Reihe sicherer Coulomb-Anregungsexperimente wurden die kollektiven Eigenschaften der beiden direkt benachbarten gerade-gerade Isotope des doppelt magischen Kerns ^{132}Sn untersucht. Ziel dieser Arbeit ist es, verlässliche Informationen über die ersten angeregten 2_1^+ -Zustände in ^{130}Sn und ^{134}Sn zu gewinnen, die empfindliche Referenzpunkte für die Entwicklung der Quadrupolkollektivität in der Nähe des Schalenabschlusses bei $Z = 50$, $N = 82$ darstellen. Zwischen November 2022 und Oktober 2025 wurden vier Experimentierkampagnen an der HIE-ISOLDE-Anlage am CERN durchgeführt. Postbeschleunigte radioaktive $^{130,134}\text{Sn}$ Ionenstrahlen wurden mit 4.4 MeV/u auf ^{206}Pb - bzw. ^{194}Pt -Targets beschleunigt. Die abregenden γ -Strahlen wurden mit dem Miniball-Spektrometer in Koinzidenz zu gestreuten Teilchen detektiert, was eine eindeutige Auswahl Coulomb-angeregter Ereignisse ermöglichte. Aus den Messdaten wurden $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ -Werte von 0.0486(57) und 0.0449(57) $e^2\text{b}^2$ für ^{130}Sn und ^{134}Sn bestimmt. Diese Ergebnisse stellen die bislang präzisesten Bestimmungen dar und lösen bestehende Widersprüche zwischen früheren experimentellen Arbeiten und theoretischen Vorhersagen. Im Kontext der Zinn-Isotopenkette bestätigen sie die erhöhte $B(E2)$ -Stärke in ^{132}Sn und liefern konsistente Hinweise auf ein lokales Maximum der Kollektivität am doppelt magischen Kern. Die aus den Daten extrahierten spektroskopischen Quadrupolmomente der 2_1^+ -Zustände in ^{130}Sn und ^{134}Sn sind innerhalb der Unsicherheiten mit null vereinbar und unterstützen somit ein nahezu sphärisches Kernmodell für die direkten Nachbarn von ^{132}Sn . Insgesamt liefern die in dieser Arbeit gewonnenen Ergebnisse gut abgesicherte experimentelle Referenzpunkte für großskalige Schalenmodellrechnungen mit realistischen Wechselwirkungen und tragen zu einem kohärenteren Verständnis der Quadrupolkollektivität und der Schalenentwicklung im $Z = 50$ -Gebiet bei und untermauern die sphärischen und doppeltmagischen Eigenschaften von ^{132}Sn .

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1. Introduction

The study of nuclear structure seeks to understand the arrangement and interactions of protons and neutrons inside the atomic nucleus. As a self-bound quantum many-body system governed by the strong interaction, the nucleus displays a rich variety of structural phenomena. These include shell effects, deformation, collective excitations, shape coexistence, and pairing correlations, which vary systematically across the nuclear chart. Understanding these features is crucial not only to describe stable nuclei, but also to predict the properties of exotic isotopes far from the valley of stability.

Nuclear structure physics connects experiment and theory through observables such as excitation energies, transition probabilities, nuclear moments, and binding energies. Over the past decades, increasingly precise measurements in nuclei far-off stability have been enabled by radioactive ion beam facilities and advanced detector arrays. On the theoretical side, models such as the nuclear shell model, collective models, and mean-field approaches have evolved into powerful tools to interpret these data within a consistent microscopic framework.

This chapter introduces the fundamental concepts used to describe the internal structure of nuclei. It begins with the nuclear shell model, which successfully explains many features of nuclear behavior based on the independent-particle motion in an average potential. The subsequent sections will highlight the role of electromagnetic observables and transition strengths in probing the structure and symmetry properties of excited nuclear states.

1.1. Nuclear Shell Model

The atomic nucleus, composed of protons and neutrons, is a strongly interacting quantum many-body system. Understanding its internal structure and stability has been a central topic in nuclear physics since the discovery of the neutron by James Chadwick in 1932 [1]. Shortly thereafter, Walter Elsasser [2, 3] proposed that nuclei exhibit shell-like structures similar to the electron shells in atoms. Empirical observations showed enhanced stability

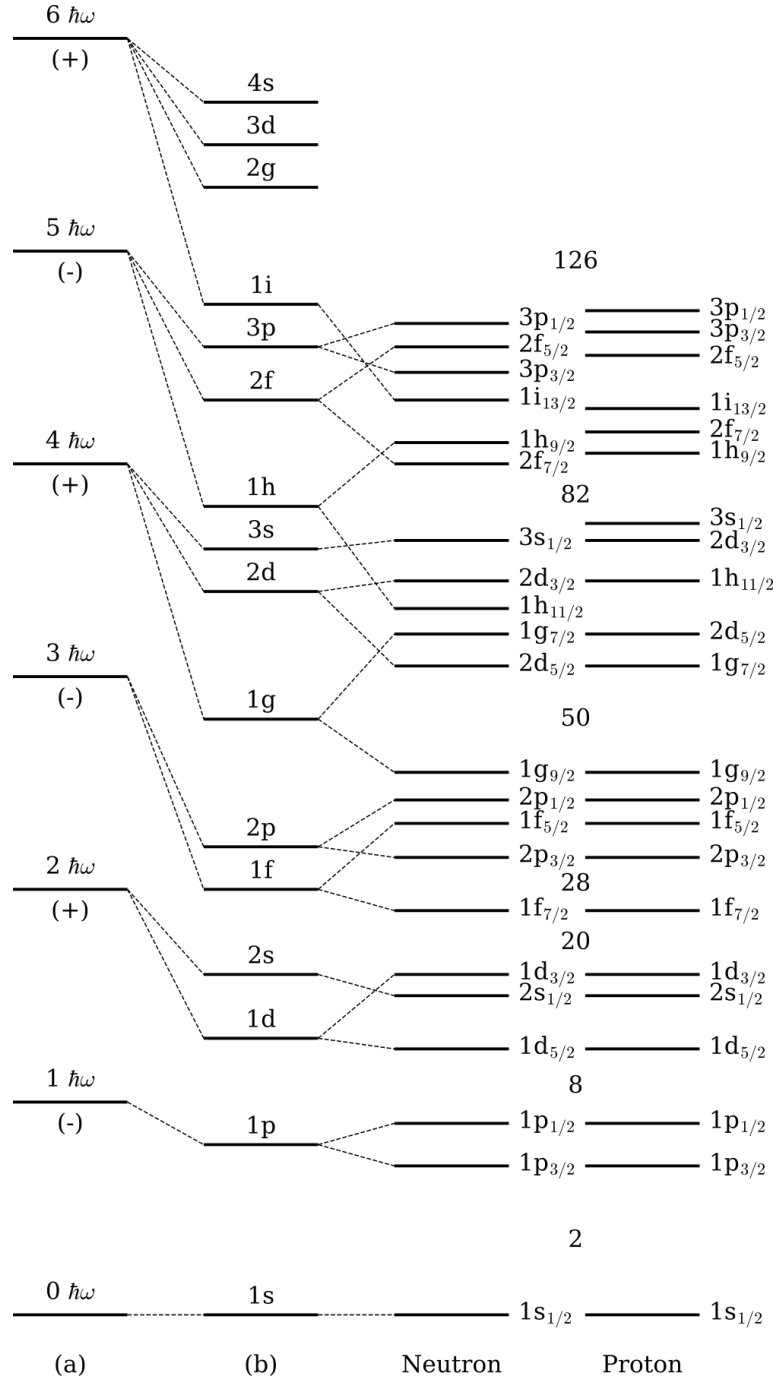


Figure 1.1. Comparison of nuclear single-particle levels in an harmonic oscillator potential (a), a spherically symmetric Wood-Saxon potential without spin-orbit coupling (b), and including spin-orbit interaction for neutrons and protons (right). The magic numbers emerge as gaps in the level density. Level positions are schematic and based on standard model calculations and experimental systematics; see, e.g., [4–6].

for nuclei with specific proton or neutron numbers—referred to as *magic numbers*—namely:

$$2, 8, 20, 28, 50, 82, 126.$$

These closed-shell nuclei are characterized by higher separation energies, lower neutron/proton capture cross sections, enhanced excitation energies for the first excited states, and reduced transition probabilities to the ground state. Nuclei with both proton and neutron numbers corresponding to magic numbers are termed *doubly magic*. A theoretical explanation for these magic numbers was independently provided in 1949 by Goeppert-Mayer [7] and by Haxel, Jensen, and Suess [8]. Their pivotal contribution was the inclusion of a strong spin-orbit interaction in the nuclear mean-field potential, analogous to the interaction observed in atomic systems. This modification allowed for a realistic description of the nuclear energy levels and accounted for the observed shell closures. The total nuclear Hamiltonian for a system of A nucleons consists of kinetic energy terms and two-body (and, in some cases, three-body) interactions. For practical reasons, especially in heavier nuclei, the problem is simplified by introducing a mean-field approximation. In this approach, each nucleon moves independently in an average potential $U(\vec{r})$ generated by all other nucleons. The total Hamiltonian is separated as:

$$H = \sum_{i=1}^A T_i + \sum_{i<j} V_{ij} = \sum_{i=1}^A (T_i + U(\vec{r}_i)) + \left(\sum_{i<j} V_{ij} - \sum_{i=1}^A U(\vec{r}_i) \right) = H_0 + H_R, \quad (1.1)$$

In this independent particle picture, the nucleons occupy discrete single-particle energy levels derived from solving the Schrödinger equation in a central potential. Due to the antisymmetric nature of fermions, the nuclear many-body wavefunction must be expressed as a Slater determinant of single-particle states. Early models such as the harmonic oscillator or Woods-Saxon potential could qualitatively reproduce some aspects of nuclear structure but failed to explain higher magic numbers. Only with the addition of the spin-orbit term,

$$V_i = V(r) + V_{ls}(r) (\vec{l} \cdot \vec{s}), \quad (1.2)$$

where $\vec{l}$ is the orbital angular momentum and $\vec{s}$ the spin, was it possible to quantitatively describe the experimentally observed magic numbers and associated nuclear properties such as spins and magnetic moments. Resulting single particle-levels using these different Hamiltonians are shown in Fig. 1.1, where also the appearance of the magic numbers is visible. While the mean-field approximation (H_0) captures the main features of the

spectrum, the residual interaction (H_R) accounts for configuration mixing and collective effects, particularly relevant in mid-shell nuclei with multiple valence nucleons.

Over the decades, the shell model has evolved into a cornerstone of nuclear structure theory, with applications ranging from the description of stable nuclei to the exploration of exotic systems far from stability. Notably, recent experimental and theoretical evidence suggests that traditional magic numbers may change in neutron-rich or proton-rich regions, revealing the dynamic and evolving nature of shell structure [9–11]. Prominent examples include the vanishing $N = 20$ shell closure in the first island of inversion around ^{32}Mg (e.g., Ref. [12]) and the emergence of new magic numbers at $N = 32, 34$ in the Ca isotope chain [13]. Even in the Sn isotope chain, *ab initio* shell-model calculations [14] as well as empirical [15] and realistic calculations including three-body forces [16] indicate a possible magic number at $N = 90$ in ^{140}Sn , which is not observed close to the line of stability.

1.2. Experimental Observables in Nuclear Physics

To understand the structure and dynamics of atomic nuclei, a broad range of experimental observables must be accessed. These observables serve as direct input for testing and refining theoretical models such as the shell model, which aims to provide a microscopic description of both ground and excited nuclear states. Especially for exotic nuclei far from stability, theoretical predictions must be validated against precise measurements to ensure the reliability of the underlying nuclear interactions.

Basic nuclear properties such as masses, excitation energies, spins, and parities of nuclear states often provide the first set of information in a spectroscopic study. These quantities are typically obtained via high-resolution particle or γ -ray spectroscopy. They allow for the identification of shell closures, energy gaps, or the occurrence of collective behavior in the nuclear excitation spectrum. However, while these basic observables are necessary, they are not always sufficient to draw firm conclusions about the internal configuration and the degree of collectivity of nuclear states. For this purpose, more sensitive quantities such as electromagnetic transition rates are required. These offer a more profound insight into the underlying wavefunctions and the configuration mixing within and across major shells [17]. Among these, the most relevant are the reduced transition probabilities $B(\sigma\lambda)$, which are directly connected to the matrix elements of the electromagnetic transition operators.

1.2.1. Electromagnetic Transition Strengths

Reduced transition probabilities $B(\sigma\lambda; I_i \rightarrow I_f)$ describe the likelihood of electromagnetic transitions between nuclear states of angular momentum I_i and I_f and provide a direct measure of the collectivity or single-particle nature of nuclear excitations. They are defined as [18]:

$$B(\sigma\lambda; I_i \rightarrow I_f) = \frac{1}{2I_i + 1} |\langle \psi_f | M(\sigma\lambda) | \psi_i \rangle|^2, \quad (1.3)$$

where $\langle \psi_f | M(\sigma\lambda) | \psi_i \rangle$ denotes the reduced matrix element of the electromagnetic multipole operator of type $\sigma\lambda$ ($\sigma = E$ for electric and M for magnetic transitions, λ being the multipolarity). The reduced transition probability directly enters the expression for the transition rate,

$$T(\sigma\lambda; I_i \rightarrow I_f) = \frac{8\pi(\lambda + 1)}{\lambda[(2\lambda + 1)!!]^2} \frac{1}{\hbar} \left(\frac{E_\gamma}{\hbar c} \right)^{2\lambda+1} \cdot B(\sigma\lambda; I_i \rightarrow I_f), \quad (1.4)$$

which describes how quickly a nuclear state decays by photon emission with energy E_γ . The lifetime τ of an excited nuclear state is inversely proportional to the total transition rate to all lower-lying states and is given by

$$\tau(I_i) = \left(\sum_{I_f} \sum_{\lambda} T(\sigma\lambda; I_i \rightarrow I_f) [1 + \alpha(\lambda)] \right)^{-1}, \quad (1.5)$$

where $\alpha(\lambda)$ denotes the internal conversion coefficient, accounting for non-radiative decay modes [19]. In practice, $B(E2)$ values between low-lying 2_1^+ and $0_{\text{g.s.}}^+$ states in even-even nuclei are among the most commonly studied transition strengths. These values reflect the quadrupole collectivity of the nucleus and are sensitive indicators for the collectivity of nuclei. Experimentally, reduced transition strengths can be determined either directly via lifetime measurements (e.g., Doppler Shift Attenuation Method (DSAM), Recoil Distance Doppler Shift (RDDS) [20]) or indirectly via Coulomb-excitation (short: Coulex) experiments. In the latter, electromagnetic matrix elements — and thus $B(\sigma\lambda)$ values — are extracted by comparing measured γ -ray yields to excitation probabilities using codes such as GOSIA [21, 22], which will be described in this thesis in Sec. 1.4. These methods are particularly valuable in regions where direct lifetime measurements are challenging due to short half-lives or low beam intensities [23]. To interpret the magnitude of $B(\sigma\lambda)$ values and assess the degree of collectivity, it is common to compare them to theoretical single-particle estimates known as *Weisskopf units* (W.u.) [19]. These provide a scale for electromagnetic transition strengths

under the assumption that only a single nucleon changes its orbit within a mean potential. The corresponding formulas for electric and magnetic transitions are:

$$B(E\lambda)_{\text{W.u.}} = \frac{1}{4\pi} \left(\frac{3}{\lambda+3} \right)^2 (1.2 A^{1/3})^{2\lambda} \quad [\text{e}^2 \text{fm}^{2\lambda}] \quad (1.6)$$

and

$$B(M\lambda)_{\text{W.u.}} = \frac{10}{4\pi} \left(\frac{3}{\lambda+3} \right)^2 (1.2 A^{1/3})^{2\lambda-2} \quad [\mu_N^2 \text{fm}^{2\lambda-2}] \quad (1.7)$$

If a measured $B(E\lambda)$ value exceeds the Weisskopf estimate significantly, it is usually interpreted as a sign of collective behavior involving coherent contributions from many nucleons. Values close to 1 W.u., on the other hand, typically reflect single-particle transitions. In this work, particular emphasis is placed on $B(E2)$ values as key observables to characterize the collective behavior and shell evolution in neutron-rich Sn nuclei. Their comparison with large-scale shell-model calculations provides insight into the effectiveness of the chosen interaction and the persistence of shell closures.

1.3. Coulomb Excitation as a Tool to Investigate Nuclear Structure

Coulomb excitation has been one of the most powerful and versatile tools to probe nuclear structure, particularly for collective excitations and electromagnetic transition strengths. Since its development in the 1950s [25], it has played a central role in determining a significant portion of known $B(\sigma\lambda)$ values and nuclear deformation parameters. Especially in regions far from stability, where alternative approaches such as lifetime measurements or inelastic scattering become increasingly difficult due to low beam intensities and short half-lives, Coulomb excitation remains a highly effective and selective method. The basic principle of Coulomb excitation is to excite a nucleus via the electromagnetic interaction alone, using the long-range nature of the Coulomb force. A schematic view of a Coulomb excitation pro-

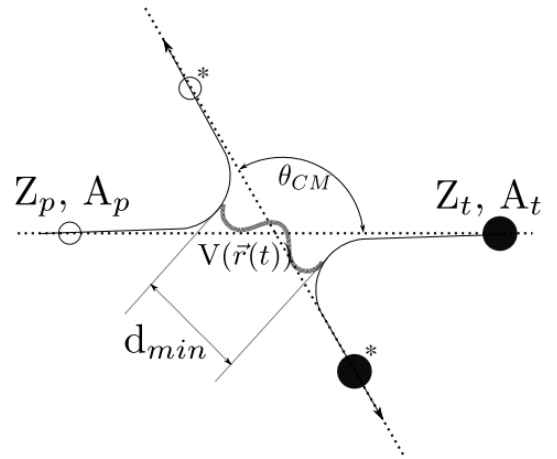


Figure 1.2. Schematic view of a Coulomb excitation process in the center-of-mass system. The interaction potential $V(\vec{r}(t))$ peaks at the point of closest approach $d_{\min}$, where the excitation probability is highest. The asterisk marks excited nuclei. Figure adopted from Ref. [24] with kind permission from D. Rosiak.

cess is shown in Fig. 1.2. A projectile nucleus with an atomic number Z_p is directed onto a target nucleus with an atomic number Z_t . If the beam energy E remains below the Coulomb barrier, the distance of closest approach $d_{\min}$ is sufficiently large to prevent any significant contribution from the nuclear force. Under these conditions, the interaction is purely electromagnetic and can be described within a well-controlled theoretical framework.

The excitation occurs through the time-dependent Coulomb field generated by the motion of the target and projectile with respect to one another. The strength of the interaction — and hence the probability of exciting a specific nuclear state — depends on several parameters, most notably the charges Z_p and Z_t , the relative energy E , and the scattering angle θ_{CM} . The maximum interaction potential is reached at the point of closest approach, where the excitation probability is highest. However, since the Coulomb interaction extends over a long range, the full trajectory of the interacting nuclei must be considered in the calculation of excitation probabilities.

This makes Coulomb excitation particularly well suited to a semi-classical treatment: the nuclear motion is approximated by classical trajectories, while the excitation process is treated quantum mechanically. In this picture, single and multiple-step excitations — both within the projectile and the target — can be described using matrix elements of the electromagnetic multipole operators. These matrix elements encode detailed structural information, such as deformation and collectivity, and can be extracted by comparing calculated and measured γ -ray yields from the de-excitation of nuclear states.

In the context of this work, Coulomb excitation is employed to determine reduced transition strengths $B(E2)$ and related electromagnetic matrix elements. These observables provide insight into the shape evolution and collectivity of neutron-rich nuclei and are particularly sensitive to the underlying shell structure. The theoretical formalism used to describe this process — including classical particle trajectories and excitation amplitudes — will be outlined in the following sections. A more detailed description is given in e.g., Refs. [26, 27] and references therein.

1.3.1. Semi-classical Particle Trajectories in Coulomb Excitation

The semi-classical approach is valid when the wavelength of the projectile is small compared to the distance of closest approach, which corresponds to a large Sommerfeld parameter. The Sommerfeld parameter η is a dimensionless quantity that characterizes the strength of

the Coulomb interaction between two nuclei with atomic numbers Z_p and Z_t , a projectile mass number A_p , and beam energy E_{CM} (in MeV):

$$\eta = \frac{Z_p Z_t e^2}{\hbar v} \approx \frac{Z_p Z_t \sqrt{A_p}}{6.32 \sqrt{E_{\text{CM}}}}. \quad (1.8)$$

For $\eta \ll 1$, the Coulomb interaction acts as a small perturbation, and first-order approximations (Born approximation) are valid. In contrast, for $\eta \gg 1$, the strong Coulomb field significantly alters the projectile's trajectory, which can then be approximated as a classical Rutherford orbit. In the experiments discussed in the present thesis using $^{130,134}\text{Sn}$ beams onto ^{206}Pb and ^{194}Pt targets, the Sommerfeld parameter amounts to $\eta_{^{130}\text{Sn}} = 308$ and $\eta_{^{134}\text{Sn}} = 293$, respectively.

The elastic scattering cross section in this regime is described by the Rutherford formula:

$$\left(\frac{d\sigma}{d\Omega} \right)_{\text{Ruth.}} = \left(\frac{Z_p Z_t e^2}{4E_{\text{CM}}} \right)^2 \frac{1}{\sin^4(\theta_{\text{CM}}/2)}, \quad (1.9)$$

where θ_{CM} is the center-of-mass scattering angle. The distance of closest approach between the projectile and target nucleus in such a trajectory is given by:

$$d_{\text{min}} = \frac{Z_p Z_t e^2}{2E_{\text{CM}}} \left(1 + \frac{1}{\sin(\theta_{\text{CM}}/2)} \right). \quad (1.10)$$

For the closest approach ($\theta_{\text{CM}} = 180^\circ$) this can be rewritten as

$$d_{\text{min}} = 2 \frac{Z_p Z_t e^2}{m_0 v^2}. \quad (1.11)$$

In inelastic processes, where energy is transferred to internal excitations, the adiabaticity of the collision becomes important. The parameter ξ is introduced to quantify the fraction of kinetic energy available for excitation:

$$\xi = \frac{Z_p Z_t e^2}{\hbar v_\infty} \cdot \frac{\Delta E}{2E}, \quad (1.12)$$

with $E = \frac{1}{2} m_0 v_\infty^2$ and ΔE the excitation energy. A related parameter is the eccentricity of

the classical orbit,

$$\varepsilon = \frac{1}{\sin(\theta_{\text{CM}}/2)}, \quad (1.13)$$

which depends on the scattering angle. The product $\xi \cdot \varepsilon$ serves as a measure for the ratio of collision time τ to the internal timescale (period) T of the nucleus, thereby indicating the degree of adiabaticity of the process. If $\xi\varepsilon \gg 1$, the excitation occurs adiabatically and is strongly suppressed. Non-adiabatic conditions ($\xi\varepsilon \lesssim 1$) favor transitions between nuclear states. In the present experiments $\xi\varepsilon \approx 1$ for all angles covered by the experimental setup and for excitation energies of ≈ 1 MeV, showing that Coulomb excitation of states around 1 MeV is feasible in the present kinematics, albeit in a borderline regime where the excitation probability remains modest and the measurement becomes experimentally demanding.

The nuclear period T is inversely proportional to the excitation frequency ω of the transition,

$$T \propto \frac{1}{\omega} = \frac{\hbar}{\Delta E}, \quad (1.14)$$

with ΔE being the energy difference between the final and initial state.

Within this semi-classical picture, the cross section for Coulomb excitation of a given nuclear state can be expressed as:

$$\left(\frac{d\sigma}{d\Omega} \right)_{\text{Coul}} = P_{if} \cdot \left(\frac{d\sigma}{d\Omega} \right)_{\text{Ruth.}}, \quad (1.15)$$

where P_{if} denotes the excitation probability for a transition from an initial state i to a final state f . This probability is given by the coherent sum over magnetic substates of the squared transition amplitudes b_{if} :

$$P_{if} = \frac{1}{2I_i + 1} \sum_{M_i, M_f} |b_{if}|^2. \quad (1.16)$$

The semi-classical formalism thus provides a practical and physically intuitive framework for calculating excitation probabilities in heavy-ion collisions at energies below the Coulomb barrier, where nuclear forces can be neglected and electromagnetic interactions dominate.

1.3.2. Coulomb Excitation as First-Order Perturbation Theory

The quantum-mechanical framework is based on the time-dependent Schrödinger equation:

$$i\hbar \frac{d}{dt} \psi(t) = (H_0 + H(t)) \psi(t), \quad (1.17)$$

where H_0 is the unperturbed Hamiltonian and $H(t)$ describes the time-dependent interaction between the nuclear charge distributions. For Coulomb excitation, $H(t)$ consists of three parts: a monopole-monopole interaction responsible for the classical Rutherford scattering, a monopole-multipole term accounting for the excitation of one nucleus by the other, and a multipole-multipole interaction representing simultaneous excitation of both collision partners. The last term is typically negligible due to the small excitation probabilities involved. In first-order per-

turbation theory, transitions between the initial state $|i\rangle$ and final state $|f\rangle$ occur directly, without involving intermediate quantum states. This is illustrated in the left part of Fig. 1.3. The corresponding transition amplitude is given by:

$$b_{if} = \frac{1}{i\hbar} \int_{-\infty}^{\infty} \langle f | H(t) | i \rangle e^{i\omega t} dt, \quad (1.18)$$

where $\omega = (E_f - E_i)/\hbar$ is the nuclear transition frequency. Since most Coulomb excitations are weak, the condition $|b_{if}| \ll 1$ is generally fulfilled, and the excitation probability is given to leading order by $P_{if} = |b_{if}|^2$.

For electric multipole transitions, the electromagnetic interaction can be expanded in terms of spherical multipoles,

$$H_E(t) = \int \rho_n(\vec{r}) \phi(\vec{r}, t) d\tau, \quad (1.19)$$

where $\rho_n(\vec{r})$ is the nuclear charge density and $\phi(\vec{r}, t)$ is the time-dependent electric potential due to the projectile:

$$\phi(\vec{r}, t) = \frac{Z_p e}{|\vec{r} - \vec{r}_p(t)|} - \frac{Z_p e}{|\vec{r}_p(t)|}. \quad (1.20)$$

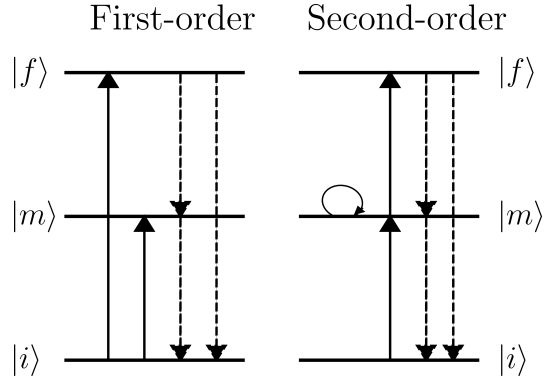


Figure 1.3. Different excitation and deexcitation paths in first-order and second order perturbation theory. For more information, see text.

This leads to a multipole expansion of the interaction Hamiltonian:

$$H_E(t) = 4\pi Z_p e \sum_{\lambda=1}^{\infty} \sum_{\mu=-\lambda}^{\lambda} \frac{Y_{\lambda\mu}(\theta(t), \phi(t))}{[r_p(t)]^{\lambda+1}} M(E\lambda, \mu), \quad (1.21)$$

where $M(E\lambda, \mu)$ denotes the electric multipole operator of order λ :

$$M(E\lambda, \mu) = \int \rho_n(\vec{r}) r^\lambda Y_{\lambda\mu}(\theta, \phi) d\tau. \quad (1.22)$$

Using these expressions, the transition amplitude becomes

$$b_{if} = \frac{4\pi Z_p e}{i\hbar} \sum_{\lambda, \mu} \frac{1}{2\lambda + 1} \langle I_f M_f | M(E\lambda, \mu) | I_i M_i \rangle \cdot S(E\lambda, \mu), \quad (1.23)$$

where the orbital integral

$$S(E\lambda, \mu) = \int_{-\infty}^{\infty} e^{i\omega t} \frac{Y_{\lambda\mu}(\theta(t), \phi(t))}{[r_p(t)]^{\lambda+1}} dt \quad (1.24)$$

contains all information about the classical trajectory.

Assuming unpolarized beams, the excitation probability is obtained by summing over all magnetic substates:

$$P_{i \rightarrow f} = \frac{1}{2I_i + 1} \sum_{M_i, M_f} |b_{if}|^2. \quad (1.25)$$

Finally, the differential Coulomb-excitation cross section can be expressed as:

$$\left(\frac{d\sigma}{d\Omega} \right)_{E\lambda} = \left(\frac{Z_p e}{\hbar v} \right)^2 \left(\frac{d_{\min}}{2} \right)^{-2\lambda+2} B(E\lambda; I_i \rightarrow I_f) \cdot f_{E\lambda}(\theta, \xi), \quad (1.26)$$

where $f_{E\lambda}(\theta, \xi)$ is a dimensionless function encapsulating the details of the trajectory and the adiabaticity of the collision. The magnetic case is treated analogously, but suppressed by a factor $\beta^2 = (v/c)^2$ for the same multipolarity λ :

$$\left(\frac{d\sigma}{d\Omega} \right)_{M\lambda} = \left(\frac{Z_p e}{\hbar c} \right)^2 \left(\frac{d_{\min}}{2} \right)^{-2\lambda+2} B(M\lambda; I_i \rightarrow I_f) \cdot f_{M\lambda}(\theta, \xi). \quad (1.27)$$

This formalism illustrates that Coulomb excitation, when treated as a first-order perturbative

process, provides direct access to electromagnetic transition matrix elements and the associated $B(\sigma\lambda)$ values, as long as the excitation remains weak and the nuclear motion is well described classically.

1.3.3. Coulomb Excitation in Second-Order Perturbation Theory

While first-order perturbation theory provides a good approximation for many Coulomb-excitation processes, it becomes insufficient in heavy-ion collisions where excitation probabilities $P_{i \rightarrow f}$ may approach unity. In such cases, higher-order effects must be considered, leading to modifications of direct excitation probabilities and enabling access to otherwise forbidden transitions via intermediate states.

The time-dependent perturbation theory can be extended to second order by including two-step excitation paths, where the system is excited from the initial state $|i\rangle$ to an intermediate state $|m\rangle$ and then to the final state $|f\rangle$. The total excitation amplitude becomes:

$$b_{if}^{(2)} = b_{if}^{(1)} + \sum_m b_{imf}^{(2)}, \quad (1.28)$$

with

$$b_{imf}^{(2)} = \frac{1}{(i\hbar)^2} \int_{-\infty}^{\infty} \langle f | H(t) | m \rangle e^{i\omega_{fm}t} dt \cdot \int_{-\infty}^t \langle m | H(t') | i \rangle e^{i\omega_{mi}t'} dt', \quad (1.29)$$

where $\omega_{xy} = (E_x - E_y)/\hbar$.

The corresponding probability for reaching the final state is then given by the square modulus of the total amplitude:

$$P_{i \rightarrow f} = \left| b_{if}^{(1)} + \sum_m b_{imf}^{(2)} \right|^2. \quad (1.30)$$

Expanding this expression, it becomes clear that not only quadratic second-order contributions appear, but also interference terms with the first-order amplitude:

$$P_{i \rightarrow f} = |b_{if}^{(1)}|^2 + 2\Re \left(b_{if}^{(1)} \cdot \sum_m b_{imf}^{(2)} \right) + \left| \sum_m b_{imf}^{(2)} \right|^2. \quad (1.31)$$

This interference is a key feature of second-order processes and can lead to observable modifications of excitation probabilities even when $|b^{(2)}| \ll |b^{(1)}|$.

Multi-Step Excitation and the Reorientation Effect

Second-order perturbation theory enables excitation pathways that are forbidden or highly suppressed in first-order processes. A prominent example is the population of the 4_1^+ state in even-even nuclei through sequential $E2$ transitions:

$$0_{\text{g.s.}}^+ \xrightarrow{E2} 2_1^+ \xrightarrow{E2} 4_1^+. \quad (1.32)$$

Due to the relatively strong coupling associated with $E2$ transitions, such two-step processes often dominate over the much weaker direct $E4$ excitation. As a result, higher-lying states in even-even nuclei are frequently populated via multi-step excitation chains. Another important second-order effect is the reorientation effect, which involves transitions between magnetic substates of the same excited level—typically the first 2^+ state in even-even nuclei:

$$|2^+, M_i\rangle \longrightarrow |2^+, M_f\rangle. \quad (1.33)$$

Although there is no energy change ($\Delta E = 0$) between the initial and final states, the interaction with the time-dependent electric field gradient can induce such transitions. The reorientation effect is particularly sensitive to the spectroscopic quadrupole moment Q_s , which is related to the diagonal $E2$ matrix element by

$$eQ_s = \sqrt{\frac{16\pi}{5}} \frac{\langle II\ 20|II\rangle}{\sqrt{2I+1}} \langle I||M(E2)||I\rangle, \quad (1.34)$$

This sensitivity allows Coulomb excitation to serve as a direct probe of both the magnitude and the sign of Q_s . The spectroscopic quadrupole moment is connected to the intrinsic quadrupole moment Q_0 through

$$Q_s = \frac{3K^2 - I(I+1)}{(I+1)(2I+3)} Q_0, \quad (1.35)$$

where I denotes the spin of the state and K is the projection of the total angular momentum I onto the body-fixed symmetry axis. For axially symmetric even-even nuclei in their ground-state band, $K = 0$ typically applies. This simplifies the formula to $Q_s = -\frac{2}{7}Q_0$. The sign of Q_0 provides information about the nuclear shape: a positive Q_0 corresponds to a prolate (rugby-ball-like) deformation, while a negative Q_0 indicates an oblate (pancake-like) shape. Under the assumption of a uniformly charged nucleus, the intrinsic quadrupole moment

can be related to the quadrupole deformation parameter β_2 via

$$Q_0 \approx \frac{3}{\sqrt{5}\pi} Z R_0^2 \beta_2 (1 + 0.36\beta_2), \quad (1.36)$$

where Z is the proton number and R_0 is the nuclear radius parameter. The equations used in this section are adopted from Refs. [22, 24–28]. These references provide comprehensive derivations and additional background on Coulomb excitation and collective quadrupole observables.

Sensitivity to Nuclear Structure

The presence of second-order processes has practical implications for the extraction of nuclear matrix elements from Coulomb-excitation data. Neglecting them can lead to systematic deviations, especially when strong low-lying transitions or close-lying excited states are involved. Therefore, accurate analysis codes such as GOSIA include the full formalism of second-order Coulomb excitation, allowing the correct treatment of interference and multi-step excitation processes.

Second-order perturbation theory thus extends Coulomb excitation by (a) accounting for interference with direct transitions, (b) enabling access to high-lying states through multi-step excitation, and (c) following the determination of spectroscopic quadrupole moments

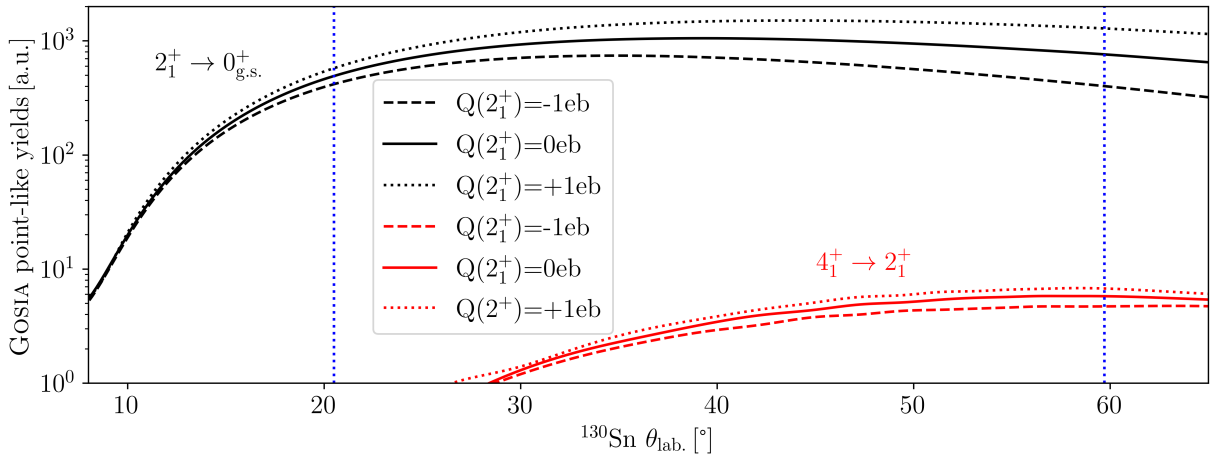


Figure 1.4. Calculated point-like yields for the $^{130}\text{Sn } 2_1^+ \rightarrow 0_{\text{g.s.}}^+$ (black) and $4_1^+ \rightarrow 2_1^+$ (red) transitions. The calculations assume a ^{206}Pb target and a beam energy of 510 MeV at the target center. The effect of a nonzero quadrupole moment on the excitation cross-sections is illustrated. The blue vertical lines indicate the angular range covered in the experiment (CD range).

via reorientation effects.

Figure 1.4 displays point-like $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ (black) and $4_1^+ \rightarrow 2_1^+$ (red) γ -ray yields in ^{130}Sn , expected for an experiment with a 510 MeV ^{130}Sn beam impinging onto a ^{206}Pb target. The effect of second-order perturbation theory is visible in the excitation of the 4_1^+ as well as in the influence of the quadrupole moment on the experimental yields. The calculations have been performed with the GOSIA code [22].

1.3.4. Safe Coulex

The concept of *safe* Coulomb excitation refers to scattering experiments in which nuclear excitation occurs exclusively via the electromagnetic interaction, with negligible contributions from the nuclear force. This condition is met when the interacting nuclei remain at a sufficiently large distance from one another, such that their nuclear surfaces do not come into contact. The minimum distance of closest approach, denoted as d_{min} , must exceed the sum of the nuclear radii of the projectile (R_p) and target (R_t), plus an additional safety margin ΔD . In Ref. [29], Cline *et al.* determined that setting $\Delta D = 5.0$ fm limits the nuclear interaction contribution to less than 0.1 %. The corresponding *safe* bombarding energy as a function of the center-of-mass scattering angle θ_{CM} is given by:

$$E(\theta_{\text{CM}}) = 0.72 \cdot \frac{Z_p Z_t}{d_{\text{min}}} \cdot \frac{A_p + A_t}{A_t} \left(1 + \frac{1}{\sin(\theta_{\text{CM}}/2)} \right), \quad (1.37)$$

where Z and A denote the atomic and mass numbers of projectile and target nuclei, respectively. The radii of the involved nuclei can be determined by assuming sphere-shaped nuclei.

$$R_i = R_0 \cdot A_i^{\frac{1}{3}} \quad (1.38)$$

For heavy nuclei, it is also possible to utilize the approach of radii of half density of the Fermi mass distribution [30]

$$R_i = r_i \left(1 - \frac{1}{r_i^2} \right) \quad (1.39)$$

with the nuclear radii as defined via the liquid drop model given as [31]

$$r_i = 1.12 \cdot A_i^{\frac{1}{3}} - 0.94 \cdot A_i^{-\frac{1}{3}} \quad (1.40)$$

For the specific case of the ^{130}Sn experiment, a safe distance of $d_{\text{min, sphere}} = 18.71$ fm and $d_{\text{min, fermi}} = 16.6$ fm have been determined. This corresponds to a safe bombarding energy of

4.0 MeV/u and 4.4 MeV/u, respectively. To enhance the Coulomb excitation cross section and ensure sufficient statistics, a higher beam energy of 4.4 MeV/u was employed, since for masses in this region, the Fermi approximation can be applied [30]. For the ^{134}Sn experiment, a ^{194}Pt target was chosen. A beam energy of 4.4 MeV/u was used again, and the analyzed scattering angles were limited to $\theta_{130\text{Sn, lab.}} 88^\circ$, since for higher scattering angles safe Coulex would no longer be fulfilled.

1.4. The Gosia code

The study of nuclear structure via Coulomb excitation has a rich history dating back to the 1950s, when early experiments used light-ion beams to probe the electromagnetic properties of nuclei. These initial investigations were relatively straightforward to interpret, relying on first- and second-order perturbation theory to describe the excitation processes [32]. The development of heavy-ion beams in the following decades enabled the population of multiple excited states through multi-step Coulomb excitation. The complexity of these new experimental scenarios necessitated more sophisticated theoretical tools for quantitative analysis. A pivotal advance came with the semi-classical theory of multiple Coulomb excitation developed by Alder and Winther [26], which was implemented in early computer codes such as that by Winther and de Boer [33]. These programs allowed the calculation of excitation amplitudes based on a set of reduced electromagnetic matrix elements provided as input.

Building on this foundation, the GOSIA code was introduced in the 1980s [21, 27] to enable a more direct and comprehensive analysis of Coulomb excitation experiments. In contrast to earlier codes that calculated excitation probabilities based on fixed matrix elements, GOSIA employs a least-squares fitting procedure wherein the matrix elements themselves are treated as adjustable parameters. By iteratively comparing calculated γ -ray yields with experimental data, the code extracts the set of matrix elements that best reproduces the observed excitation and de-excitation patterns. A key feature of GOSIA is that unknown matrix elements are normalized to independently known experimental quantities. These may include absolute transition probabilities derived from measured lifetimes, cross sections obtained from Rutherford scattering, or previously established matrix elements. This normalization provides an absolute scale for the extracted matrix elements, ensuring physical consistency and compatibility with other observables.

Crucially, GOSIA calculates both the excitation of nuclear states and their subsequent γ -ray

de-excitation within a unified framework, using the same electromagnetic matrix elements. The description of γ -ray emission is based on the CEGRY code [34] and incorporates essential effects such as angular correlations, nuclear deorientation, recoil corrections, Jacobian transformations to a common reference frame, and integration over realistic detector geometries. The first successful application of GOSIA demonstrated the uniqueness of the extracted matrix element set for ^{110}Pd [35], a result subsequently confirmed by recoil-distance Doppler shift lifetime measurements [36]. Since then, GOSIA has been validated through extensive analyses of isotopic chains such as W, Os, and Pt [37–40]. It has played a critical role in elucidating nuclear shape evolution and shape coexistence across a wide range of nuclei, including transitional regions [41, 42], rare earth elements [43, 44], actinides [45], and nuclei exhibiting exotic octupole deformations [46].

In addition to the standard version, a variant known as GOSIA2 is used to perform the minimization using *target normalization*. In this approach, electromagnetic matrix elements of the projectile nucleus are determined relative to known matrix elements of the target nucleus. This is particularly useful when absolute cross sections or lifetime information for the projectile are not available, but well-established electromagnetic properties of the target can serve as a reference. This is usually the case for experiments using rare isotope beams, like radioactive ion beams. GOSIA2 allows both excitation and de-excitation processes in the target and projectile to be treated simultaneously within the same fitting routine, ensuring a consistent relative normalization. In the context of this work, GOSIA2 is employed to extract matrix elements of the projectile by referencing them to the known properties of the target nucleus. A more detailed description of GOSIA is given in Ref. [22, 27].

2. Physics Case

The ten doubly magic nuclei (${}^4_2\text{He}$, ${}^{16}_8\text{O}$, ${}^{40}_{20}\text{Ca}$, ${}^{48}_{20}\text{Ca}$, ${}^{48}_{28}\text{Ni}$, ${}^{56}_{28}\text{Ni}$, ${}^{78}_{28}\text{Ni}$, ${}^{100}_{50}\text{Sn}$, ${}^{132}_{50}\text{Sn}$, ${}^{208}_{82}\text{Pb}$) form indispensable reference points for nuclear-structure theory. Among them, tin occupies a special place: with $Z = 50$ it is the heaviest element that hosts two doubly magic isotopes, ${}^{100}\text{Sn}$ at the proton drip line and ${}^{132}\text{Sn}$ on the neutron-rich side. This unique situation turns the Sn isotopic chain into a stringent laboratory for testing microscopic descriptions of shell structure across a wide range of isospin, from the vicinity of $N = Z$ to the region beyond the major neutron shell closure at $N = 82$.

The interest in nuclei around ${}^{132}\text{Sn}$ extends beyond nuclear structure physics. The neighborhood of the heaviest unstable doubly magic nucleus is intimately connected to rapid neutron-capture (r) process nucleosynthesis: the persistence of the $N = 82$ shell closure influences reaction-path trajectories and contributes to shaping the abundance distribution near the second r -process peak [47–49]. In this context, recent progress has been driven by the interplay of multimessenger observations and advances in precision nuclear data. JWST (James Webb Space Telescope) observations of the kilonova associated with Gamma-Ray-Burst 230307A provided direct evidence for heavy-element production, in-

54			${}^{108}\text{Xe}$	${}^{110}\text{Xe}$	${}^{112}\text{Xe}$	${}^{114}\text{Xe}$	${}^{116}\text{Xe}$	${}^{118}\text{Xe}$	${}^{120}\text{Xe}$	${}^{122}\text{Xe}$	${}^{124}\text{Xe}$	${}^{126}\text{Xe}$	${}^{128}\text{Xe}$	${}^{130}\text{Xe}$	${}^{132}\text{Xe}$	${}^{134}\text{Xe}$	${}^{136}\text{Xe}$	${}^{138}\text{Xe}$
52		${}^{104}\text{Te}$	${}^{106}\text{Te}$	${}^{108}\text{Te}$	${}^{110}\text{Te}$	${}^{112}\text{Te}$	${}^{114}\text{Te}$	${}^{116}\text{Te}$	${}^{118}\text{Te}$	${}^{120}\text{Te}$	${}^{122}\text{Te}$	${}^{124}\text{Te}$	${}^{126}\text{Te}$	${}^{128}\text{Te}$	${}^{130}\text{Te}$	${}^{132}\text{Te}$	${}^{134}\text{Te}$	${}^{136}\text{Te}$
50	${}^{100}\text{Sn}$	${}^{102}\text{Sn}$	${}^{104}\text{Sn}$	${}^{106}\text{Sn}$	${}^{108}\text{Sn}$	${}^{110}\text{Sn}$	${}^{112}\text{Sn}$	${}^{114}\text{Sn}$	${}^{116}\text{Sn}$	${}^{118}\text{Sn}$	${}^{120}\text{Sn}$	${}^{122}\text{Sn}$	${}^{124}\text{Sn}$	${}^{126}\text{Sn}$	${}^{128}\text{Sn}$	${}^{130}\text{Sn}$	${}^{132}\text{Sn}$	${}^{134}\text{Sn}$
48	${}^{98}\text{Cd}$	${}^{100}\text{Cd}$	${}^{102}\text{Cd}$	${}^{104}\text{Cd}$	${}^{106}\text{Cd}$	${}^{108}\text{Cd}$	${}^{110}\text{Cd}$	${}^{112}\text{Cd}$	${}^{114}\text{Cd}$	${}^{116}\text{Cd}$	${}^{118}\text{Cd}$	${}^{120}\text{Cd}$	${}^{122}\text{Cd}$	${}^{124}\text{Cd}$	${}^{126}\text{Cd}$	${}^{128}\text{Cd}$	${}^{130}\text{Cd}$	${}^{132}\text{Cd}$
46	${}^{96}\text{Pd}$	${}^{98}\text{Pd}$	${}^{100}\text{Pd}$	${}^{102}\text{Pd}$	${}^{104}\text{Pd}$	${}^{106}\text{Pd}$	${}^{108}\text{Pd}$	${}^{110}\text{Pd}$	${}^{112}\text{Pd}$	${}^{114}\text{Pd}$	${}^{116}\text{Pd}$	${}^{118}\text{Pd}$	${}^{120}\text{Pd}$	${}^{122}\text{Pd}$	${}^{124}\text{Pd}$	${}^{126}\text{Pd}$	${}^{128}\text{Pd}$	${}^{130}\text{Pd}$
	50	52	54	56	58	60	62	64	66	68	70	72	74	76	78	80	82	84

Figure 2.1. Even-even nuclear chart from $Z = 46$ – 54 and $N = 50$ – 84 , highlighting the magic numbers $Z = 50$, $N = 50$ and $N = 82$ (black rectangles). Stable isotopes are indicated in black, unstable ones in white.

cluding ejecta components with nuclei in the vicinity of $A \approx 130$ [50]. Complementarily, high-precision mass measurements of neutron-rich tin isotopes beyond $N = 82$ revealed a reduced neutron-pairing gap and a softened two-neutron separation-energy trend, with direct consequences for r -process network flows near $A \sim 135$ [14]. In parallel, β -decay and β -delayed neutron-emission studies of very neutron-rich Sn isotopes demonstrated that one- and two-neutron emission probabilities are highly sensitive to the underlying shell structure and to the degree to which the $N = 82$ closure persists towards more neutron-rich systems [51].

From the perspective of nuclear structure, the ^{132}Sn region provides a particularly clean environment to isolate key ingredients of effective interactions. Because $Z = 50$ is closed, low-lying excitations in the even-even Sn isotopes are dominated by neutron configurations, while proton degrees of freedom remain largely inert at low excitation energy. As a consequence, spectroscopy along the Sn chain constitutes an excellent testing ground for the neutron-neutron component of shell-model interactions, and for the evolution of shell structure as neutrons are added across major gaps. At the same time, nearby odd-mass and odd-odd nuclei offer sensitive probes of single-particle character and the robustness of shell closures. Laser spectroscopy of the $Z = 49$ indium isotopes — one proton hole below the $Z = 50$ shell — revealed an abrupt change in magnetic-dipole and quadrupole moments at $N = 82$, indicating that a nearly pure single-particle character is recovered only at the doubly magic core [52]. Consistently, recent spectroscopy of the odd-odd nucleus ^{128}Ag shows a well-defined seniority structure and no excitations across the $N = 82$ and $Z = 50$ gaps, underscoring the robustness of these closures in the region [53].

While shell closures provide a natural organizing principle, the emergence of collectivity in their vicinity is a central theme of this thesis. A key observable in this context is the reduced electric quadrupole transition probability, $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ which typically exhibits a pronounced minimum at doubly magic nuclei and increases in neighboring open-shell systems. For ^{132}Sn , however, available evidence suggests that this simple expectation may not apply. A pioneering Coulomb-excitation experiment, supported by large-scale shell-model (LSSM) and Monte Carlo shell-model (MCSM) calculations, reported an enhanced quadrupole collectivity and attributed it to increased proton excitations across the $Z = 50$ shell gap [23, 54]. Although ^{132}Sn itself exhibits a pronounced shell closure [23, 55], the extent to which quadrupole collectivity is enhanced specifically at ^{132}Sn , and how it evolves in the adjacent isotopes, remains experimentally unresolved: transition probabilities in ^{130}Sn and ^{134}Sn are still largely unmeasured, leaving the systematics around the shell closure

an open question.

On the theoretical side, the situation is similarly unsettled. MCSM calculations by Togashi *et al.* predict an enhancement by about a factor of two in ^{132}Sn relative to its neighbors, whereas LSSM calculations by Naïdja, Nowacki and Sieja [23, 56] suggest an even stronger increase by roughly a factor of four. At the same time, the MCSM framework reproduces the overall systematics of $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values across the Sn isotopic chain, yet it disagrees with the only existing estimates for $^{130,134}\text{Sn}$ reported in conference proceedings [57–60]. Those tentative values — $B(E2) = 0.023(5)$ and $0.029(5) e^2\text{b}^2$ for ^{130}Sn and ^{134}Sn , respectively — are consistent with LSSM predictions of 0.028 and $0.027 e^2\text{b}^2$, but deviate substantially from the MCSM results of 0.055 and $0.044 e^2\text{b}^2$. Resolving this discrepancy is essential for establishing a consistent picture of quadrupole collectivity in and around the doubly magic nucleus ^{132}Sn .

In addition to even–even Sn, theoretical studies emphasize the role of proton–neutron correlations in shaping quadrupole response near the closed shells. Shell-model calculations employing different nucleon–nucleon interactions suggest that enhanced quadrupole collectivity in the ^{132}Sn region can arise from constructive proton–neutron quadrupole coherence, such that adding or removing only a single nucleon from the doubly magic core may induce a pronounced collective response [61]. Testing this mechanism in the odd-proton nucleus ^{131}Sb requires, in particular, precise knowledge of electromagnetic matrix elements in its even–even core nucleus ^{130}Sn , thereby providing a direct structural information for the onset of collectivity in the region [62].

The following sections summarize key experimental developments along the Sn isotopic chain, starting from the doubly magic ^{100}Sn and progressing towards increasingly neutron-rich nuclei. Special attention is devoted to the region around ^{132}Sn , where the most stringent constraints on shell evolution and quadrupole collectivity emerge, and where the focus of this thesis is set.

2.1. The Sn Isotope Chain Between the Doubly-Magic Nuclei ^{100}Sn and ^{132}Sn

The even-even isotopes of the nuclear chart between $N = 50 - 84$ and $Z = 46 - 54$ are shown in Fig. 2.1. The two doubly-magic nuclei, ^{100}Sn and ^{132}Sn , highlighted here by the crossing magic numbers, are of enduring experimental and theoretical interest. The region around the $N = Z$ nucleus ^{100}Sn has only recently become experimentally accessible [63].

Considerable effort has been devoted to verifying its doubly-magic character, although excited states of this "holy grail of nuclear structure research" [64] have yet to be observed.

A breakthrough came with the observation of superallowed α decays from ^{108}Xe to ^{104}Te and finally to ^{100}Sn [65], with reported half-lives of $t_{1/2}(^{108}\text{Xe}) = 58_{-23}^{+106} \mu\text{s}$ and $t_{1/2}(^{104}\text{Te}) < 18 \text{ ns}$. The half-life of ^{100}Sn itself was first determined in 2012 at GSI [66]. In 2019, a tenfold increase in ^{100}Sn production was achieved at RIKEN [67], allowing a more precise measurement of $t_{1/2}(^{100}\text{Sn}) = 1.18(8) \text{ s}$. For a comprehensive review about the ongoing research around ^{100}Sn , see Ref. [87].

Starting from the doubly-magic ^{100}Sn to more neutron-rich isotope excitation energies of the first 2^+ states along the tin isotopic chain are remarkably constant between ^{102}Sn and ^{130}Sn , as shown in the left plot of Fig. 2.2. This trend is well described within the seniority scheme [88, 89], where the ground state is approximated by a condensate of $J^\pi=0$ neutron pairs [18]. A slight increase around $N = 64$ is attributed to a weak sub-shell closure. However, the overall constancy suggests that pairing correlations dominate and dilute shell effects, resulting in similar 2_1^+ configurations across the chain. Because of the large and nearly constant $Z = 50$ shell gap, proton excitations require significant energy, which

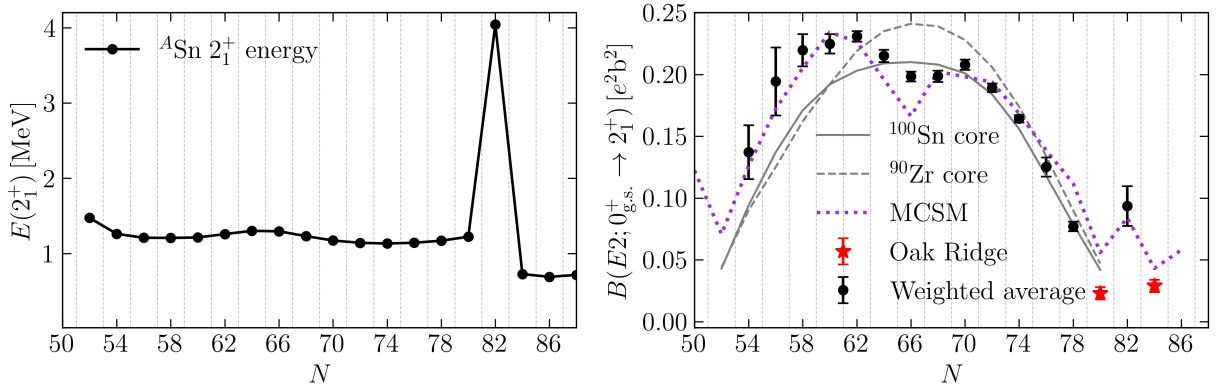


Figure 2.2. Left: Experimental excitation energies of the first 2^+ state along the Sn isotope chain from ^{100}Sn to ^{138}Sn . Values are taken from Ref. [68]. For ^{100}Sn , the excitation energy is still unknown. Between the doubly-magic nuclei, the excitation energy of the first 2^+ state shows a remarkable constant trend, with an enhanced value at doubly-magic ^{132}Sn and a subsequent lowering of the 2^+ energies. Right: Experimental $B(E2; 0_{g.s.}^+ \rightarrow 2_1^+)$ values for Sn isotopes (top). MCSM calculations adapted from Ref. [54] and LSSM from Ref. [69] with a ^{100}Sn and ^{90}Zr core are added. Experimental values used for the weighted average are from Ref. [70–86]. Only values published in this century have been considered. Values for $^{130,134}\text{Sn}$ from conference proceedings [57–59] and ^{132}Sn from Ref. [23].

explains the ~ 1.2 MeV excitation energy of the first 2^+ state throughout the chain. Unlike the excitation energies, the $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values exhibit a more complex behavior. The weighted mean of recent experimental results are shown in the right plot of Fig. 2.2 together with different theoretical approaches with (MCSM [54]) and without (marked as ^{100}Sn and ^{90}Zr core [69]) excitation above the $Z = 50$ and $N = 82$ shell closure. A symmetric, parabolic trend between the two doubly-magic nuclei is expected, if seniority was conserved. This is indeed shown by the theory excluding core excitation. However, instead, enhanced $B(E2)$ values are observed in the lighter Sn isotopes, and a sudden drop at $N = 66$ is still discussed [73, 78]. Recent large-scale MCSM calculations [54] using an extended model space with eight orbits for protons and neutrons reproduce this trend well. These calculations explain the long-standing puzzle of the enhanced collectivity in neutron-deficient Sn isotopes [74–76]. For $N > 66$, the $B(E2)$ values revert to a parabolic trend, except for ^{132}Sn , which is again described by LSSM and MCSM calculations.

2.1.1. The Special Role of ^{132}Sn and its even-even Neighbors

Since the early 1980s, the nucleus ^{132}Sn has been regarded as one of the most promising candidates for exhibiting pronounced shell closures at both proton and neutron numbers. In fact, as early as 1981, J. Blomqvist stated that “double-magic ^{132}Sn exhibits the strongest shell closure of any observed nucleus” [90] despite the fact, that “nuclear reaction studies are basically excluded“. This bold conclusion was drawn at a time when detailed nuclear reaction studies around ^{132}Sn were essentially unreachable due to the short half-life and low production rates of isotopes in this region. Nuclei around ^{132}Sn are furthermore challenging to access, since evaluated production mechanisms such as fusion-evaporation reactions predominantly populate proton-rich rather than neutron-rich nuclei.

45 years later, ^{132}Sn remains one of the most interesting yet experimentally challenging nuclei. It is the heaviest known unstable doubly-magic nucleus and sits at the intersection of several active research frontiers, including the evolution of shell structure far from stability, the r-process in nuclear astrophysics, the role of tensor forces in shell evolution, and the benchmarking of modern shell-model interactions. Despite impressive progress in radioactive beam production and detection techniques, fundamental properties of ^{132}Sn and its surrounding nuclei — such as the structure of excited states, transition strengths, and static moments — are still only partially known. A first milestone for proofing the magic character of ^{132}Sn came from single-particle transfer reactions on ^{133}Sn [55], which reveal the purity of neutron single-particle states beyond the shell closures. Only in 2018, a modern

Coulomb excitation experiment using the Miniball [91] array at the HIE-ISOLDE [92, 93] facility provided the first direct experimental confirmation of the spherical shape and doubly-magic nature of ^{132}Sn [23].

The enhanced $B(E2)$ value in ^{132}Sn was accurately determined here, providing direct evidence for its double magicity. Due to the closed $Z = 50$ and $N = 82$ shells, both proton and neutron excitations are hindered, resulting in a sudden increase of the 2_1^+ energy. The 2_1^+ state exhibits a mixed proton-neutron character, which explains the locally enhanced $B(E2)$ value. The influence of proton excitation leads to the enhanced transition probability. Early measurements published in a conference proceeding [59] yielded values compatible with the more precise results of Ref. [23].

In contrast to ^{132}Sn , the neighboring isotopes $^{130,134}\text{Sn}$ remain even more elusive. Despite their proximity to the doubly-magic nucleus and their relevance for understanding the evolution of shell structure toward neutron-rich systems, no precise transition matrix elements or electromagnetic observables of first excited states have been published to date. The only available information come from conference proceedings, which highlight the experimental challenges associated with this nucleus [57–60]. Measuring transition strengths in ^{130}Sn and ^{134}Sn allows for direct comparison with the doubly-magic reference nucleus ^{132}Sn . In particular, deviations from expected shell-model systematics may reveal subtle effects such as core excitations, proton-neutron correlations, or the onset of collectivity. Thus, ^{130}Sn and especially ^{134}Sn serve as key nuclei for testing the predictive power of modern nuclear structure models and for testing different approaches in terms of the neutron–neutron interaction and the choice of effective charges. This benchmark role is particularly important for ^{134}Sn : as experiments become increasingly difficult toward the neutron-rich frontier, well-validated theoretical descriptions are essential for extrapolating to even more exotic nuclei, and therefore require stringent experimental reference points.

Another interesting feature of ^{132}Sn is its remarkable structural analogy to the other well-established doubly-magic nucleus, ^{208}Pb . Both nuclei possess the same number of valence nucleons beyond the major closed shells at $Z = 50$, $N = 82$ for ^{132}Sn and $Z = 82$, $N = 126$ for ^{208}Pb . Consequently, their low-lying level structures exhibit striking similarities: both show a first excited 2_1^+ state at high energy, weak collective features, and clear single-particle character in nearby odd-mass isotopes [55]. These analogies offer a unique testing ground for nuclear models, especially shell-model interactions that aim to describe the evolution of structure toward the neutron drip line.

While this summary does not claim completeness, it highlights the progress made in exploring the Sn isotopes over the last two decades.

2.2. Previous Experiments in ^{130}Sn

As discussed in the previous section, the region surrounding the doubly-magic nucleus ^{132}Sn has long attracted considerable interest due to its significance for testing nuclear-structure models. However, its location far from the valley of stability on the neutron-rich side of the nuclear chart makes it difficult to access with conventional techniques such as single-nucleon transfer or fusion-evaporation reactions. Consequently, experimental data on neighboring nuclei such as ^{130}Sn remain scarce. The first experimental observations in ^{130}Sn date back to the 1970s. Fogelberg *et al.* [94] identified a long-lived 10_1^+ isomeric state using β -delayed γ -ray spectroscopy and determined a half-life of $t_{1/2} = 1.61(15) \mu\text{s}$. Decades later, Kameda *et al.* [95] studied isomeric decays following the in-flight fission of ^{238}U at 345 MeV/u at RIKEN, obtaining a shorter half-life of $t_{1/2} = 1.14_{-0.21}^{+0.27} \mu\text{s}$. More recently, Dunlop *et al.* [96] performed β -decay and β -delayed neutron spectroscopy of ^{131}In at TRIUMF-ISAC using the GRIFFIN spectrometer and reported a refined value of $t_{1/2} = 1.52(9) \mu\text{s}$.

The first measurements of electromagnetic transition strengths of low-lying states in ^{130}Sn were conducted in the early 2000s at the Holifield Radioactive Ion Beam Facility (HRIBF) at Oak Ridge National Laboratory (ORNL). There, a radioactive ^{130}Sn beam produced via the Isotope Separation On-Line (ISOL) method was directed onto a carbon target. Scattered carbon recoils were detected using the HyBall array, while emitted γ rays were recorded with the CLARION spectrometer. The resulting spectra are shown in Fig. 2.3 for three different beams: a pure ^{130}Sn and ^{128}Sn beam, and a $A = 128$ cocktail beam. From this measurement, a preliminary $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ value of $0.023(5) e^2 \text{b}^2$ was extracted [58, 60]

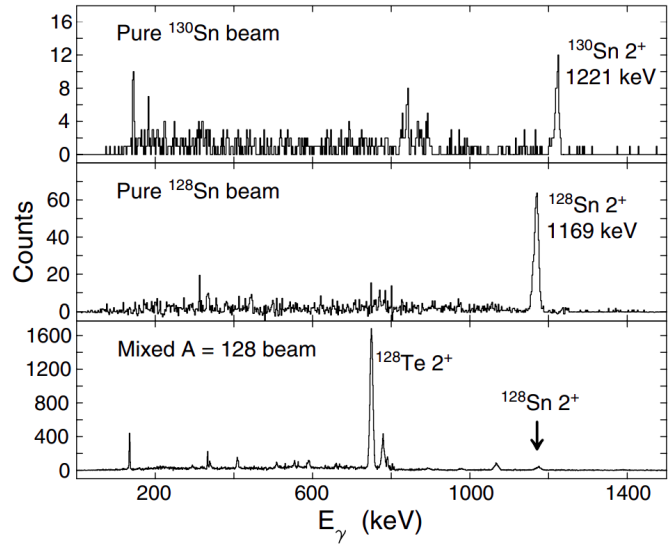


Figure 2.3. Gamma-ray spectra from inelastic scattering of ^{130}Sn on carbon. Adopted from Ref. [57] with kind permission from Elsevier.

based on 66(10) counts in the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition, significantly lower than the $0.055 e^2 b^2$ given by MCSM calculations discussed in the previous section. The technique used to extract these values — inelastic scattering on a ^{12}C target — is described in detail in Ref. [97]. Importantly, this method is highly sensitive to the assumed target thickness, which directly affects the normalization of the extracted matrix elements. For example, in ^{132}Te , a previously underestimated target thickness led to a revision of the $B(E2)$ value from $0.172(17) e^2 b^2$ to $0.216(22) e^2 b^2$ [98]. Furthermore, the exact beam composition is crucial in determining unknown transition strength values, since the obtained γ -ray yields are compared with the number of scattered carbon nuclei, which reacted with the beam of interest.

To date, no direct measurement of $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ in ^{130}Sn has been published in peer-reviewed literature. This makes ^{130}Sn a prime candidate for dedicated Coulomb excitation studies, which would allow one to validate theoretical models in the region just below ^{132}Sn and close the experimental gap between the mid-shell Sn nuclei and doubly-magic ^{132}Sn .

In this context, it is relevant to note that approaches which effectively reduce the neutron effective charge can reproduce the older experimental systematics: Kumar *et al.* [99] reported lifetime measurements in the isobaric nucleus ^{130}Te and the $N = 82$ isotone ^{132}Te , and extended their analysis to ^{130}Sn using shell-model calculations based on the *sn100pn* interaction [100]. Their theoretical results show good agreement with earlier values from ORNL by Radford *et al.* [58] by reducing the effective neutron charge, but deviate significantly from predictions of MCSM calculations by Togashi *et al.* [54].

A dedicated remeasurement of $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ in ^{130}Sn is therefore crucial to discriminate between competing theoretical descriptions and to establish a reliable experimental value for electromagnetic collectivity just below ^{132}Sn .

2.2.1. The isomeric 7_1^- and 5_1^- States in ^{130}Sn

In addition to the 10_1^+ isomers discussed earlier, ^{130}Sn also exhibits a long-lived 7_1^- isomeric state and 5_1^- state located above it. The level scheme, including these isomeric states, is shown in Fig. 2.4. These states provide valuable insight into the structure of neutron-rich nuclei near the doubly-magic ^{132}Sn , since they are expected to be nearly pure two neutron-hole states [101]. While the 7_1^- state decays only via β decay into ^{130}Sb , the 5_1^- state decays via an $E2$ γ -ray transition and an $E1$ transition into the 7_1^- and 4_1^+ state respectively.

The 5_1^- state was first investigated by Fogelberg *et al.* [94], who determined its lifetime to

be $t_{1/2} = 52(3)$ ns using β -delayed γ -ray spectroscopy.

A later study by Pietri *et al.* [102] also addressed this state, though the results of 63(4) ns differed slightly. Recently, a new lifetime measurement of the 5_1^- state was published by Piau *et al.* [103], reporting a refined value of $t_{1/2} = 50(5)$ ns. This result is in tension with earlier findings by Pietri *et al.*, emphasizing the need for further clarification. The quadrupole moment of the 7_1^- isomer has been a subject of active study, particularly regarding its deformation properties. Gorges *et al.* [104] performed laser-spectroscopic measurements of isotope shifts at ISOLDE to extract charge radii and quadrupole moments. However, the extracted spectroscopic quadrupole moment for the 7^- state was ambiguous, yielding two possible solutions:

- $Q_s(7^-) = -0.426(13)_{\text{stat}}(18)_{\text{sys}}$ eb,
- $Q_s(7^-) = +0.452(4)_{\text{stat}}(19)_{\text{sys}}$ eb.

Earlier this century, Le Blanc *et al.* [101] determined a value of $Q_s(7_1^-) = -0.36(11)$ eb, consistent only with the negative solution reported by Gorges *et al.*. This result suggests a moderate oblate deformation of the 7_1^- state. Based on these findings, Le Blanc and collaborators proposed a $\nu(d_{3/2}^{-1}h_{11/2}^{-1})$ two-hole configuration as the dominant structure of the 7_1^- isomer.

Similar long-lived isomeric 7_1^- and 5_1^- states have also been observed in neighboring $N = 80$ isotones with $Z < 50$, such as ^{128}Cd [105] and ^{126}Pd [106, 107], suggesting a systematic appearance of such configurations in this region of the nuclear chart. More recently, new isomeric states have been identified in ^{130}Cd [108]. The observed lifetimes of these states are valuable for constraining shell-model predictions and refining the parameters used in the description of electromagnetic transition strengths.

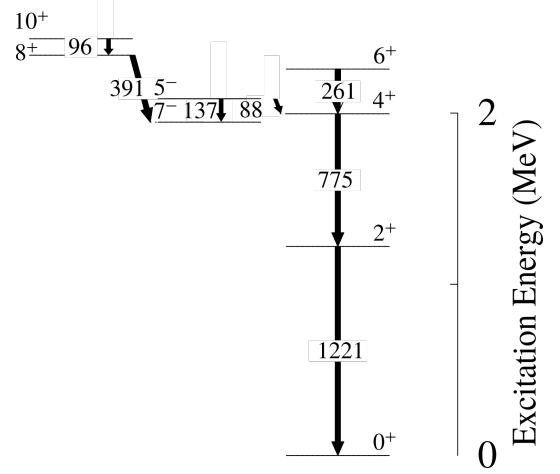


Figure 2.4. Levelscheme of ^{130}Sn . The isomeric 7_1^- state decays via β decay. Transition energies are given in keV.

2.3. Previous experiments in ^{134}Sn

For the even heavier isotope ^{134}Sn , the available experimental information becomes increasingly scarce. The first excited states in this nucleus were identified in 1997 by Zhang *et al.* [109] and were later refined by Korgul *et al.* [110]. In the latter work, the half-life of the long-lived 6_1^+ state was also determined to be 80(15) ns. Precise β -decay half-lives of ^{134}Sn and neighboring nuclei were measured at RIKEN in 2015 [111], while charge radii in this region were obtained at ISOLDE, CERN [104].

In 2021, the first lifetimes of states below the 6_1^+ level were published [112], measured with the fast-timing method at the ISOLDE Decay Station. This work yielded the first experimental value for the lifetime of the 4_1^+ state and the most precise value to date for the 6_1^+ state. They also give an estimate for the half-life of the 2^+ state with $t_{1/2}(2_1^+) = 53(30)$ ps.

A comparison of their 4_1^+ result with various theoretical approaches showed that none of them is reproducing the experimental $B(E2; 4_1^+ \rightarrow 2_1^+)$ value satisfactorily. The only further available information on the electromagnetic properties of the first excited 2^+ state comes from the conference proceedings introduced earlier. Unlike their ^{130}Sn experiment, in this Coulomb-excitation measurement γ rays were detected using an array of 150 BaF_2 crystals, significantly increasing the total detection

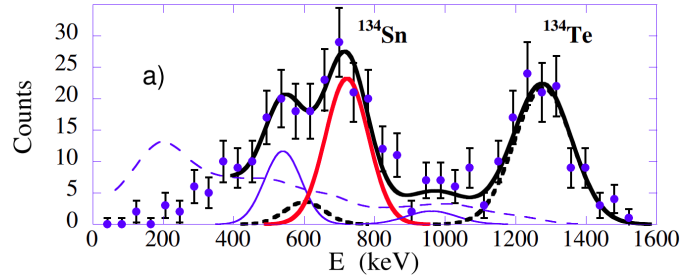


Figure 2.5. Spectra obtained from Ref. [59] with kind permission from Elsevier. The decay of the 725 keV $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition following Coulomb excitation is visible, along with contaminating transitions from ^{134}Ba and ^{134}Te . The absence of an HPGe array in this experiment resulted in reduced energy resolution.

efficiency but at the cost of energy resolution compared to the setup used for the ^{130}Sn experiment. The experimental spectrum is shown in Fig. 2.5. Although the authors report the observation of the 2_1^+ de-excitation in ^{134}Sn , the limited resolution complicates the spectrum considerably, as the beam contained substantial fractions of ^{134}Ba and ^{134}Te . Transitions from these nuclei overlap with the region of interest and strongly affect the spectral shape. Based on these data, the group extracted a value of $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.029(5) e^2 b^2$, which, similar to the value in ^{130}Sn , deviates significantly from the predictions of MCSM calculations.

To bring the previously discussed information together: the nuclei around the ^{132}Sn

shell closure remain experimentally poorly explored. Transition strengths of the lowest excited states have so far only been reported in conference proceedings, and these values show significant deviations from the most recent state-of-the-art MCSM calculations. A comparison between theoretical and experimental results is shown in Fig. 2.6. The figure displays the relative deviations of various theoretical predictions from the experimental values, with the experimental uncertainty indicated by the grey shaded band. Shown are results from MCSM, and LSSM [69] calculations.

While the LSSM approaches show systematic deviations from the experiment, the MCSM generally reproduces the measured values well. For ^{130}Sn , all models overpredict the experimental $B(E2)$ strength by roughly a factor of two. The doubly magic nucleus ^{132}Sn is again well described, but for ^{134}Sn the MCSM result deviates significantly from the experiment, while the LSSM calculations is limited to nuclei below ^{132}Sn . Understanding the origin of this discrepancy is an open question in nuclear structure physics, with implications for modelling the nuclear physics input of astrophysical r-process calculations.

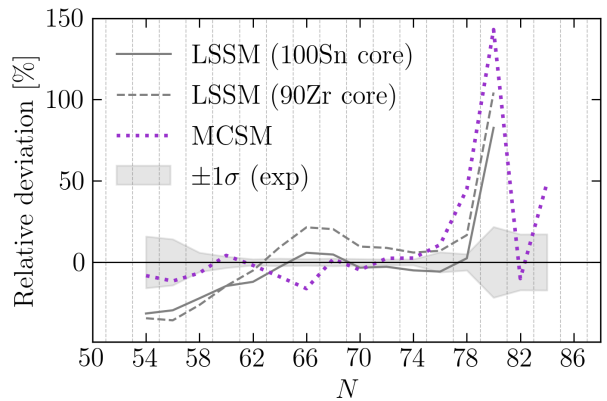


Figure 2.6. Deviations of calculated $B(E2)$ values along the Sn isotope chain relative to the experimental results. Except for ^{130}Sn and ^{134}Sn , all isotopes are well reproduced by the theoretical models, in particular by the MCSM.

3. The Radioactive Ion Beam Facility ISOLDE, CERN

Understanding the structure of atomic nuclei, particularly those far from the valley of stability, remains one of the central challenges in modern nuclear physics. While ground-state properties of more than 3400 bound nuclei have been determined, only 253 are stable, and merely 33 additional radioactive nuclei occur naturally on Earth. Consequently, over 90 % of all known nuclides are short-lived and must be produced artificially through nuclear reactions [113]. These systems are of key interest not only for fundamental studies of nuclear structure, but also for astrophysical processes such as the rapid neutron-capture process (*r*-process) and the proton-capture process (*p*-process), which govern the nucleosynthesis of the heavy elements [114]. To access such exotic nuclei, various nuclear reactions are employed. Neutron-rich isotopes are typically produced via fission of heavy elements, whereas proton-rich nuclei are often formed in fusion–evaporation reactions. One approach relies on *in situ* measurements, in which a stable beam is directed onto a stable target, and the reaction products are studied directly. Although this method benefits from comparatively large production cross sections, it is constrained by competing reaction channels and limited beam–target combinations.

A powerful complementary technique is the use of post-accelerated radioactive ion beams (RIBs). Here, exotic nuclei are first produced in a primary target, ionized, mass-separated, and subsequently delivered to an experimental station. This two-step process provides access to isotopically pure beams, reduces background contributions, and allows tuning of the beam energy to probe a broad range of nuclear observables. The method is technically demanding due to often low production yields and short half-lives of the species of interest. Efficient release from the target material, selective ionization, and high-resolution mass separation are essential to obtain beams of sufficient purity and intensity. Two principal RIB production techniques are currently in use: the In-Flight (IF) and the Isotope Separation On-Line (ISOL) methods [115]. The IF technique relies on high-energy reactions to produce fast-moving nuclei, which are separated in flight using magnetic spectrometers. In contrast, the ISOL method — implemented at ISOLDE (CERN) [93] — produces low- to medium-



Figure 3.1. Schematic layout of the ISOLDE hall, highlighting the GPS and HRS mass separators connected to two separate target stations. Components discussed in this thesis are marked in orange-red. Adapted and modified from Ref. [118] with kind permission of CERN.

energy beams with high selectivity and often superior intensities, making it particularly suitable for detailed nuclear-structure investigations.

At ISOLDE, radioactive isotopes are generated by bombarding thick target materials (e.g., UC_x) with 1.4 GeV protons from the Proton Synchrotron Booster (PSB). The reaction products diffuse out of the heated target material, are ionized, mass-separated, and subsequently delivered either at low energy (≈ 60 keV) or post-accelerated by the HIE-ISOLDE linear accelerator [92, 116]. The efficiency of the ISOL technique depends on several factors: (i) the chemical properties of the element, (ii) diffusion and effusion times within the target, (iii) the applied ionization method (surface, plasma, or laser ionization), and (iv) the performance of the mass separators. In practice, the ISOL technique is well suited for nuclei with half-lives exceeding approximately 4 ms [117]. A schematic layout of the ISOLDE hall is shown in Fig. 3.1. The individual parts highlighted in red will be discussed in the following section.

3.1. Target Design and Materials

The target system is central to the production of radioactive ion beams (RIBs). It must withstand high proton-beam intensities, efficiently produce isotopes, and enable rapid isotope release from the target material. ISOLDE receives 1.4 GeV protons from the PSB with intensities of up to 3.2×10^{13} protons per pulse (4 μ A peak), limited to 2 μ A average due to safety regulations. Two targets, each connected to an independent mass separator (GPS or HRS), can operate simultaneously.

A typical ISOLDE target consists of a cylindrical container of approximately 20 cm length and 2 cm diameter, made from graphite or tantalum. The interior is filled with material in the form of powders, pellets, molten metals, or pressed fibers. Proton-induced reactions (spallation, fragmentation, fission) generate a wide variety of isotopes inside the target matrix. For neutron-rich isotopes, UC_x is widely used due to its high fission yield and fast release properties, and it was employed in all experiments conducted for this work.

3.2. Ion Sources at ISOLDE

ISOLDE employs several types of ion sources, each optimized for specific elements based on their physical and chemical characteristics. The choice of ion source is essential for achieving high ionization efficiency while minimizing isobaric and isomeric contamination.

The most basic form of ionization is provided by the *surface ion source*. Surface ionization relies on the interaction between hot metal surfaces (typically tungsten or tantalum) and atoms with low ionization potentials. Upon contact, electrons are transferred from the atom to the surface, resulting in efficient ionization for alkali and alkaline-earth elements. However, this method is ineffective for elements such as tin, which has an ionization potential of approximately 7.3 eV [119].

For improved isotopic and isobaric selectivity, *resonance ionization* can be employed. The Resonance Ionization Laser Ion Source (RILIS) uses precisely tuned lasers to excite and ionize atoms through multi-step resonant transitions. This technique enables element-selective ionization and is well suited for elements with higher ionization potentials, including Sn, In, and Hg [120]. Although RILIS provides high beam purity for many elements, it is less effective for molecular species, which may dissociate under laser irradiation.

A third option available at ISOLDE is the *plasma ion source*. The Forced Electron Beam Induced Arc Discharge (FEBIAD) ion source is a central component of ISOLDE's plasma-

based ionization scheme and is particularly effective for reactive and refractory elements. It generates a confined plasma by accelerating electrons between a heated cathode and an anode. Collisions between electrons and neutral atoms or molecules lead to ionization in the plasma volume. The FEBIAD source is especially advantageous for ionizing molecular beams, such as tin sulfide (SnS), which provide high beam purity by chemically binding tin to sulfur and thereby suppressing isobaric contaminants such as cesium and xenon [121].

Before the establishment of SnS^+ beam production, RILIS provided the highest tin beam purity [122]. RILIS suppresses heavier isobars such as Sb, Te, I, and Xe, whose ionization potentials exceed 8.6 eV. However, surface-ionizable elements with low ionization potentials, such as Cs (3.89 eV), Ba (5.21 eV), and In (5.78 eV), are also efficiently ionized in the hot Nb or W cavity used by RILIS, resulting in significant background. This contamination posed challenges for decay spectroscopy of very neutron-rich Sn isotopes. For example, while β -delayed neutron emission allowed the observation of ^{137}Sn , precision studies of ^{136}Sn were hindered by background from the long-lived isomer $^{136\text{m}}\text{Cs}$ [123].

To address these limitations, an alternative approach using molecular ion beams was developed. The extraction of SnS^+ proved highly efficient and selective [124]. SnS molecules form reliably across a wide range of sulfur-containing compounds and target temperatures. A $\text{SnS}^+ / (\text{Sn}^+ + \text{SnS}^+)$ ratio of up to 58 % was achieved at 1700 K, decreasing only slightly at higher temperatures. In addition, the desorption time of SnS is approximately two orders of magnitude shorter than that of atomic Sn, significantly improving release efficiency, particularly for short-lived isotopes [121].

The FEBIAD MK5 ion source used for SnS^+ ionization employs noble gases (e.g., Ar or Xe) to initiate and sustain the plasma. Typical operating parameters include electron-acceleration voltages of 100–200 V and pressures between 10^{-4} and 10^{-5} mbar. The plasma temperature is adjusted to satisfy $k_{\text{B}}T_{\text{p}} \geq W_{\text{SnS}}$, where W_{SnS} is the ionization energy of SnS . An additional transverse magnetic field increases the residence time of molecules in the plasma, enhancing ionization efficiency. These features enable FEBIAD sources to reach ionization efficiencies of up to 80 % in the 7–10 eV ionization-potential range and to deliver beam currents of 1–10 μA , making them a robust option for producing high-purity beams of refractory and short-lived isotopes [121, 124].

3.2.1. Mass Separation at ISOLDE

Once ionized, the singly charged ions are extracted at energies between 30 and 60 keV and directed to one of the two mass separators: the General Purpose Separator (GPS) or the High-Resolution Separator (HRS). These devices separate ions according to their mass-to-charge ratio (A/q), ensuring that only the isotopes of interest are delivered to the experimental stations.

The GPS is designed for versatility and high throughput. It provides moderate mass resolution and is typically used when beam purity is less critical or when the isotope of interest is produced with high yield. Its compact design and rapid tuning capability make it well suited for exploratory measurements and routine beam delivery.

In contrast, the HRS offers significantly higher mass resolution of up to $M/\Delta M \approx 15\,000$ when higher-order optical corrections are applied. Its optical system consists of two magnetic separation stages interleaved with electrostatic quadrupoles and multipole lenses for the correction of aberrations. The use of double-focusing dipole magnets enables precise selection of the desired isotope while efficiently suppressing isobaric contaminants. The HRS is therefore essential for high-precision applications such as collinear laser spectroscopy and decay spectroscopy, where beam purity is critical. The availability of both separators—each connected to an independent target station—allows ISOLDE to operate multiple experiments in parallel while maintaining flexibility and high performance. This dual-separator configuration is a key element of ISOLDE’s capability to deliver more than 1300 isotopes from over 70 elements to a broad international user community. During the experimental campaign for this work, both separators and their associated target stations were employed on two separate occasions.

3.3. High-Energy Beam Preparation at ISOLDE

After mass separation, the ionized beams are injected into the Penning trap REXTRAP [125], where they are accumulated, cooled, and subsequently extracted as compact ion bunches. For molecular beams, dissociation typically occurs either within REXTRAP or during the following stages of beam preparation, ensuring that only the desired atomic ions continue along the beam line. The cooled and bunched ions are then transferred to the electron-beam ion source REXEBIS [126], where they are confined for several tens of milliseconds and exposed to an intense electron beam. In this environment, successive electron-impact ionization increases their charge state in a controlled manner. This charge-breeding process

is essential for matching the A/q requirements of the downstream linear accelerator. Once the desired charge state distribution is achieved, the ions are extracted from REXEBIS and injected into the HIE-ISOLDE accelerator.

The HIE-ISOLDE post-acceleration system consists of a normal-conducting RFQ followed by a sequence of superconducting quarter-wave resonator (QWR) cavities arranged in cryomodules. In total, the linac comprises four superconducting cryomodules, each containing five superconducting cavities that provide continuous acceleration of the charge-bred beams up to energies of 10 MeV/u [92, 116, 117]. A final A/q selection using a dipole magnet downstream of the linac ensures that only ions with the correct charge state are transmitted to the experimental area.

At the high-energy end of ISOLDE, three dedicated experimental stations are available. The Scattering Experiments Chamber (SEC) offers a flexible environment for a wide range of reaction studies, including elastic and inelastic scattering as well as transfer reactions on light and heavy targets. The ISOLDE Solenoidal Spectrometer (ISS) is optimized for precision measurements of transfer reactions and inelastic-scattering processes, exploiting a superconducting solenoid to provide high-resolution particle tracking far from stability. The high-purity germanium detector array Miniball, used in the experiments presented in this thesis, enables high-efficiency γ -ray spectroscopy and is particularly suited for Coulomb excitation studies with low-intensity radioactive ion beams. Together, the REX/HIE-ISOLDE post-acceleration chain and the associated experimental stations constitute the core of ISOLDE's high-energy physics program, enabling detailed investigations of nuclear structure and reactions across a wide range of exotic isotopes. In the course of this work, all measurements were performed at the Miniball station, following beam preparation through the full REXTRAP–REXEBS–HIE-ISOLDE sequence.

3.4. The Miniball Spectrometer

The Miniball spectrometer is a state-of-the-art γ -ray detection array that was specifically developed to meet the demanding conditions of experiments with radioactive ion beams (RIBs), such as those delivered at ISOLDE. Studies of exotic, often short-lived nuclei place stringent requirements on the detection system: in addition to high photopeak efficiency and excellent energy resolution, advanced capabilities for event reconstruction are essential. A further key aspect of many RIB experiments is that the γ -emitting nuclei typically recoil with velocities of a few percent of the speed of light. As a consequence, Doppler shifts

and Doppler broadening can become dominant contributions to the observed line shapes. Accurate Doppler correction therefore relies on a precise determination of the emission angle and, accordingly, requires high angular granularity of the detector array. Miniball was designed to address these requirements in an optimized yet highly flexible configuration, enabling high-sensitivity and high-precision γ -ray spectroscopy with radioactive beams.

At its core, Miniball consists of an array of high-purity germanium (HPGe) detectors, each segmented into six individual volumes to precisely determine the interaction point of incoming γ rays. This segmentation, combined with modern digital signal processing, enables accurate Doppler correction of γ -ray energies—an essential feature when studying fast-moving ions in post-accelerated beams. A major strength of Miniball is its compact and modular design. Unlike larger, fixed detector arrays, Miniball can be reconfigured to match the geometry and physics goals of each experiment. This flexibility allows for optimized coverage around the target, improving detection efficiency and angular resolution, both of which are crucial for extracting transition matrix elements and electromagnetic transition strengths.

In combination with ancillary detectors—such as silicon detector arrays for charged-particle identification—Miniball enables powerful particle- γ coincidence measurements. This makes the system particularly effective for studying collective excitations via Coulomb excitation as well as single-particle structure through transfer reactions. Integration with the HIE-ISOLDE post-accelerator further enhances its capabilities, enabling experiments across a wide range of beam energies and reaction mechanisms. Miniball has played a key role in advancing our understanding of nuclear structure far from stability, yielding results ranging from precise measurements of $B(E2)$ transition strengths to the development of detailed level schemes of exotic nuclei. Notable achievements using Miniball include the first observation of pear-shaped nuclei [46] and the confirmation of the doubly magic nature of ^{132}Sn [23].

Refurbished Miniball detector array

During CERN’s second long shutdown (LS2), when no proton beam was available, the Miniball array was temporarily relocated to RIKEN to participate in the HIKARI campaign (see, e.g., Ref. [128]). In parallel, an extensive refurbishment and redesign of the detector system was undertaken at the University of Cologne with the goal of improving modularity, maintainability, and long-term performance.

One of the key upgrades was the development of a new encapsulation concept. In contrast to the previous welded design, the new housing employs a screw-based assembly, providing significantly easier access to the cold electronics inside the cryostat. This greatly improves maintainability and modularity, allowing individual components to be serviced or replaced with minimal effort. The cryostat system itself was also redesigned. Previously, each triple cryostat was connected to the liquid-nitrogen dewar via a long, mechanically sensitive neck. Although suitable for cooling, this configuration complicated mounting and introduced mechanical stress due to the substantial mass of the germanium crystals. The new cryostat features a more compact and robust coupling, improving both stability and handling. The refurbishment further included a complete upgrade of the preamplifier electronics. Inspired by developments for the AGATA HPGe array [129], the new electronics are more compact, modular, and thermally optimized. This results in improved signal quality,

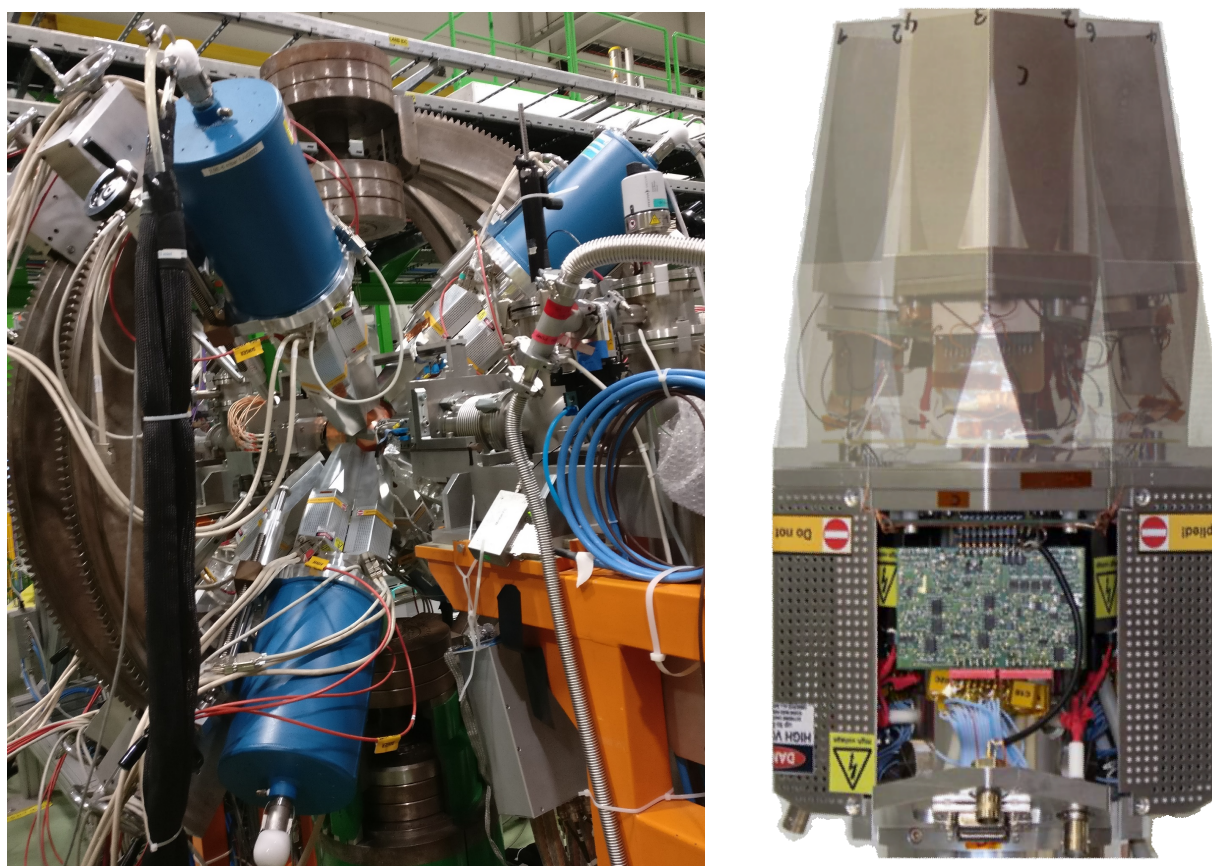


Figure 3.2. Left: Fully equipped Miniball spectrometer surrounding the target chamber at ISOLDE. Right: One of the triple cryostats together with the new digital electronics based on the AGATA-like preamplifier.

reduced noise, and easier integration—critical factors for high-precision spectroscopy with low-intensity radioactive ion beams. With these upgrades, the Miniball spectrometer has been transformed into a more versatile and user-friendly system, optimized for modern RIB experiments. The refurbished array was shipped back to ISOLDE in 2022. The fully equipped setup is shown on the left side of Fig. 3.2, while the right side illustrates one of the triple cryostats together with its new digital electronics. For all experiments performed in this campaign, the Miniball array was operated in combination with a double-sided silicon strip detector (DSSSD), a CD-shaped array used to measure scattered ejectiles and recoils and to determine their energies and emission angles. Using the velocity vectors reconstructed from the DSSSD and the interaction positions measured in the Miniball crystals, an accurate Doppler correction of the γ rays can be performed. The design of the CD-detector is shown in Fig. 3.3. A CD with 16 rings (p-strip) and 24 stripes (n-side) has been used in all experiments. In the first three experiments, two n-side strips have been connected to one to reduce charge sharing.

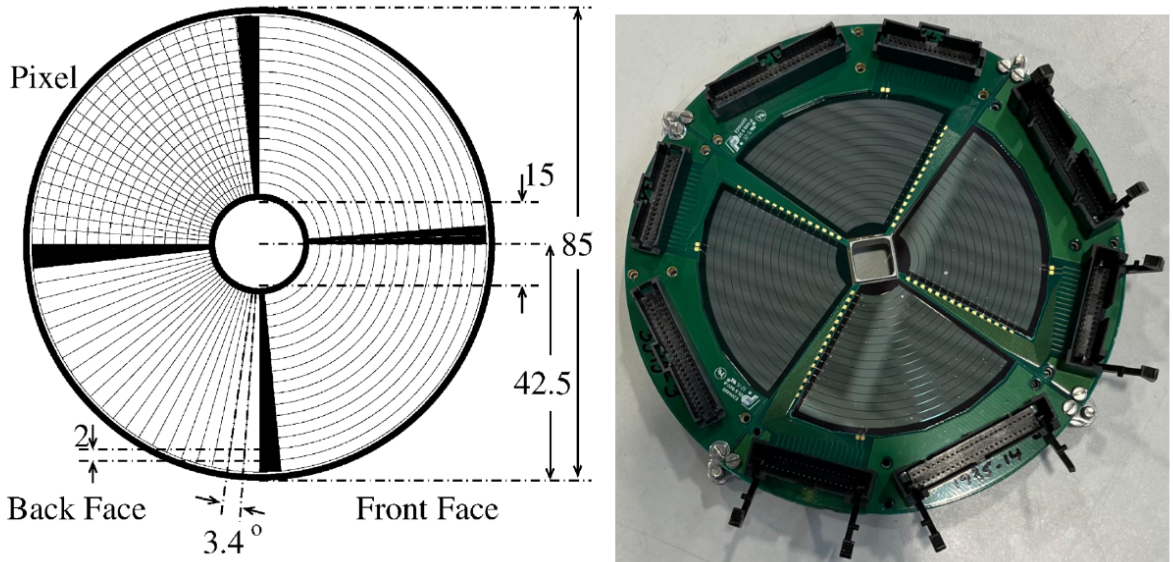


Figure 3.3. Left: Schematic drawing of a CD detector used in the present experiment (adapted from Ref. [127] with kind permission of Elsevier.). Right: Photograph of the CD detector as operated during the campaign.

4. Experimental Details

During the course of the thesis, four experiments were conducted at ISOLDE using the Miniball detector array. Here, the basic information about the performed experiments are explained. Only the differences between the experimental setups employed in the respective experiments is discussed, since all experiments used the Miniball detector array in combination with a DSSSD detector for particle- γ coincidence measurements.

4.1. First Beamtime: Coulex of ^{130}Sn in November 2022

The first experiment of the renewed Miniball campaign at ISOLDE following CERN's Long Shutdown 2 was carried out in November 2022. Later measurements in the campaign targeted gamma-electron coincidence spectroscopy. The presented experiment did not rely on such coincidences. Nevertheless, the SPEDE chamber equipped with the SPEDE electron spectrometer [130] was already installed for the subsequent program. To ensure safe Coulomb-excitation conditions, a beam energy of 4.4 MeV/u was chosen, impinging

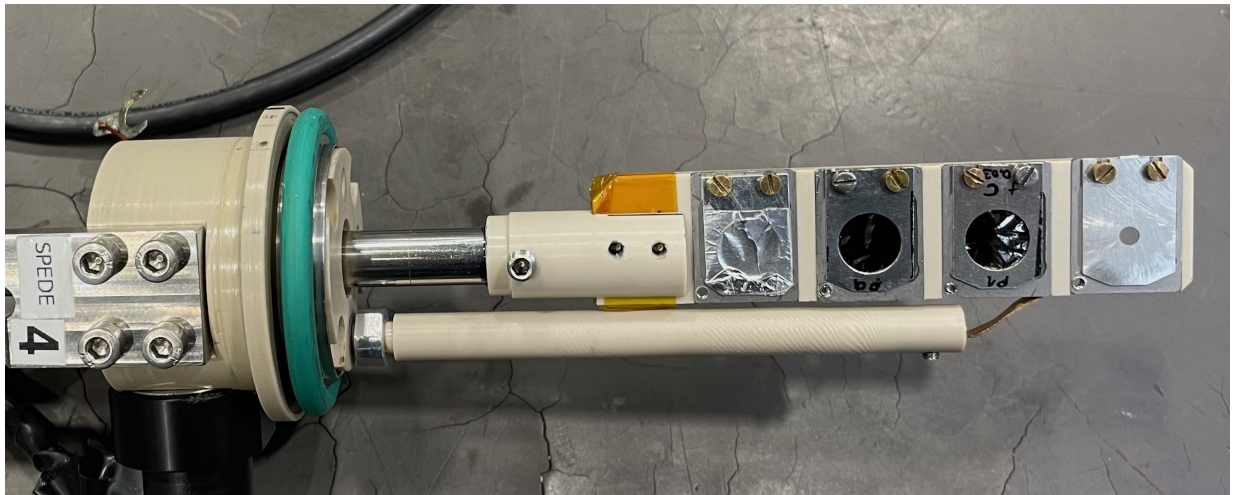


Figure 4.1. Target ladder used during the November 2022 experiment. From right to left, the ladder includes a 3 mm aperture for beam alignment and rate diagnostics, two ^{206}Pb targets of different thicknesses backed by a thin layer of natural carbon, and a beam catcher.

onto a ^{206}Pb target. The target thickness had been determined beforehand in Cologne via α -energy-loss measurements. The target ladder configuration is shown in Fig. 4.1. From right to left, it contained:

- Position 1: 3 mm aperture
- Position 2: $2.0(2) \text{ mg/cm}^2 \text{ } ^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$
- Position 3: $3.1(2) \text{ mg/cm}^2 \text{ } ^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$
- Position 4: Beam catcher

For particle detection, the DSSSD was mounted about 30 mm downstream of the target. The beam time commenced on 3rd November following calibration runs with a ^{152}Eu source and the stable-beam reactions $^{22}\text{Ne}(\text{d,p})^{23}\text{Ne}$ and $^{22}\text{Ne}(\text{d,n})^{23}\text{Na}$, which were used to determine the Miniball detector geometry. Miniball was operated with seven of its eight triple cryostats. The beam was expected to contain a significant population of the long-lived isomeric $^{130}\text{Sn}_{7-}$. To monitor the population ratio between the ground state and the isomer, a single HPGe detector was positioned on the long-energy side of the accelerator.

Once the radioactive beam was delivered, it became apparent that the count rates were too high for the data-acquisition system: the Miniball DAQ suffered from an almost 100% dead time. The beam was therefore attenuated by a factor of 100, which reduced the overall rates but did not sufficiently restore the DAQ performance. This was the first in-beam operation of the FEBEX modules from GSI with a radioactive ion beam, and the system had been installed only a few days before the experiment. During the measurement, it was not possible to directly monitor the individual counting rates in the HPGe detectors. Instead, the acquisition was observed to be dominated by dead time, with long periods in which pile-up was so severe that essentially no events were written to disk. A post-experiment assessment of the FEBEX4 system, including an estimate of the trigger rates, showed that such conditions already occurred at a global trigger rate of the order of 5 kHz. Thus, even at what would normally be considered low counting rates for Miniball experiments, the combination of long $\sim 50 \mu\text{s}$ signal tails and the limitations of the online processing led to very high pile-up probabilities and large effective dead times.

Operational interruptions related to the linac and PS Booster as well as a failure of a magnet further reduced the available beam time by approximately 2.5 days. The experiment concluded on 8th November 2022. Out of the allocated 120 h, about 75 h were ultimately delivered as radioactive ion beam. Additional beam time was therefore granted, and a

second experiment was conducted in October 2023.

4.2. Second Beamtime: Coulex of ^{130}Sn in October 2023

In the second beam time, the plunger (Coulex) chamber was used instead of the SPEDE chamber. For the purpose of this experiment, this mainly affects the upstream part of the setup, since no SPEDE detector was installed. Furthermore, a different target wheel was mounted, providing six possible target positions. The target wheel is shown on the left in Figure 4.2, while the plunger chamber in front of half of the Miniball spectrometer is shown on the right. The following targets were mounted on the wheel, where position “6” corresponds to the slot at the top, and the numbering increases clockwise:

- Position 1: $1.8 \text{ mg/cm}^2 \text{ }^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$
- Position 2: 3 mm aperture
- Position 3: 5 mm aperture
- Position 4: $4.2 \text{ mg/cm}^2 \text{ }^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$
- Position 5: $3.8 \text{ mg/cm}^2 \text{ }^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$
- Position 6: $2.2 \text{ mg/cm}^2 \text{ }^{206}\text{Pb} + 0.03 \text{ mg/cm}^2 \text{ nat. C}$

The experiment was scheduled directly after an experiment employing an ^{132}Sn beam. Owing to the high quality of the data obtained in that experiment, one additional day of beam time was allocated to our measurement from this proposal [131]. The long-lived granddaughter nucleus of ^{132}Sn — ^{132}Te — with a half-life of 3.2 days, was present as background activity during our run. The DAQ system employed for the 2022 experiment operated with a common trigger and a shared dead time, which led to severe pile-up and substantial overall dead time losses. This was changed for this experiment, where every channel is read out individually, or stores information into a buffer until it gets a readout signal, resulting in a dead time per channel. Using this improvement, a count rate of 10 kHz per crystal, as observed in the present experiment, was easily manageable for the DAQ.

Energy and efficiency calibration measurements with various radioactive sources were carried out prior to the ^{132}Sn experiment and repeated after our ^{130}Sn beam time. Position calibration measurements were also performed before the ^{132}Sn beam time, and the Miniball detector positions were not modified thereafter. Again, seven of the eight triple clusters were used for the experiment. The RIB experiment for ^{130}Sn started on 11th October

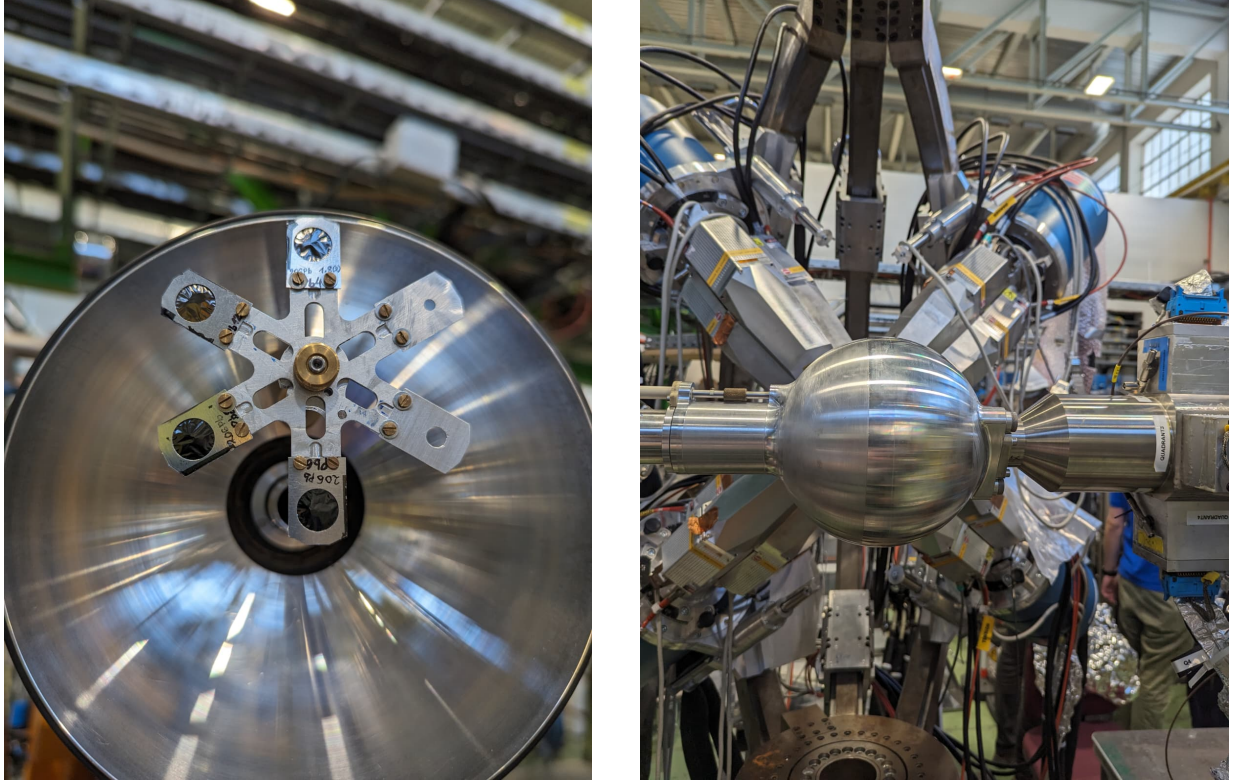


Figure 4.2. Target wheel (left) equipped with the ^{206}Pb targets and apertures of different diameters. The right-hand side shows the plunger (Coulex) chamber mounted in front of half of the Miniball spectrometer.

2023 and ended on 16th October 2023. For the main physics run, the 4.2 mg/cm^2 ^{206}Pb target (position 4) was used to increase the statistics compared to the first experiment, at the expense of a reduced accuracy of the Doppler correction and a slightly degraded CD detector resolution due to the increased energy loss and straggling in the thicker target. In total, the experiment lasted for 108 h, with short interruptions caused by maintenance work and occasional losses of the proton beam roughly 45 h after the start of the run. Accounting for these interruptions, the effective in-beam time amounts to about 100 h.

4.3. Third Beamtime: Coulex of ^{134}Sn in July 2024

Following the measurements with the radioactive ^{130}Sn beam, the focus of the experimental program was shifted to ^{134}Sn . With four additional neutrons compared to ^{130}Sn , this nucleus lies further from stability, and the expected beam intensities are considerably lower, by about a factor of 150 according to the ISOLDE yield database. The experiment was again prepared using measurements with stable beams. In addition to runs with a CD_2 target, an auxiliary measurement with a ^{196}Pt target was performed to obtain an additional constraint on the azimuthal (ϕ) angles of the detectors. These data also provided information on the distance between the CD detector and the target and were used for the calibration of the

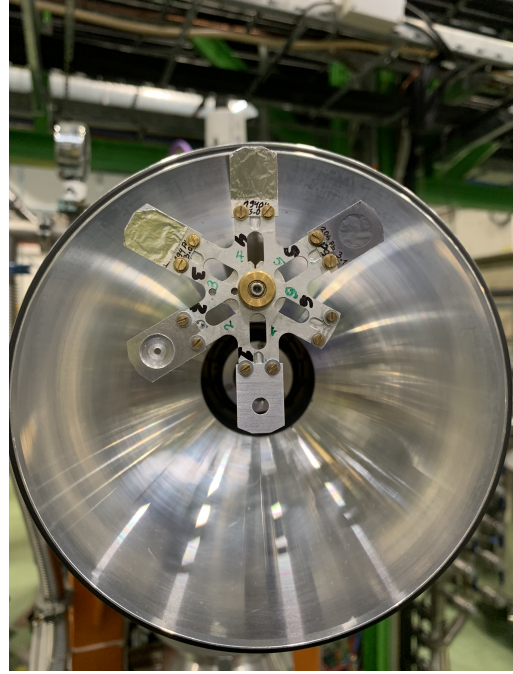


Figure 4.3. Target wheel used during the ^{134}Sn experiment in 2025. The main part of the measurement was carried out with ^{194}Pt targets, while an additional ^{206}Pb target was also mounted.

CD detector. Since these preparatory measurements could already be carried out in June 2024, while the ^{134}Sn beam time was scheduled to start only on 19th July 2024 as the first experiment with Miniball in 2024, sufficient time was available to complete the calibrations, such that the experiment was optimally prepared from the analysis perspective. On the day before the scheduled start of the experiment, first tests with the primary UC_x production target were performed. During these tests, it was found that the heating element for the sulfur supply did not remain stable under proton irradiation, which prevented the formation of the desired molecular beam. For this target unit, no second heating element and no backup target had been installed. As a consequence, despite the extensive preparatory work, the planned ^{134}Sn Coulomb-excitation experiment could not be carried out.

4.4. Fourth Beamtime: Coulex of ^{134}Sn in October 2025

An additional Coulomb-excitation experiment on ^{134}Sn was therefore granted and performed in September/October 2025, as the penultimate Miniball campaign before the Long

Shutdown 3. For this beam time, the General Purpose Separator (GPS) was employed. A breeding time of 230 ms in the EBIS was applied, which has a direct impact on the beam composition: with a half-life of $T_{1/2} = 1.05(1)$ s, a non-negligible fraction of the ^{134}Sn ions decays to ^{134}Sb during the charge-breeding process. To obtain detailed information on the time-dependent beam composition, an ionization chamber was installed downstream of the target. On four occasions during the experiment, the target was retracted, and dedicated ionization-chamber runs were carried out to monitor the relative contributions of the different isotopes in the beam.

The beam energy was again chosen as 4.4 MeV/u. In contrast to the previous experiments, a ^{194}Pt target was adopted for the main part of the measurement instead of ^{206}Pb . This choice was motivated by the proximity of the first 2_1^+ states in ^{134}Sn (725 keV) and ^{206}Pb (803 keV), which would complicate the γ -ray analysis if a ^{206}Pb target were used as the primary Coulomb-excitation partner. The following targets were mounted on the target wheel:

- Position 1: 5 mm aperture
- Position 2: 2 mm aperture
- Position 3: 3.0 mg/cm^2 ^{194}Pt (96.45 %)
- Position 4: 3.0 mg/cm^2 ^{194}Pt (96.45 %)
- Position 5: 3.1 mg/cm^2 ^{206}Pb (99.8 %) + 0.03 mg/cm^2 nat. C
- Position 6: empty.

The experiment started on 25th September 2025 and finished on 3rd October 2025. In total, 185 h of beam time were delivered, with only short interruptions. Of this, approximately 140 h correspond to radioactive-ion-beam data taken on the ^{194}Pt targets. Additional beam time was devoted to the ionization-chamber runs. Furthermore, the ^{206}Pb target was used for about one day to investigate Coulomb excitation of low-lying transitions in odd-odd isobars. For the first time within the experimental campaign underlying this work, the fully equipped Miniball spectrometer with all eight triple-cluster detectors was available for this measurement.

5. Data Analysis of 2022 Experiment

This chapter presents the analysis of the experiment conducted in 2022. As the first experiment performed with the Miniball spectrometer at ISOLDE following the Long Shutdown 2 and the detector refurbishment campaign, it naturally encountered some of the typical initial challenges associated with recommissioning complex equipment. As will become evident in the course of this chapter, the data acquisition setup at the time was not yet fully optimized for the specific requirements of the experiment. Furthermore, already during the beam tuning phase, an intensity was used that was too high for stable operation. This led to elevated count rates in the germanium detectors, which in turn resulted in (a) significant dead times and (b) a degradation in peak quality in the recorded spectra. Additionally, a magnet failure of the HIE-ISOLDE accelerator caused a two-day interruption of the beam time, further reducing the available statistics. The focus of this work is on the subsequent experiment conducted in 2023, which will be discussed in detail in the following chapter. Nevertheless, the main steps of the analysis of the 2022 data are briefly outlined here. Since the procedures for data preparation and analysis closely resemble those applied in the 2023 experiment, only a concise summary is given at this point. A more detailed description can be found in the analysis section of the following chapter.

5.1. Miniball Calibration

Prior to the shipment of the refurbished Miniball array to CERN, a characterization of the energy resolution of the individual detector segments was performed at the University of Cologne. For this purpose, a ^{60}Co calibration source was used. The full width at half maximum (FWHM) of the 1.33 MeV γ -ray line was determined for each segment. The resulting FWHM values are displayed on the right-hand side of Fig. 5.1. Most segments had an energy resolution well below 2.5 keV, with only a few detectors slightly exceeding this value. This corresponds to a typical relative energy resolution of approximately 0.19 % at 1.33 MeV, which confirms the excellent performance of the Miniball detectors after refurbishment.

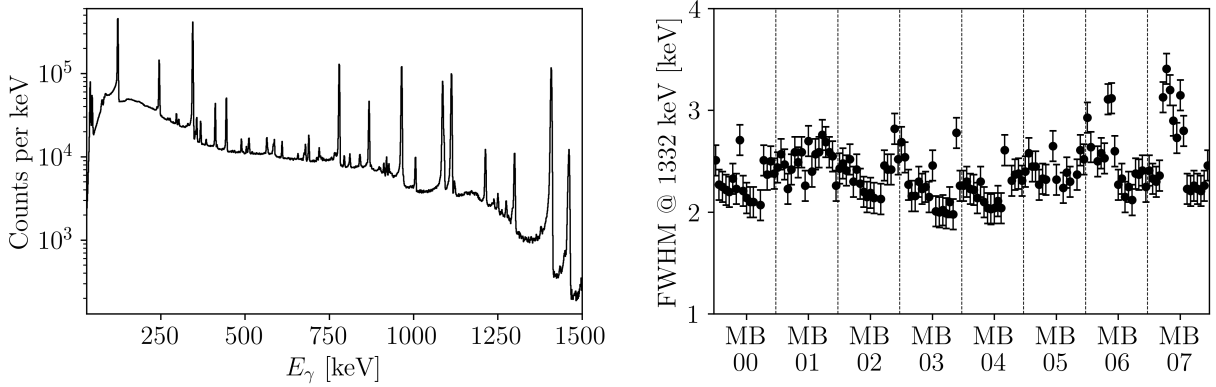


Figure 5.1. Left: Calibrated ^{152}Eu spectrum measured at ISOLDE. The spectrum shows the sum of all segments after energy calibration. Right: Energy resolution (FWHM) of each Miniball segment measured at 1.33 MeV using a ^{60}Co source at the University of Cologne prior to installation at ISOLDE.

Following the installation of Miniball at ISOLDE, a calibration measurement was carried out using a ^{152}Eu source. The well-known γ -ray transition energies of this source served as reference points for a segment-wise energy calibration. A second-order polynomial function was fitted to the measured channel-to-energy relations for each segment individually. The resulting energy-calibrated spectrum, shown on the left-hand side of Fig. 5.1, represents the sum spectrum of all calibrated segments. The calibration measurement was further used to provide information for the relative efficiencies for different energies of the setup. For the present experiment, Cluster 3 was excluded, leaving seven of the eight triple clusters. Accounting for one non-operational crystal in Cluster 1, the measurement was performed with 20 of the 24 crystals. Due to limitations in the quality of background spectra acquired during the experiment, a time-dependent energy shift tracking (e.g., via periodic gain stabilization) could not be performed. As will be discussed in a later section, the resolution of the peaks in the background data was insufficient, and the dead-time too enhanced to enable reliable peak position monitoring over time.

5.2. CD Calibration

For the calibration of the DSSSD, a two-point energy calibration was performed using the radioactive ion beam data and the kinematic lines from the scattered Sn and Pb nuclei. The right-hand side of Fig. 5.2 displays the spectra measured at four different scattering angles. Two distinct peaks can be observed in each spectrum, allowing for a separate two-point calibration of each ring and sector. The expected energies of scattered ejectiles and recoils

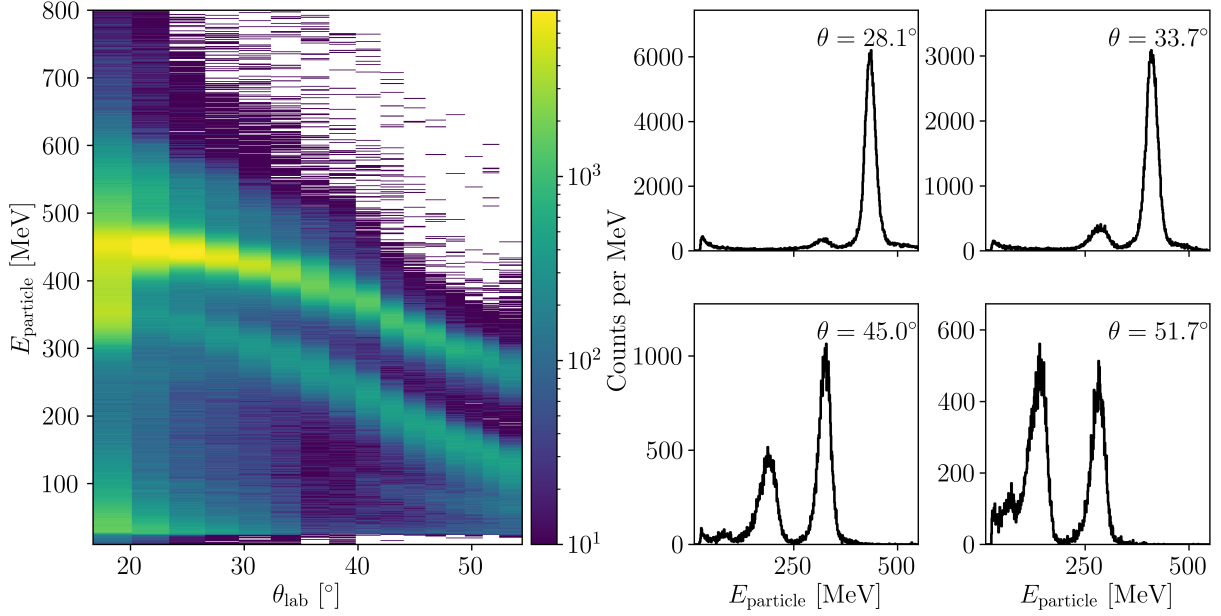


Figure 5.2. Left: Calibrated particle energy plotted as a function of scattering angle θ_{lab} . The z-axis (intensity) is shown on a logarithmic scale. The scattered Sn nuclei are visible in the higher-lying kinematic line, Pb on the lower one. For large scattering angles, a component arises at lower energies below the kinematic line of the Pb nuclei. This comes from ^{12}C , which was used as a fronting on the target. Right: Calibrated energy spectra for four different scattering angles. A clear separation between scattered ejectiles and recoils is visible.

in vacuum were calculated using the LISE++ code [132]. The calibrated spectra shown here represent the sum over all four quadrants.

On the left side of Fig. 5.2, the final energy vs. scattering angle (θ_{lab}) matrix is shown. A clear separation between recoils and ejectiles is visible. The energy resolution of the DSSSD in this data set was excellent, except for the first sector. Due to the large Rutherford scattering cross section at low angles, this sector suffered from considerable radiation damage during the experiment. The CD was placed at a distance of 30 mm to the target in the SPEDE chamber, covering scattering angles of 16.1 – 52.7° in the laboratory system.

5.3. Miniball Angle Determination

A stable beam experiment was performed to determine the geometrical parameters of the Miniball detector array — namely the angles ϕ , θ , and α (the latter describing the rotation of a cluster around its long central axis), as well as the distance R from the target position to the detector. A ^{22}Ne beam was directed onto a CD_2 target, inducing (d,p) and (d,n) reactions. Due to the mass asymmetry between projectile and recoil, these reactions are

strongly forward-focused.

As a result, the γ -ray energy emitted from excited states populated in the reaction depends predominantly on the angle of the segment in which the γ ray is detected, rather than on the scattering angle of the recoiling nucleus. The observed energy is Doppler shifted and varies as a function of the emission angle. By measuring the Doppler-shifted energy of the well-known 440.2 keV transition in ^{23}Na for each segment and comparing it with the known rest energy, the detector geometry can be reconstructed.

Figure 5.3 shows the measured γ -ray energy per segment before (red) and after (black) Doppler correction. The red horizontal line indicates the rest energy of the 440.2 keV transition in ^{23}Na . Cluster 3 was not operational during the experiment. Crystal 0 in Cluster 1 was malfunctioning and also not applied in the in-beam experiment. Additionally, Crystal 1 of Cluster 1 and Crystal 2 of Clusters 2 and 5 were not used for the angle determination but were employed in the subsequent radioactive ion beam experiment. The position calibration could still be performed for all these detectors. Since the Doppler shift depends on the four geometrical parameters mentioned above, as well as on the velocity of the decaying nucleus, measurements in five segments are sufficient to determine a detector's position. Even for Cluster 1, where only one crystal was used, the available statistics were sufficient for a reliable position determination.

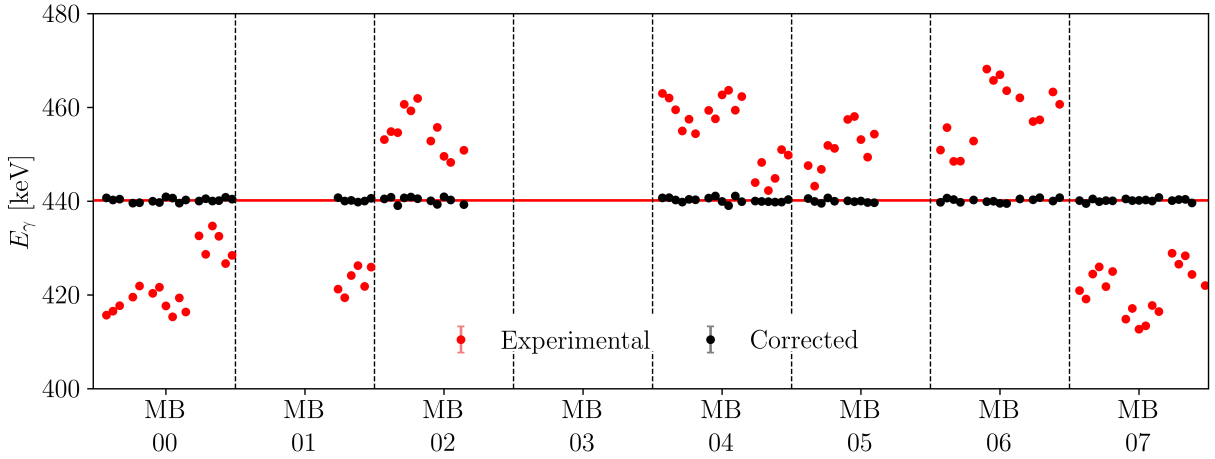


Figure 5.3. Measured Doppler-uncorrected γ -ray energy for each segment (red). The black points show the energy after Doppler correction, based on the extracted detector positions. The red line marks the known transition energy of 440.2 keV in ^{23}Na .

5.4. Beam Composition

The radioactive ion beam used in this experiment was produced using the molecular beam technique. While the employed separators are highly effective in discriminating between different isotopes, they do not provide separation of isobars. To minimize isobaric contamination, a molecular beam of $^{130}\text{Sn}^{34}\text{S}$ was chosen. This approach exploits the fact that tin is one of the few isobars capable of forming stable molecular bonds with sulfur, thereby significantly reducing isobaric impurities in the beam. However, minor contamination from other species cannot be entirely excluded. In addition, ^{130}Sn exhibits a long-lived 7_1^- isomeric state with a half-life of several minutes, which decays via β decay to ^{130}Sb . Determining the contribution of this isomeric component in the beam is also of critical importance for the interpretation of the observed decay patterns.

To obtain an initial estimate of the beam composition, the full set of detected γ rays was analyzed. As the data in this experiment were collected with a particle trigger, all these γ rays were recorded within a time window from -3 to $+1\ \mu\text{s}$ relative to the particle detection. The resulting γ -ray spectrum is shown in Fig. 5.4. Peak positions corresponding to transitions following the β decay of ^{130}Sn in its ground and isomeric states are shown in green and orange, respectively. Transitions from the daughter nucleus ^{130}Sb , which decays into excited states of the stable nucleus ^{130}Te , are marked in blue. The observed

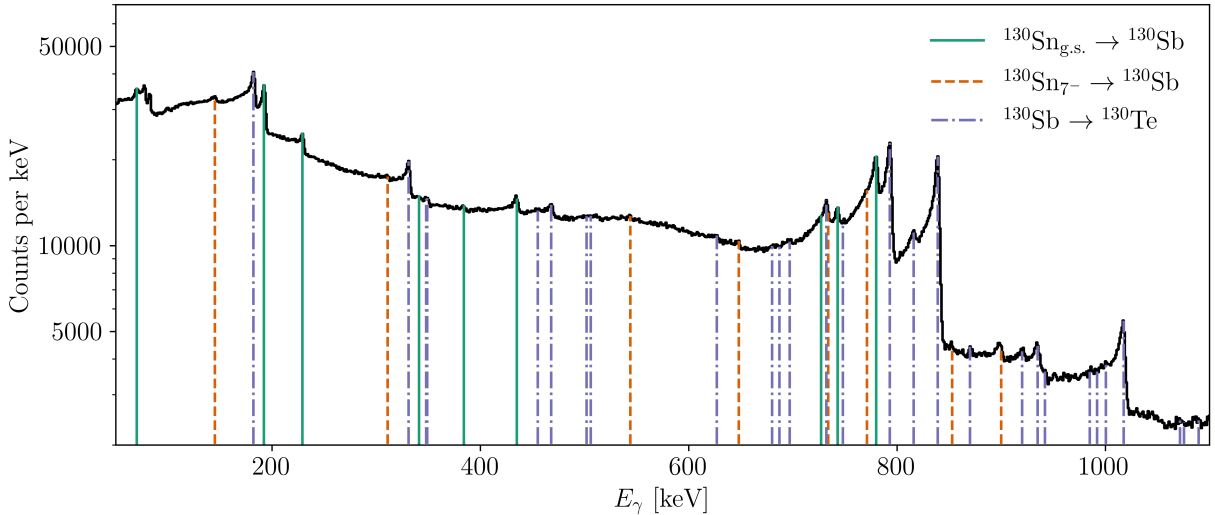


Figure 5.4. γ -ray spectra summed over all crystals and the entire beam time. Expected γ -ray energies from transitions following the β decay of ^{130}Sn and its daughter nuclei are indicated. No additional peaks are observed in the spectra. The peaks deviate significantly from a Gaussian-like shape.

spectral features are consistent with the presence of ^{130}Sn in the beam. However, the peak shapes deviate significantly from the expected Gaussian form. Instead, they exhibit pronounced low-energy tails—sometimes resembling triangular structures. This distortion is attributed to excessively high counting rates during data acquisition, which resulted not only in substantial dead time but also in the observed non-Gaussian peak shapes. At certain times during the experiment, pile-up effects in the HPGe detectors became so severe that the recorded count rate dropped to 0 Hz, as virtually all detected γ rays were affected by pile-up.

As a consequence, a detailed determination of the beam composition was not feasible with these data. Nevertheless, the fact that all prominent lines in the spectrum can be attributed to the decay of ^{130}Sn or ^{130}Sb strongly suggests that these isotopes were the only radioactive beam constituents.

To obtain a more precise characterization of the radioactive component of the beam, an additional single-crystal HPGe detector was installed upstream of the beamline capacitors for the beam composition study [133]. The detected radiation originates from radioactive beam nuclei that accumulate along the beamline and decay in situ. γ -ray spectra were recorded in intervals of 30 minutes.

The spectra from three consecutive 30-minute intervals recorded before the beam interruptions resulting from problems with the PS Booster and a magnet failure of the HIE-ISOLDE accelerator, are shown in the left panel of Fig. 5.5, in black, blue, and red, respectively. Transitions associated with the decay of $^{130}\text{Sn}_{\text{g.s.}}$ (780 keV), $^{130}\text{Sn}_{\text{is.}}$ (899 keV), and ^{130}Sb (816 keV and 840 keV) are marked. ^{130}Sb also has an isomeric state, which is fed by the decay of $^{130}\text{Sn}_{\text{g.s.}}$, whereas the isomeric ^{130}Sn decays mostly into the ground-state of Sb. The shown spectra exhibit no significant differences over time, indicating that the activity had reached equilibrium before the interruption. The right panel displays three spectra measured after the beam stop. Unfortunately, these exhibit clear double-peak structures that vary in shape and intensity depending on the energy, rendering them unsuitable for a reliable beam composition analysis. The 899 keV transition from $^{130}\text{Sn}_{\text{is.}}$ has a known absolute intensity of 8.1(9) %, and the corresponding peak is therefore expected to be weak. However, since the spectra before the beam interruption represent an equilibrium state, they can be used to infer the beam composition, provided that intensity and detection efficiency corrections are applied. A detailed justification for this approach, as well as a decay scheme showing all relevant γ -ray transitions, is presented in Section 6.8 and Appendix A.

One significant complication is that an absolute efficiency calibration of this detector was not feasible. The origin of the γ rays detected in the HPGe detector is not point-like but distributed along the beamline, and no dedicated source measurements were performed. To still enable a meaningful estimation of the composition—specifically the relative contribution of $^{130}\text{Sn}_{\text{g.s.}}$ and $^{130}\text{Sn}_{\text{is.}}$ —two γ -ray transitions with closely spaced energies were used: the 780 keV line (g.s. decay) and the 899 keV line (isomeric decay). The relative detector efficiency for these energies was taken from the Miniball efficiency curve, yielding a correction factor of 1.08. For the reasons outlined above, systematic uncertainties of the efficiencies can only be estimated. Because an efficiency increase towards higher energies is not physically expected over such a small energy interval, an uncertainty of 0.08 to the correction factor. The total intensities of the 780 keV and 899 keV transitions were obtained from the spectra in Fig. 5.5, corresponding to efficiency- and absolute-intensity-corrected counts of 49266(3154) and 30986(6107), respectively. The uncertainty estimate accounts for the unknown spatial distribution of activity, the absence of a source calibration, and the statistical uncertainties in the weak 899 keV peak.

Based on this correction, the beam composition was estimated to be 61(7) % $^{130}\text{Sn}_{\text{g.s.}}$ and 39(8) % $^{130}\text{Sn}_{\text{is.}}$. The sum of these two transitions amounts to 80234(6874), which agrees well with the activity of ^{130}Sb , determined from the 840 keV $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in ^{130}Te , yielding 79200(1200). This 840 keV line has an absolute intensity of 100 % for the β decay

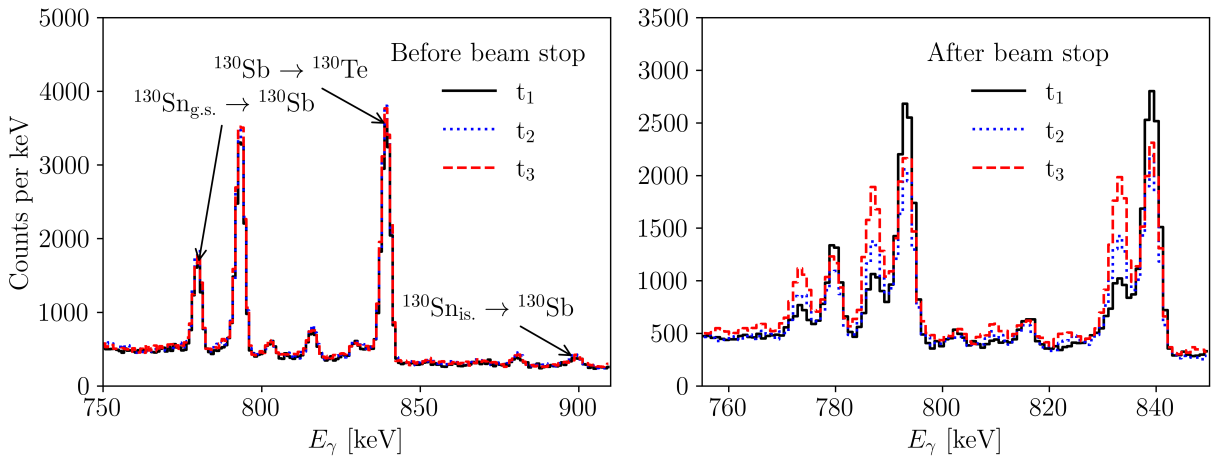


Figure 5.5. Left: Three spectra of 30-minute measurement of decaying particles in the beam line. A constant rate is given, indicating that the activity is in an equilibrium. Right: After the beam stop, the three peaks marked on the left side show a double-peak structure in the spectra, hindering the analysis of the beam composition.

of $^{130}\text{Sb}_{\text{g.s.}}$ and $^{130}\text{Sb}_{\text{is.}}$, respectively, and therefore serves as an excellent reference transition. All β -decays along the $^{130}\text{Sn} \rightarrow ^{130}\text{Sb} \rightarrow ^{130}\text{Te}$ cascade accumulate in the 2_1^+ state of ^{130}Te , which decays via this 840 keV transition. A significant component of ^{130}Sb in the beam could therefore not be identified.

5.5. Experimental Result of 2022 Experiment

Figure 5.6 shows Doppler-corrected γ -ray spectra in the region of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transitions in ^{130}Sn (left) and ^{206}Pb (right). The spectra are gated either on ejectile particles (top) or on recoils (bottom). Each spectrum was fitted with a Gaussian peak shape on top of a linear background. The total fit is shown in red, while the Gaussian and background components are indicated by dashed and dotted curves, respectively. All spectra are binned in 8 keV. The extracted peak intensities, together with their statistical uncertainties, are annotated in the individual panels. In all cases, these uncertainties exceed 12 % and represent only the statistical contribution; systematic effects are not included at this stage. The fitted FWHM ranges from 20 keV to 68 keV. A low-energy tail component was not included in the fit model. Nevertheless, the spectrum in the lower right panel shows indications of a tail towards lower energies, which likely influences the extracted intensity. Since a reliable parametrization of this feature was not feasible, the corresponding uncertainty was increased for the subsequent analysis.

These data were used to perform a GOSIA analysis, employing the extracted peak areas as input yields. The data set was divided into two experiments: one gated on scattered Sn nuclei and one on Pb recoils. Internal conversion coefficients were calculated using the BRICC code [134]. Stopping powers for the integration of yields over the full target thickness and angular coverage of the CD detector were obtained from the SRIM code [135]. Due to the limited statistics, the diagonal matrix element — related to the quadrupole moment of the 2_1^+ state — was neglected. This approximation is justified to first order, as ^{130}Sn lies near the doubly magic nucleus ^{132}Sn , whose spherical character has been demonstrated in earlier studies (e.g., [23, 54]). A 4^+ state was included in the level scheme as a buffer state, although its decay could not be observed in the present data set. Using the measured input values, together with the determined beam composition and setup efficiency, a reduced transition probability of

$$B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.054_{-14}^{+10} (\text{stat.}) e^2 \text{b}^2$$

was extracted. It should be noted, however, that the quoted uncertainty is purely statistical. A reliable estimate of the systematic is challenging, primarily due to the poorly constrained isomeric composition of the beam, which — as discussed in the previous chapter — currently represents the dominant source of uncertainty. Including this contribution as a sole systematic uncertainty, the transition probability is expressed as

$$B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.054_{-14}^{+10} (\text{stat.}) (7)_{\text{syst.}} e^2 \text{b}^2.$$

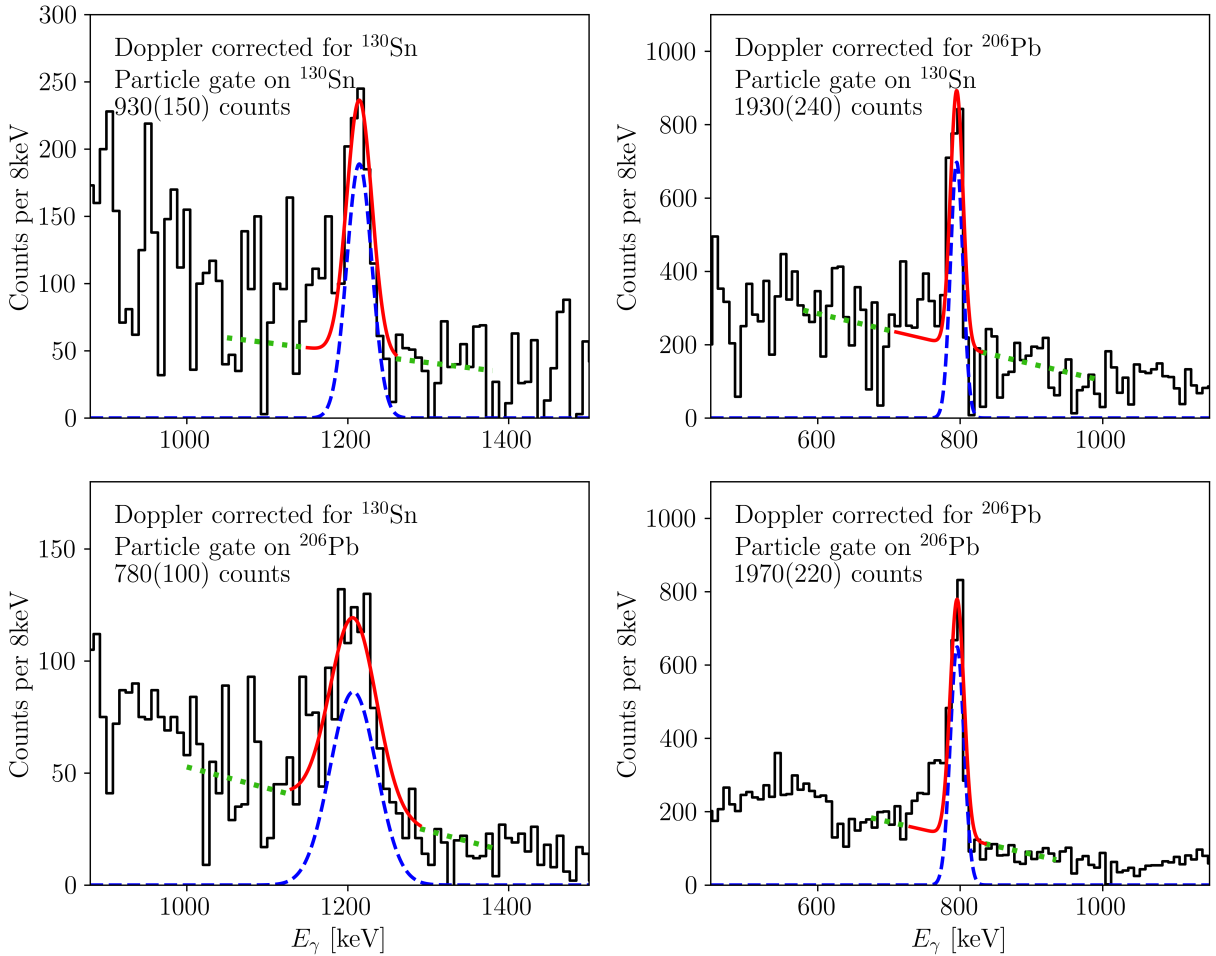


Figure 5.6. Doppler-corrected γ -ray spectra around the $2^+ \rightarrow 0^+$ transitions in ^{130}Sn (left) and ^{206}Pb (right), gated on ejectiles (top) and recoils (bottom). Each spectrum was fitted with a Gaussian peak on a linear background (total fit in red); the Gaussian and background contributions are shown as dashed and dotted curves, respectively. The spectra are displayed with 8 keV binning, and the resulting peak intensities with their statistical uncertainties are given in each panel. For more information, see text.

This result is particularly noteworthy as it shows good agreement with the theoretical prediction by Togashi *et al.* [54] ($0.055 e^2 b^2$) and exceeds the value reported in the conference proceedings by Radford *et al.* [57, 58] by a factor of approximately 2.5. However, potential systematic contributions, such as the exact choice of the fitted lineshape, uncertainties in beam energy, target thickness, reorientation effect, or the distance between target and CD detector — affecting the effective scattering angle θ — are not discussed here. Their impact on the overall uncertainty is considered negligible in light of the already dominant uncertainty arising from the beam composition. With a relative uncertainty of 31 – 39 %, this result does not yet permit a definitive interpretation. The large overall uncertainties stem from several experimental limitations encountered during the initial data acquisition phase. Early technical challenges, such as an imprecise beam focus, the provisionally assigned data acquisition system, and an excessively high proton current, prevented reliable monitoring and control of the instantaneous Miniball count rate, which reached particularly high values during the experiment. Furthermore, an unexpected multi-day interruption reduced the available beam time. As a result, only 75 h of data were collected, compared to the originally approved 120 h.

In order to improve the results of the extracted $E2$ strength, a follow-up experiment was approved. This second measurement of ^{130}Sn Coulomb excitation is described in the following chapter.

6. Data Analysis of 2023 Experiment

This chapter describes the data analysis procedures applied in the Coulex experiment performed in October 2023. It begins with a description of the data acquisition and preprocessing steps, followed by the calibration of the DSSSD and germanium detectors. The subsequent sections address the characterization of the beam composition and the optimization of the Doppler correction applied to the γ -ray spectra. Based on the Doppler-corrected spectra, the experimental yields for the excitation of the 2_1^+ state (and the corresponding $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ decay) are extracted and used as input for the GOSIA2 analysis. Finally, the data are examined for signatures of excitation originating from an isomeric component in the beam.

6.1. Data Acquisition and Preprocessing

During the experiment, FEBEX modules developed at GSI were used to digitize the signals from the various detector components, including the HPGe Miniball detectors, the DSSSD, and the timing signal from the EBIS. The digitized data were continuously recorded and stored to tape for subsequent offline analysis. The offline data processing was carried out using the MB-SORT framework [136], a modular software package originally developed for the analysis of Miniball experiments at ISOLDE. During this thesis, the code was extended and adapted to meet the specific requirements of the present experiment, including changes in detector configuration and calibration procedures.

The first step in the preprocessing chain consisted of converting the raw binary data into the ROOT file format. This provided a structured and flexible foundation for all subsequent stages of the analysis. In the next step, physical events were built by correlating signals from different detector systems that occurred within a coincidence time window of $\pm 3 \mu\text{s}$. An event was defined as any instance in which either a charged particle or a γ ray was detected by the setup. The event-building algorithm grouped all time-correlated signals, allowing for the reconstruction of the full interaction based on information from the Miniball and DSSSD. Following event construction, detector-specific calibrations were applied. These included

energy and timing calibrations, software thresholds, and corrections for configuration-specific effects — such as the re-assignment of crystal segments necessitated by the

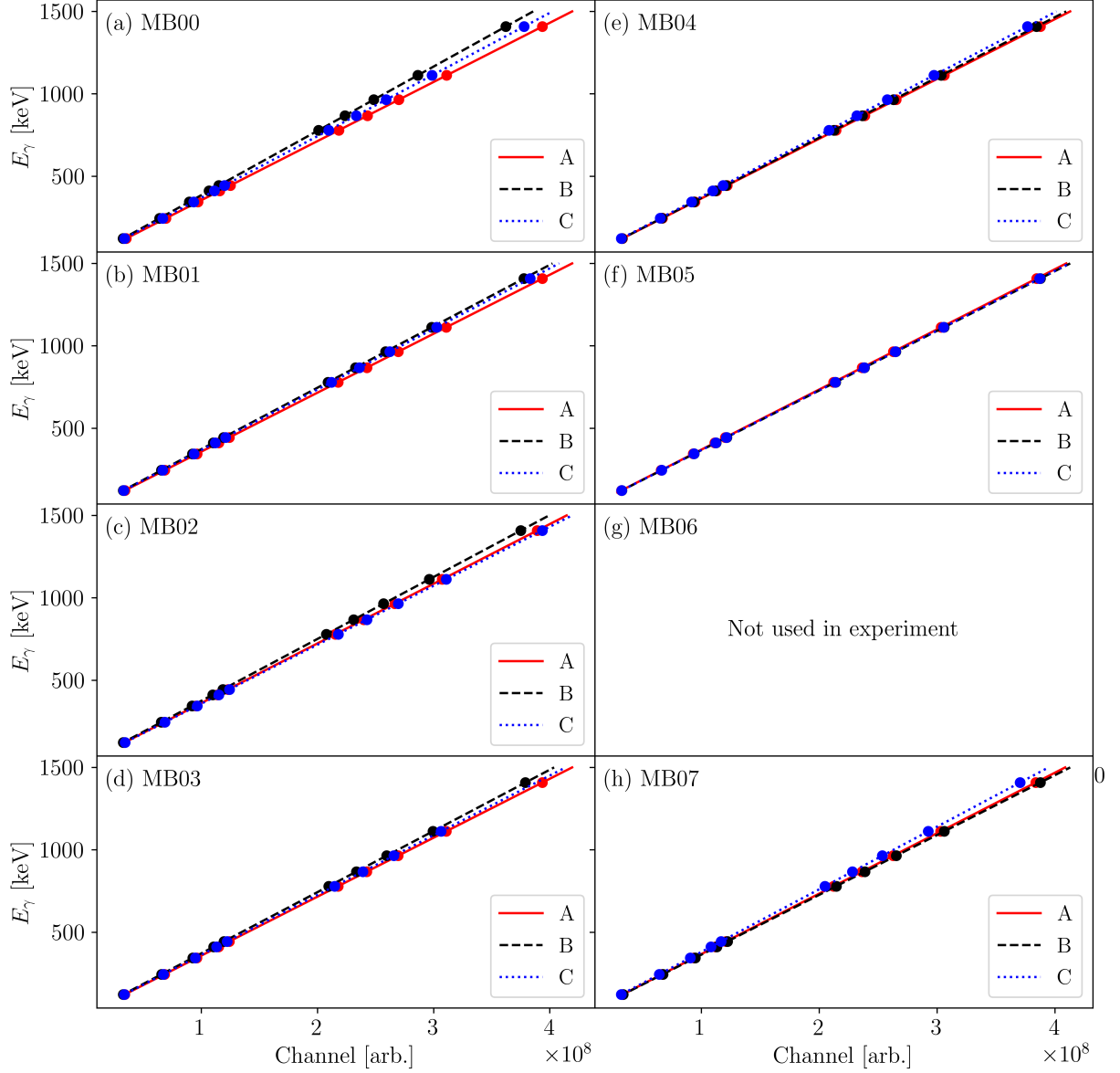


Figure 6.1. Quadratic calibration functions converting ADC channel numbers into γ -ray energies for all Miniball clusters. The plots show the fitted calibration curves for the core signals of each of the three crystals per cluster. Seven out of the eight Miniball clusters were used in the present experiment; cluster 6 was excluded due to previously observed peak distortions. For each included cluster, the energy calibration was based on a set of well-known γ -ray transitions from a ^{152}Eu source, and the resulting fits reflect the segment-wise energy response of the individual detector cores.

recent re-encapsulation of the Miniball detectors. In the final stage of data pre-processing, calibrated events were used to generate the histograms that form the basis of the data analysis. These include timing spectra between detector systems, particle energy spectra, and Doppler-corrected γ -ray energy spectra.

6.2. Miniball Calibration

To calibrate the HPGe detectors of the Miniball setup, a measurement with a radioactive ^{152}Eu calibration source was conducted overnight. Since each segment of the detector array requires a dedicated energy calibration, an automated procedure was implemented. In the first step, a peak-finding algorithm was applied to the raw, uncalibrated spectra to identify the two most prominent γ -ray lines. These were unambiguously assigned to the well-known transitions at 121.78 keV and 344.28 keV. A linear two-point calibration was then

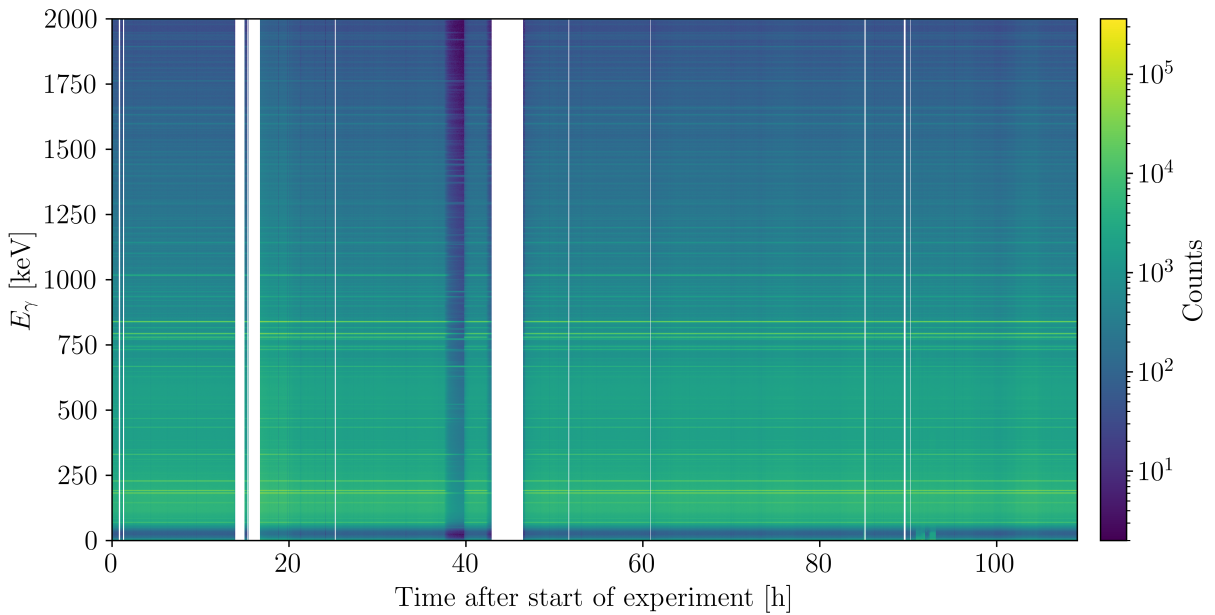


Figure 6.2. The γ -ray energy is plotted as a function of time after the start of the experiment. Prominent horizontal lines indicate intense background γ -ray transitions that are consistently visible throughout the run. Vertical white bands correspond to periods without data acquisition due to beam interruptions. The most pronounced gap, occurring at approximately 45 hours of beam time, was caused by an emergency shutdown of the PS Booster. Additional vertical structures stem from shorter interruptions, such as the loss of individual accelerator cavities. In addition, after approximately 35 hours of beam time, the beam was deliberately stopped while data acquisition continued to record background spectra. As a result, this time interval exhibits reduced statistics.

performed based on these peaks, providing an initial energy calibration for each individual spectrum. Subsequently, the calibrated spectra were used to fit additional γ -ray peaks corresponding to known transitions at 121.78, 244.69, 344.29, 411.12, 443.98, 778.90, 867.34, 964.13, 1112.12, and 1408.01 keV. These peak centroids served as input for a quadratic energy calibration, allowing for a more accurate conversion from ADC channel numbers to physical γ -ray energies, especially at higher energies. This procedure was repeated for all detector segments. For the core signals of each crystal, the calibration was cross-checked and, where necessary, manually refined to ensure consistency and accuracy across the setup. Figure 6.1 shows the resulting calibration curves for all used Miniball triple clusters. Each panel displays the fitted calibration functions and the corresponding peak positions for the three core segments of a cluster. The quadratic contribution to the calibration function amounts to approximately 0.03 % at an energy of 1000 keV. It should be noted that cluster 6 was excluded from the experiment, as it exhibited a double-peak structure in previous measurements, indicating unreliable performance of the corresponding electronic channels. Using the summed spectra and the same peaks as used for the energy calibration, a relative efficiency analysis was performed. These efficiencies will be discussed later in section 6.11.1

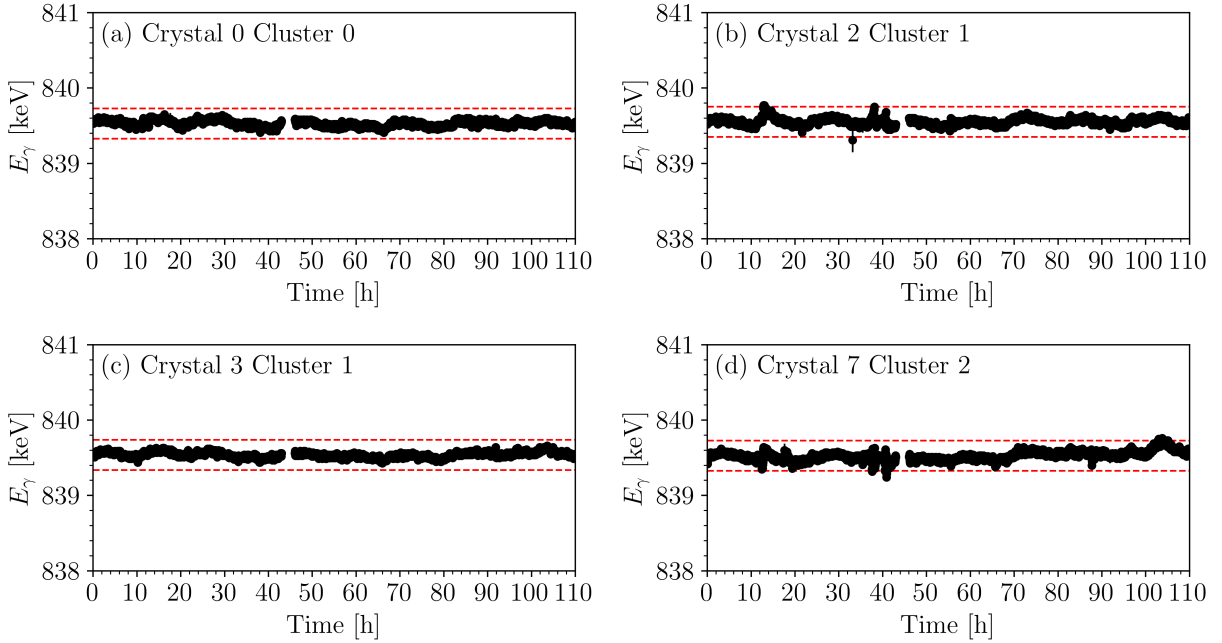


Figure 6.3. Fitted peak position of the 839.5 keV background line for four different core signals, drawn against the time after the beginning of the experiment. γ -ray is emitted after β -decay of ^{130}Sb . The peak position is well within 0.2 keV for all times and all cores. The ± 0.2 keV interval is highlighted as a red dashed line.

in more detail.

HPGe detectors can exhibit slight energy drifts over time, which potentially affect the precision of energy calibration in long experiments. To account for such effects, a time-resolved analysis was performed. Figure 6.2 shows a two-dimensional histogram of γ -ray energy as a function of time since the beginning of the experiment detected with the Miniball array. Prominent horizontal structures correspond to background γ -ray lines, predominantly originating from β -decay of ^{130}Sn and its daughter nuclei. Vertical features in the matrix reflect interruptions in the beam delivery: for example, the gap around 45 hours after the start of the experiment marks a shutdown of the PS Booster, while shorter interruptions are attributed to temporary instabilities such as tripping cryostats.

To investigate the energy stability of the detector system, the matrix was produced for every core signal of each crystal and divided into time slices of 20 seconds. For each time slice, a projection onto the energy axis was performed. Gaussian fits were applied to several full-energy peaks. The fitted peak position of the 839.5 keV line is shown as a function of time in Fig. 6.3 for four different crystal core signals. The observed variations remained below 0.2 keV (shown as a dashed red line) over the entire runtime, demonstrating excellent energy stability. To verify the reliability and quality of the calibration across the entire detector system, the procedure was repeated for each individual segment. No significant energy drift was observed. Consequently, no time-dependent correction of the energy calibration was required.

6.3. CD Calibration

For the initial energy calibration of the CD detector, a two-point procedure was employed. The known energies of scattered ^{130}Sn and ^{206}Pb nuclei served as reference values and were used to calibrate all rings and segments of the detector. While this approach enables a basic alignment of the energy response across the detector, it does not account for nonlinearities arising from particle-dependent effects. In particular, the so-called pulse height deficit (PHD) constitutes a significant source of systematic uncertainty in silicon detectors. This effect originates from the reduced charge collection efficiency for particles with high specific energy loss (dE/dx), leading to a lower signal amplitude than expected from the actual kinetic energy deposited in the material [137]. Since the magnitude of the PHD depends on the atomic number Z of the detected ion, heavier particles are affected more strongly than lighter ones. As a consequence, a universal calibration function cannot adequately describe

both recoils and ejectiles, which differ substantially in Z . This leads to unphysically large offsets in the calibration. While this is sufficient for single-hit coincidence events—where only one particle is detected on each side of the CD—it fails in events involving multiple hits, especially when one particle is detected in two neighboring strips. In such cases, the summed calibrated energy of both strips deviates from the true total energy, as the Z -dependent signal suppression is not properly corrected.

To address this issue, a more refined calibration procedure was adopted, following the method described in Ref. [138]. This method exploits the fact that, for events with a strip multiplicity of one — i.e., exactly one p-strip and one n-strip register a signal so that the full particle energy is deposited in a single strip on each side — the energies measured

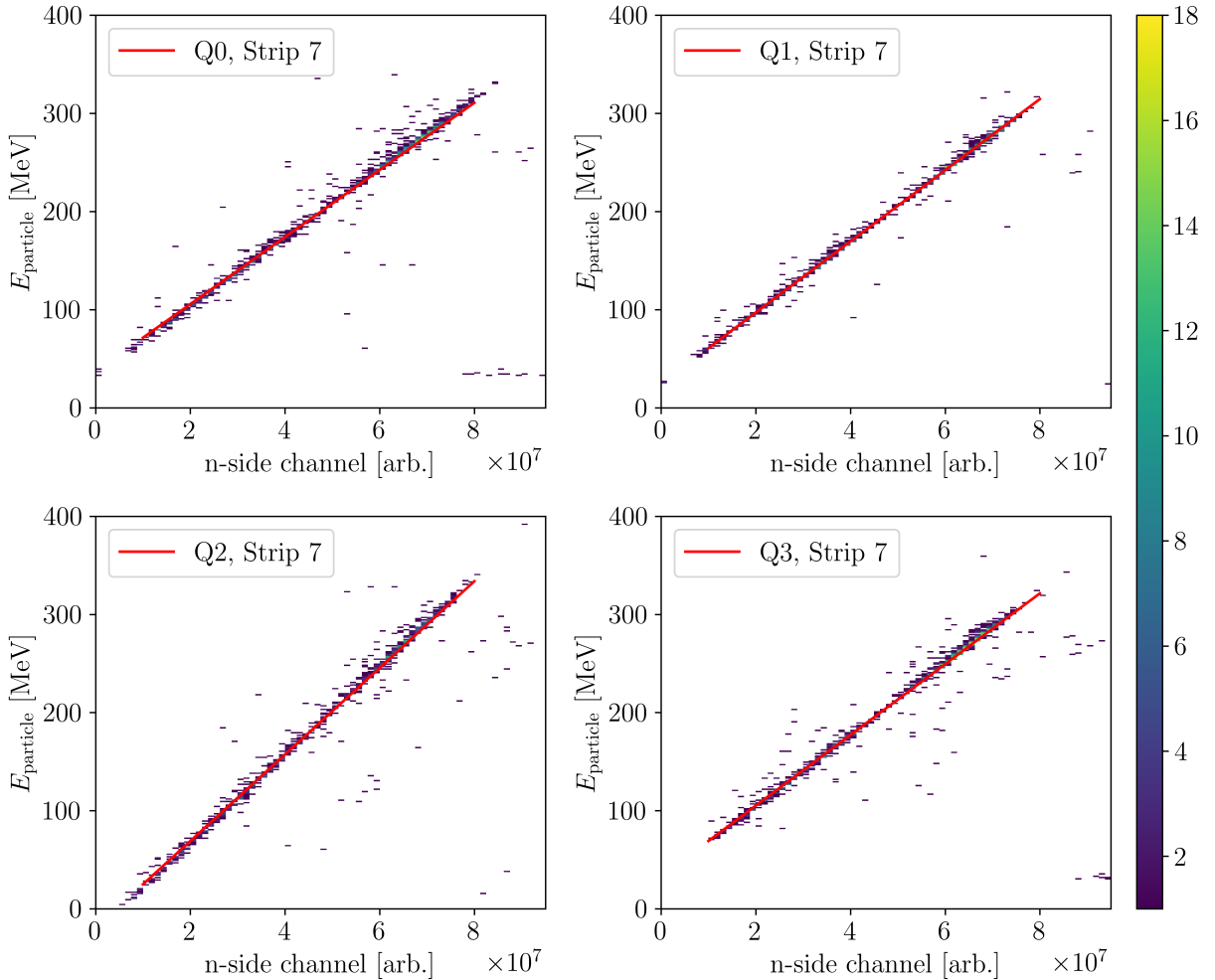


Figure 6.4. Calibration of the n-side strips. Shown is the linear correlation between the calibrated p-side energy and the raw n-side channel for strip number 7 in all four quadrants.

on the p - and n -sides must be identical. In the present analysis, ring 9 was chosen as the starting point for the iterative calibration procedure. A dedicated two-point calibration was first applied to the p -side of this ring, using the two well-separated energy peaks as reference points. Subsequently, two-dimensional histograms were constructed for each of the twelve n -side strips in a given quadrant. The energy measured on the p -side (in keV) is plotted along one axis, while the uncalibrated ADC value of the corresponding n -strip was plotted along the other. Events were selected such that exactly one hit occurred on a given n -strip and one on the already calibrated p -strip. A linear correlation is observed, and a linear fit was used to derive the calibration function for each n -strip. An example of this calibration process is shown in Fig. 6.4, illustrating the fit for n -strip number 7 across all four quadrants.

Once all n -strips had been calibrated using this method, the procedure was repeated for the p -side strips. In this step, one calibrated n -strip was fixed, and the coincident p -side signals were analyzed analogously. Again, linear fits were applied to derive the final calibration parameters for all p -side strips.

The result of this iterative calibration procedure is presented in Fig. 6.5, where the experimentally determined particle energies are compared to simulations carried out using the KINSIM code [139]. The upper kinematic line comes from the scattered ^{130}Sn , while the

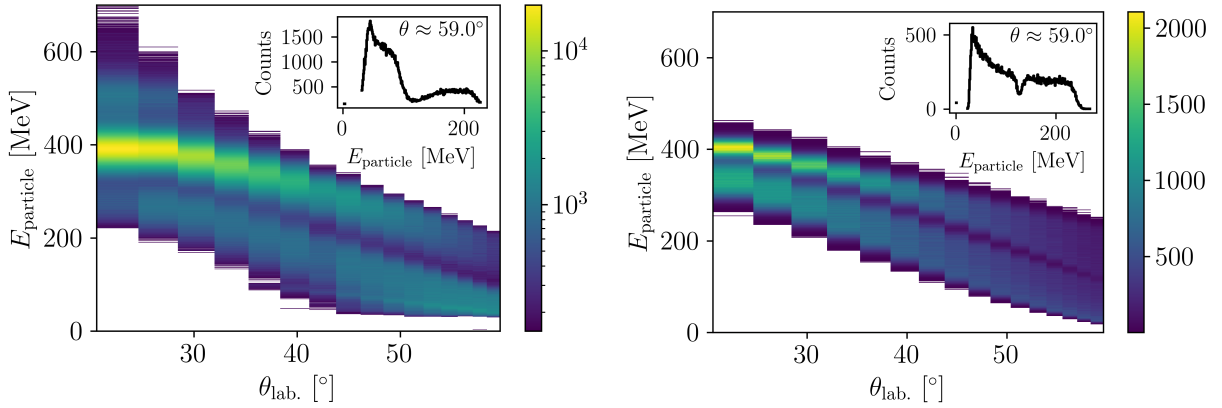


Figure 6.5. The particle energy is shown as a function of the scattering angle in the laboratory system. On the left, the calibrated experimental data is shown; on the right, a simulated spectra using KINSIM [139]. In the experimental data, the calibration of the first two detector rings deviates from the simulated trend due to beam-induced damage. Note that a logarithmic color scale is used for the experimental spectrum. In the inner panels, the projection onto the y-axis is shown for ring 15. The ^{130}Sn intensity is enhanced in the simulated spectra since the cross-section was increased to reduce the simulation time. For more information, see text.

lower line is from the recoiling ^{206}Pb nuclei. The experimental spectra exhibit an enhanced yield at lower recoil energies, as highlighted in the inset. This feature is reproduced by the simulations, supporting its interpretation as a physical effect. Another plausible origin is target contamination, e.g., by ^{12}C or ^{16}O , which may contribute predominantly in this energy range. Such an effect has also been observed in Ref. [140]. Moreover, a small ^{12}C contribution is expected because the ^{206}Pb layer was deposited onto a $^{\text{nat.}}\text{C}$ substrate during target production.

6.3.1. Threshold and Multiplicity

In the next step, software thresholds were implemented individually for each channel of the CD detector, including both the n-side and p-side, to suppress events with artificially high multiplicities. Figure 6.6 shows the resulting multiplicity distributions for the n-side versus the p-side for quadrant 0, exemplarily before (left) and after (right) applying these thresholds. Initially, multiplicities of up to nine are observed. This is partly caused by electronic noise but also by the high beam intensity used during the experiment. The elevated beam current was necessary because of the inherently low Coulomb-excitation cross sections in the investigated system, which lies close to the doubly magic nucleus ^{132}Sn —only two neutrons away—and thus exhibits limited collectivity. Furthermore, the excitation energy was relatively high, which further reduced the Coulomb excitation probability. To compensate for these limitations and to achieve meaningful statistics, a high-intensity beam

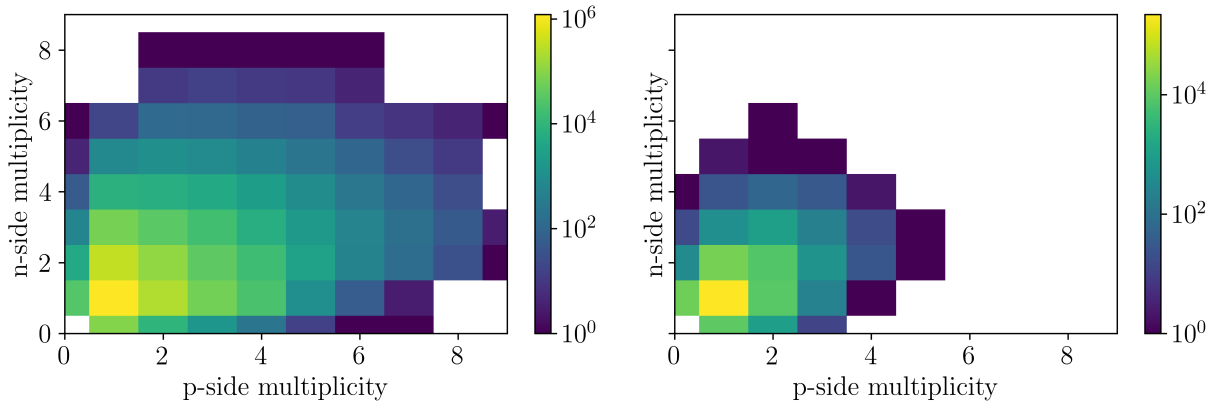


Figure 6.6. Left: Multiplicity correlation between the n-side and p-side of the DSSSD in quadrant 0 prior to calibration. At this stage, no threshold corrections were applied to individual strips and rings, leading to an increased number of spurious multiplicity combinations and electronic noise artifacts. Right: The same multiplicity matrix after applying calibrated thresholds and individual adjustments for each segment. The cleaned-up distribution shows a realistic event distributions and a significant reduction of random background coincidences.

had to be employed, inevitably leading to higher event rates and consequently increased multiplicities.

However, high multiplicity events significantly complicate the identification of the actual interaction point within the detector. When multiple strips on both the n- and p-sides fire in coincidence, it becomes difficult to reliably reconstruct the energies and spatial information associated with the true physical event. Many of these coincidences are not genuine physical events but instead stem from electronic noise or beam-induced background. Therefore, software thresholds were applied in a channel-by-channel manner to reject events dominated by noise or signal pile-up.

To validate the effectiveness of these thresholds, additional correlation plots were analyzed. These included cases where only one strip on each side of the detector fired in coincidence, as well as situations in which one side registered one strip and the other two, or where two strips fired simultaneously on both sides. In ideal conditions, events with one strip on each side appear as a narrow diagonal line in the correlation matrix. Events with two strips on one side and one on the other appear slightly above or below the diagonal, while true two-vs-two coincidences are located further off-diagonal on both sides. These diagnostic patterns confirmed that the chosen thresholds were sufficient to suppress background contributions while preserving real interaction events.

The final multiplicity distributions after applying the optimized thresholds, as shown in Fig. 6.6, demonstrate that most events now correspond to clean 1 vs. 1 coincidences. The number of events with multiplicities higher than two was reduced drastically. It should be noted that the innermost ring of the detector was significantly affected by radiation damage from previous experiments, during which the detector was exposed to high beam intensities. This damage led to a degradation in energy resolution, making it impossible to distinguish between beam and target particles in that region. To minimize the impact of this degraded ring on surrounding channels, an adequate higher threshold was applied to it. In addition, some detector strips and rings were nonfunctional during the experiment. Particles hitting these dead regions could not be properly identified and were therefore assigned a multiplicity of zero.

6.3.2. Position of the CD Detector

The CD detector was originally assumed to be mounted 26 mm downstream of the target position. This value was initially determined using two independent methods: an α -source

calibration and kinematic constraints from the known reaction kinematics of ^{130}Sn and ^{206}Pb . The α -calibration was carried out using a ^{226}Ra source placed at the target position. The intensity of the emitted α particles was measured across the rings of the CD to extract the relative solid angles and infer the source-to-detector distance. This measurement, performed shortly after the experiment by Liam Gaffney, resulted in an average distance of 26.18 ± 0.10 mm [141]. However, during the Doppler-correction analysis, it became evident that this value did not yield optimal correction. Two main factors were identified: first, the PCB holding the CD detector was not tightly fixed during the experiment, possibly shifting its effective position; second, the α -source used in the calibration measurement had been placed upstream of the target ladder, whereas the target in the radioactive beam experiment was mounted downstream. This change in geometry results in a systematic offset in the CD-to-target distance. To accurately determine the effective distance during the experiment, the kinematics of the reaction were examined. For binary reactions occurring at the center of the target material, the following conditions are fulfilled:

$$\phi_{\text{lab.}}^{\text{ejectile}} + \phi_{\text{lab.}}^{\text{recoil}} = 360^\circ, \quad \theta_{\text{CM}}^{\text{ejectile}} + \theta_{\text{CM}}^{\text{recoil}} = 180^\circ.$$

If both reaction partners (e.g., ^{206}Pb and ^{130}Sn) are detected in coincidence, the sum of their reconstructed center-of-mass scattering angles can be calculated as a function of the

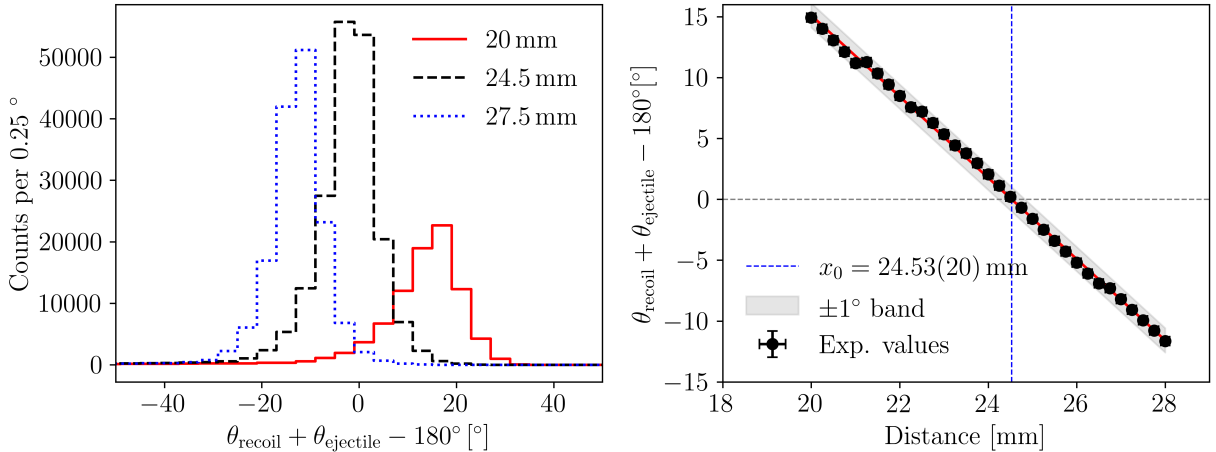


Figure 6.7. Left: Sum of center-of-mass scattering angles θ for ejectile and recoil, subtracted by 180° , shown for three assumed CD-to-target distances between 20 mm and 27.5 mm. For the correct distance, the distribution is expected to be symmetric around zero. A Gaussian fit with tailing was applied to each distribution. Right: Extracted mean values from the fits as a function of assumed distance. The grey band represents an uncertainty of $\pm 1^\circ$. The zero crossing of the linear fit (blue line) defines the best estimate for the actual CD position.

assumed CD-to-target distance d . The left panel of Fig. 6.7 shows this angular sum, offset by 180° , for three exemplary distances. A distribution symmetric around zero is observed at $d = 24.5$ mm, while for smaller distances the peak shifts to positive angles. To quantify this trend, the full data set was analyzed for distances between 20 mm and 28 mm in 0.25 mm steps. For each value, a Gaussian fit with left-tailing was applied to the angular sum distribution, and the resulting mean value was extracted. These values are shown in the right panel of Fig. 6.7. A clear linear correlation is observed. A linear fit was applied to the points, and the intersection with zero defines the distance at which the angular condition is fulfilled. To estimate the uncertainty, an angular deviation of $\pm 1^\circ$ was translated into a distance uncertainty, visualized as a gray band in the figure. This yields a final value for the effective CD-to-target distance of $d = 24.53(20)$ mm. From this value, the angular coverage of each CD ring in the laboratory frame was calculated. The laboratory scattering angles θ_1 and θ_2 , which correspond to the inner and outer boundaries of a given ring n , can be derived from geometry using the following expressions:

$$\theta_1 = \tan^{-1} \left(\frac{c + n \cdot \Delta_{\text{dead}} + n \cdot w_{\text{ring}}}{d} \right), \quad \theta_2 = \tan^{-1} \left(\frac{c + n \cdot \Delta_{\text{dead}} + (n + 1) \cdot w_{\text{ring}}}{d} \right),$$

Table 6.1. Calculated laboratory angles θ_1 and θ_2 for each CD ring for a detector distance of $d = 24.53(20)$ mm, including a dead layer thickness of 0.1 mm.

n	θ_1 [$^\circ$]	θ_2 [$^\circ$]
0	20.15(15)	23.96(17)
1	24.15(17)	27.74(19)
2	27.92(19)	31.28(21)
3	31.45(21)	34.56(22)
4	34.72(22)	37.61(23)
5	37.76(23)	40.43(23)
6	40.57(23)	43.03(23)
7	43.16(23)	45.43(23)
8	45.54(23)	47.64(23)
9	47.74(23)	49.68(23)
10	49.77(23)	51.56(23)
11	51.65(23)	53.29(22)
12	53.38(22)	54.90(22)
13	54.98(22)	56.39(22)
14	56.46(22)	57.76(21)
15	57.83(21)	59.05(21)

where $c = 9$ mm is the distance from the detector surface to the first ring, $\Delta_{\text{dead}} = 0.1$ mm is the dead layer thickness, $w_{\text{ring}} = 1.9$ mm is the ring width, and $d = 24.53(20)$ mm is the target-to-detector distance. The results for all 16 rings ($n = 0$ to 15) are summarized in Table 6.1, showing the corresponding angular ranges in the laboratory system. It should be noted that, for optimal Doppler correction, the azimuthal angle ϕ of the individual CD strips is also a critical parameter. However, this angle cannot be constrained using the previously described method, as variations in the target-to-detector distance have no influence on ϕ . Since the CD is mounted using three screws located at fixed positions in ϕ , only three discrete rotational alignments of the detector are mechanically possible. All three configurations were tested by evaluating the quality of the Doppler correction. A rotation angle of -118° was found to yield the best agreement with expected peak shapes and energies and was therefore adopted for the final analysis.

6.4. Particle- γ Coincidences

The time correlation between detected particles and coincident γ -rays has to be analyzed to distinguish prompt events from random coincidences. This distinction is essential for background suppression and the reliable identification of γ -ray transitions following nuclear reactions. In Fig. 6.8, the measured γ -ray energy is plotted against the time difference

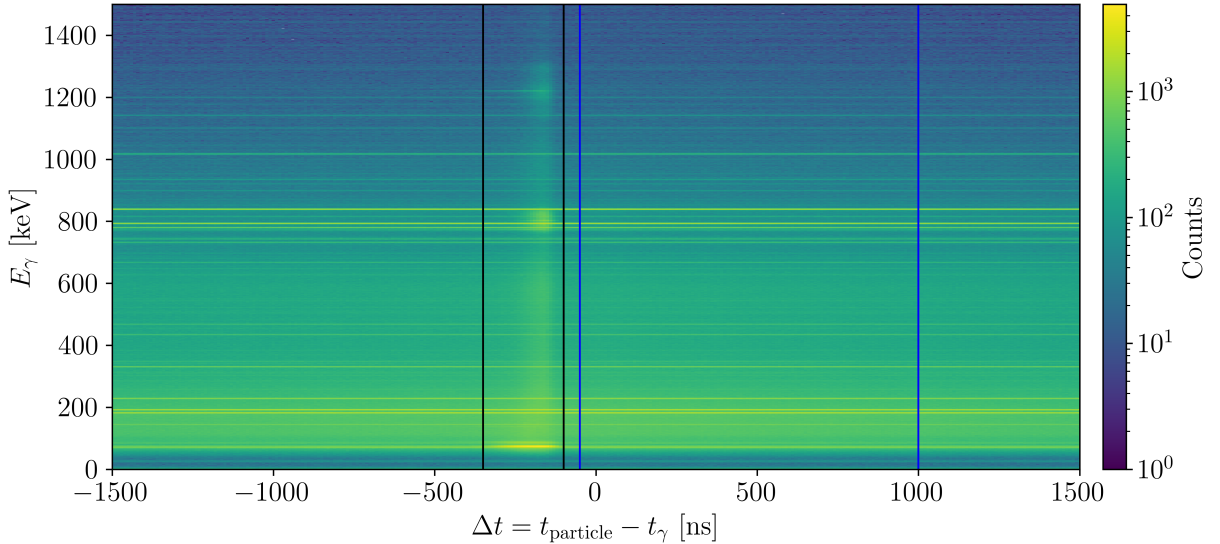


Figure 6.8. γ -ray energy plotted against the time difference between the γ ray and the detected particle. The defined prompt time window of 250 ns is indicated in black, while the random time window is shown in blue.

between the detected particle and the corresponding γ -ray. Since particle detectors exhibit response times faster than those of HPGe detectors, the time difference is not centered around zero. For smaller γ -ray energies, a left tail of the prompt peak is observed due to the time-walk effect. In experiments in which multiple prompt γ -rays of different energies are observed, a time-walk correction can be applied to align the timing differences across the time spectrum. However, in the present experiment, the majority of the observed γ rays are not prompt but rather occur randomly in time, and insufficient peaks were detected prompt. As a result, a time-walk calibration was not feasible. This does not compromise the quality of the data but leads to a broader prompt time window. In Fig. 6.8, this prompt window is indicated by black vertical lines, while the blue-shaded area denotes the random-timing region.

A straightforward background subtraction using a simple prompt-to-random ratio was not possible, due to the non-constant nature of the background across the time range. In particular, prominent background peaks arise from γ -rays following β -decay of ^{130}Sn and its daughter nuclei. The lifetime of those decays ranges from a few minutes of ^{130}Sn up to half an hour in ^{130}Sb , leading to a massive background after reaching an equilibrium. To accurately account for this, characteristic peaks were fitted in both the prompt and random spectra. The ratio of the fitted peak intensities was then used to scale the random spectrum before subtraction.

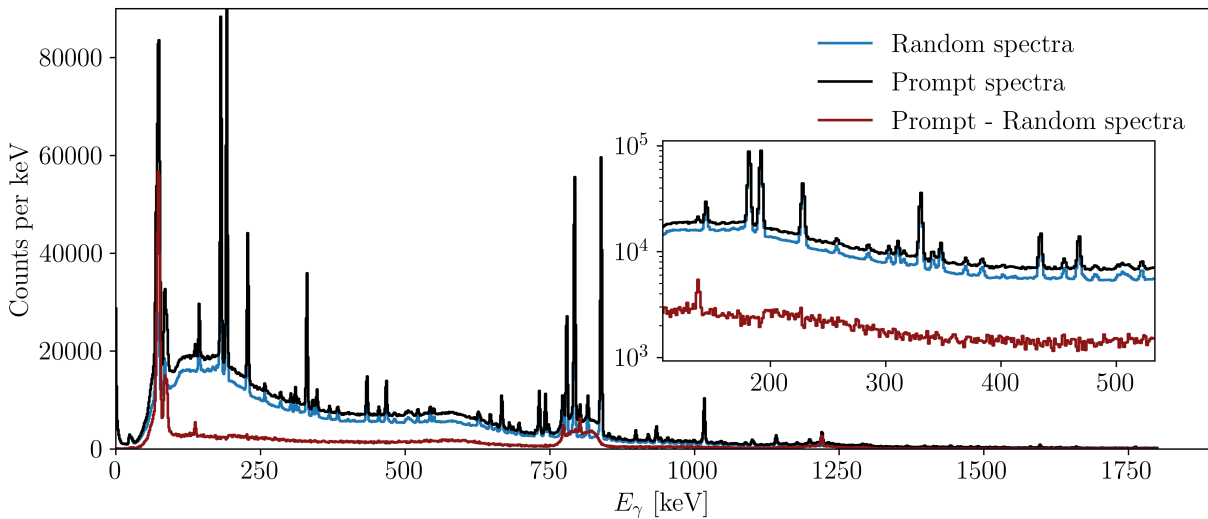


Figure 6.9. The prompt, random, and background-subtracted γ -ray spectra are shown in black, blue, and red, respectively. The inset displays the energy region from 100 to 500 keV on a logarithmic scale. For further details, see text.

The resulting prompt and random spectra are shown in Fig. 6.9, again in black and blue, respectively. The random spectrum was scaled by a factor of 0.29 before subtraction. The final background-subtracted spectrum is displayed in red. The absence of any residual background peaks confirms that the chosen scaling factor of 0.29 accurately reflects the true prompt-to-random ratio. It is noteworthy that the background subtraction worked remarkably well, even though the random coincidences exceeded the true prompt peak by an order of magnitude. The inset in Fig. 6.9 highlights the energy region from 100 to 500 keV. Here, a distinct peak is visible at 139 keV. This peak originates from the excitation of a $^{130}\text{Sn}_{7-}$ beam component into a long-lived 5_1^- state. Since this state decays after the nucleus has come to rest, the corresponding γ -ray is not Doppler-shifted. This particular feature will be discussed in more detail later in this thesis.

6.5. Miniball Angle Determination

In experiments with fast radioactive ion beams, excited nuclei typically decay in flight in the case of experiments at HIE-ISOLDE with approximately 10 % of the speed of light. As a result, the observed γ -ray lines are Doppler shifted and significantly broadened. Moreover, the compact geometry of the Miniball array, with detectors placed close to the target, enhances the Doppler broadening. The Doppler-shifted photon energy $E_{\gamma,\text{lab.}}$ in the laboratory frame is given by the relativistic expression:

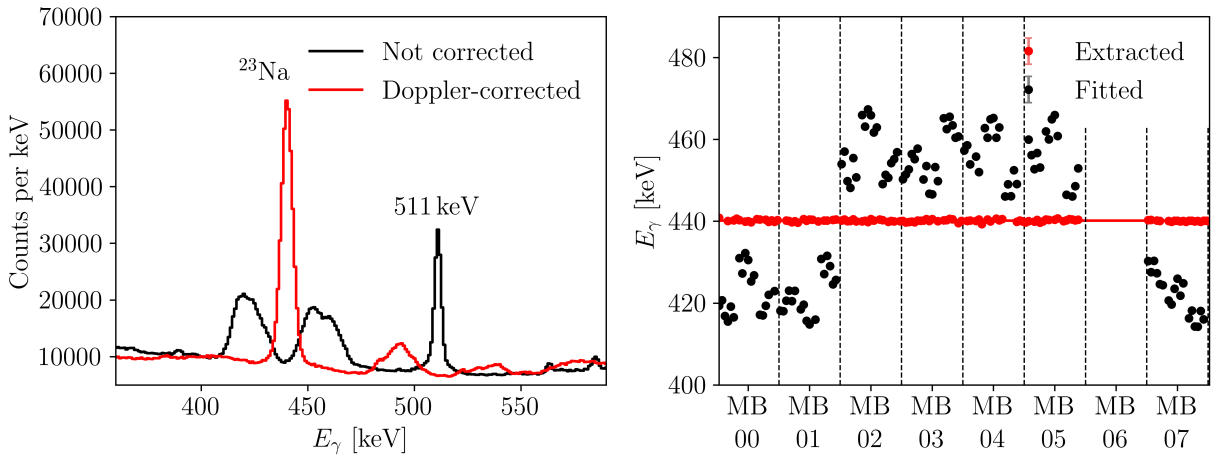


Figure 6.10. Left: The measured spectrum is shown in black without any Doppler correction. In red, the Doppler-corrected spectrum is displayed, using the final values for α , θ , ϕ , and r . Right: Fitted peak positions of the 440 keV transition are shown in black, and the corresponding Doppler-corrected positions are shown in red. Miniball clusters 6 was not used in this experiment.

$$E_{\gamma,\text{lab.}} = \frac{\sqrt{1 - \beta^2}}{1 - \beta \cos \theta_{p\gamma}} \cdot E_0, \quad (6.1)$$

where E_0 is the photon energy in the rest frame of the emitting nucleus, $\beta = v/c$ is the recoil velocity in units of the speed of light, and $\theta_{p\gamma}$ is the angle between the particle and the emitted γ ray. The angle $\theta_{p\gamma}$ can be calculated from the particle and photon directions in the laboratory system:

$$\cos \theta_{p\gamma} = \sin \theta_\gamma \sin \theta_p \cos(\phi_p - \phi_\gamma) + \cos \theta_\gamma \cos \theta_p. \quad (6.2)$$

To perform the Doppler correction, both the velocity β and the angle $\theta_{p\gamma}$ must be known.

The velocity can either be estimated from the measured energy in the particle detector using $\beta \approx \sqrt{2E/m_0c^2}$ or obtained from kinematic calculations under the assumption that the reaction occurs at the target center. Another way is to determine the velocity at which the Doppler-correction works the best; this is described later in section 6.9.1.

The angle $\theta_{p\gamma}$ requires precise knowledge of both the particle direction (measured in the CD detector) and the γ -ray direction, which is determined by the position of the Miniball crystal that registered the photon. To determine these detector angles accurately, a dedicated calibration measurement using the reaction $^{22}\text{Ne}(d,p\gamma)^{23}\text{Na}$ was performed. A ^{22}Ne beam at 3.5 MeV/u was directed onto a deuterated polyethylene (CD_2) target of 3.2 mg/cm^2 thickness. The kinematics of this reaction favor forward scattering, and the recoil angle is negligible. Consequently, the Doppler shift of the deexciting γ -ray depends predominantly on the HPGe-angle, while the particle angle can be treated as constant. The 440.2 keV transition in ^{23}Na , depopulating the $5/2^+ \rightarrow 3/2^+$ ground-state transition, serves as

a reference. Figure 6.10 shows the summed γ -ray spectrum of all Miniball detectors before (black) and after (red) Doppler correction. The 511 keV annihilation line remains sharp

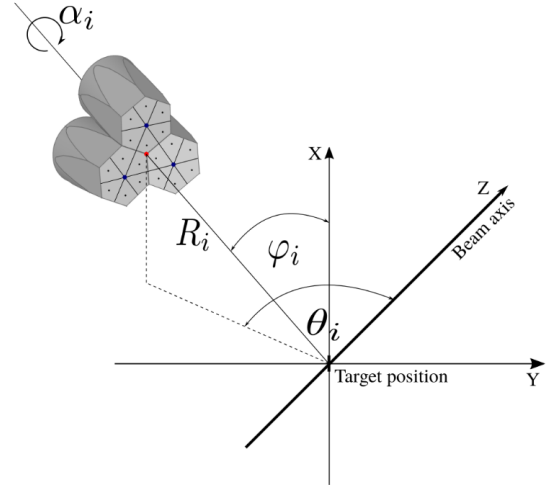


Figure 6.11. Definition of the defining quantities of the Miniball detector position, the angles θ and ϕ , the rotation of the cryostat α and the distance of the cryostat to the target R . Figure adapted from Ref. [24] with kind permission of D. Rosiak.

in the uncorrected spectrum but is misaligned after correction, as expected. In contrast, the 440 keV line appears broadened without correction and becomes sharp once the correct angles and velocities are used.

The position of each Miniball triple cluster is defined by four parameters: the distance from the target center to the crystal front face R_i , the polar and azimuthal angles θ_i and ϕ_i , and the in-plane rotation angle α_i around the cluster axis. These parameters, illustrated in Fig. 6.11, are used to compute the γ -ray emission angle. The internal segmentation and crystal arrangement of the triple cluster are predefined in the analysis framework.

For each segment, the 440 keV peak is fitted individually. A minimization algorithm is applied to find the optimal detector geometry that reproduces the known transition energy after Doppler correction. As starting values, the angles from the mechanical detector frame are used. The sixfold segmentation of each crystal provides 18 data points per cluster to determine the four geometric parameters, which ensures reliable fits even if individual segments are missing. All fits are visually inspected and repeated if necessary. Segments that were inactive or whose spectra could not be fitted reliably are excluded from the final result. The resulting energy spectra for each segment and cluster reveal the expected energy

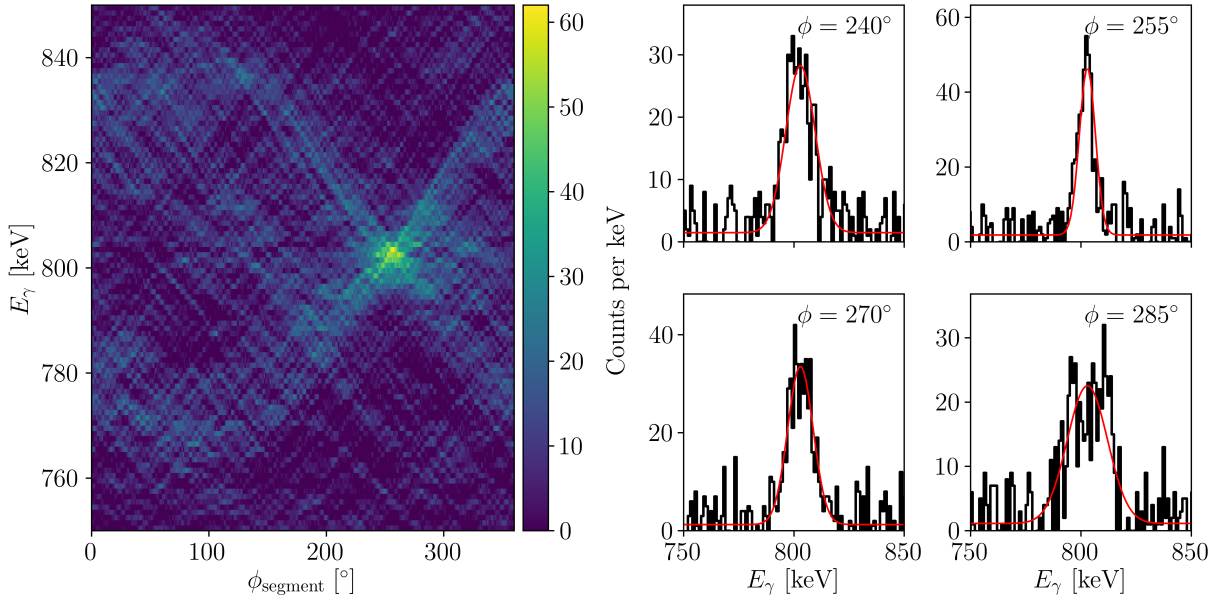


Figure 6.12. Left: Doppler-corrected energy of spectra measured for segment 3 of crystal 1 in cluster 3 depending on the ϕ angle of this segment. A minimum is visible at around 250° . Right: Projections onto four different ϕ s are shown, showing the dependence of the FWHM of the peak from the ϕ angle of the segment.

shifts depending on the detector placement. Crystals positioned in the forward direction show a shift toward higher energies, while backward-facing crystals are shifted to lower energies. These relative shifts confirm the consistency of the extracted angles and validate the Doppler correction procedure.

6.5.1. ϕ Determination

However, as seen from equation 6.1, the applied Doppler correction depends only weakly, to first order, on the azimuthal angle ϕ of the HPGe detectors if $\theta_p = 0^\circ$. Consequently, only the extracted parameters θ , α , and R can be considered reliable. Using the values deduced from the stable beam data, an FWHM of the 803 keV $2^+ \rightarrow 0^+$ transition in ^{206}Pb , obtained by gating on the recoil kinematics, of 13 keV was determined. This value is slightly higher than the 10 keV resolution reported in the work of David Rosiak [24] for a comparable setup.

To further constrain the azimuthal detector angle ϕ , an additional optimization approach was developed in this thesis by making use of the Coulomb-excitation data itself. A particle gate was applied to select the kinematic line of recoiling ^{206}Pb nuclei. For each Mini-ball segment, the Doppler-corrected energy of the 803 keV transition was evaluated as a function of different assumed ϕ -values. To this end, three-dimensional histograms were created, in which the Doppler-corrected energy was plotted against the assumed ϕ angle and the segment number.

Figure 6.12 (left) shows a two-dimensional slice of this three-dimensional histogram, illustrating the dependence of the Doppler-corrected energy of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in ^{206}Pb on the assumed emission angle ϕ for segment number 6 of crystal 1 in cluster 3. A pronounced dependence is visible, with a clearly defined minimum around $\phi = 255^\circ$, indicating the best alignment between the assumed and true

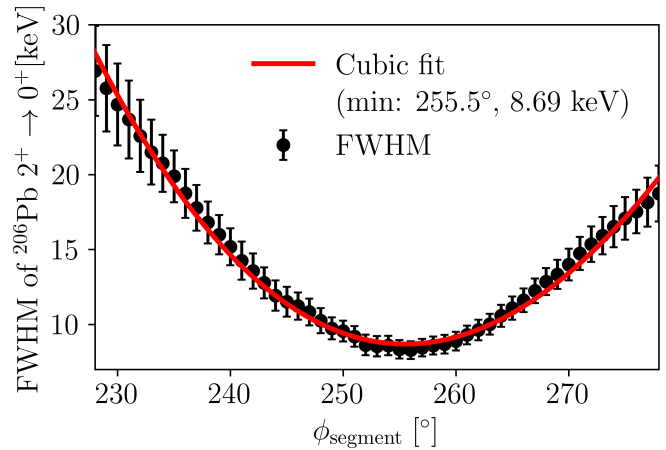


Figure 6.13. Exemplary representation of the ϕ determination for segment 4 of cluster 0 crystal 0. A Gaussian function was fitted by projecting the matrix shown in Fig. 6.12 for each ϕ . The FWHM and its error were determined. A quadratic function was fitted around the area where the fit converges and the minimum of this function was determined. This calculated FWHM and the parabolic function are shown here in black and red respectively.

detector angle. On the right-hand side of Fig. 6.12, projections of the Doppler-corrected spectra for several individual ϕ -values are shown. For each angle, a Gaussian function was fitted to the resulting photopeak. The quality of the Doppler correction was quantified by evaluating the FWHM of these peaks. To identify the optimal ϕ -value, the FWHM was plotted as a function of the assumed angle (Fig. 6.13). A third-order polynomial was fitted to the data points, and its minimum was taken as the most probable value for ϕ .

This procedure provided a refined estimation of the azimuthal orientation of the segments. This method of optimization was only feasible due to the exceptionally high statistics of the present experiment, which allowed for a segment-by-segment analysis of the Doppler-corrected energy distributions. Using the previously determined ϕ -values and the known energy of the uncorrected $5/2^+ \rightarrow 3/2^+$ transition measured in the stable beam experiment, a combined minimization method was introduced. In this approach, the difference between the Doppler-corrected transition energy and the known reference energy was minimized simultaneously with the deviation between the experimentally determined ϕ -value and the geometrically assigned segment angle. This allowed for a further refinement of the segment geometry by including both experimental observables in the fit procedure, ensuring an optimal and consistent alignment of the detector angles with respect to the observed Doppler shift.

The resulting values of ϕ after minimization are shown in Fig. 6.14 in black compared to measured values for each segment in red. The inner panel shows the differences of calculated and determined ϕ values. The resulting angles are shown in Table 6.2. Detector 6 was excluded in the given experiment due to double-peaks in the spectra. In addition, only segment signals (no core signals) were used in the present analysis. Segments that were non-operational in some crystals were excluded, resulting in a total of 115 segments

Table 6.2. Geometrical parameters of the Miniball clusters after optimization.

ID	θ [°]	ϕ [°]	α [°]	R [mm]
00	138.52	309.62	176.88	108.48
01	134.92	236.18	341.36	113.54
02	49.52	238.93	173.45	106.65
03	49.19	306.43	344.72	110.49
04	49.61	54.64	21.33	105.42
05	46.13	128.48	332.21	107.40
07	134.64	55.78	23.15	111.77

considered.

This improved angle determination results in a Doppler-corrected FWHM of 8.6(2) keV at 803 keV for gating onto the whole recoil kinematics.

6.6. Kinematic Gates

The use of a thick target, the resulting energy straggling of the scattered nuclei, and the finite energy resolution of the CD detector system lead to substantially broadened particle-energy distributions. In the one-dimensional projection of the particle-identification matrix, the recoil and ejectile contributions therefore overlap over a wide energy range. A clean separation based on simple thresholds or rectangular cuts is not feasible and would introduce significant cross-contamination. To minimize systematic misidentification in the extraction of particle-resolved observables, kinematic gates must be defined in regions where the two particle populations are sufficiently well separated.

To establish analysis gate conditions for the overlapping peaks, the observed peak structures were parameterized using analytical functions that approximate the experimental line shapes. In Fig. 6.5 the measured particle energy is shown in black for ring 10 of the CD.

As illustrated here, for this larger scattering angles ($49.77\text{--}51.56^\circ$) the spectra deviate markedly from a Gaussian shape and instead exhibit a broad maximum with pronounced

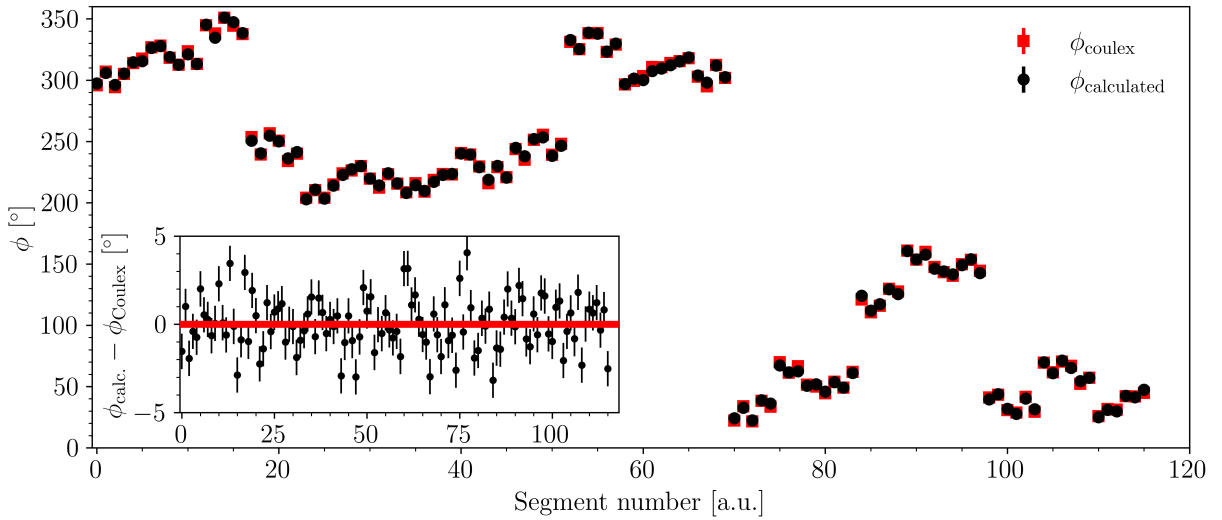


Figure 6.14. ϕ values deduced from Coulomb excitation spectra for each segment are shown in red. In black, the final phi values for each segment are shown. Inner plot: Residuals between measured and determined values.

exponential-like tails. The peak shapes were therefore modeled with a bifurcated Gaussian, which allows for different widths on the low- and high-energy sides of the centroid:

$$f_{\text{Asym}}(x) = A \cdot \begin{cases} \exp\left(-\frac{(x - \mu)^2}{2\sigma_L^2}\right), & x < \mu, \\ \exp\left(-\frac{(x - \mu)^2}{2\sigma_R^2}\right), & x \geq \mu, \end{cases} \quad (6.3)$$

where A denotes the amplitude, μ the centroid, and σ_L and σ_R the widths on the low- and high-energy sides, respectively. For each CD ring, the experimental spectrum was fitted with the sum of two such components, representing the recoil-like and ejectile-like contributions.

In Fig. 6.15 the result of this fit is shown for ring 10, including the individual components and their sum overlaid on the measured spectrum. The fit is restricted to energies above 100 MeV. At lower energies, electronic noise and additional contributions (e.g. scattering of ^{12}C in the target, cf. Fig. 6.5) distort the spectral shape such that the adopted parameterization does not provide an adequate description of the full distribution. Since the present fit model is

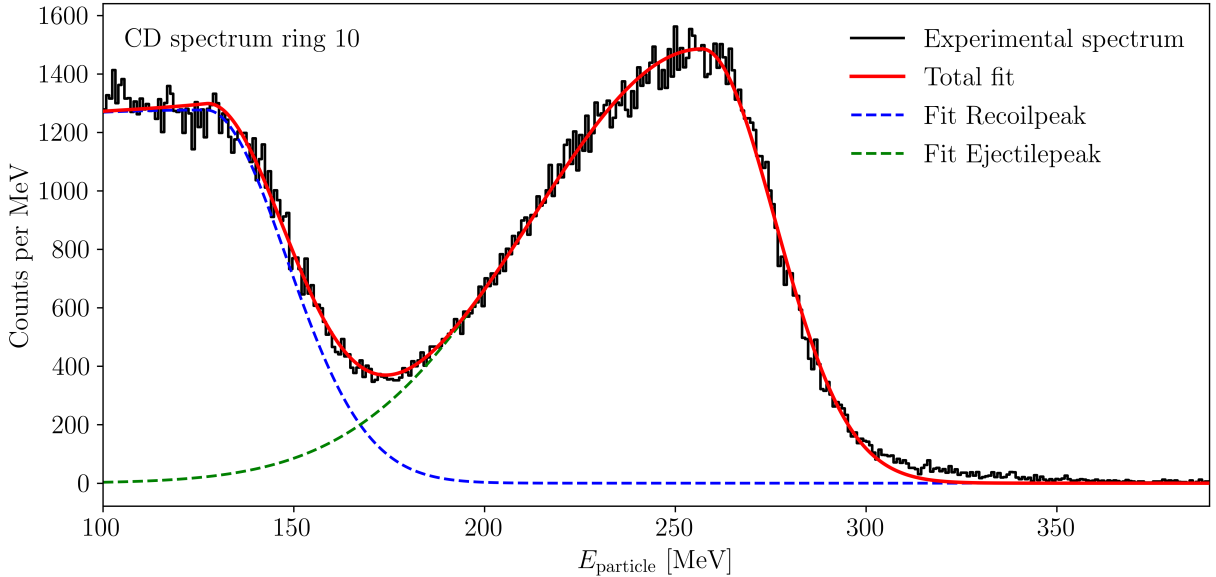


Figure 6.15. Experimental particle-energy spectrum for ring 10 of the CD detector (all sectors combined). The recoil-like (blue) and ejectile-like (green) contributions are described by bifurcated Gaussians; the red curve shows their sum. The parameterisation enables the definition of high-purity kinematic gates even in the overlap region.

only required to define gates in the overlap region between the two particle populations, the low-energy part is not considered further.

As a practical gating criterion, energy intervals were selected such that the contribution of the opposite component is below 1 % of its respective maximum. This procedure yields high-purity recoil and ejectile selections and forms the basis for the subsequent extraction of particle-resolved observables.

6.7. Beam Composition

For ^{130}Sn , neither transitional nor diagonal matrix elements of low-lying states are known, and no lifetimes have been reported for the states populated in the present experiment. The analysis technique of the measurement is therefore based on a normalization with respect to well-known transition probabilities of the target nucleus ^{206}Pb . The matrix elements are, in this case, deduced by calculating normalization constants based on the target excitation and then minimizing the difference of calculated and experimental yields of target and beam excitation by varying the unknown matrix elements. Thus, the amount of target excitation induced by the nuclei of interest has to be known precisely, and hence, the beam composition has to be determined accurately.

Potential beam contamination along the isobaric $A = 130$ chain must be considered, as the mass separation process is only sensitive to the mass-to-charge ratio. In previous experiments employing molecular $^{132}\text{Sn } ^{34}\text{S}$ beams, contamination with mass number $A = 166$ were observed [24]. Indications of possible contaminants in the present experiment can be obtained by analyzing the background γ -ray spectra. These spectra, which predominantly contain γ -rays following the β -decay of implanted isotopes and their daughter nuclei, are shown in Fig. 6.16. Transitions from the decay of $^{130}\text{Sn}_{\text{g.s.}}$ are highlighted in red, those from $^{130}\text{Sn}_{\text{is.}}$ in blue, and those from ^{130}Sb in green. Additional transitions corresponding to nuclei with mass number $A = 132$ stem from long-lived decay products of a previous experiment, in which ^{132}Sn was used as the radioactive ion beam. This decays into ^{132}Te with a half-life of 3.204(13) days [68], resulting in persistent background contributions.

No evidence for decays of $A = 166$ isotopes was found, and no γ -ray transitions from other $A = 130$ nuclei were identified beyond those of ^{130}Sn and ^{130}Sb . However, stable beam components cannot be detected through the background spectra alone. To further investigate this, particle-gated and Doppler-corrected γ -ray spectra were analyzed. These are shown in Fig. 6.17, Doppler-corrected for ejectile- (red) and recoil-like particles (black).

Apart from known transitions in ^{130}Sn and ^{206}Pb , no additional peaks appear sharp in the spectra. Given that ^{130}Sn is expected to be the least collective nucleus along the $A = 130$ isobaric chain, the excitation probability into 2_1^+ states should be higher in other isotopes. The absence of de-exciting γ -rays from these isotopes suggests that contributions from the three stable $A = 130$ isotopes — ^{130}Te , ^{130}Xe , and ^{130}Ba — are negligible.

This interpretation is supported by chemical considerations: ^{130}Xe , being a noble gas, does not form stable molecules with sulfur; ^{130}Ba is not produced in the primary target used, according to the ISOLDE Yield Tables; and ^{130}Te is not expected to form molecular bonds with sulfur either. The Doppler correction was performed for the second time assuming a ^{162}Dy beam, the only stable $A = 162$ isotope. Peaks at expected energies of 80 keV ($2_1^+ \rightarrow 0_{\text{g.s.}}^+$) and 180 keV ($4_1^+ \rightarrow 2_1^+$) cannot be observed. Since ^{162}Dy is a midshell nucleus with high collectivity, excitation of the 2_1^+ and 4_1^+ states would have been expected, if the beam consisted partly of this isotope. At least the decay of the 4_1^+ should be visible in the spectra, whereas the 2_1^+ state decays most prominently by inner conversion. According to the ISOLDE yield tables, ^{162}Dy is not expected from a primary uranium target, which matches the experimental observations.

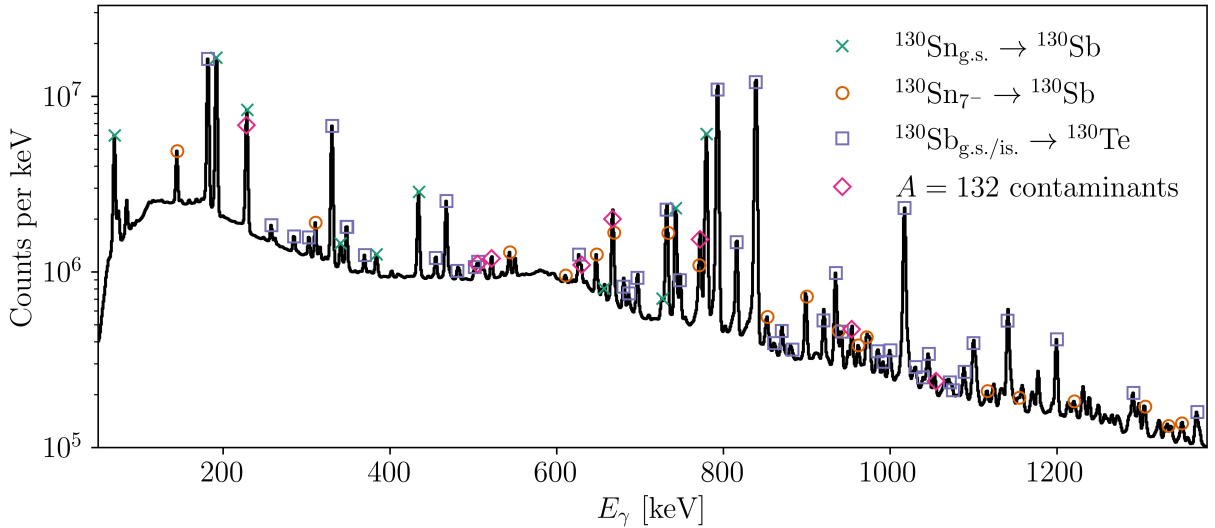


Figure 6.16. Summed γ -ray single spectra over the full duration of the experiment and all detectors. Transitions following the β -decay of ^{130}Sn and its daughter nuclei are indicated. Additional γ -ray lines from residual activity of a previous experiment using a ^{132}Sn beam and the same target chamber are also visible. These originate from long-lived decay products with half-lives of several days, resulting in persistent background contributions.

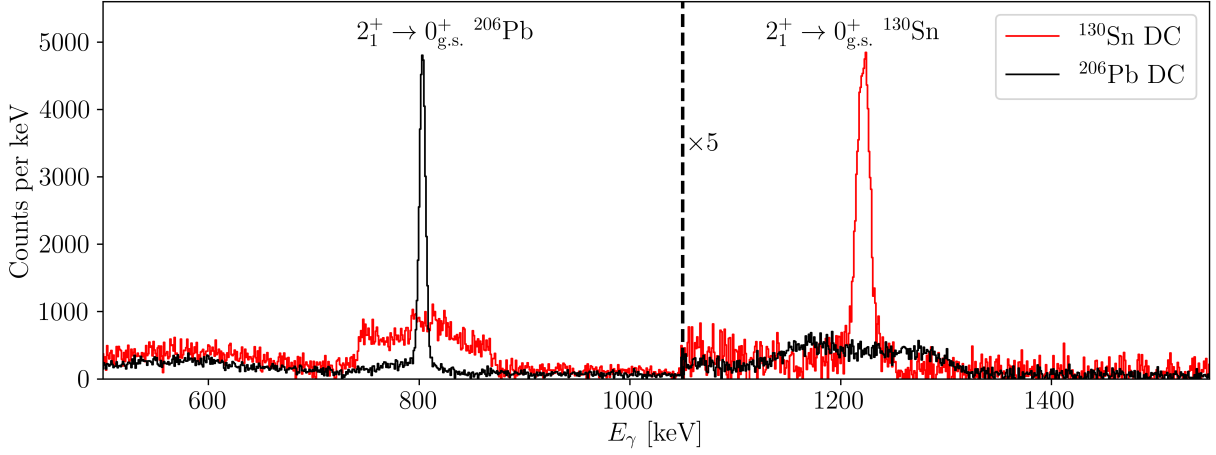


Figure 6.17. Doppler-corrected and background-subtracted γ -ray spectra. Gate on target- and beam-like nuclei are added up to maximize the statistics. Only the two $2_1^+ \rightarrow 0_{g.s.}^+$ transitions of ^{130}Sn and ^{206}Pb appear sharp. No excitation of any isobaric $A = 130$ contaminant is visible.

6.8. Ratio of ^{130}Sn and ^{130}Sb in Beam

Following the arguments given in the previous chapter, the remaining beam components are ^{130}Sn and ^{130}Sb . Both isotopes have an isomeric state, with half-lives in the range of minutes; thus, the possible beam components are $^{130}\text{Sn}_{g.s.}$, $^{130}\text{Sn}_{is.}$, $^{130}\text{Sb}_{g.s.}$, and $^{130}\text{Sb}_{is.}$. Since all of these isotopes are unstable, and their decays are connected, their β -delayed γ -ray yields can be used to determine the relative contribution. In Fig. 6.18 the decay scheme of these isotopes is displayed. In Appendix A, it is shown, that in the equilibrium state, the activity of the isotopes simplifies to the ion currents, and the beam composition can be analyzed by the activity of the isotopes. Since the absolute efficiency and dead time of the setup are unknown, only relative beam ratios can be determined, which is sufficient for the performed experiment.

6.8.1. Decay Spectra During Experiment

The composition of the radioactive beam was investigated by analyzing the time dependence of characteristic γ transitions during the experiment. As described in the previous chapter, the intensity ratios of selected transitions, corrected for detection efficiency, allow for a quantitative determination of the relative contributions of individual isotopes in the beam.

The time-resolved γ -ray spectra were obtained by projecting the energy-versus-time matrix in intervals of 20 s. For each resulting spectrum, Gaussian functions were fitted to extract the peak areas of transitions associated with the decay of ^{130}Sn and ^{130}Sb in their ground

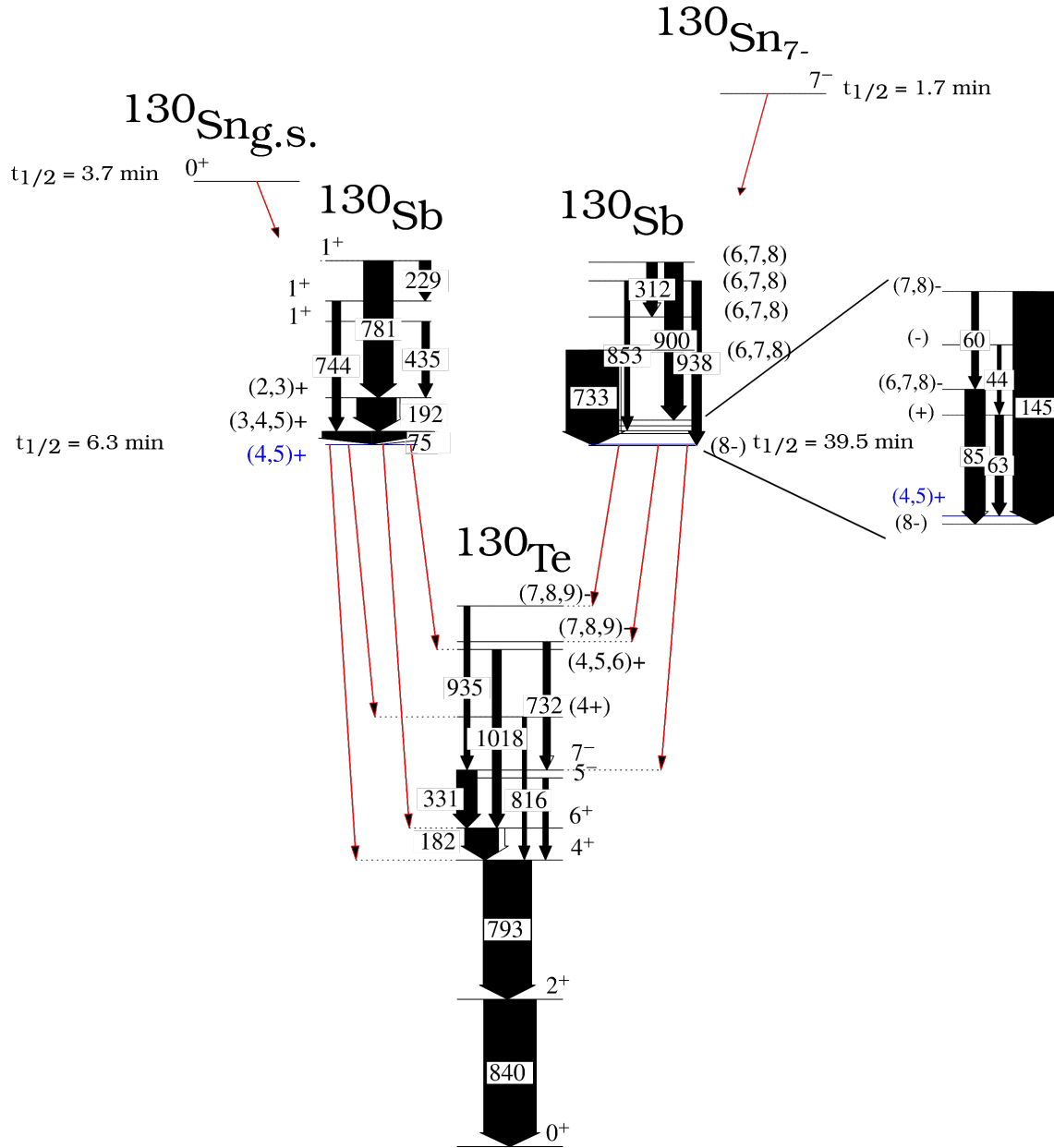


Figure 6.18. Decay scheme of $^{130}\text{Sn}_{\text{g.s.}}$, $^{130}\text{Sn}_{\text{is.}}$ and their daughter nuclei. γ -ray transitions with relative intensity $>10\%$ are shown. Arrow width indicates the relative intensities of γ rays with respect to the most prominent γ rays and are taken from [142] for excited states in ^{130}Sb and from [143] for excited states of ^{130}Te . Energies are taken from the NNDC data center [68]. β -decay lines (in red) indicate which states in ^{130}Te are populated by the two different decay branches. Note that the set of β -decay transitions shown is not exhaustive. It is intended only to highlight those states that are populated exclusively via either the left or the right decay branch. The isomeric $(4,5)^+$ state in ^{130}Sb is highlighted in blue to visualize the population by both, the decay of $^{130}\text{Sn}_{\text{g.s.}}$ and $^{130}\text{Sn}_{\text{is.}}$.

and isomeric states. The corresponding decay schemes are shown in Fig. 6.18, and the energy-time matrix is discussed in detail in chapter 6.2. Unlike the calibration study, the focus here lies on the evolution of peak intensities rather than peak positions.

The ground state of ^{130}Sn ($^{130}\text{Sn}_{\text{g.s.}}$) undergoes β -decay to excited states in ^{130}Sb , which decay further into the $(4,5)^+$ isomer. This isomer in turn decays to excited states in ^{130}Te . The isomeric 7^- state in ^{130}Sn populates negative-parity states in ^{130}Sb or directly feeds its (8^-) ground state. Only a small branch, $y = 0.045(5)\%$, connects the negative-parity cascade to the positive-parity isomeric $(4,5)^+$ state of ^{130}Sb , which is highlighted in blue in Fig. 6.18.

Both, the ground and isomeric states of ^{130}Sb , decay further into excited states of ^{130}Te . As all decay chains ultimately populate the first 2^+ state in ^{130}Te , the 839.5 keV transition

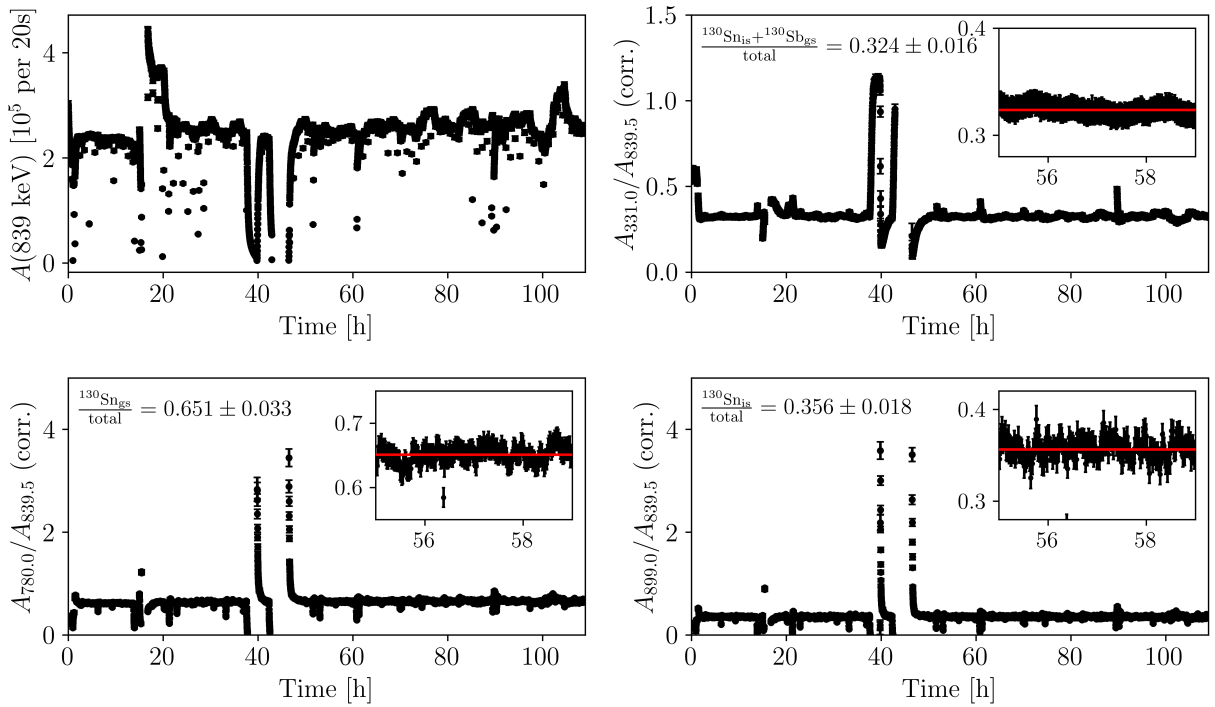


Figure 6.19. Top left: Fitted peak intensity of the 839.5 keV transition in ^{130}Sn as a function of time. The effects of extended beam interruptions are visible as reduced intensities. Approximately 20 h after the start of the experiment, the proton current was increased and subsequently reduced again due to elevated rates in the HPGe detectors. Top right, bottom left, bottom right: Ratios of the 839.5 keV intensity to the 331 keV, 780 keV, and 899 keV transitions, respectively. In each panel, the inset highlights the time interval between 55 and 59 h after the beginning of the experiment, where the intensity ratios appear stable. The resulting average values, including their uncertainties, are indicated in the main plots. For further details, see text.

serves as a proxy for the total activity of all beam components. The fitted intensity of the 839.5 keV line, displayed in the top-left panel of Fig. 6.19, shows characteristic fluctuations due to changes in beam conditions. Notably, a significant increase in intensity is observed around 20 h into the experiment, corresponding to a temporary rise in proton current, which was later reduced due to high detector rates. The remaining panels in Fig. 6.19 display the intensity ratios of the 331 keV ($^{130}\text{Sb}_{\text{is.}} \rightarrow ^{130}\text{Te}$), 780 keV ($^{130}\text{Sn}_{\text{g.s.}} \rightarrow ^{130}\text{Sb}$), and 899 keV ($^{130}\text{Sn}_{\text{is.}} \rightarrow ^{130}\text{Sb}$) transitions to the 839.5 keV line. Short beam interruptions lead to temporary shifts in these ratios, but in all cases, the values tend toward an equilibrium over time. A detailed view of the quasi-stationary regime between 55 and 59 h is shown in the inset panels. By using these isolated transitions that are uniquely associated with a specific parent isotope and applying efficiency corrections, the relative beam composition can be deduced. The mean values in this interval, including the squared sum of their statistical and systematic uncertainties, are indicated in the main panels. Systematic uncertainties arise primarily from efficiency corrections and the adopted absolute γ -ray intensities taken from Ref. [68].

The lower spectra give the relative ratio of $^{130}\text{Sn}_{\text{g.s.}}$ (left) and $^{130}\text{Sn}_{\text{is.}}$ (right). Similarly to the experiment conducted in 2022, this already amounts to 100.7(36) % of the beam, showing no signs of antimony in the beam. As further proof, the 331 keV ratio compared to the 839.5 keV line is shown on the top right, which is populated by the ground-state decay of ^{130}Sb . The latter is fed by $^{130}\text{Sn}_{\text{is.}}$ plus the relative amount of $^{130}\text{Sb}_{\text{g.s.}}$ in the beam. The deduced value for $^{130}\text{Sn}_{\text{is.}} + ^{130}\text{Sb}_{\text{g.s.}}$ shown in the top right panel is lower than the value for $^{130}\text{Sn}_{\text{is.}}$ alone. This stems from the value of $y = 0.045(5)$, which is the ratio of $^{130}\text{Sn}_{\text{is.}}$ decaying into $^{130}\text{Sb}_{\text{is.}}$. According to this analysis, the beam consists of 65(3)% $^{130}\text{Sn}_{\text{g.s.}}$ and 35(2)% $^{130}\text{Sn}_{\text{is.}}$. An ^{130}Sb component in the beam could thus not be identified.

6.8.2. Decay Spectra After Experiment

After the experiment was completed, the decay of radioactive isotopes remaining in the target chamber was monitored as a function of time. This provides an independent characterization of the beam composition by analyzing the time dependence of the emitted γ -ray transitions. The left panel of Fig. 6.20 shows the γ -ray energy plotted against time after the end of the experiment. The last in-beam run is also shown to the left of $t = 0$, which marks the time when the beam was stopped. A short measurement gap around $t = 0$ appears as a white band in the plot, reflecting the time between the last beam event and the start of the decay spectroscopy.

For several transitions from nuclei with a short lifetime — such as the line at 1017 keV, fed by $^{130}\text{Sb}_{\text{g.s.}}$ with $T_{1/2} = 6.3$ min — a rapid decay is clearly visible. The right panel of Fig. 6.20 shows the time evolution of four selected γ -ray peaks. For each time, a Gaussian function was fitted to the respective energy peak, and the peak area was extracted as a function of time.

As shown in the level scheme in Fig. 6.18, the 899 keV transition is fed by the decay of $^{130}\text{Sn}_{\text{is.}}$. By fitting a single exponential function to the decay curve, the activity at $t = 0$ as well as the decay constant—and thus the half-life—can be determined. As shown in Chapter 6.8, the activity at $t = 0$ is directly proportional to the amount of the isomeric isotope present in the beam. The 780 keV line is populated via the decay of $^{130}\text{Sn}_{\text{g.s.}}$, and was likewise fitted with a single exponential decay function. The 331 keV line is emitted in the decay of $^{130}\text{Sb}_{\text{g.s.}}$, which itself is populated via the decay of $^{130}\text{Sn}_{\text{is.}}$. To describe this decay chain, the following two-step decay model is used:

$$A_{\text{Sb}}(t) = \lambda_{\text{Sb}} \left[N_{\text{Sb},0} \cdot e^{-\lambda_{\text{Sb}}t} + \frac{\lambda_{\text{Sn}} N_{\text{Sn},0}}{\lambda_{\text{Sb}} - \lambda_{\text{Sn}}} (e^{-\lambda_{\text{Sn}}t} - e^{-\lambda_{\text{Sb}}t}) \right], \quad (6.4)$$

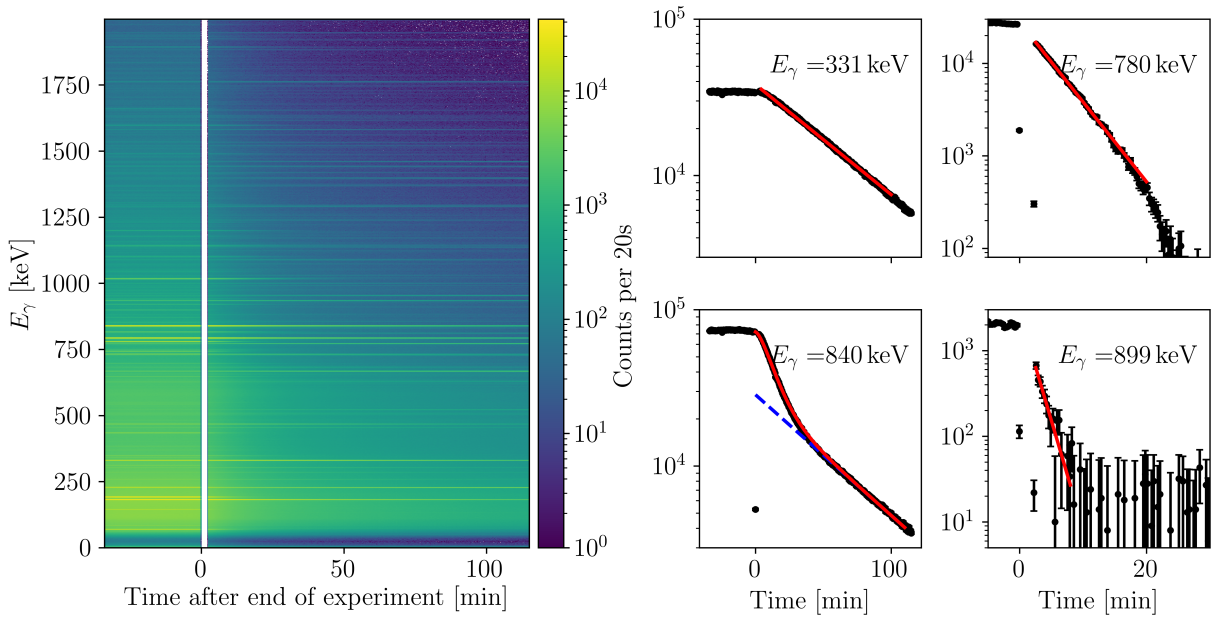


Figure 6.20. Left: γ -ray energy as a function of time after the end of the experiment. The last run before the beam was stopped is shown to the left of $t = 0$. The gap around zero reflects a short interruption between the last in-beam run and the beginning of the decay spectroscopy. Right: Time-dependent decay curves for four selected peaks. See text for details. The deviating points at $t \approx 0$ arise from finishing and starting the respective runs and were excluded in the fit.

Table 6.3. Results of the fits are shown in Fig. 6.20. For each isotope, the activity at $t = 0$ was determined as well as the half-life. Since the decay chain is known precisely, the ratios can be deduced by looking at the activities at $t = 0$. For more information, see text.

Isotope	$A(t = 0)$ [1/20s]	λ [1/min]	$T_{1/2,\text{exp.}}$ [min]	$T_{1/2,\text{lit.}}$ [min]	Ratio in beam
$^{130}\text{Sn}_{\text{g.s.}}$	60600(900)	0.192(5)	3.61(8)	3.72(4)	0.67(4)
$^{130}\text{Sn}_{\text{is.}}$	30120(1200)	0.0170(5)	40.3(12)	39.5(8)	0.33(4)
$^{130}\text{Sb}_{\text{g.s.}}$	28560(800)	0.45(5)	1.54(19)	1.7(1)	-0.003(20)
$^{130}\text{Sb}_{\text{is.}}$	62700(1000)	0.115(5)	6.02(10)	6.3(2)	0.008(20)

where λ_{Sn} and λ_{Sb} denote the decay constants of ^{130}Sn and ^{130}Sb , respectively, $N_0 = A(0)/\lambda$ is the number of nuclei at $t = 0$, and $A(t)$ is the measured activity. The first term accounts for directly present ^{130}Sb , while the second term describes the contribution from the decay of ^{130}Sn .

Since the decay constant and initial activity of $^{130}\text{Sn}_{\text{is.}}$ are already known from the fit to the 899 keV decay, the function can be fitted to the data with only two free parameters: the activity and decay constant of $^{130}\text{Sb}_{\text{g.s.}}$. As the influence of $N_{\text{Sn},0}$ and λ_{Sn} on the shape of the decay curve is weak—clearly visible in the nearly pure exponential form of the 331 keV decay—the scaling factor y introduced in section 6.8 is neglected here. The resulting fit provides the activity of $^{130}\text{Sb}_{\text{g.s.}}$ at $t = 0$ and its decay constant.

The 839.5 keV transition receives contributions from both $^{130}\text{Sb}_{\text{g.s.}}$ and $^{130}\text{Sb}_{\text{is.}}$. However, since the ground-state decay is significantly slower, the late-time behavior (i.e., at $t > 70$ min) is dominated by $^{130}\text{Sb}_{\text{g.s.}}$. Therefore, an exponential function was first fitted to the late-time tail using the decay constant previously obtained from the 331 keV fit. This contribution is shown as a dashed blue line in Fig. 6.20. After subtracting this component, Equation 6.4 was fitted to the remaining early-time data ($t < 70$ min). Here, both the decay constant and activity of $^{130}\text{Sn}_{\text{g.s.}}$ are known from the 780 keV fit, such that only the activity and decay constant of $^{130}\text{Sb}_{\text{is.}}$ remain as free parameters.

The extracted values for the activities at $t = 0$, the decay constants, and the corresponding half-lives $T_{1/2} = \ln(2)/\lambda$ are listed in Table 6.3. The beam composition was deduced as follows: As discussed in the previous chapter, the activity of the 839.5 keV transition at $t = 0$ corresponds to the total γ -ray activity resulting from all beam components contributing to this line. By determining the individual contributions of each isotope to this total activity, the relative fractions of isotopes in the beam can be quantified.

The activities of $^{130}\text{Sn}_{\text{g.s.}}$ and $^{130}\text{Sn}_{\text{is.}}$ directly reflect the respective beam intensities, as both decay predominantly via well-separated transitions. The activity of $^{130}\text{Sb}_{\text{is.}}$, on the other hand, consists of three components: a direct contribution from $^{130}\text{Sb}_{\text{is.}}$ present in the beam, the decay of $^{130}\text{Sn}_{\text{g.s.}}$, and a fraction $y = 0.045(5)$ of the decay of $^{130}\text{Sn}_{\text{is.}}$ feeding into $^{130}\text{Sb}_{\text{is.}}$. Similarly, the measured activity of $^{130}\text{Sb}_{\text{g.s.}}$ at $t = 0$ originates from the remaining fraction $(1 - y)$ of the $^{130}\text{Sn}_{\text{is.}}$ decay as well as from a possible direct beam component of $^{130}\text{Sb}_{\text{g.s.}}$.

Based on the fitted values of all four activities at $t = 0$, a second, independent determination of the beam composition can be performed by solving the corresponding system of contributions. The resulting isotope fractions are listed in the last column of Table 6.3. Again, it is evident that the contribution from ^{130}Sb is negligible, while ^{130}Sn clearly dominates the beam composition with a fraction of 0.67(4). The final ratios of $^{130}\text{Sn}_{\text{g.s.}}$ and $^{130}\text{Sn}_{\text{is.}}$ in the beam were deduced from the weighted average of the two presented methods to 66(2) % and 34(2) %, respectively. This value matches ratios found in previous works, e.g. in Ref. [101].

6.9. $0_{\text{g.s.}}^+ \rightarrow 2_1^+$ Excitation in ^{130}Sn

Based on the data obtained through the steps outlined in the previous sections, experimental results can now be derived. The following section begins with a discussion of how the optimal Doppler-correction velocity was determined. Since a small contribution from two-step excitation into the 4_1^+ state is expected—whose 775 keV decay lies close to the 803 keV transition of the 2_1^+ state in ^{206}Pb —an accurate Doppler correction is essential.

6.9.1. Velocity Determination of Ejectiles and Recoils

In addition to the emission angle, the velocity v of the fast moving reaction partners plays a crucial role in the Doppler shift of the emitted γ ray. An accurate determination of the velocity is therefore essential for a proper Doppler correction. Various approaches have been tested in this work to optimize the correction and minimize the full width at half maximum (FWHM) of Doppler-shifted peaks.

Two standard methods are commonly used: extracting the velocity from the energy measured in the CD detector or calculating it from the known reaction kinematics. In the latter

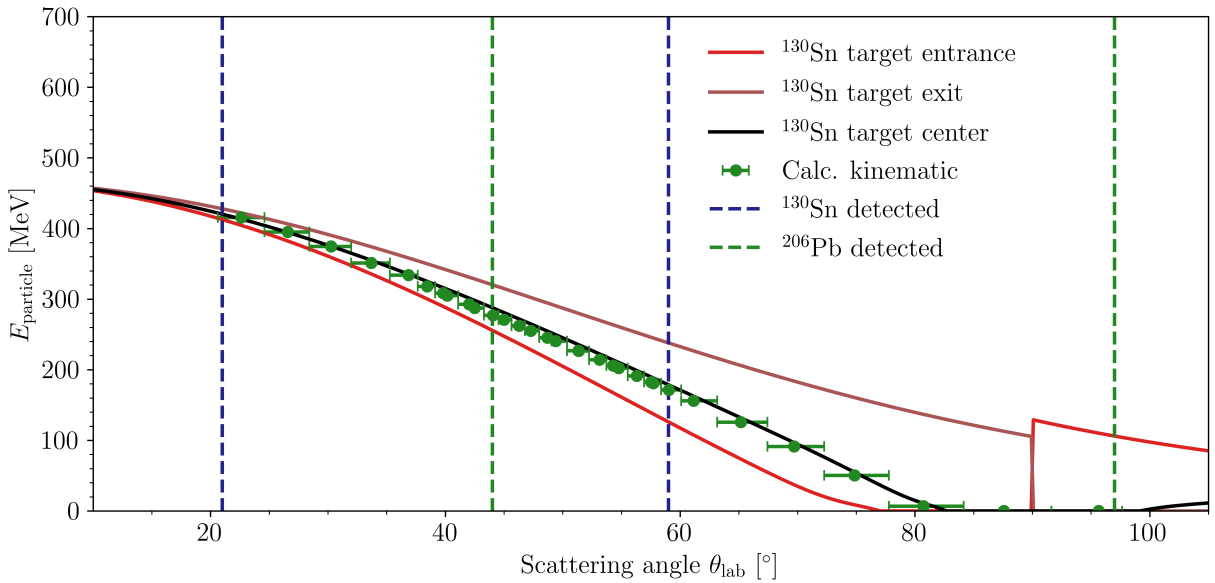


Figure 6.21. Energy of scattered ^{130}Sn nuclei after leaving the target as a function of the laboratory scattering angle. The red bands indicate the energy range for reactions occurring at the entrance (light red) and exit (dark red) of the target. The black line corresponds to reactions at the target center, as assumed by the Doppler-correction sorting code. Green points are values calculated by the code and agree well with the LISE-based kinematics. The blue and green shaded angular regions indicate where the CD detects the Sn ejectiles and Pb recoils, respectively.

case, the velocity can be deduced from the measured scattering angle in the CD and the known beam energy, assuming a two-body reaction. Since the lifetimes of the excited states are expected to exceed the flight time of the nuclei through the target, the γ -ray emission predominantly occurs in vacuum after the particles have exited the target.

Figure 6.21 shows the calculated energy of scattered ^{130}Sn nuclei as a function of the laboratory scattering angle θ_{lab} . The red-shaded areas represent the range of energies for reactions occurring at the target entrance (upstream side of the target) in light red and target exit (the downstream side of the target) in dark red. The variation arises due to differences in energy loss, which depends on both the initial energy and the effective path length through the target material. Reactions occurring near the entrance of the target result in larger energy losses compared to those occurring near the exit. The black line shows the energy for reactions assumed to take place in the center of the target. The energy of scattered Sn was calculated using the LISE [132] simulation code. The green points represent the energies calculated by the sorting code for the angles given from the individual rings of the CD. These agree well with the black line.

The blue and green vertical bands mark the angular ranges in which the CD detects ^{130}Sn ejectiles and ^{206}Pb recoils, respectively. In the latter case, the kinematics of the undetected Sn nuclei can be inferred from the measured recoil properties.

Despite the good agreement between the sorting code and the central kinematics, it is evident that these calculated values do not represent the average kinetic energy of the ejectiles after the target. This is due to two main effects: first, the energy loss in the target is not linear with depth but depends on the actual energy of the particle. Second, the effective path length through the target does not scale linearly with the reaction position. These effects become particularly pronounced at large scattering angles, especially around 90° , where particles originating from the center of the target are stopped completely, while those from the back side may exit with energies above 100 MeV.

Additionally, as discussed in Section 6.6, parts of the kinematic lines are truncated due to the overlap of the observed recoil and ejectile peaks in the CD spectra, which introduces a systematic shift in the observed scattered energies towards higher values.

An alternative and often more reliable approach is to use the energy measured directly in the CD, which naturally accounts for the effective energy loss without requiring assumptions on the reaction point inside the target. However, as discussed in Section 6.3, the energy

calibration of the CD detector is not straightforward. The main focus of the calibration procedure was to minimize deviations between the four quadrants and to enable a reliable treatment of charge-sharing events between neighboring rings and strips. A precise and absolute calibration that yields the correct particle energy was not achievable.

Several effects, such as energy loss in the detector's dead layer and the PHD, have so far been neglected, although they would need to be accounted for in a fully accurate calibration. In addition, as shown in Fig. 6.15, the observed peaks are very broad. This broadening arises both from the relatively thick target foil used in the experiment (4.2 mg/cm^2) and from the intrinsic energy resolution of the CD detectors.

Moreover, the measured energy can exclusively be used for Doppler correction of the particle that was actually detected. It provides no direct information about the velocity of the undetected reaction partner. Since target-based normalization is applied in the present analysis, it is essential to perform proper Doppler correction for both reaction products. This underscores the need for a consistent and accurate velocity determination strategy.

An alternative approach was adopted in this work, relying directly on the experimental data. By evaluating particle- γ coincidences for each individual ring of the CD detector, two-dimensional matrices were generated that display the Doppler-corrected γ -ray energy as a function of an assumed velocity of the scattered particle in vacuum. Four representative matrices are shown in Fig. 6.22, each corresponding to a kinematic gate on ^{130}Sn . The top row illustrates Doppler corrections performed under the assumption that the recoiling particle is ^{206}Pb , while the bottom row assumes the ejectile itself to be ^{130}Sn .

In all cases, a distinct minimum in energy broadening is observed, reflecting the optimal velocity for the Doppler correction. For each velocity step, a spectrum was extracted, and a Gaussian function was fitted to the photopeak. The FWHM of the fit function was then used as a metric to determine the best estimate for the velocity v at which the nucleus decays.

The best fits obtained in this procedure are presented in Figs. 6.23, 6.24, 6.25, and 6.26, representing Doppler corrections for both beam-like and target-like nuclei under kinematic gates on ^{130}Sn and ^{206}Pb , respectively. The resulting minimal FWHMs are summarized in the bottom-right panel. For the ^{130}Sn ejectile, a resolution of 14 keV was achieved when gated on its kinematics and between 15 and 20 keV when gated on the ^{206}Pb kinematics. The 803 keV transition in ^{206}Pb yielded a remarkably narrow FWHM of 4–8.5 keV under

a self-gated Doppler correction, which is an exceptional result for a radioactive ion beam experiment with Miniball. Even when gated on the ^{130}Sn kinematics, the resolution remained below 10 keV in most cases.

All three explained methods for the velocity determination have been tested; the last one showed overall the best results for the FWHM and was adopted for the whole analysis.

For small scattering angles of both Sn and Pb nuclei, the undetected particle in the laboratory frame receives a large momentum transfer. As a result, nuclei excited near the entrance of the target may be stopped before decaying, whereas those scattered near the back of the target might exit before emitting γ rays. This effect, previously illustrated in

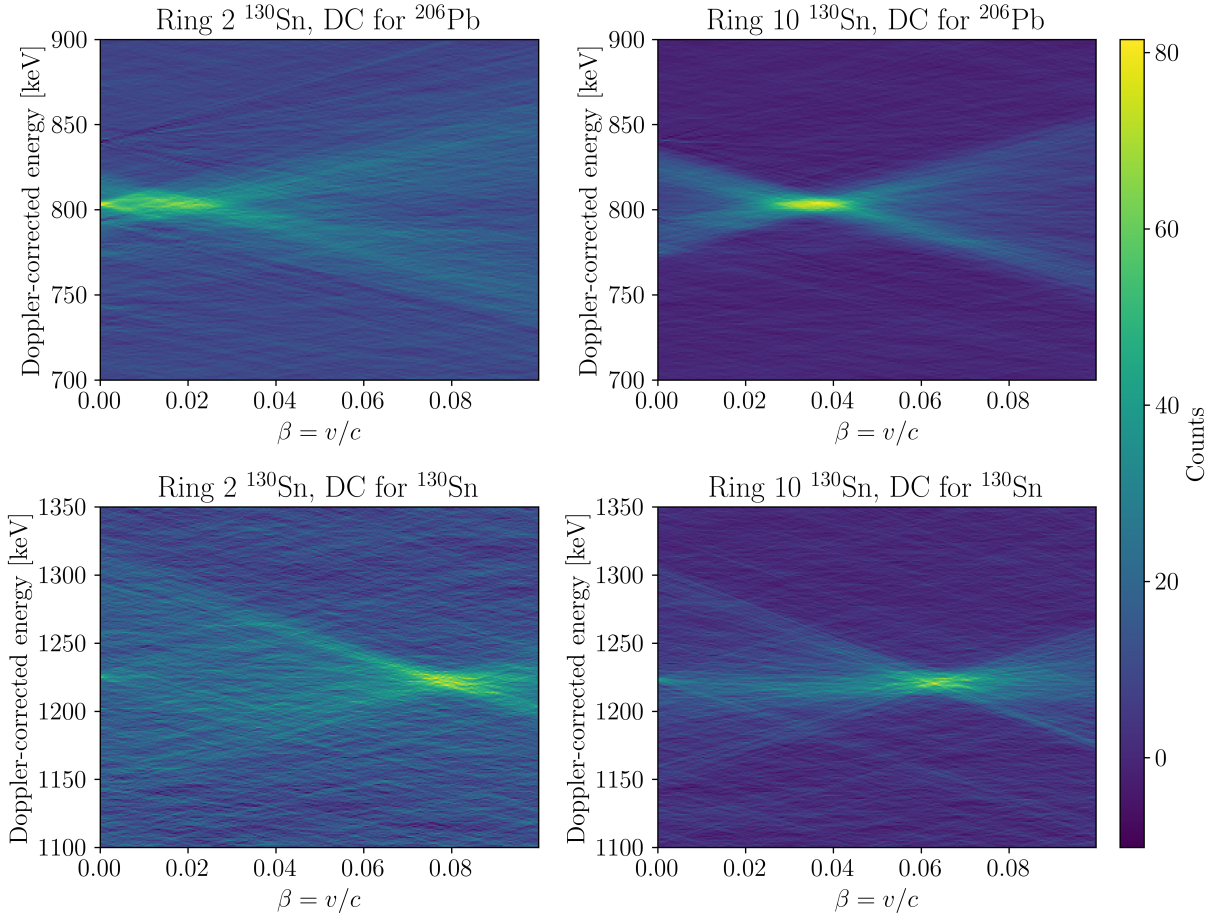


Figure 6.22. Doppler-corrected γ -ray energy for different assumed velocities. Left: Gate was placed onto the kinematics of ^{130}Sn detected in ring 2 of the CD detector. Doppler-correction performed for ^{206}Pb (top) and ^{130}Sn (bottom), respectively. On the right, the same matrices are shown with a gate on ring 10 of the CD.

Fig. 6.21, causes a spread in the actual recoil velocities ranging from 0 to 2 % of the speed of light, which prevents the formation of a clear minimum in the velocity matrix of Fig. 6.22.

Consequently, these CD rings were excluded from the final analysis, as a reliable determination of the peak integrals was not feasible. The same effect is visible in the FWHMs extracted from Gaussian fits, which exhibit abrupt variations between neighboring CD rings.

For instance, in Fig. 6.25, the fitted FWHM of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in ^{130}Sn at 1220 keV is only 5 keV for small scattering angles — an unrealistically narrow value — because some nuclei decay at rest. However, the fitted Gaussian did not reproduce the actual Doppler-broadened line shape. A closer inspection of the lower matrices in Fig. 6.22 reveals a subtle structure at zero velocity in both rings. This can be attributed to the population of the isomeric 7^- component of the ^{130}Sn beam, which is excited in the experiment into a 5_1^- state. This state decays via a cascade involving the 4_1^+ and 2_1^+ states. Due to its relatively long lifetime, it decays well after the time-of-flight to the particle detector, producing sharp, unbroadened γ -ray peaks in the laboratory frame.

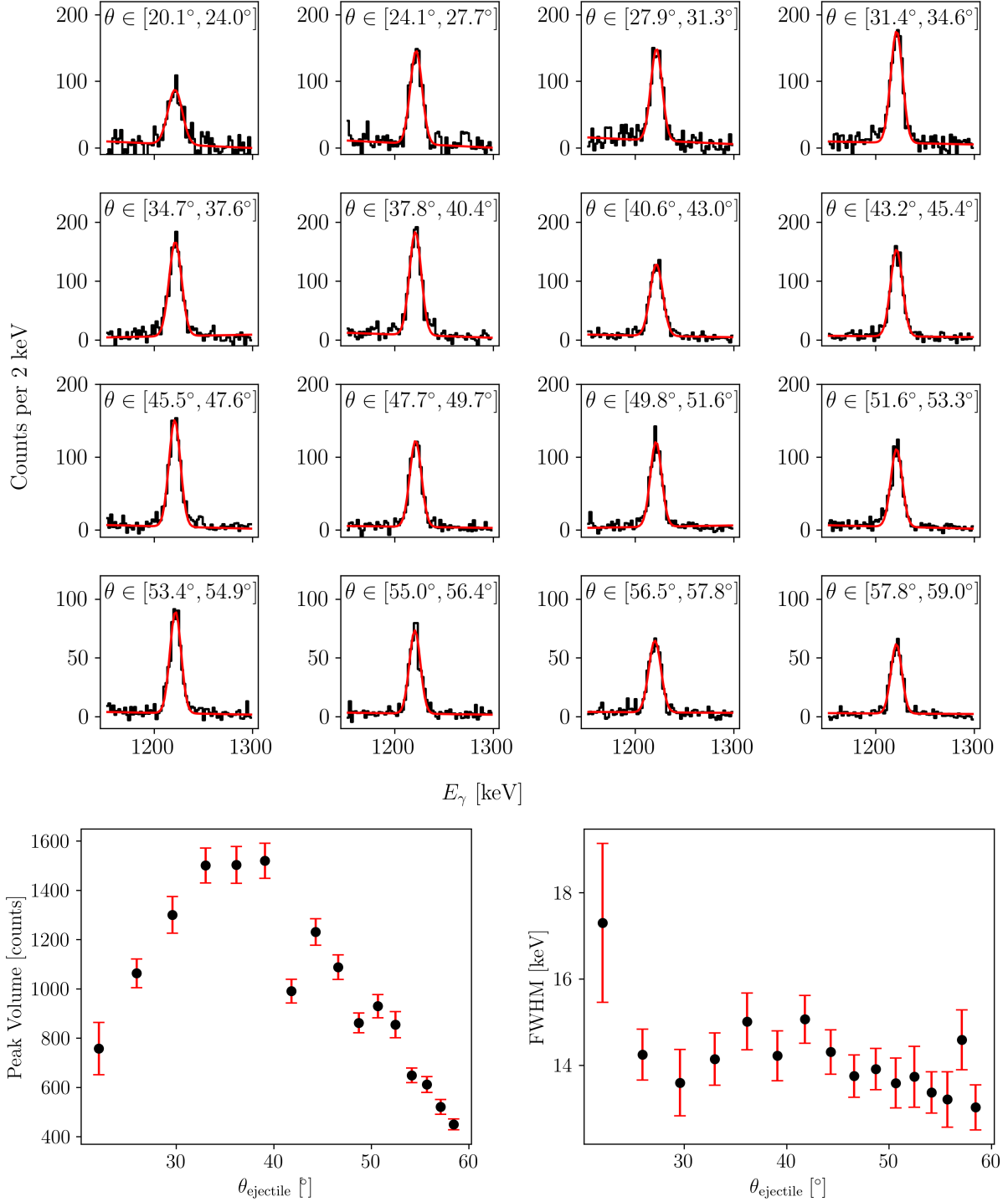


Figure 6.23. Doppler-corrected γ -ray spectra gated on beam-like reaction products and corrected for ^{130}Sn . Events were selected by applying an angular gate to a single ring of the CD detector, thereby restricting the ejectile emission angles to a narrow range. The red curve represents a Gauss fit plus a linear background applied to the most prominent photopeak. The resulting peak volume and full width at half maximum (FWHM) extracted from the fit are shown in the two lower panels as a function of the ejectile angle.

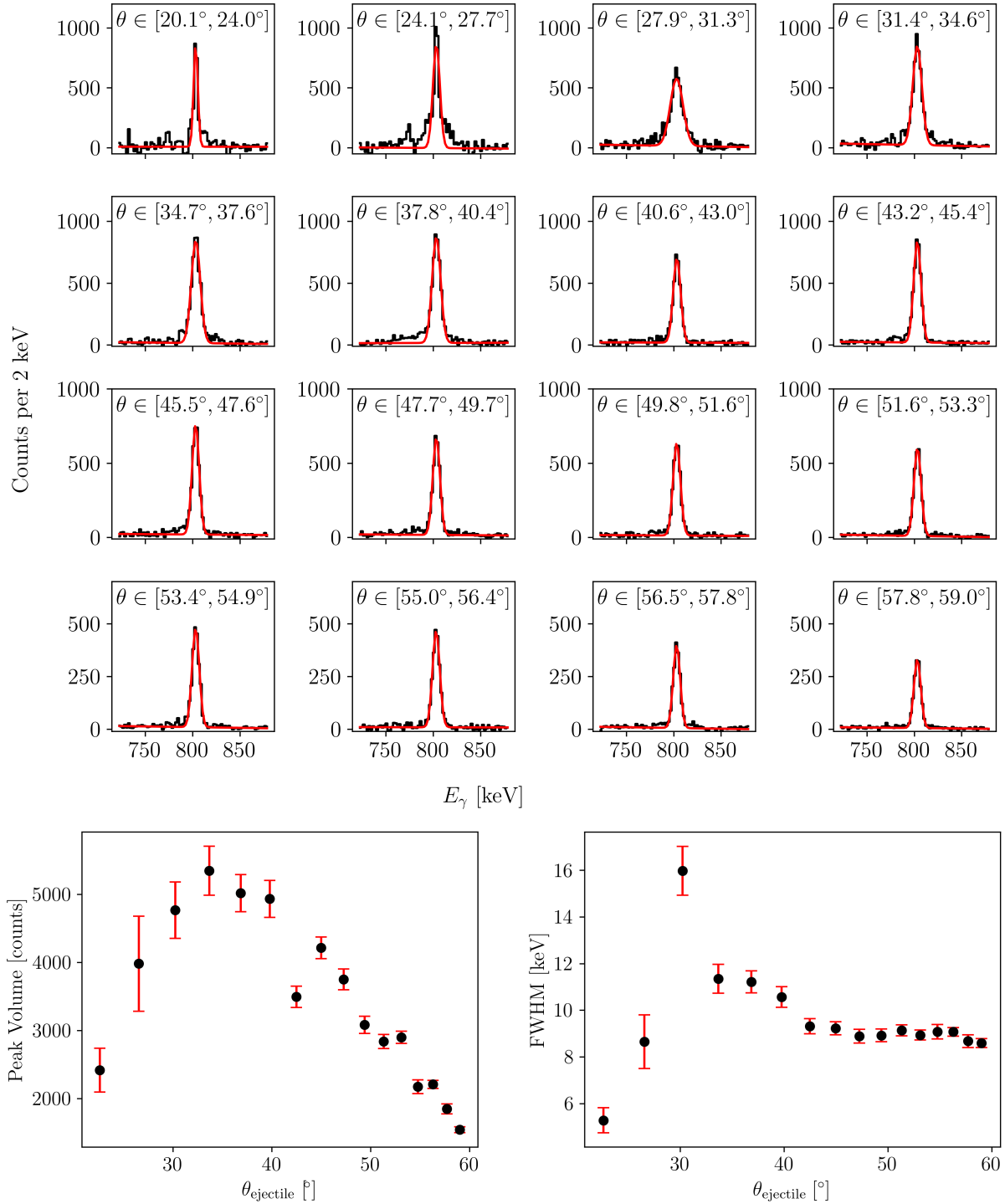


Figure 6.24. Doppler-corrected γ -ray spectra gated on beam-like reaction products and corrected for ^{206}Pb . Events were selected by applying an angular gate to a single ring of the CD detector, thereby restricting the ejectile emission angles to a narrow range. The red curve represents a Gauss fit plus a linear background applied to the most prominent photopeak. The resulting peak volume and full width at half maximum (FWHM) extracted from the fit are shown in the two lower panels as a function of the ejectile angle.

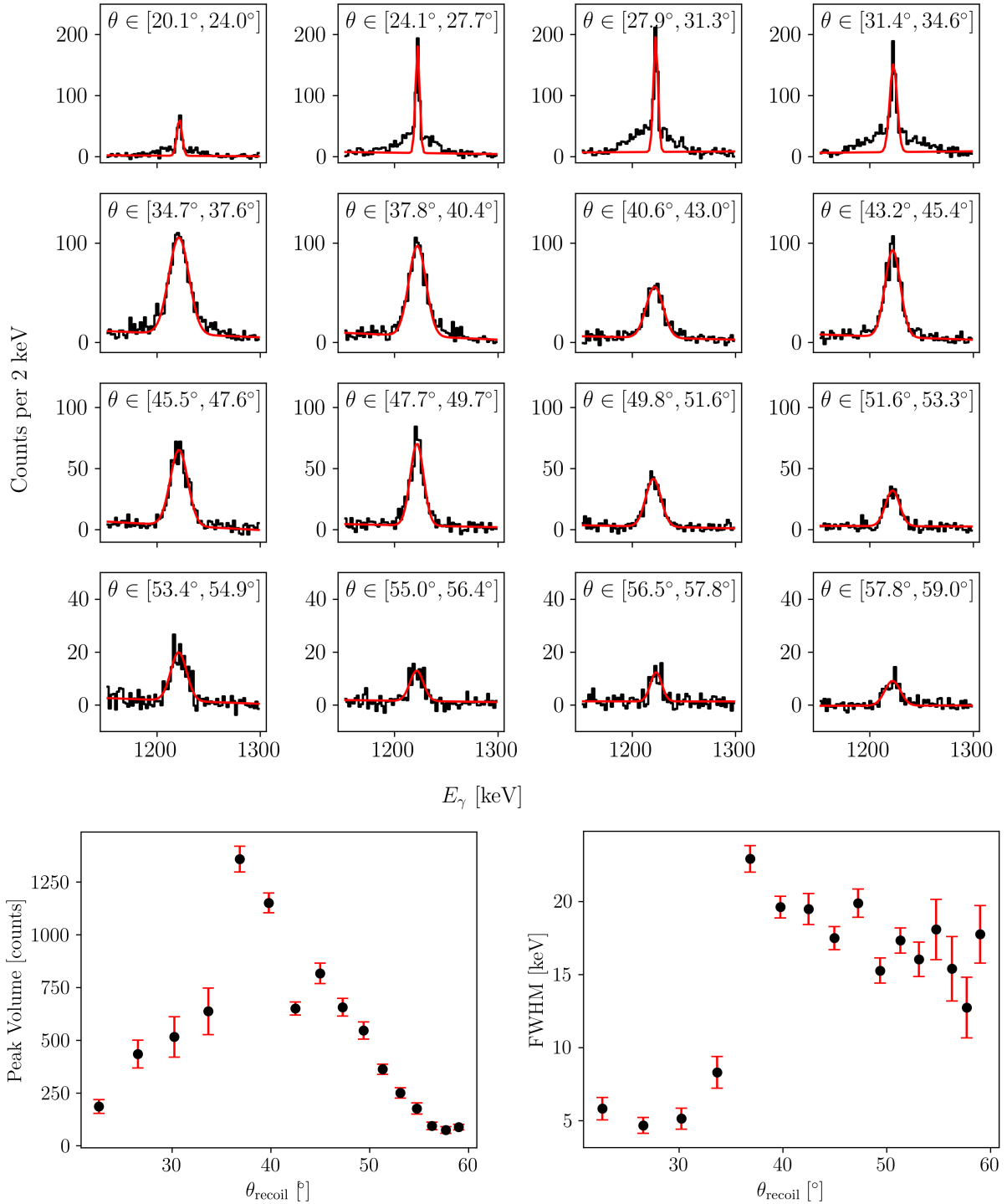


Figure 6.25. Doppler-corrected γ -ray spectra gated on target-like reaction products and corrected for ^{130}Sn . Events were selected by applying an angular gate to a single ring of the CD detector, thereby restricting the ejectile emission angles to a narrow range. The red curve represents a Gauss fit plus a linear background applied to the most prominent photopeak. The resulting peak volume and full width at half maximum (FWHM) extracted from the fit are shown in the two lower panels as a function of the ejectile angle.

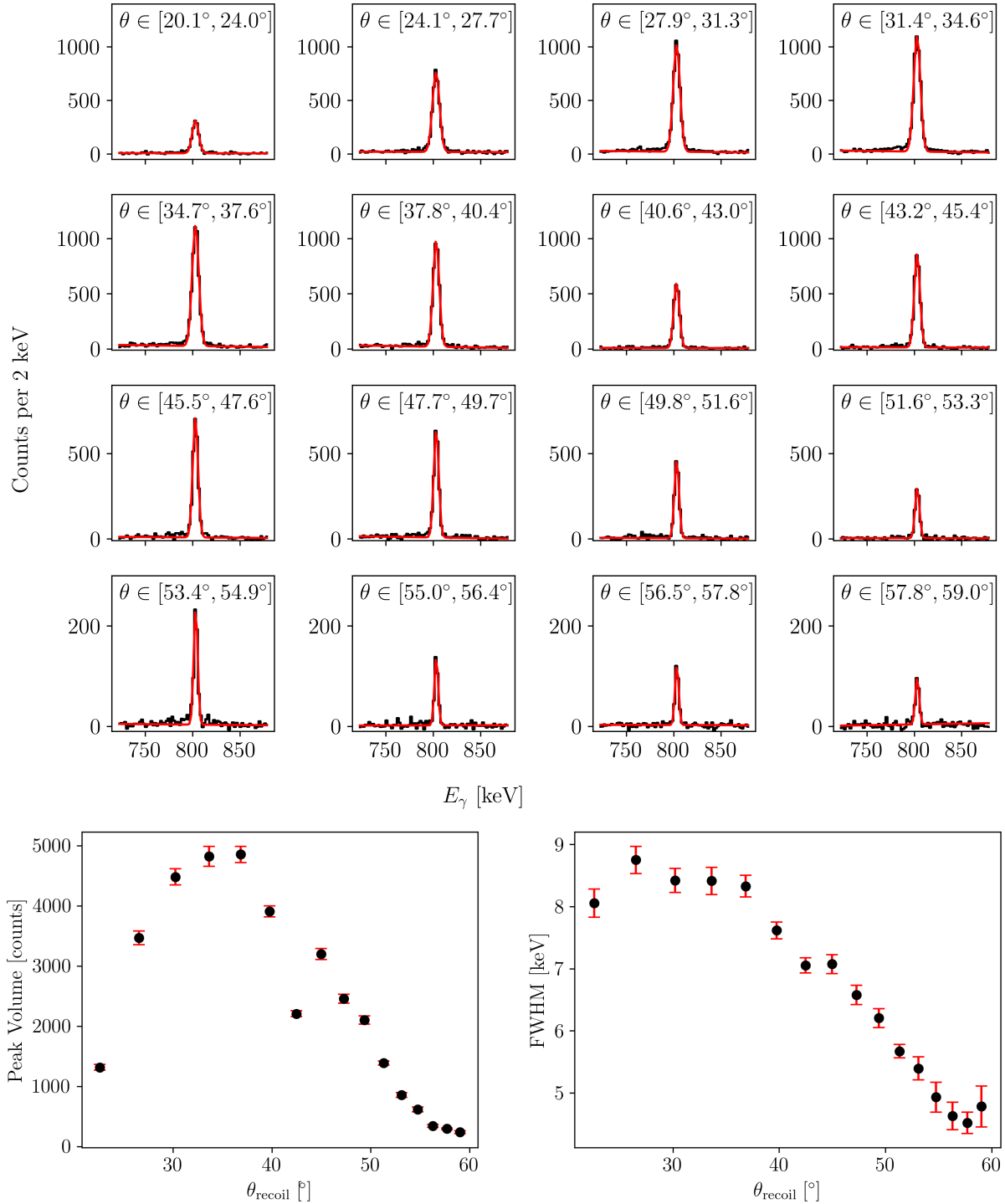


Figure 6.26. Doppler-corrected γ -ray spectra gated on target-like reaction products and corrected for ^{206}Pb . Events were selected by applying an angular gate to a single ring of the CD detector, thereby restricting the ejectile emission angles to a narrow range. The red curve represents a Gauss fit plus a linear background applied to the most prominent photopeak. The resulting peak volume and full width at half maximum (FWHM) extracted from the fit are shown in the two lower panels as a function of the ejectile angle.

6.9.2. Determination of γ -ray yields

In the lower left panels of Figs. 6.23, 6.24, 6.25, and 6.26, the fitted peak volumes of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transitions in ^{130}Sn and ^{206}Pb are shown as a function of the scattering angle of the detected particle. In total, 32 individual fits were performed—16 for each isotope per gate. However, several fits were not included in the final analysis for the following reasons.

In the case of spectra where the ejectile was detected, the first three rings were excluded. At small laboratory scattering angles of the ejectile, the corresponding recoil angles are large. As previously discussed, this leads to strong variations in energy loss: some recoils are stopped in the target, while others pass through it. This behavior results in pronounced discrepancies in the FWHM of the γ -ray peaks, particularly for the first three rings. Only those rings for which a reliable recoil velocity could be determined—based on the method described in the preceding chapter—were included in the analysis. For the remaining 13 rings, the fits show consistent FWHM values and smoothly varying intensities as a function of angle. A notable exception is ring 7, where a sudden drop in intensity is observed due to a non-functional strip in one of the CD detector sectors. Nevertheless, since all angular bins are independently treated in the GOSIA2 analysis, this deviation does not preclude the inclusion of this data point.

The situation is more complex when ^{206}Pb recoils are detected instead of the ejectile. Again, small laboratory angles of the recoil correspond to large angles of the ejectile. As illustrated in Fig. 6.21, for scattering angles $\theta_{^{130}\text{Sn}} > 75^\circ$ (i.e., $\theta_{^{206}\text{Pb}} < 33.8^\circ$) some ^{130}Sn ejectiles are already fully stopped in the target material. For $\theta_{^{130}\text{Sn}} > 80^\circ$ most ejectiles are stopped before they can exit the target. In these cases, the velocity used for the Doppler correction is effectively zero, which does not allow to distinguish whether the observed 2_1^+ γ -ray decay follows direct population or originates from a cascade via the 5_1^- state populated by excitation of the 7_1^- beam component. As will be explained in section 6.11, the 5_1^- state is long-lived, and its own and all following decays appear unshifted in the spectra. Therefore, all data points beyond $\theta_{^{130}\text{Sn}} > 80^\circ$ in this angular region were excluded from the analysis.

At larger scattering angles of ^{206}Pb recoils (i.e., $\theta_{^{206}\text{Pb}} > 44^\circ$), the coincident ^{130}Sn ejectiles are also within the acceptance of the CD detector. In this range, both reaction partners can be detected simultaneously. Since each γ -ray can only be attributed to one kinematic configuration—either that of the recoil or the ejectile—to avoid double counting, these coincident events are consistently assigned to the ejectile channel.

The final intensities used in the analysis are listed in Table 6.4 as a function of the scattering angle in the laboratory system. For the measured $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transitions in ^{130}Sn and ^{206}Pb , the quoted uncertainties are purely statistical. The ^{206}Pb intensities have been corrected for the detection efficiency, which varies with energy. In addition, the values were scaled by a factor of 0.66(2), reflecting the fact that the beam consisted of 66(2) % ground-state ^{130}Sn . Since Coulomb excitation of ^{206}Pb is independent of whether the tin nucleus is in its ground or isomeric state, this factor can be applied directly. The uncertainties of the energy-dependent efficiency and the beam composition were adopted here.

6.9.3. Two-step Excitation

Coulomb excitation is not limited to the population of the first excited states, higher-lying states can also be populated via multi-step excitation processes. In the present experiment, the second excited state of interest is the 4_1^+ state in ^{130}Sn at an excitation energy of 1991 keV. This state decays exclusively to the 2_1^+ state via a 775 keV γ -ray transition.

In the Doppler-corrected γ -ray spectra, no peak corresponding to the $4_1^+ \rightarrow 2_1^+$ transition

Table 6.4. Determined counts in the given angular range for the $2^+ \rightarrow 0^+$ transitions in ^{130}Sn and ^{206}Pb . For the latter, the measured counts as well as counts corrected for the beam composition and efficiency are shown. Only intensities used in the analysis are given here.

Number	$\theta_{^{130}\text{Sn, lab.}}$	^{130}Sn counts	^{206}Pb measured	^{206}Pb corrected
1	31.45°–34.56°	1501(71)	5345(360)	2822(220)
2	34.72°–37.71°	1503(74)	5017(273)	2649(178)
3	37.76°–40.43°	1520(71)	4932(271)	2604(176)
4	40.57°–43.03°	991(48)	3494(157)	1845(110)
5	43.16°–45.43°	1231(54)	4214(160)	2225(122)
6	45.54°–47.64°	1088(50)	3750(153)	1980(112)
7	47.74°–49.68°	862(40)	3083(126)	1628(92)
8	49.77°–51.56°	930(47)	2837(103)	1498(80)
9	51.65°–53.29°	855(53)	2898(90)	1530(77)
10	53.38°–54.90°	649(29)	2174(101)	1148(70)
11	54.98°–56.39°	612(32)	2208(60)	1166(56)
12	56.46°–57.76°	521(30)	1848(74)	976(55)
13	57.83°–59.05°	450(22)	1542(44)	814(39)
14	60.67°–64.32°	648(30)	2030(102)	1072(68)
15	64.53°–68.62°	1151(50)	3818(191)	2016(128)
16	68.70°–73.48°	1358(65)	4703(235)	2483(158)

was identified at the expected energy. Due to its proximity to the strong 803 keV transition from the Coulomb excitation of the target nucleus, which dominates the spectrum Doppler-corrected for ^{130}Sn , a reliable extraction of the much weaker 775 keV transition is hindered. To suppress this background and investigate the possible intensity of the $4_1^+ \rightarrow 2_1^+$ decay, events with for recoil kinematic Doppler-corrected energies within ± 10 keV of 803 keV were excluded from the analysis.

The resulting spectrum is shown in Fig. 6.27. The inset highlights the energy region around the expected 775 keV transition from the 4_1^+ state. No significant peak could be observed. Since the two-step excitation pathway contributes to the sensitivity on the quadrupole moment of the 2_1^+ state, and to constrain the transition strength $B(E2; 4_1^+ \rightarrow 2_1^+)$, an upper limit for this transition intensity was extracted using the method proposed by O. Helene [144]. Following this paper, the determination of statistically meaningful upper limits for signal intensities in the presence of background and limited counting statistics is achieved. It considers both the number of observed events in the signal region and the expected background to derive a confidence interval. The method involves iteratively evaluating the probability that a given signal strength is consistent with the number of observed counts in the signal region, assuming Poisson statistics. This is repeated for

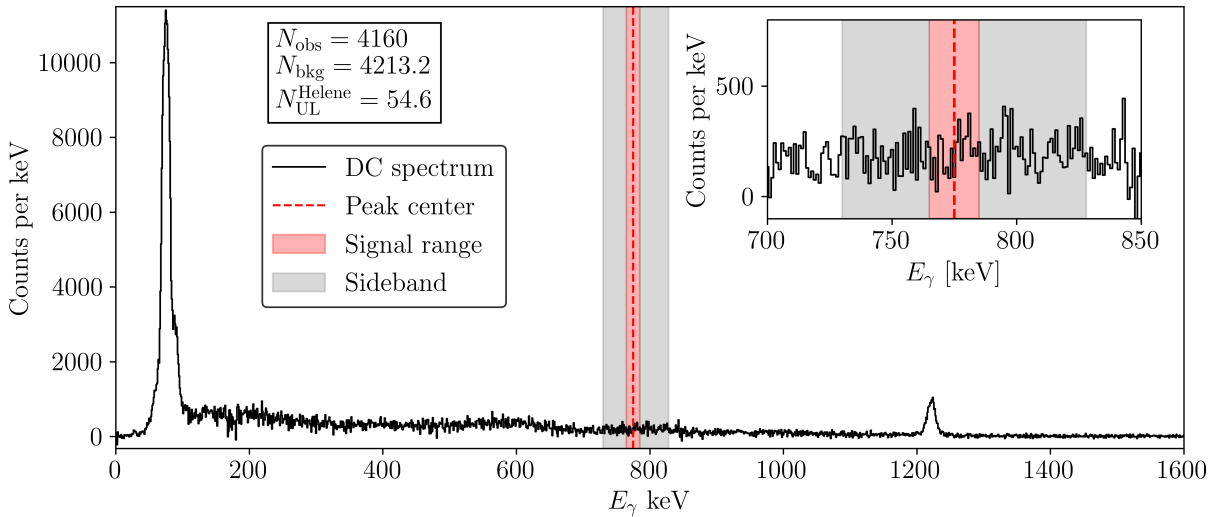


Figure 6.27. Doppler-corrected γ -ray spectrum for ^{130}Sn , after excluding events within ± 10 keV of the 803 keV transition from the ^{206}Pb target. The inset highlights the energy region around the expected $4_1^+ \rightarrow 2_1^+$ transition at 1991 keV. No peak could be identified in this region. An upper limit for the transition intensity was determined using the method of Helene [144], based on the number of counts in the signal and sideband regions. The resulting upper limit corresponds to 55 counts at the 95 % confidence level.

increasing signal strengths until the desired confidence level is reached. In this analysis, the upper limit for the $4_1^+ \rightarrow 2_1^+$ transition intensity was determined to be 55 counts at the 95% confidence level.

6.10. Gosia2 Analysis of 2_1^+ State

With the experimentally extracted yields of the observed γ -ray transitions in ^{130}Sn and ^{206}Pb , the previously unknown $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ value in ^{130}Sn can be investigated. The value is determined by performing χ^2 minimization with the GOSIA2 code.

In GOSIA, the dataset is separated into “experiments” by grouping events according to different ranges of laboratory scattering angles. Since Coulomb-excitation cross sections depend on the scattering angle, and the influence of the reorientation effect (e.g. through $Q_s(2_1^+)$) becomes stronger towards larger laboratory angles (see Sec. 1.3.3), subdividing the data into multiple experiments improves the accessibility of the spectroscopic quadrupole moment in the GOSIA2 analysis.

GOSIA2 evaluates the measured γ -ray yields by calculating Coulomb-excitation yields for each experiment at the specified mean beam energy and mean scattering angle. The calculated yields are scaled with experiment-specific normalization constants and compared to the experimental values. The electromagnetic matrix elements are varied to obtain the best overall agreement, quantified by the minimum of a global χ^2 function constructed from the yield differences. In the present case, the normalization is constrained by the known reduced $E2$ matrix element in ^{206}Pb , $\langle 0_{\text{g.s.}}^+ || E2 || 2_1^+ \rangle = 0.3178(47) \text{ eb}$ [24, 68].

Since the experimental yields are not obtained at a single, fixed beam energy and scattering angle but rather over finite kinematic ranges — due to energy losses in the target and the angular coverage of the CD — the corresponding calculated yields have to be integrated over the relevant intervals in beam energy and scattering angle. These beam-energy and angular ranges are therefore required inputs to represent each experiment consistently at its mean kinematic conditions in the GOSIA2 analysis.

The energy loss of a 572 MeV ^{130}Sn beam traversing the full target thickness was calculated with LISE++ to be 111 MeV. In the particle gates, parts of the measured ejectile spectrum are cut off at low energies due to the overlap region of recoil and target-like peaks, as discussed in Sec. 6.3. Events corresponding to particles scattered from the upstream front side of the target experience the largest energy loss and are therefore suppressed by the

applied lower energy gate; consequently, the maximum beam energy used in the GOSIA integration had to be adjusted. Back-calculating the lower energy threshold of the scattered ejectiles yields an upper beam energy of 559 MeV. Accordingly, only events with ^{130}Sn beam energies between 461 MeV and 559 MeV were selected in the analysis and this interval was adopted for the GOSIA analysis.

The target thickness was determined with an uncertainty of 0.2 mg/cm^2 , corresponding to an uncertainty of approximately 5 MeV in the reaction energy at the target exit. To account for uncertainties from both the target thickness and the energy resolution of the silicon detectors, the upper energy limit was varied by the same amount.

The scattering angles and their associated uncertainties were determined in Sec. 6.3. Internal conversion coefficients were calculated using the BRICC code [134]. The stopping powers required for the integration over the full target thickness and the angular acceptance were taken from SRIM [135]. For each ring of the CD detector and for each detected particle type, an individual experiment was defined, resulting in a total of 32 datasets. Each

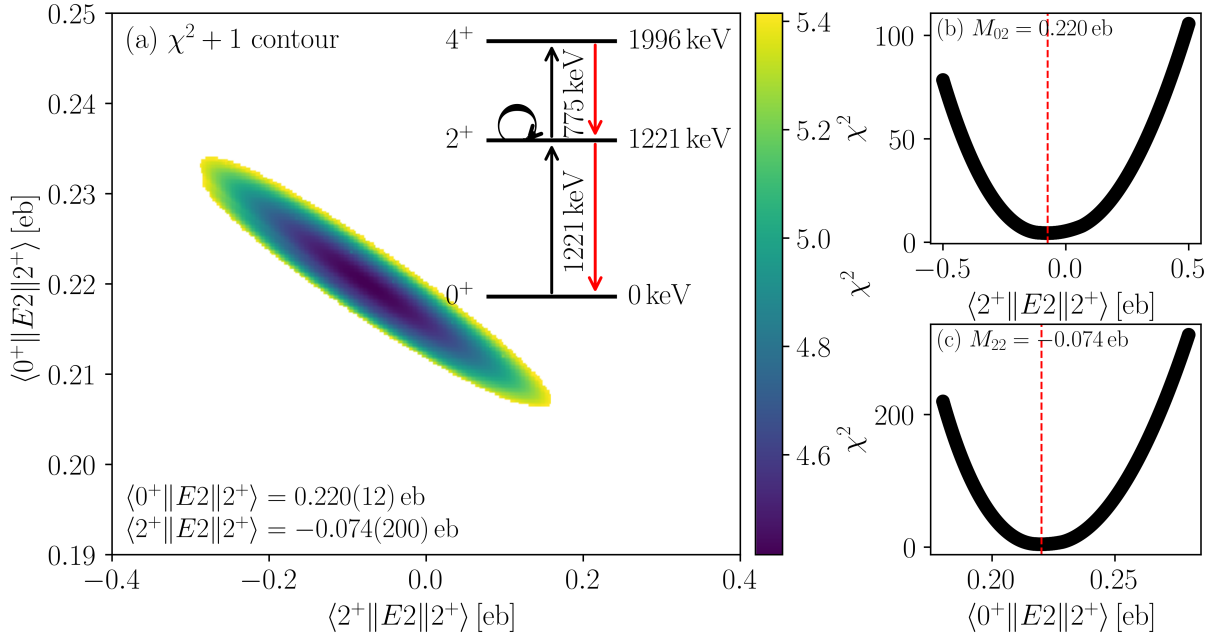


Figure 6.28. Left: $\chi^2 + 1$ surface plot based on fixed values of the diagonal and transitional $E2$ matrix elements. For each pair of values, a GOSIA2 calculation was performed, and the resulting χ^2 value is plotted. The shown uncertainty region corresponds to the area where $\chi^2 \leq \chi^2_{\min} + 1$. The level scheme used in the analysis is given. Results are presented as well. Right: The projection onto the resulting transitional (b) and diagonal (c) matrix elements and the χ^2 distribution around the minima, which is given as a red, vertical line.

experiment was normalized independently. For the final analysis, 16 of these experiments were selected, as discussed in the previous section.

The error estimation method implemented in GOSIA via the `OP,ERRO` option is incompatible with GOSIA2. Therefore, the uncertainties of the transitional and diagonal $E2$ matrix elements were determined from two-dimensional χ^2 -surface scans. For this purpose, both matrix elements were fixed on a two-dimensional grid, and a full GOSIA2 minimization was performed at each grid point while allowing all remaining fit parameters, including additional matrix elements and normalization constants, to vary freely. The resulting χ^2 surface was used to determine the region satisfying $\chi^2 < \chi_{\min}^2 + 1$. The reported error intervals were then obtained from the projection of this joint allowed region onto the respective matrix-element axes.

The resulting χ^2 surface is shown in Fig. 6.28(a), where the uncertainty region corresponds to the area enclosed by the condition $\chi^2 \leq \chi_{\min}^2 + 1$. The level scheme used in the GOSIA2 analysis is displayed as an inset on the right side of the figure. Black arrows indicate the excitation paths, including the possible two-step excitation into the 4^+ state, which was not seen in the spectra, but serves as a buffer state for possible unobserved feeding. Red arrows correspond to the de-excitation paths that were assumed for the decay scheme and included in the analysis.

The projections onto the transitional $\langle 0_{\text{g.s.}}^+ \| E2 \| 2_1^+ \rangle$ (b) and diagonal $\langle 2_1^+ \| E2 \| 2_1^+ \rangle$ (c) matrix elements are also shown, confirming the unambiguous location of the minimum. The uncertainty of the transitional matrix element is largely driven by the inclusion of the reorientation effect. If the quadrupole moment is neglected, i.e. for a vanishing diagonal matrix element, a projection of the surface onto the $\langle 0_{\text{g.s.}}^+ \| E2 \| 2_1^+ \rangle$ axis yields $\langle 0_{\text{g.s.}}^+ \| E2 \| 2_1^+ \rangle = 0.216(4) \text{ eb}$. Including the reorientation effect allows to determine the $Q_s(2_1^+)$ values results in a slightly increased transitional matrix element, while adding an additional uncertainty to the transitions strength.

To account for systematic effects associated with the uncertainties of the scattering angles and the beam-energy range, the $\chi_{\min}^2 + 1$ analysis was repeated for varied values of these inputs. The results, including statistical and systematic contributions, are

$$\langle 2_1^+ \| E2 \| 0_{\text{g.s.}}^+ \rangle = 0.220(13) \text{ eb}, \quad Q_s(2_1^+) = -0.06(15) \text{ eb}.$$

Here, the spectroscopic quadrupole moment was obtained from the fitted diagonal matrix

element using the standard relation $Q_s(2_1^+) = \sqrt{\frac{16\pi}{5}} \frac{1}{\sqrt{2J+1}} \langle 2_1^+ || E2 || 2_1^+ \rangle$ with $J = 2$. Using the fitted transition matrix element, the corresponding reduced transition probability is obtained as

$$B(E2; 0_{g.s.}^+ \rightarrow 2_1^+) = 0.0486(57) e^2 b^2.$$

For comparison, setting $Q_s(2_1^+) = 0$ (i.e. neglecting the diagonal matrix element and assuming a spherical 2_1^+ state) results in $B(E2; 0_{g.s.}^+ \rightarrow 2_1^+) = 0.0466(17) e^2 b^2$.

6.11. $7_1^- \rightarrow 5_1^-$ Excitation in ^{130}Sn

During the analysis of the excitation of the 2_1^+ state, it was observed that a fraction of the emitted $2_1^+ \rightarrow 0_{g.s.}^+$ γ rays does not originate from in-flight decay (and is therefore not Doppler shifted), but instead appears unshifted, i.e. emitted after the recoil has been stopped, even when applying a particle gate on the ejectiles.

Figure 6.29 (left panel) shows the measured γ -ray energy as a function of $\cos(\theta_{\text{particle}-\gamma})$, where $\theta_{\text{particle}-\gamma}$ denotes the angle between the detected particle and the γ -ray emission direction. If all decays occurred in flight, a linear correlation is expected; this behavior indeed dominates the matrix. In addition, however, a pronounced vertical structure is

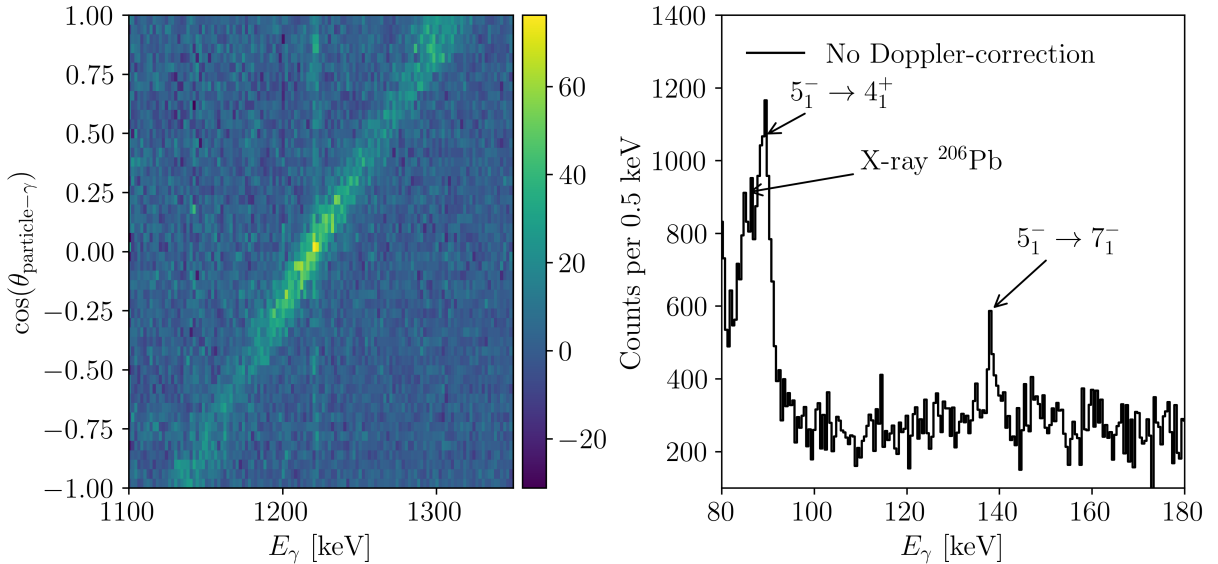


Figure 6.29. Left: Emission angle $\theta_{\text{particle}-\gamma}$ of the γ -ray transition shown against the measured, non Doppler-corrected γ -ray energy. The vertical structure indicates, that sum states decay, while the nuclei is at rest. A particle gate was placed on the ejectile kinematics. Right: The γ -ray spectra for energies below 200 keV without Doppler-correction. The decay of the isomeric 5_1^- state into the 4_1^+ and 7_1^- state are visible as well as X-ray lines from the ^{206}Pb target.

visible at the rest energy of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition at 1221 keV, indicating a component of γ -ray emission from stopped nuclei.

Since the spectrum is gated on the ejectile kinematics, this unshifted contribution cannot arise from events in which the ejectile is stopped in the target. Instead, it is attributed to feeding from the long-lived isomeric branch: the 7_1^- isomer is Coulomb excited to the 5_1^- state, which subsequently decays after the recoil has been stopped. The relevant level scheme is shown in Fig. 6.30.

The half-life of the 5_1^- state was first determined by Fogelberg *et al.* in 1981 [94] from γ -ray spectroscopy following the β decay of ^{130}I , yielding 52(3) ns. A later measurement within the RISING campaign at GSI reported a longer value of 62(5) ns [145], in disagreement with the earlier result. More recently, Piau *et al.* obtained 50(5) ns, again inconsistent with the RISING result. Despite these discrepancies, all measurements clearly indicate that the 5_1^- state does not decay in flight in the present experiment, since the time of flight between target and stopper is below 10 ns for all scattering angles. Consequently, the decay transitions associated with the 5_1^- state appear unshifted in particle-gated γ -ray spectra.

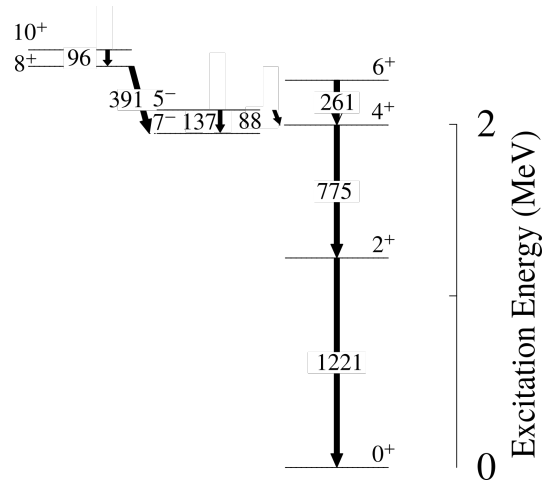


Figure 6.30. Levelscheme of ^{130}Sn . The isomeric 7_1^- state decays via β decay. Transition energies are given in keV.

This interpretation becomes more apparent in the right plot of Fig. 6.29. The uncorrected (i.e. not Doppler-corrected) γ -ray spectrum gated on the ejectiles is shown. Besides Pb X-ray lines, two distinct peaks are observed at 89 keV and 137 keV. These energies correspond to the decay of the 5_1^- state to the 4_1^+ and 7_1^- states, respectively. Lastly, Fig. 6.31 displays γ rays detected in coincidence with the 89 keV line. The subsequent cascade $4_1^+ \rightarrow 2_1^+ \rightarrow 0_{\text{g.s.}}^+$ is clearly visible and likewise appears unshifted, consistent with emission after stopping. This provides strong evidence that the isomeric 7^- component is indeed Coulomb excited in the present measurement, namely via excitation to the 5_1^- state. This constitutes the first observation within the Sn isotopic chain of Coulomb excitation involving an isomeric component in this way. Transition strengths associated with the $7^- \rightarrow 5_1^-$ excitation have so far been inferred from lifetime information. In the present work, this picture can be

complemented by direct Coulomb-excitation data.

For the following analysis, the long lifetime of the 5_1^- state has two important implications. First, the corresponding energy uncertainty of the lines is significantly smaller than Doppler-corrected peaks, since the majority of the decays of the 5_1^- state appear unshifted in the particle-gated spectra. Second, the GOSIA analysis faces efficiency-related complications. For the 2_1^+ state it is typically assumed that recoils and ejectiles decay effectively at the target position, which is valid for lifetimes on the order of picoseconds. In the case of the 5_1^- state, however, the decay occurs predominantly at the position of the CD detector, leading to a γ -ray efficiency that differs from that at the target. Since a calibration source cannot be distributed over the CD detector surface, GEANT4 simulations [146, 147] were performed to quantify this effect. These simulations are described

in the following subsection. Subsequently, the procedure for extracting γ -ray yields from the 5_1^- decay is outlined, followed by a lifetime determination from particle- γ time differences. Finally, the experimentally determined lifetime and the measured yields are used as input for a GOSIA analysis, which is presented in the last part of this chapter.

6.11.1. Miniball Efficiency at CD Position

The absolute photopeak efficiency of the Miniball spectrometer is routinely determined for γ rays emitted at the target position using radioactive sources with well-known γ -ray emission intensities. For the present work, a ^{152}Eu source was employed for this purpose. In this experiment, however, the decay of the long-lived 5_1^- state occurs essentially at rest at the position of the CD detector rather than at the target. Consequently, an absolute efficiency is required both at the target position and at the CD position.

As it is not feasible to distribute a calibration source uniformly over the CD detector surface, the efficiency at the CD position was determined from simulations. The target chamber was implemented in a dedicated GEANT4 framework using CAD drawings [148]. The Miniball spectrometer was added around the chamber. For simplicity, the segmentation of the setup

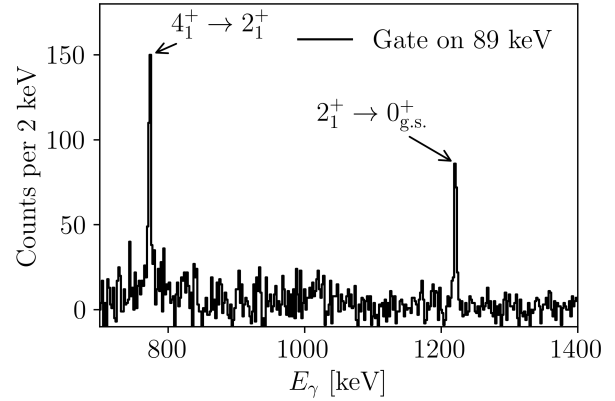


Figure 6.31. Non Doppler-corrected γ -ray spectrum gate on ejectiles and on 89 keV $5_1^+ \rightarrow 4_1^+$ transition. The cascade down to the ground-state is seen.

was neglected in the simulations; likewise, the electronics and dewar behind the cryostats were not included. A visualization of the simulated geometry is shown in Fig. 6.32.

To benchmark the simulation, a ^{133}Ba source was placed in the experimental setup at three positions inside of the plunger chamber (top, right, and bottom) and compared to spectra recorded at the target position. At $E_\gamma = 356\text{ keV}$, the measured relative efficiencies at these locations were found to be 89 % (right), 92 % (top), and 87 % (bottom). The observed differences arise because detector 6 was not used, rendering the effective geometry non-symmetric. The corresponding simulated relative efficiencies amount to 88 % (right), 93 % (top), and 89 % (bottom), in good agreement with the measurements. No explicit uncertainties are quoted here, since these values were obtained solely from count-rate ratios between spectra taken at the target position and at the respective chamber positions; the associated statistical uncertainties are negligible. The dominant systematic contribution stems from the limited accuracy with which the source position can be reproduced in the simulation. Nevertheless, the level of agreement indicates that the simulation provides a sufficient description of the spectrometer geometry and its attenuation properties with respect to γ -ray detection.

In the next step, a source was simulated at the position of the CD detector. Since the γ -ray emission is not expected to be uniformly distributed over the CD surface — owing to angular-dependent Coulomb-excitation cross sections — the emission points were sampled according to the angular distribution inferred from the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ intensities determined in the analysis of the 2_1^+ state. As the population of the 5_1^- state proceeds via the excitation

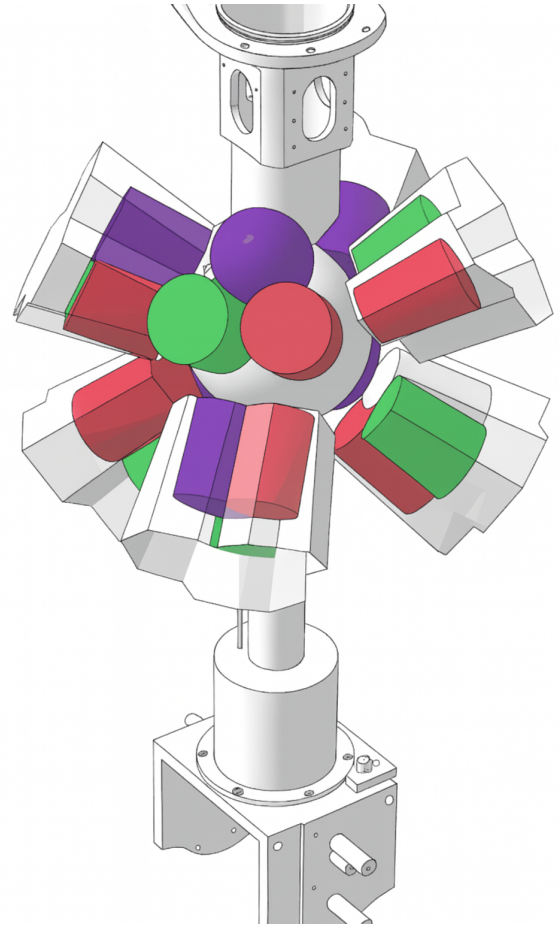


Figure 6.32. GEANT4 visualization of the Mini-ball spectrometer surrounding the plunger chamber. The latter is based on CAD drawings. The Miniball spectrometer is limited to the cryostats, while the Dewar and electronics on the back of the setup are neglected.

of the 7_1^- state followed by an $E2$ transition into the 5_1^- state, this procedure constitutes a reasonable approximation for the spatial distribution of the here relevant decays.

Figure 6.33 compares three efficiency curves: (i) the simulated absolute efficiency for a source at the target position, (ii) the simulated absolute efficiency for a source distributed over the CD according to the procedure described above, and (iii) the experimentally determined efficiency at the target position. For the simulated spectra, the absolute efficiency can be extracted directly from the known emitted intensities. For the measured ^{152}Eu spectrum, an absolute normalization was not available in the present configuration; therefore, the experimental curve was normalized to the simulated absolute efficiency at $E_\gamma = 1408 \text{ keV}$.

Forth-order Debertin fit functions (c.f. Ref. [149]) are added for all three data sets. The simulated and measured efficiencies at the target position show excellent agreement overall. A small discrepancy is observed for the 121 keV line, which is plausibly attributed to simplifications in the simulated plunger geometry (e.g. omitted cabling, the target wheel, and other small components). Due to the larger interaction probability at low energies, such details affect low-energy γ rays more strongly than higher-energy transitions. For this reason, the simulation is considered reliable above approximately $E_\gamma 120 \text{ keV}$. Deviations are expected to be significant below this energy value.

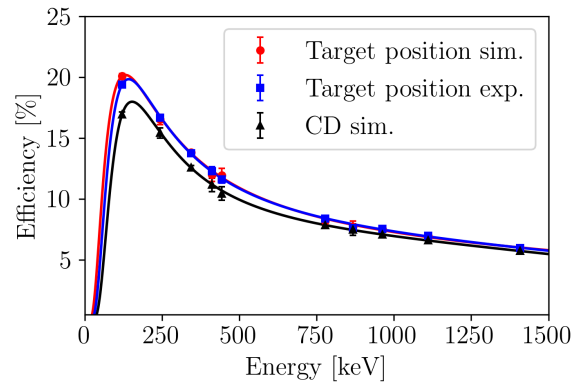


Figure 6.33. Comparison of experimental efficiency (blue) measured at target position and efficiencies determined from simulated γ -ray spectra. In the simulations the source was placed once at target position (shown in red) and once spread over the CD (black). Fits employing a fourth order Debertin function are added for all data points.

As expected, the efficiency curve obtained for the CD-distributed source is slightly lower than the corresponding efficiency at the target position, reflecting the different emission location and the additional material traversed by the γ rays before entering the detectors. At the energies most relevant for the 5_1^- -state analysis, i.e. for the 137 keV ($5_1^- \rightarrow 7_1^-$) and 775 keV ($4_1^+ \rightarrow 2_1^+$) transitions, the efficiency at the CD position amounts to 92 % and 94 %, respectively, relative to the efficiency at the target position. To account for residual systematic effects of the simplified geometry and the source-position modeling in the simulation, an additional relative systematic uncertainty

of 2% was assigned to the CD-to-target efficiency ratio. This value is guided by the study with the ^{133}Ba source discussed above, where the simulated relative efficiencies at different chamber positions reproduce the measured ratios at the level of a few percent. The quoted 2% uncertainty therefore represents a conservative estimate of the remaining modeling uncertainty relevant for the γ -ray efficiency difference between γ -ray emission from the target position and emission from the area of the CD Si detector.

6.11.2. Determination of γ -ray Yields

The level scheme relevant for the population and decay of the 5_1^- state is shown in Fig. 2.4. The 5_1^- state decays to the 7_1^- state via a 137 keV $E2$ transition and to the 4_1^+ state via an 89 keV $E1$ transition, with relative intensities of 50(4)% and 100(10)%, respectively [68]. The 4_1^+ state subsequently feeds the 2_1^+ state and finally the ground state through the 775 keV ($4_1^+ \rightarrow 2_1^+$) and 1221 keV ($2_1^+ \rightarrow 0_{\text{g.s.}}^+$) transitions.

For the GOSIA2 analysis, neither the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ at 1221 keV intensity nor the 88 keV

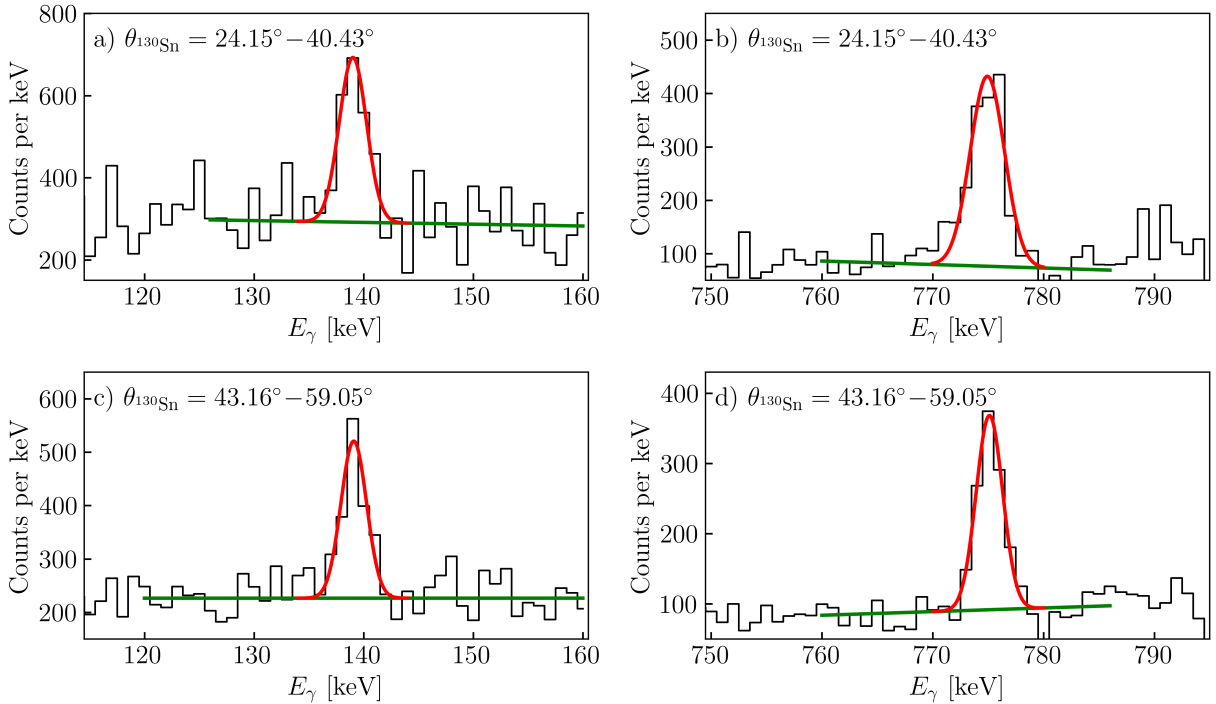


Figure 6.34. Non-Doppler-corrected γ -ray spectra gated on events registered in the CD detector, separated into two angular intervals: 24.15° – 43.04° (top) and 43.16° – 59.05° (bottom). The left panels show the $5_1^- \rightarrow 7_1^-$ transition at 137 keV, the right panels the 775 keV $4_1^+ \rightarrow 2_1^+$ transition. Gaussian peak-shapes (red) are fitted on top of a linear background (green).

$5_1^- \rightarrow 4_1^+$ transition can be used to quantify the population of the 5_1^- state. The $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ line contains a contribution from the direct Coulomb excitation of the 2_1^+ state by the ground-state component of the beam, which cannot be disentangled on an event-by-event basis. The 88 keV transition is likewise unsuitable: as discussed in the previous sections, its low-energy lies in a region where the CD-position efficiency cannot be determined with sufficient accuracy, and experimentally the peak is additionally obscured by the intense Pb $K\beta$ X-ray line at 84 keV originating from the target, preventing a reliable yield extraction.

Instead, the population of the 5_1^- state was constrained via the decay of the 4_1^+ state, using the 775 keV $4_1^+ \rightarrow 2_1^+$ transition. This approach is justified because the 4_1^+ state is not expected to be significantly populated by multi-step Coulomb excitation under the present conditions. Any potential contribution from two-step feeding can be assessed directly in the non-Doppler-corrected spectra: owing to the long lifetime of the 5_1^- state, the relevant decays occur at rest, and the corresponding lines appear narrow without Doppler broadening. A dedicated determination of the 5_1^- lifetime, supporting this assumption, is presented in the following subsection.

For the yield extraction, the CD detector was subdivided into two angular intervals, 24.15° – 40.43° and 43.16° – 59.05° . The seventh ring of the CD was not used, since one of the quadrants had a dead strip. The corresponding non-Doppler-corrected spectra are shown in Fig. 6.34. Photopeak areas were determined from least-squares fits using a Gaussian peak-shape superimposed on a linear background, as illustrated in the figure. The extracted peak yields for all relevant transitions in both angular intervals are summarized in Table 6.5. The quoted statistical uncertainties (given for the uncorrected yields) are dominated by residual background fluctuations after random subtraction.

Table 6.5. Uncorrected yields and efficiency-corrected intensities for the 137 keV ($5_1^- \rightarrow 7_1^-$) and 775 keV ($4_1^+ \rightarrow 2_1^+$) transitions in the two CD angular intervals. Uncertainties on the uncorrected yields are statistical (dominated by residual background fluctuations after random subtraction), while uncertainties on the corrected intensities include the quadratic sum of statistical and efficiency-related systematic contributions.

Energy [keV]	Transition	Angular coverage [°]	$I_{\text{Uncorrected}}$	$I_{\text{Corrected}}$
137	$5^- \rightarrow 7^-$	24.15–40.43	1093(160)	532(80)
		43.16–59.05	937(95)	456(49)
775	$4^+ \rightarrow 2^+$	24.15–40.43	1248(90)	1328(100)
		43.16–59.05	993(78)	1056(92)

The uncorrected yields were subsequently corrected for the relative photopeak efficiency. This correction includes (i) the energy-dependent efficiency relative to the 803 keV transition from recoil excitation and (ii) the geometric and attenuation differences between emission at the target position and at the CD position, as discussed in the previous section. The GEANT4 simulations were further adjusted to account for the different angular coverage of the two CD intervals. The corrected intensities are quoted with uncertainties obtained from the quadratic sum of statistical and systematic contributions. The systematic component, originating from the efficiency determination, is small compared to the statistical uncertainty and therefore has only a minor impact on the total uncertainty.

Only ejectile kinematics were used for the yield extraction. For recoil-gated data, the decay position of the long-lived 5_1^- state cannot be reconstructed with sufficient precision.

6.11.3. Lifetime Analysis

The lifetime of the long-lived 5_1^- state can be extracted from the time difference between a detected particle and a coincident γ ray, as the decay time is comparable to the timing resolution of the experimental setup. In principle, the most direct signature would be provided by the prompt transitions depopulating the 5_1^- state, in particular the 88 keV and 137 keV lines. In practice, however, the detection in the low- γ -ray energy regions are strongly affected by Compton-scattered γ rays and by intense β -delayed activity, which

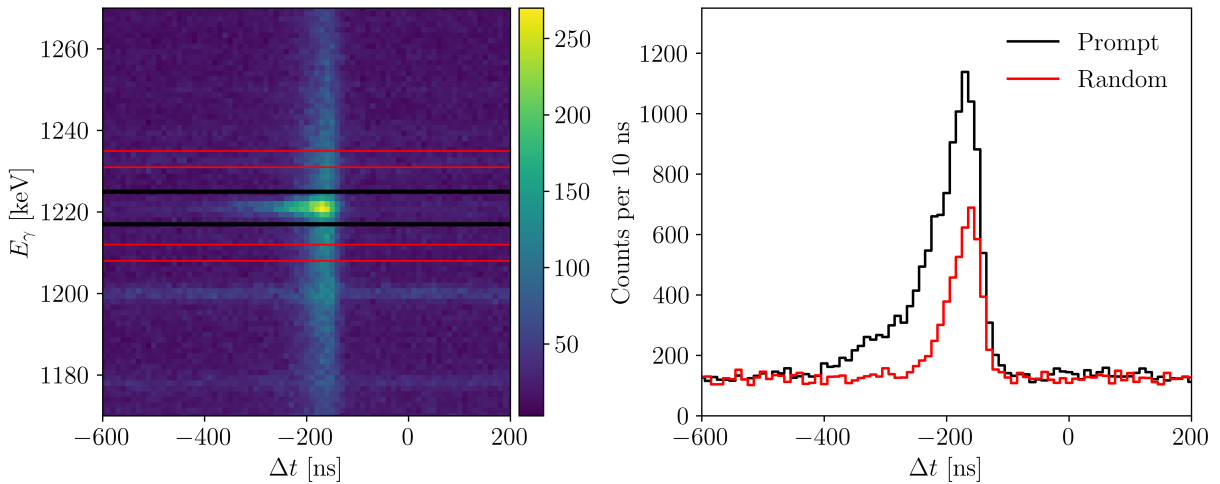


Figure 6.35. Left: Uncorrected γ -ray energy E_γ versus time difference Δt between the detected γ ray and the particle. The prompt gate around the stopped 1221 keV transition is indicated in black, while the random background regions are marked in red. Right: Corresponding Δt spectra for the prompt (black) and random (red) selections.

prevents a reliable timing analysis. Likewise, the 775 keV transition is not suitable due to contributions from the nearby 803 keV target-excitation line and from prominent β -delayed γ rays, such as the 780 keV line following the decay of $^{130}\text{Sn}_{\text{g.s.}}$.

For the lifetime determination, the 1221 keV $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition was therefore employed. Figure 6.35 shows on the left the measured γ -ray energy against the time difference of detected γ rays and particles. To isolate the component fed by the 5_1^- decay from the direct $0_{\text{g.s.}}^+ \rightarrow 2_1^+$ excitation, the Doppler-broadened contribution arising from the direct Coulomb excitation of the 2_1^+ state was subtracted together with the random-coincidence background from the prompt time region. The regions defined as prompt and random are highlighted in the plot in black and red, respectively. The resulting time-difference spectrum is shown in Fig. 6.36 and shows predominantly the delayed population of the 2_1^+ state originating from the decay sequence via the 5_1^- state.

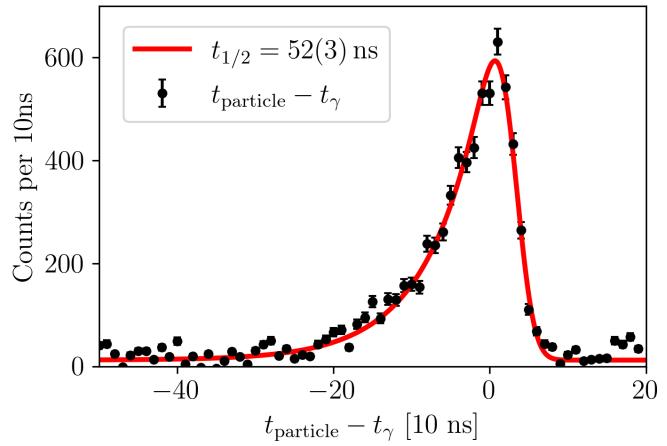


Figure 6.36. Time difference between the detected ejectile and a coincident γ ray gated on the 1221 keV $2^+ \rightarrow 0_{\text{g.s.}}^+$ transition. The solid line shows a fit obtained from the convolution of a Gaussian prompt response with an exponential decay component.

The spectrum was fitted using a response function consisting of a Gaussian prompt peak, representing the instrumental timing resolution, convolved with an exponential decay component accounting for the delayed emission. The fit result is included in Fig. 6.36.

Since the lifetimes of the intermediate positive-parity states in the cascade are expected to be very short compared to the extracted timescale, the obtained effective lifetime can be attributed to the 5_1^- state. The resulting lifetime is found to be $\tau(5_1^-) = 52(3)$ ns.

6.11.4. Gosia2 Analysis of the 5_1^- State

Using the experimental γ -ray yields extracted in Sec. 6.11.2, the lifetime of the 5_1^- state determined in Sec. 6.11.3 and the established branching ratios from NNDC [68], a dedicated GOSIA2 analysis was performed. The aim of this analysis is threefold. First, the combined information from Coulomb-excitation yields and the independently determined lifetime provides improved constraints on the transition strengths associated with the isomeric decay.

In particular, literature values are presently limited by sizable uncertainties in the adopted branching ratios, resulting in transition-strength uncertainties that are large compared to the precision of the lifetime. Second, the simultaneous description of Coulomb-excitation observables and the decay time offers an important consistency check of the assumed beam composition, since the population of the isomeric sequence is sensitive to the relative contributions of the different beam components. Finally, the GOSIA2 framework allows the spectroscopic quadrupole moment of the 5_1^- state to be investigated through its impact on the Coulomb-excitation dynamics.

For the GOSIA2 calculation, a modified level scheme had to be employed, shown in Fig. 6.37. In GOSIA, the state from which the excitation is initiated must be the lowest-lying state included in the calculation. In the present case, however, the experimentally observed intensities used to constrain the 5_1^- population are associated with transitions within the positive-parity cascade ($4_1^+ \rightarrow 2_1^+$) that is fed by the decay of the isomeric sequence rather than by direct Coulomb excitation from the ground state. Such “downstream” levels cannot be implemented straightforwardly in an unmodified scheme without violating the ordering requirement. To retain sensitivity to the experimentally accessible decay yield while respecting the GOSIA2 formalism, the level scheme was adapted by introducing an auxiliary 3^+ state above the 7_1^- state. This fictitious level was chosen such that it does not couple to the 7_1^- state and therefore does not affect the excitation path.

Instead, it acts as a sink (buffer) for the intensity leaving the 4_1^+ state. Practically, the measured yield of the $4_1^+ \rightarrow 2_1^+$ transition was mapped onto an effective $M1$ transition $4_1^+ \rightarrow 3^+$, ensuring that the observed depopulation of the 4_1^+ state is accounted for in the intensity balance of the calculation.

In neighboring Sn isotopes, the 7_1^- state is located above the 4_1^+ state, and the two levels may be connected by a $E3$ component. While such a decay branch is not established

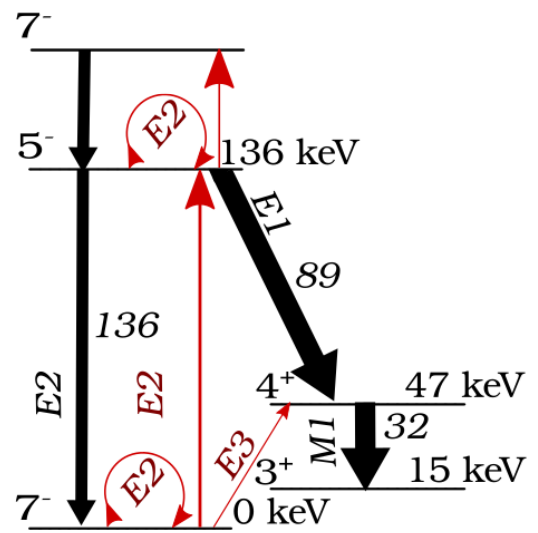


Figure 6.37. Level scheme adopted for the GOSIA2 analysis of the isomeric sequence. A fictitious 3^+ state was introduced as a buffer level (see text). Red arrows indicate matrix elements relevant for the excitation, black arrows those relevant for the de-excitation and intensity flow.

in ^{130}Sn , it cannot be excluded a priori. An additional $E3$ link between the 7_1^- and 4_1^+ states was therefore included in the GOSIA2 model to test its potential impact on the extracted matrix elements. Furthermore, an additional buffer level above the 5_1^- state was introduced and coupled to it via an $E2$ transition. These changes in the adopted level-scheme ensure numerical stability and provides a controlled intensity flow in the de-excitation part of the scheme without altering the excitation process. The spectroscopic quadrupole moment of the 7_1^- state was fixed to a literature-guided value of $Q_s(7_1^-) = 0.36(11) \text{ eb}$ [101]; although it has no significant influence on the excitation of the isomeric sequence under the present conditions, it enters the calculation formally and was therefore constrained to avoid unphysical solutions.

In addition to the experimental information discussed above, target normalization was employed for the isomeric component to provide an absolute scale for the calculated yields. For each CD angular interval, the total number of observed $^{206}\text{Pb } 2_1^+ \rightarrow 0_{\text{g.s.}}^+$ counts was scaled by the isomeric beam fraction of $^{130}\text{Sn}_{7^-}$, amounting to 34(2) %. This procedure yields an effective normalization condition for the analyzes and thereby increases the constraining power of the data set. All remaining input parameters, in particular stopping powers and internal conversion coefficients, were adopted analogous to the analysis of the ground-state excitation of the $^{130}\text{Sn}_{\text{g.s.}}$ beam.

The outcome of the GOSIA2 fit is summarized in Fig. 6.38, where projections of the $\chi^2 + 1$

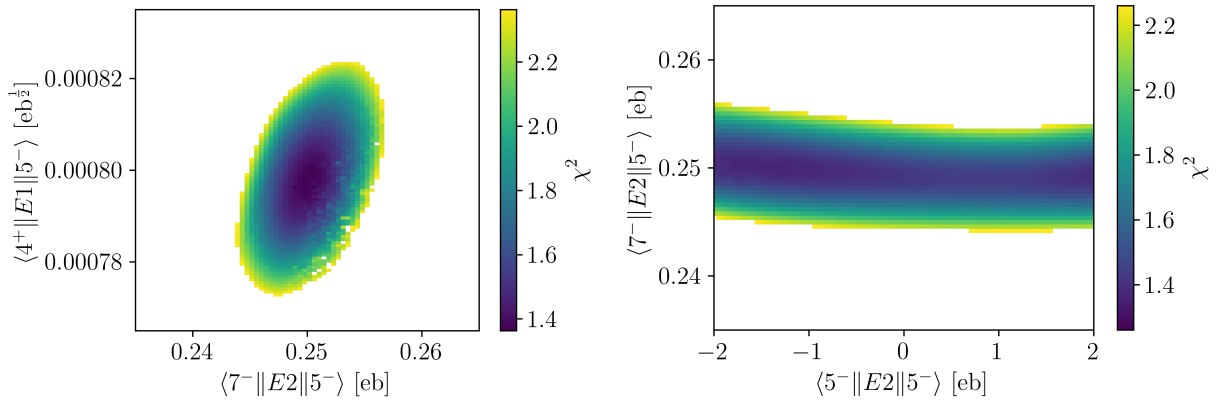


Figure 6.38. Two-dimensional projections of the $\chi^2 + 1$ region obtained from the GOSIA2 analysis of the 5_1^- isomeric sequence in ^{130}Sn . Left: correlation between the transitional matrix elements $\langle 4_1^+ \| E1 \| 5_1^- \rangle$ and $\langle 7_1^- \| E2 \| 5_1^- \rangle$, exhibiting a well-defined minimum. Right: dependence of $\langle 7_1^- \| E2 \| 5_1^- \rangle$ on the diagonal matrix element $\langle 5_1^- \| E2 \| 5_1^- \rangle$ (related to $Q_s(5_1^-)$), for which no clear minimum is observed.

region are shown. The left panel displays the correlation between the transitional matrix elements $\langle 4_1^+ \| E1 \| 5_1^- \rangle$ and $\langle 7_1^- \| E2 \| 5_1^- \rangle$. A well-defined minimum is obtained at

$$\langle 4_1^+ \| E1 \| 5_1^- \rangle = 7.96(25) \times 10^{-4} e b^{1/2}, \quad \langle 7_1^- \| E2 \| 5_1^- \rangle = 0.249(6) e b,$$

demonstrating that the combination of Coulomb-excitation yields and the independently determined lifetime provides strong constraints on the isomeric decay strengths. The right panel shows the transitional $\langle 7_1^- \| E2 \| 5_1^- \rangle$ matrix element as a function of the diagonal $\langle 5_1^- \| E2 \| 5_1^- \rangle$ matrix element, which is directly related to the spectroscopic quadrupole moment of the 5_1^- state. In contrast to the left panel, no pronounced minimum is observed within the explored parameter range, indicating that the present data set does not provide a meaningful constraint on the static quadrupole moment of the 5_1^- state. The extracted transitional matrix elements can be converted into transition strengths, yielding $B(E2; 7_1^- \rightarrow 5_1^-) = 0.0042(2) e^2 b^2$ and $B(E1; 4_1^+ \rightarrow 5_1^-) = 71(3) \times 10^{-6} e^2 b$. These values are in excellent agreement with those derived solely from the newly determined lifetime, namely $B(E2; 7_1^- \rightarrow 5_1^-) = 0.0040(11) e^2 b^2$ and $B(E1; 4_1^+ \rightarrow 5_1^-) = 72(10) \times 10^{-6} e^2 b$, respectively.

7. Data Analysis of 2025 Experiment

Following the successful experiment on ^{130}Sn , the subsequent study aimed at determining transition strengths in ^{134}Sn , located just beyond the doubly-magic nucleus ^{132}Sn . The experiment was originally scheduled to take place in July 2024 over the course of one week, as the first Miniball experiment at CERN in 2024. In preparation, extensive calibration measurements were performed, including source runs and experiments with stable beams to establish the energy and position calibrations of the individual detectors, similar to what was described for the 2023 data. These calibrations had already been evaluated in advance, ensuring that the ^{134}Sn campaign was well-prepared from the analysis perspective. The initial difficulties encountered during the ^{130}Sn analysis—such as problems with the data acquisition and the sorting code—had been resolved meanwhile, as demonstrated in the last chapter, so that the new experiment could proceed under the best possible conditions.

Unfortunately, the 2024 run could not start as intended. The heating element in the primary UC_x target had not been sufficiently tested and proved unstable under proton irradiation. As a result, molecular formation could not be achieved, and the experiment had to be terminated before it even began. However, the experiment was successfully rescheduled and repeated in October 2025. The following chapter describes the data preparation and analysis procedures of this experiment. Since most of the analysis steps follow those already outlined for the ^{130}Sn data, they are briefly summarized here, focussing on analysis steps that differ from those used in the 2023 dataset. Nevertheless, all necessary analysis methods and calibrations are explained.

7.1. Miniball Calibration

Similarly to the 2023 experiment, the energy calibration of all Miniball detectors was performed using γ -ray lines from a ^{152}Eu source measured prior to the experiment. The well-known transition energies of the ^{152}Eu source, covering a wide range up to about 1.4 MeV, allowed for a precise determination of both the linear and non-linear components of the energy response for each crystal. The energy calibration functions and the used

E_γ -Channel pairs are plotted in Fig. 7.1. The stability of the calibration was again observed by looking at γ rays after β decay for each segment and for the total setup over time.

The absolute full-energy peak efficiency of the Miniball array was determined using a set of standard calibration sources (^{133}Ba , ^{152}Eu , and ^{226}Ra), which were measured under identical

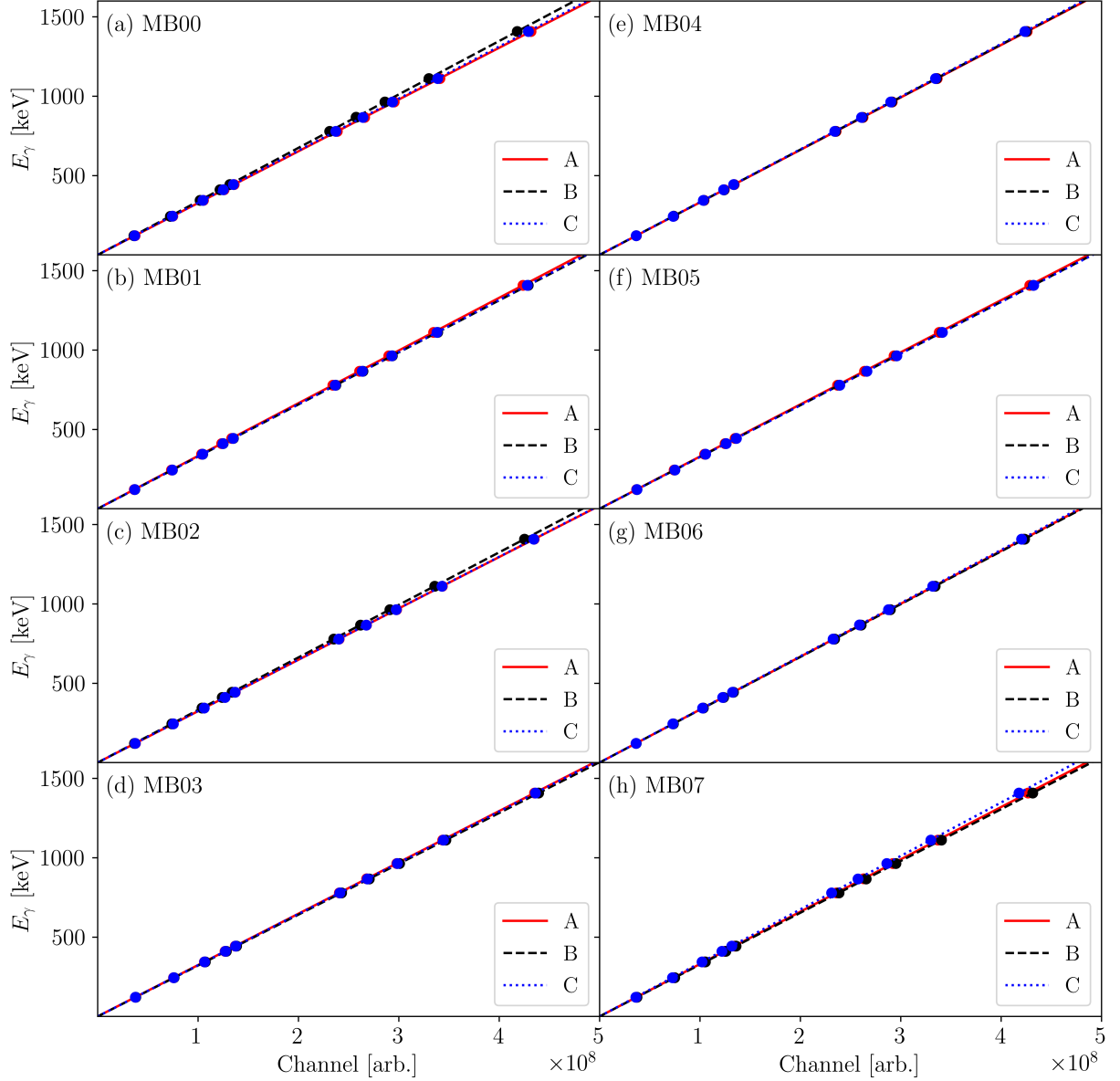


Figure 7.1. Energy calibration of the Miniball detectors using γ rays from a ^{152}Eu source measured prior to the experiment. For all eight Miniball triple clusters, Segment 1 of each cluster is shown. A parabolic calibration function is fitted to the data and accurately reproduces the known transition energies over the full energy range.

geometrical conditions as during the experiment. These measurements were carried out in the week following the experiment, with acquisition times ranging from about 5 h for ^{133}Ba to nearly one day for ^{152}Eu , ensuring sufficient statistics for all relevant γ -ray transitions. The relative γ -ray intensities of the calibration lines were adopted from Refs. [150, 151]. The resulting efficiency data points were then fitted using the empirical function employed by Debertin and Helmer [149], which parameterizes the logarithm of the photopeak efficiency $\varepsilon(E_\gamma)$ as a fourth-order polynomial in the logarithm of the γ -ray energy E_γ :

$$\ln \varepsilon(E_\gamma) = a_0 + a_1 \ln(E_\gamma) + a_2 [\ln(E_\gamma)]^2 + a_3 [\ln(E_\gamma)]^3 + a_4 [\ln(E_\gamma)]^4. \quad (7.1)$$

The coefficients a_i were obtained from a least-squares fit to the measured data points. For the determination of the absolute normalization of the efficiency curve, a low-activity ^{60}Co source was used, providing two well-known high-energy γ lines at 1173 keV and 1332 keV. The absolute efficiency at 1173 keV was determined to be $(6.62 \pm 0.15)\%$, which served as the reference point for scaling the entire efficiency curve. The resulting parametrization accurately describes the detector response over the full energy range relevant for the experiment and was subsequently applied in the analysis of the in-beam γ -ray spectra.

7.2. Position of the CD Detector

The distance of the CD to the target provides the information about the scattering angles θ that are covered by the individual rings of the CD. To obtain a precise value, an α -source measurement was carried out. A triple α source was placed at the target position and measured for 1.5 h. The number of particles detected in each strip is proportional to the solid angle covered by that strip. This solid angle depends on the distance between the

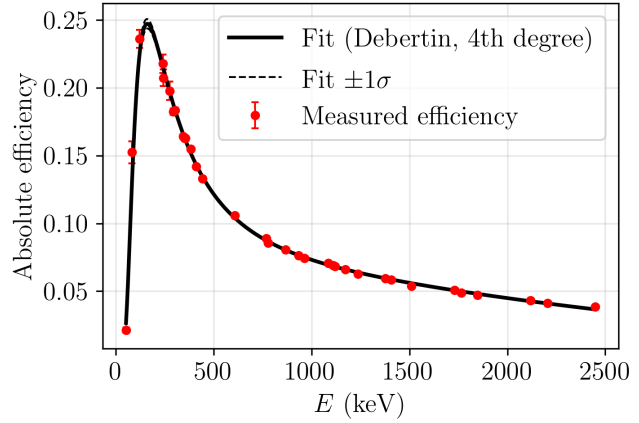


Figure 7.2. Absolute full-energy peak efficiency of the Miniball detector array. The efficiency curve was obtained using γ rays from ^{133}Ba , ^{152}Eu , and ^{226}Ra calibration sources. Relative intensities were taken from Ref. [150] for ^{226}Ra and from Ref. [151] for ^{133}Ba and ^{152}Eu . The absolute normalization was determined using a low-activity ^{60}Co source.

detector and the target. By studying how the number of detected decays varies with the strip index, the distance can therefore be determined. The procedure and the results are illustrated in Fig. 7.3.

The fit model applied is the same as the one implemented in the analysis code. For each quadrant, the measured strip intensities $I(s)$ as a function of strip index s are described by

$$I(s) = A \Delta\Omega(s; d), \quad (7.2)$$

where A is a free normalization constant and d is the distance to be extracted. The strip solid angle $\Delta\Omega(s; d)$ is determined from the CD geometry using the inner and outer strip radii,

$$R_b(s) = 9.05 + (15 - s) \cdot 2.0 \text{ mm}, \quad R_a(s) = R_b(s) + 1.9 \text{ mm},$$

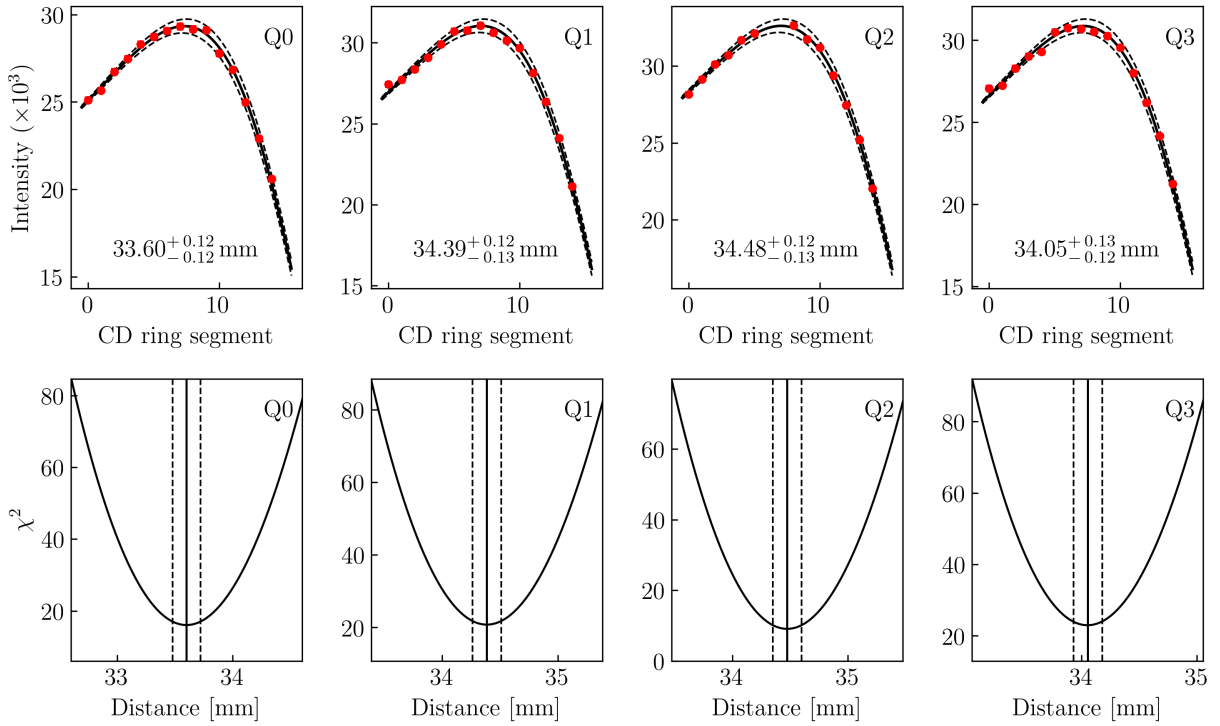


Figure 7.3. Top: Strip-wise intensities for the four CD quadrants (red points) together with the best-fit solid-angle model (black) and the 1σ envelopes (black dashed). The extracted distance d with asymmetric 1σ uncertainties is annotated in each panel. Bottom: χ^2 profiles as a function of d for the same quadrants. The best-fit distance is marked by a solid line, and the 1σ crossings by dashed lines.

leading to

$$\Delta\Omega(s; d) = 2\pi \left[\frac{R_a^2 + d^2 - d\sqrt{R_a^2 + d^2}}{R_a^2 + d^2} - \frac{R_b^2 + d^2 - d\sqrt{R_b^2 + d^2}}{R_b^2 + d^2} \right] \phi(s). \quad (7.3)$$

An empirical azimuthal coverage factor $\phi(s)$ is included, parametrized as a cubic polynomial [141],

$$\phi(s) = \frac{-0.00625 s^3 + 0.07056 s^2 - 0.54542 s + 77.66278}{360}. \quad (7.4)$$

The uncertainty on d is obtained from a χ^2 profile. For a dense grid of fixed d values, the normalization A is refitted, which yields $\chi^2(d)$. The 1σ interval corresponds to the distances where χ^2 increases by one unit with respect to the minimum.

Figure 7.3 shows for each quadrant in the top row the strip intensities together with the best-fit model in black. Dashed black lines indicate the 1σ envelopes. The extracted distance with its asymmetric 1σ uncertainty is given inside each panel. The bottom row

Table 7.1. Angular coverage of the 16 rings for a detector distance of $d = 33.8(5)$ mm. Uncertainties stem solely from $\sigma_d = 0.5$ mm.

Ring	$\theta_{\min}$ [deg]	θ_{mean} [deg]	$\theta_{\max}$ [deg]
1	14.91(22)	16.40(23)	17.87(25)
2	18.03(25)	19.47(27)	20.89(28)
3	21.04(28)	22.43(30)	23.79(32)
4	23.93(32)	25.26(33)	26.57(33)
5	26.70(33)	27.97(35)	29.21(37)
6	29.34(37)	30.55(37)	31.73(38)
7	31.85(38)	33.00(38)	34.12(40)
8	34.23(40)	35.32(40)	36.38(40)
9	36.49(40)	37.52(42)	38.51(42)
10	38.62(42)	39.59(42)	40.53(42)
11	40.63(42)	41.54(42)	42.43(42)
12	42.53(42)	43.39(42)	44.23(42)
13	44.31(42)	45.13(42)	45.92(42)
14	46.00(42)	46.77(42)	47.51(42)
15	47.59(42)	48.31(42)	49.01(42)
16	49.09(42)	49.77(42)	50.43(42)

displays the corresponding χ^2 profiles as a function of distance. The minimum is marked with a solid vertical line and the 1σ crossings with dashed lines.

From the four quadrants, an average CD distance of 34.02(6) mm was obtained. Using measurements performed with a laser a distance of 33.8 mm was found. An analysis similar to the one presented in section 6.3.2 was performed as well resulting in a distance of 33.5(5) mm, agreeing with the values shown here. However, the latter is not sensitive to different distance of the individual quadrants. Since the value determined with the laser does not depend on the thickness of the α -source, this value was adopted in the present work. An uncertainty of 0.5 mm was applied. The angular coverage of the individual segments determined from these distances is given in Table 7.1.

7.3. Energy Calibration of the CD Detector

The energy calibration of the CD detector was performed similarly as in the 2023 experiment. Initially, one inner p -side strip was calibrated absolutely by comparison with kinematics calculations using the KINSIM code [139]. Subsequently, all n -side strips were calibrated *relative* to this reference p -strip using only events with a single hit on both detector sides (later called multiplicity 1v1). In a second step, one of the calibrated n -side strips is then used as reference to calibrate all p -side strips across the detector.

This procedure is illustrated in Fig. 7.4, which shows four example rings of quadrant 1 and n -strip 15. For each ring, a linear function was fitted to the raw p -side energy spectrum to match the corresponding n -side energy scale. The same approach was applied to all rings of all quadrants to ensure a consistent energy calibration across the entire CD detector.

After the calibration procedure, the multiplicity distributions were examined to verify the detector performance. It was observed that events with one p -strip and multiple n -side hits were enhanced compared to the 2023 data, even after increasing the detector thresholds. This increase can be explained by the detector readout configuration: in the 2023 experiment, two n -side strips were electronically combined into one channel, whereas in the present setup, all 24 n -side strips of each quadrant were read out individually. Consequently, the total number of possible coincidences on the n -side was higher, leading to a larger fraction of events with multiplicities greater than one.

7.3.1. Correction for Events with two n -side Hits

During the analysis, it was observed that events with a multiplicity of two on the n -side (n mult. = 2) not only contributed to the total n -side energy, but also affected the corresponding energy measured on the p -side. In the data-sorting procedure, energies detected simultaneously in two neighboring n -side strips are automatically summed to form a single event. However, when comparing these $2n$ events to $1v1$ events, a systematic deviation in the p -side energy became apparent.

This is illustrated in the left diagram of Fig. 7.5. The expected kinematic line for $1v1$ events is indicated in black. If the $2n$ events did not influence the p -side energy, the

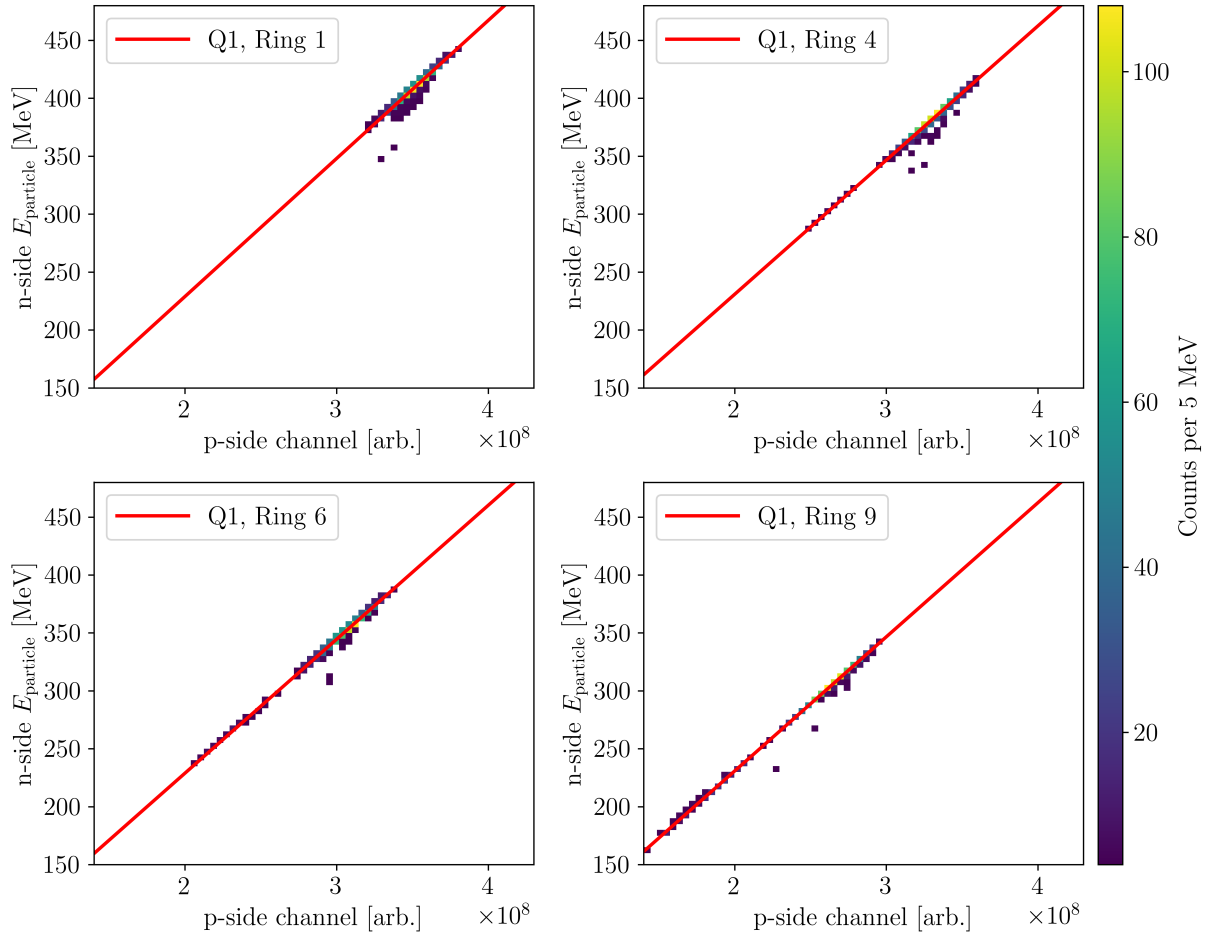


Figure 7.4. Linear energy calibration of the CD detector. Shown are four representative rings of quadrant 1. The calibrated n -side energies were used to determine the calibration of the corresponding p -side strips. A multiplicity of one on both the n - and p -side of the DSSSD was required.

corresponding data points would align along this linear correlation. Instead, a clear parabolic-like deviation is observed, indicating that the summed n -side signal modifies the baseline or charge collection on the p -side. Since the total energy used in the subsequent analysis corresponds to the projection onto the p -side energy axis, this distortion produces a double-peak structure, as shown in the right panel of Fig. 7.5 (black curve) for high energies.

To correct for this effect, a quadratic function was fitted to the data points in the left panel, describing the deviation as a function of the relative energy sharing between the two coincident n -side hits. The derived correction function has the form

$$\text{cor}\left(\frac{E_n}{E_{n,\text{sum}}}\right) = -1.22\left(\frac{E_n}{E_{n,\text{sum}}}\right)^2 + 1.22\frac{E_n}{E_{n,\text{sum}}} + 1, \quad (7.5)$$

where E_n is the energy deposited in one n -side strip and $E_{n,\text{sum}}$ is the summed energy of both coincident hits. This correction factor was determined individually for each quadrant

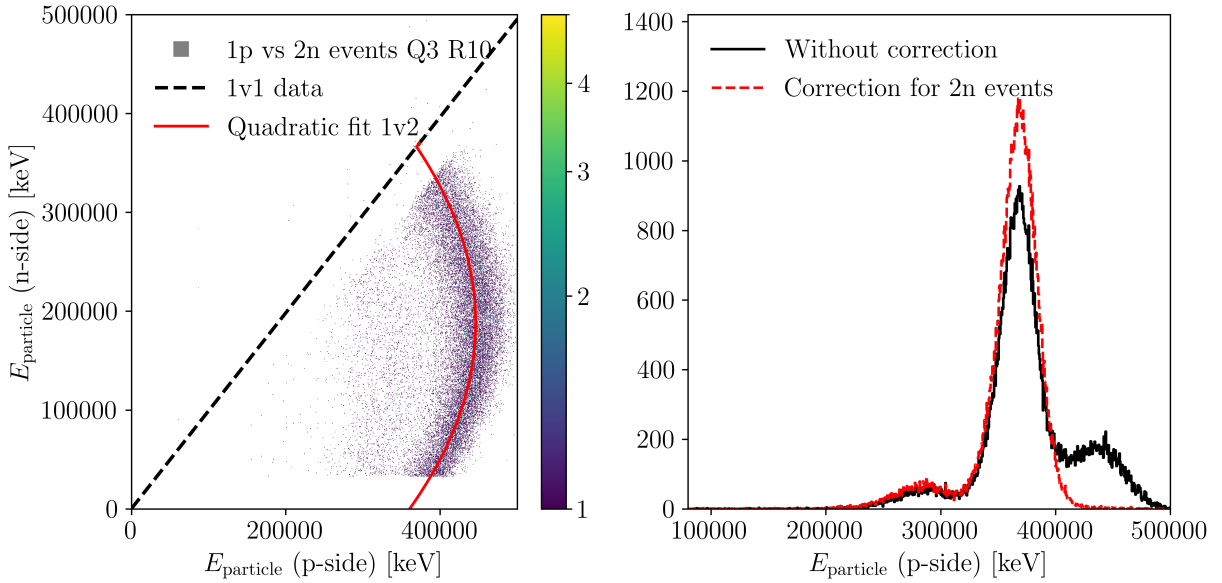


Figure 7.5. Correction for events with multiplicity 2 on the n -side. The left panel shows the correlation between p - and n -side energies for ring 10 of quadrant 3. The black dotted line indicates the linear calibration function derived from multiplicity 1 events, corresponding to equal energies on both sides. The red line represents a quadratic fit applied to the multiplicity 2 events on the n -side, which is used to correct their energies. The right panel shows the projected p -side energy spectrum before (black) and after (red) the correction. At 280 MeV the recoiling ^{194}Pt nuclei are visible, while the peak at 360 MeV shows the scattered ejectiles. For further details, see text.

and ring of the CD detector. The effect of the correction is illustrated in Fig. 7.5 (right), where the red spectrum shows the result after applying the correction. The single-particle peak is clearly recovered with improved statistics, while the overall energy resolution of the CD detector remains unaffected. The physical origin of this effect cannot be fully explained here. A possible contribution from a baseline shift induced by the coincident $2n$ events has been discussed in Ref. [24].

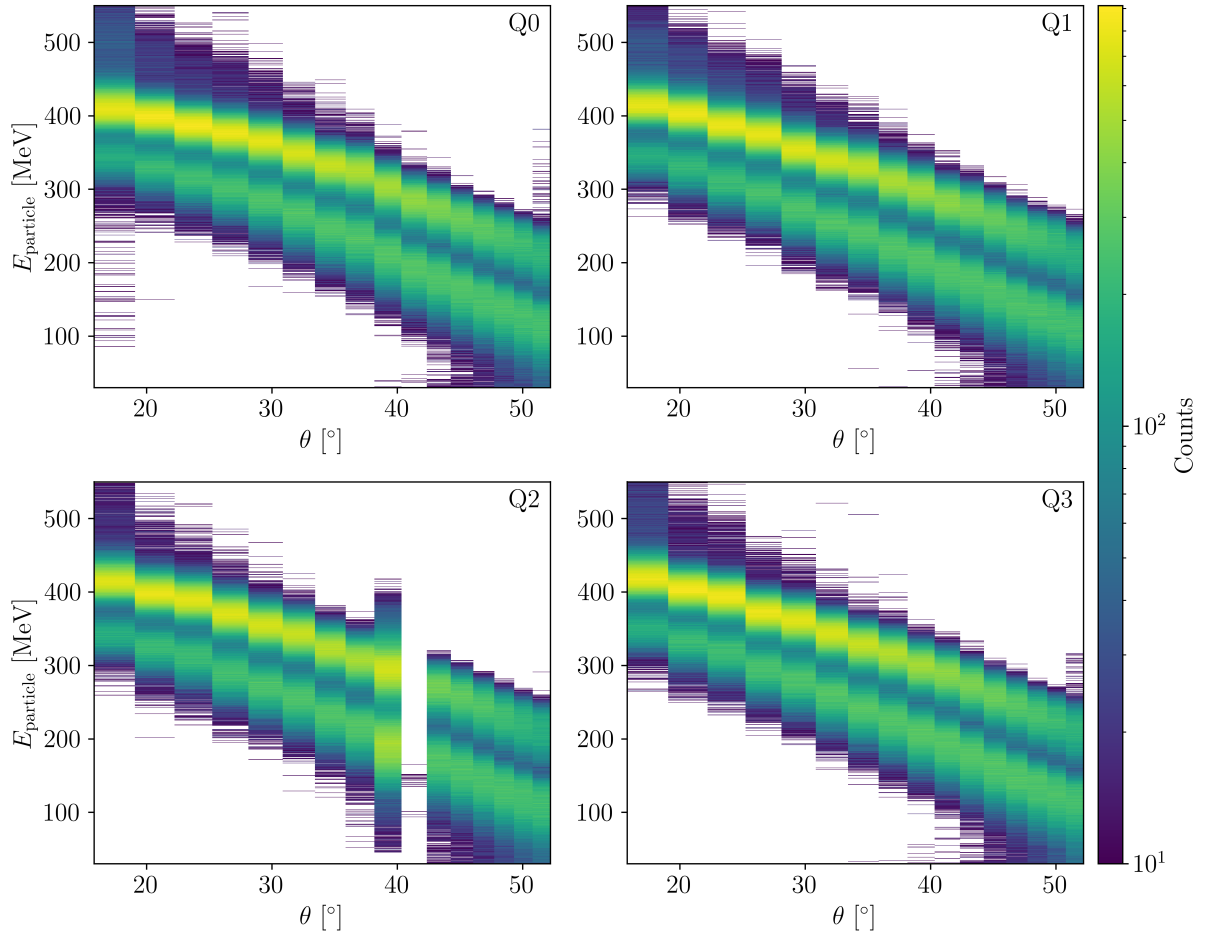


Figure 7.6. Energy-calibrated kinematic spectra for all four quadrants of the CD detector are shown on a logarithmic scale. Only particles detected in coincidence with a γ ray are included. All four quadrants show good agreement, and a clear separation between recoil- (bottom) and ejectile-like (top kinematic line) nuclei is visible. The double-peak structure observed at higher energies was corrected by applying the multiplicity 2 n -side event correction. Rings 9 and 10 of quadrant 2 were electronically shortened and are therefore combined into a single ring.

7.4. Kinematic Gates

The final energy-calibrated spectra, corrected for $2n$ events, are shown in Fig. 7.6 for all four quadrants of the CD detector. Only particles detected in coincidence with a γ ray are included. The energy response of the individual quadrants is in excellent agreement, demonstrating the uniformity of the calibration. In quadrant 2, rings 9 and 10 were electrically shorted, and therefore all events from this region were assigned to ring 9, leaving ring 10 empty. Despite the rather thick 3.0 mg/cm^2 ^{194}Pt target, a clear separation between ejectile and recoil bands is observed in all quadrants, although some overlap occurs at the smallest scattering angles.

To avoid applying gates on regions of overlap, Gaussian peak functions were fitted to both the ejectile and recoil lines in the kinematic spectra. Only regions that could be clearly attributed to either ejectile or recoil events were included in the gates. Since, in the subsequent γ -ray analysis, spectra from several θ -angles are summed, it is essential to remove a consistent fraction of ejectile and recoil events across all rings. To achieve this, the expected relative number of detected ejectiles was compared to the Rutherford differential cross section, calculated for each ring geometry. Only particle events were included without the condition of a coincident γ ray. The ratio between the measured and calculated ejectile yields should remain approximately constant for all scattering angles if the assumed beam intensity is correct. Figure 7.7 shows this comparison, where the calculated Rutherford yields are normalized to an assumed beam intensity of 6×10^4 particles per second.

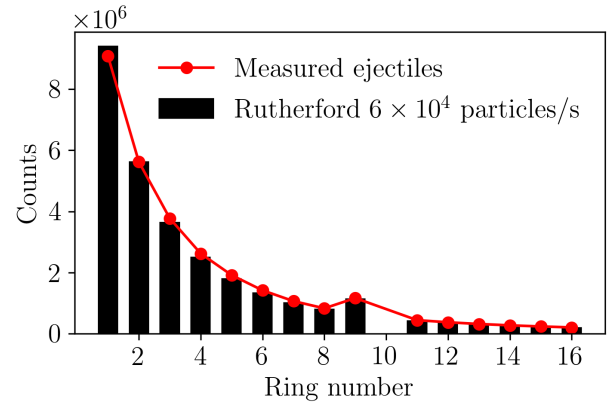


Figure 7.7. Measured and calculated ejectile yields for all CD rings. The black bars show the expected yields from the Rutherford cross section assuming a beam intensity of 6×10^4 particles per second, while the red points represent the measured ejectile counts. Rings 9 and 10 were combined due to the electrical short between them. The good agreement between data and calculation confirms the consistency of the assumed beam intensity.

The projectile isotope was taken to be ^{134}Te , consistent with the dominant composition of the beam, which also contained a significant fraction of ^{134}Ba , as discussed later. Rings 9

and 10 were combined for the reasons mentioned above. The measured yields follow the expected Rutherford trend very well, and the ratio of measured to calculated ejectiles is close to unity across all angles, confirming the applied gates.

7.5. Particle- γ Coincidences

The γ -ray energies measured in coincidence with a detected particle as a function of the time difference Δt are shown in Fig. 7.8. The prompt coincidence region is indicated in black, while the random time region is shown in blue. The ratio of prompt to random events, amounting to 26.8, is significantly enhanced compared to the ^{130}Sn case. This improvement mainly arises from the strongly reduced contribution of β -delayed γ rays. Both ^{134}Sn and its daughter nucleus ^{134}Sb predominantly decay directly to the ground state of their respective daughter nuclei, resulting in a much lower number of γ rays from delayed decays. In addition, the beam current during the ^{134}Sn experiment was reduced by approximately a factor of ten compared to the ^{130}Sn run. As discussed in the following chapter, roughly half of the incoming beam consisted of stable ^{134}Ba , which does not contribute to random γ -ray background, but, as a very collective nucleus, very much contributes to the prompt events. These factors combined lead to a much cleaner particle- γ coincidence spectrum.

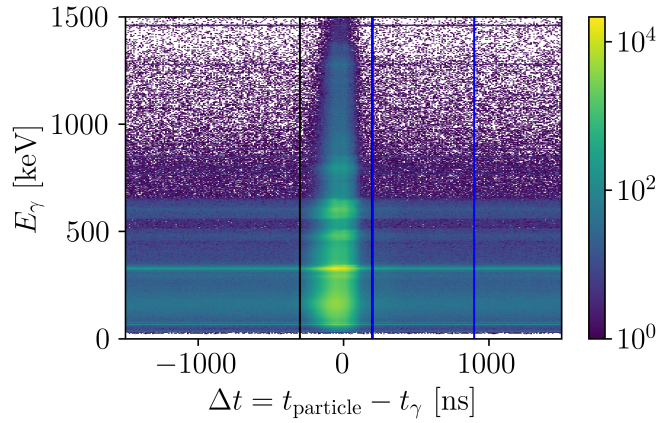


Figure 7.8. γ -ray energy as a function of the time difference between the detected particle and the γ ray. β -delayed γ rays populate the entire matrix, while prompt signals are concentrated around $\Delta t = -50$ ns.

7.6. Miniball Angle Determination

As already discussed for the previous experiment, the stable-beam reaction $^{22}\text{Ne}(d,n)^{23}\text{Na}$ does not provide sufficient information on the ϕ angles of the Miniball clusters. Since the beam intensity in the radioactive beam experiment was expected to be significantly lower than in the ^{130}Sn experiment, it was assumed, prior to the experiment, that the RIB data cannot be used to determine the ϕ angles in the same way as for the 2023 data.

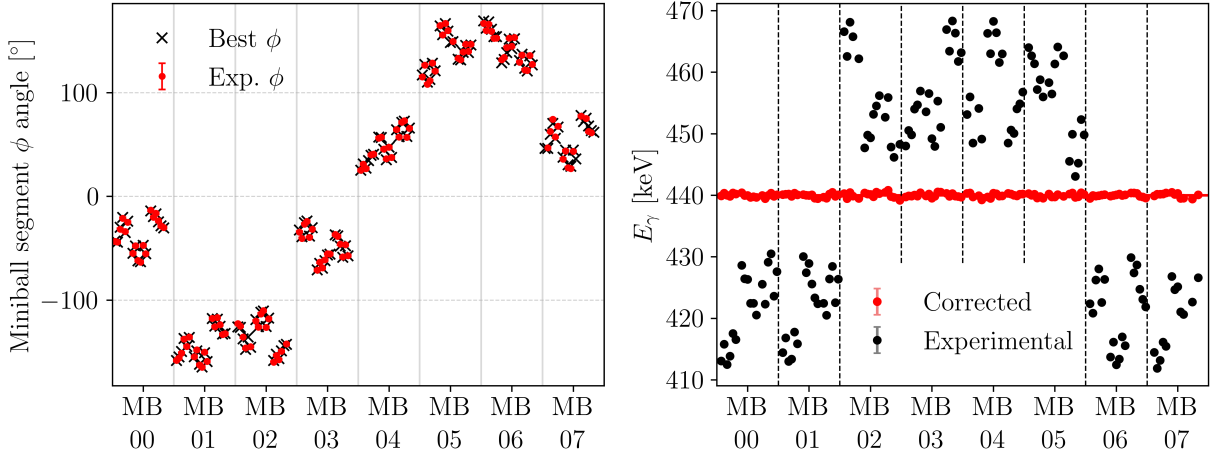


Figure 7.9. Left: Extracted ϕ angles of the Miniball segments with their statistical uncertainties (black), and values after minimization. Right: measured Doppler-shifted energy of the 440 keV transition in ^{23}Na (black) and corrected energies after minimization (red). In both cases, the minimization procedure yields good agreement with the expected values.

A Coulomb-excitation experiment with ^{22}Ne on a ^{181}Ta target was therefore carried out prior to the RIB experiment, but the resulting yields were not sufficient to reliably extract the ϕ angles of each segment. Therefore, data from an earlier RIB experiment investigating the Coulomb excitation of ^{162}Yb were analyzed. This nucleus exhibits strong collective behavior, as expected for a midshell nucleus. As a result, not only the first 2^+ state but also higher-lying states were excited. The 2_1^+ state has a lifetime of about 600 ps and therefore decays not at the target position but in flight between the target and the CD detector. Using this transition for the determination of ϕ angles is not suitable, as the angles must be referenced to the target position.

Instead, the decay of the 4_1^+ state was employed for the determination of the ϕ angles. Figure 7.9 (left) shows the extracted ϕ angles (black) with their statistical uncertainties, along with the values obtained after the minimization procedure. On the right-hand side, the Doppler-shifted energy of the 440 keV transition in ^{23}Na is shown in red, together with the corrected energies after minimization in black. Both plots demonstrate good agreement between the expected values and the minimized results. In the present experiment, a significant amount of ^{134}Ba was observed in the beam. Owing to its strong collectivity, substantial excitation of the 2_1^+ state was observed and used to verify the ϕ angles determined from the Yb calibration data.

7.7. Target Excitation

The γ -ray spectrum gated on coincident detection of ejectiles and recoils and Doppler-corrected for target-like nuclei is shown in Fig. 7.10. Several transitions in ^{194}Pt appear as sharp peaks, indicating that the target was excited up to at least the 8_1^+ and 6_2^+ states. Apart from the prominent ^{194}Pt lines, only a weak 355 keV transition from ^{196}Pt is visible, consistent with a target composition dominated by the desired isotope. However, the target excitation was mainly caused by Coulomb excitation of isobaric components present in the beam. A detailed analysis of the beam composition is therefore required to quantify their respective contributions and to properly account for target-excitation effects in the subsequent interpretation.

7.8. Beam Composition

Using the angles and distances determined in the previous sections, Doppler-corrected γ -ray spectra were reconstructed. In Fig. 7.11, the spectrum is shown, Doppler-corrected for the ejectiles. The $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ decay at 725 keV of ^{134}Sn is clearly visible and marked in red. However, several additional lines from transitions corresponding to deexcitation of other $A = 134$ isobars are also present, most prominently from ^{134}Te and ^{134}Ba . The much stronger intensities of these transitions suggest that the majority of the beam consisted of

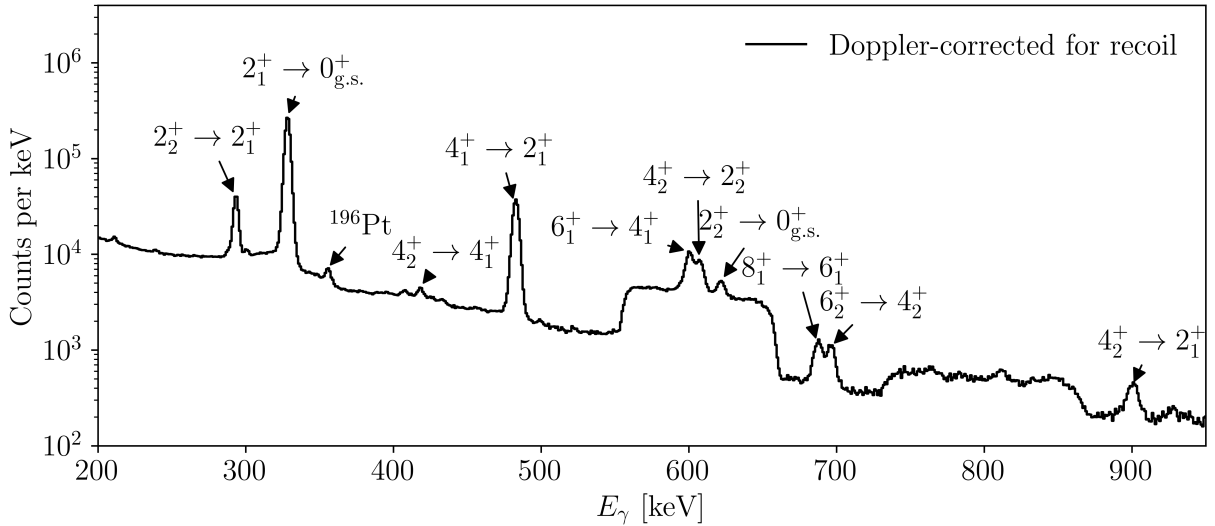


Figure 7.10. Doppler-corrected γ -ray spectrum of target-like recoils, gated on ejectiles and recoils. The visible peaks are predominantly assigned to known transitions in ^{194}Pt . A weak contribution from the 355 keV transition in ^{196}Pt is also observed.

these contaminants.

Given that the half-life of ^{134}Sn ($t_{1/2} = 1.04\text{ s}$) is comparable to the charge-breeding time ($\Delta t = 0.23\text{ s}$), a considerable amount of ^{134}Sb is expected to be present in the beam, produced by the β decay of ^{134}Sn during the breeding process. Therefore, the beam can be inferred to contain at least four components: ^{134}Sn , ^{134}Sb , ^{134}Te , and ^{134}Ba . The enhanced background at around 600 keV comes from the target excitation.

While the presence of odd-odd isotopes between these even-even nuclei cannot be excluded, no evidence for their excitation was observed during the experiment. Since γ transitions in the odd-odd isotopes typically occur at energies below 500 keV, a region dominated by the strong collective transitions from the ^{194}Pt target, the target material was changed to ^{206}Pb for a period of 16 h. This allowed for a more sensitive search for low-energy transitions, in particular from possible odd-odd beam contaminants. The spectra is added in Appendix B.1. No decays of excited states from ^{134}I or ^{134}Cs were observed. One peak in each, the uncorrected and the Doppler-corrected spectra may be attributed to excitations in ^{134}Sb .

The background spectrum shown in Fig. 7.12 displays γ rays emitted following the β

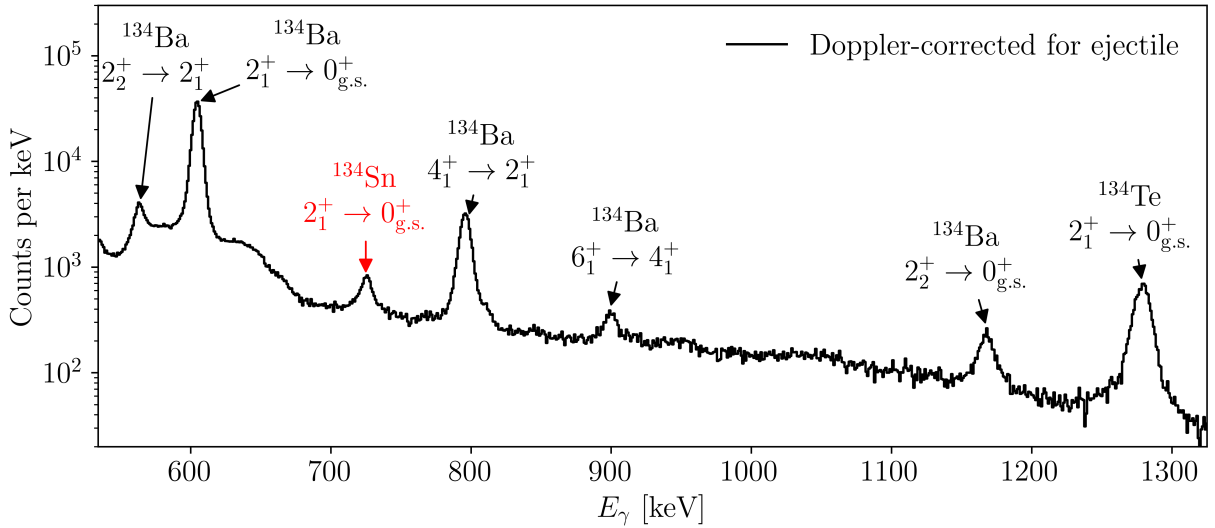


Figure 7.11. γ -ray spectra Doppler-corrected for the ejectiles, gated on all detected ejectiles and recoils. The $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition of ^{134}Sn , which is of primary interest for this work, is highlighted in red. In addition, several sharp transitions originating from other $A = 134$ isobars, in particular ^{134}Te and ^{134}Ba , are clearly visible. The relative intensities of these lines indicate that the beam contained a significant fraction of these isotopes, rather than being composed purely of ^{134}Sn . Further details on the identification and relative contributions of the individual components are discussed in the text.

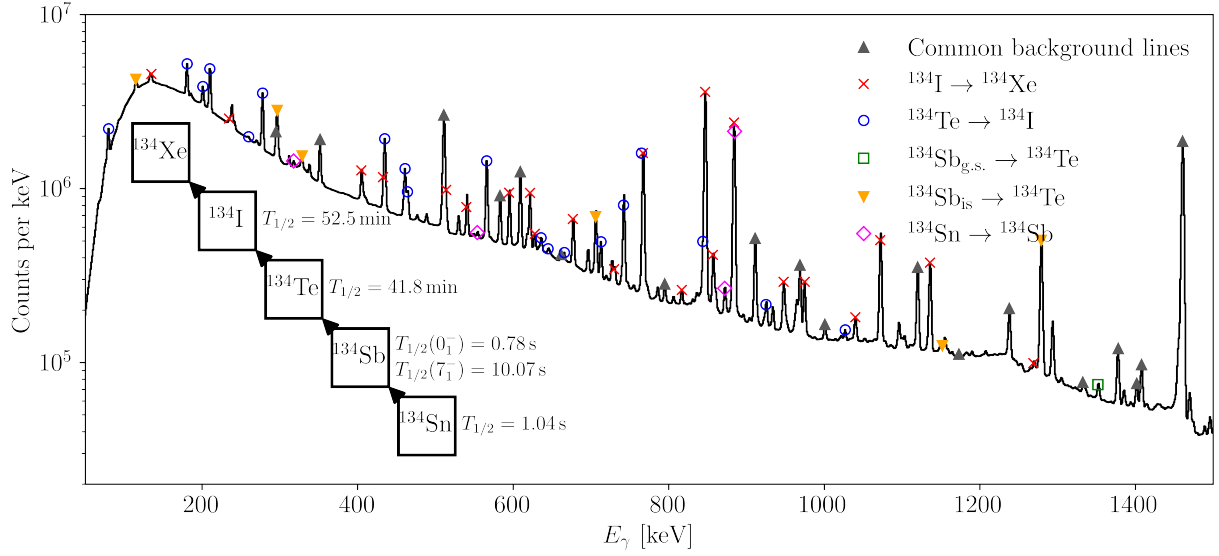


Figure 7.12. Background γ -ray spectrum measured without particle gating. Transitions assigned to the β -decay chain along the $A = 134$ isobars — from ^{134}Sn through ^{134}Sb to ^{134}Te — are marked, as well as typical background lines originating from environmental radiation. No transitions associated with unexpected isotopes were observed, confirming that the beam consisted solely of isobars with mass number $A = 134$.

decay of nuclei in the beam. Transitions corresponding to the decay chain within the $A = 134$ isobaric family are clearly identified and indicated in the figure, along with typical environmental and target-related background lines. No additional transitions from isotopes outside the $A = 134$ chain were observed, confirming that the beam composition was limited to isobars with $A = 134$.

The method of estimating the relative beam composition by comparing γ -ray intensities in the background spectra, as successfully applied in the case of ^{130}Sn , cannot be directly transferred to the present data. On the one hand, ^{134}Ba is stable, and its abundance in the beam therefore cannot be deduced from γ -ray decay signatures. On the other hand, the majority of ^{134}Sn nuclei decay directly to the ground state of ^{134}Sb , while only 6(3) % [68] populate excited states. This results in limited statistics for the corresponding γ peaks and introduces a relative uncertainty of approximately 50 % in their intensity determination. To obtain a more reliable estimate of the contributions from the individual beam components, the excitation of the target induced by each isotope was studied using a combination of in-beam Coulomb-excitation data, background spectra, and complementary measurements with an ionization chamber. The ionization-chamber measurements were carried out at four different times during the experiment to monitor potential temporal variations in

the beam composition. The methodology used to determine the beam composition from these combined observables is described in detail in the following sections. First, the total contributions of the long-lived ^{134}Te and stable ^{134}Ba components are determined. This is followed by an analysis of the ^{134}Sn -to- ^{134}Sb ratio.

7.8.1. ^{134}Te and ^{134}Ba

The Doppler-corrected, ejectile-gated γ -ray spectra are dominated by Coulomb excitation of ^{134}Te and ^{134}Ba . The level scheme including Coulomb-excited states observed in these beam-like isobars, as well as in the target-like isotope ^{194}Pt , is shown in Fig. 7.13. In all cases, an additional buffer state was included to account for possible unobserved feeding.

To quantify the respective contributions of ^{134}Ba and ^{134}Te to the excitation of the ^{194}Pt target, the yields of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transitions were analyzed. The underlying idea is to determine the amount of ^{194}Pt target excitation by providing the measured total yields from beam excitation (e.g. ^{134}Ba and ^{134}Te), together with the known matrix elements of beam and target isotopes, as input to GOSIA2, and then inferring the corresponding target-excitation yield.

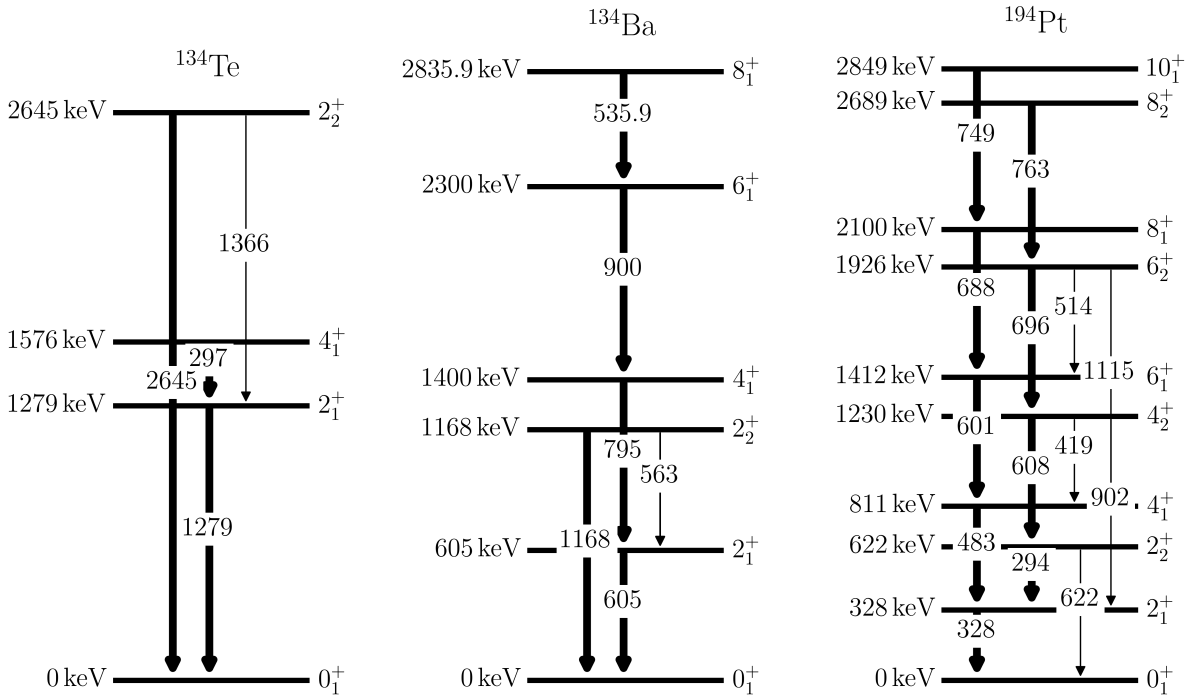


Figure 7.13. Level scheme of ^{134}Te , ^{134}Ba , and ^{194}Pt as observed in the Coulomb-excitation data. A buffer state was added at the top of each scheme to account for unobserved feeding.

The matrix elements used for the 2_1^+ excitations are $\langle 2_1^+ || E2 || 0_{g.s.}^+ \rangle = 0.827(7) \text{ eb}$ for ^{134}Ba , $\langle 2_1^+ || E2 || 0_{g.s.}^+ \rangle = 0.323(6) \text{ eb}$ for ^{134}Te , and $\langle 2_1^+ || E2 || 0_{g.s.}^+ \rangle = 1.285(7) \text{ eb}$ for ^{194}Pt . Matrix elements, lifetimes, branching ratios, and multipolarities of higher-lying transitions in these isotopes are listed in Tab. B.1 in Appendix B.2. Because feeding from higher-lying states affects the population of the 2_1^+ levels, these quantities were included in the GOSIA2 calculations.

For the determination of γ -ray yields, Gaussian functions were fitted to the corresponding $2_1^+ \rightarrow 0_{g.s.}^+$ photopeaks. To maximize statistics, the data were summed over all CD rings while applying the kinematic gate on the ejectiles. In total, 132 289(500) and 4 168(100) counts were observed for the $2_1^+ \rightarrow 0_{g.s.}^+$ transitions in ^{134}Ba and ^{134}Te , respectively, and 694 002(1 000) counts for the corresponding transition in the ^{194}Pt target.

The ^{194}Pt target enrichment of 96.45 % was taken into account by scaling the ejectile-induced yields accordingly. After correcting for detector efficiencies, normalized yields for the ejectile excitations were deduced.

For the GOSIA2 analysis, stopping powers and internal conversion coefficients were calculated using SRIM and BRICC, respectively. The ^{194}Pt target thickness of 3 mg/cm^2 results in an energy loss of 80 MeV for the beam as it traverses the target. As in the ^{130}Sn analysis, the lower-energy gates applied to the beam-like particles suppress ejectiles scattered on the

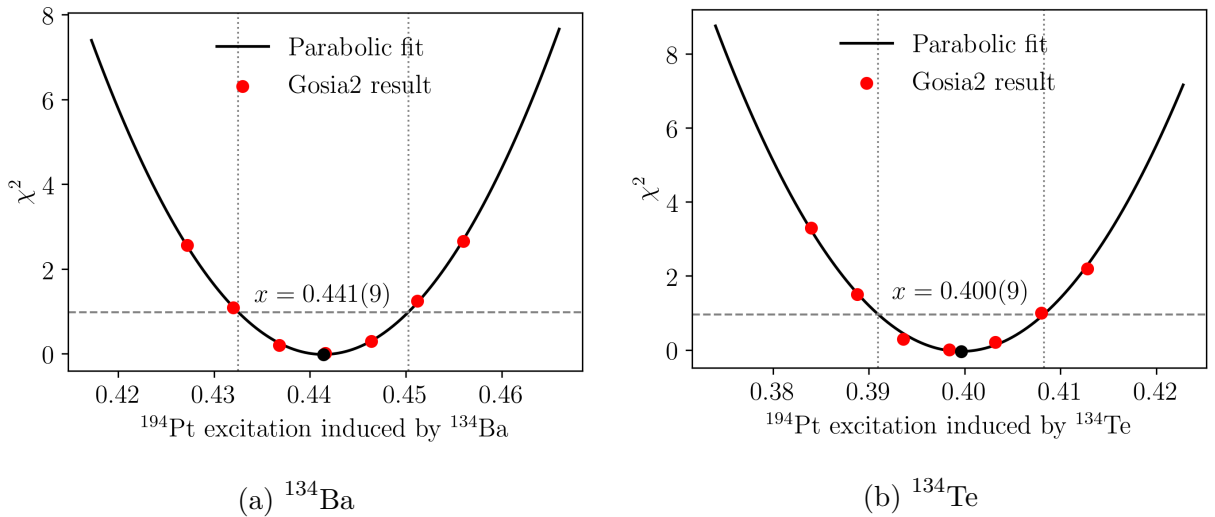


Figure 7.14. χ^2 dependence of the GOSIA2 fit on the assumed fraction of ^{194}Pt target excitation induced by (a) ^{134}Ba and (b) ^{134}Te . The solid curve shows a parabolic fit around the minimum; the quoted uncertainty is defined by the $\chi^2_{\min} + 1$ criterion.

upstream side of the target. For the present analysis, this effectively limits the maximum beam energy to 585 MeV. Consequently, an energy interval of 510–585 MeV was adopted, together with the full angular coverage of the CD, corresponding to scattering angles of 14.9° – 50.4° .

The fraction of target excitation induced by a given isotope is here defined as x_{isotope} . The measured (efficiency-corrected) ^{194}Pt target yield and its uncertainty were scaled by trial values of x_{isotope} and used as input for the GOSIA2 calculations. GOSIA2 was then used to integrate the beam-like and scaled target-like intensities over the ranges of beam energies and scattering angles. Based on these integrated yields, a χ^2 minimization was performed with GOSIA2, resulting in a χ^2 value for each assumed x_{isotope} .

The determined χ^2 data points are shown in Fig. 7.14 for ^{134}Ba and ^{134}Te with respect to x_{isotope} . A parabolic function was fitted in the vicinity of the minimum, and the uncertainty was extracted using the $\chi^2_{\text{min}} + 1$ criterion. This procedure yields target-excitation fractions of $x_{^{134}\text{Te}} = 40.0(9)\%$ and $x_{^{134}\text{Ba}} = 44.1(9)\%$. The $15.8(13)\%$ target excitation not attributed to either ^{134}Ba or ^{134}Te is assigned to ^{134}Sn and ^{134}Sb .

The $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition yields, integrated over 4 h intervals, are shown on the left-hand side of Fig. 7.15. Applying the procedure described above, the time-dependent beam

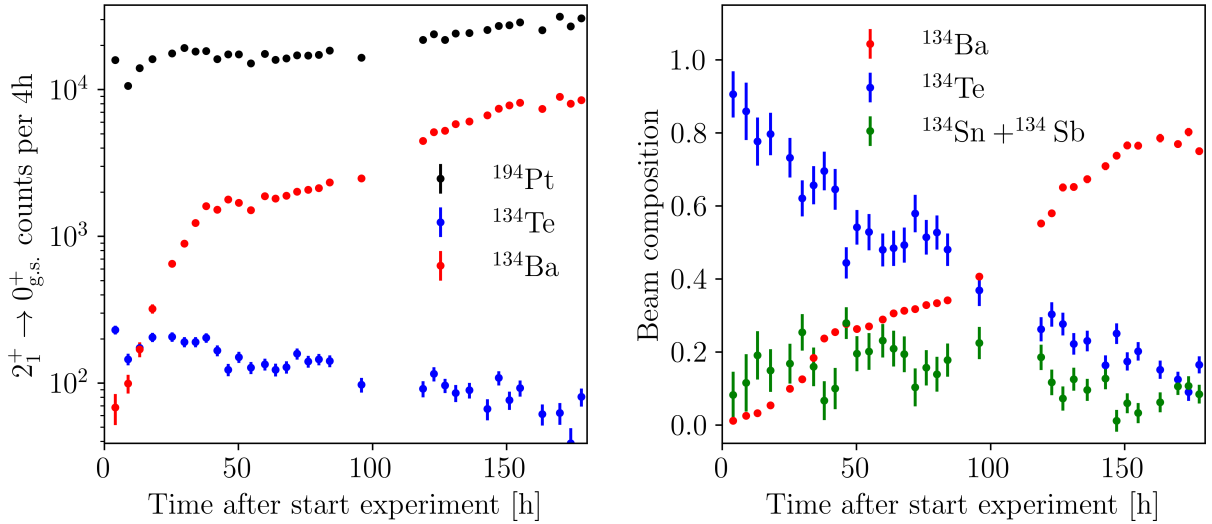


Figure 7.15. Time evolution of the beam composition deduced from Coulomb-excitation data integrated over four-hour intervals. The ^{134}Te fraction decreases with time, while the ^{134}Ba component increases. The remaining intensity is attributed to $^{134}\text{Sn} + ^{134}\text{Sb}$, which stays approximately constant at the 10–20 % level over the course of the experiment.

composition was deduced and is shown on the right-hand side of the same figure.

The relative contributions of ^{134}Ba and ^{134}Te in the beam exhibit a pronounced time dependence: at the beginning of the measurement, the beam consisted almost exclusively of ^{134}Te , whereas the stable contaminant ^{134}Ba accumulated continuously during operation. This led to an increasing Ba fraction and a corresponding reduction of the ^{134}Te component. In contrast, the remaining $^{134}\text{Sn} + ^{134}\text{Sb}$ component stays nearly constant over the full running period.

For the subsequent GOSIA2 analysis of ^{134}Sn , the beam composition does not need to be known at every point in time. The decisive quantity is the amount of target excitation induced by ^{134}Sn , which serves as the normalization reference. Variations in the $^{134}\text{Te}/^{134}\text{Ba}$ ratio therefore do not compromise the validity of the ^{134}Sn analysis.

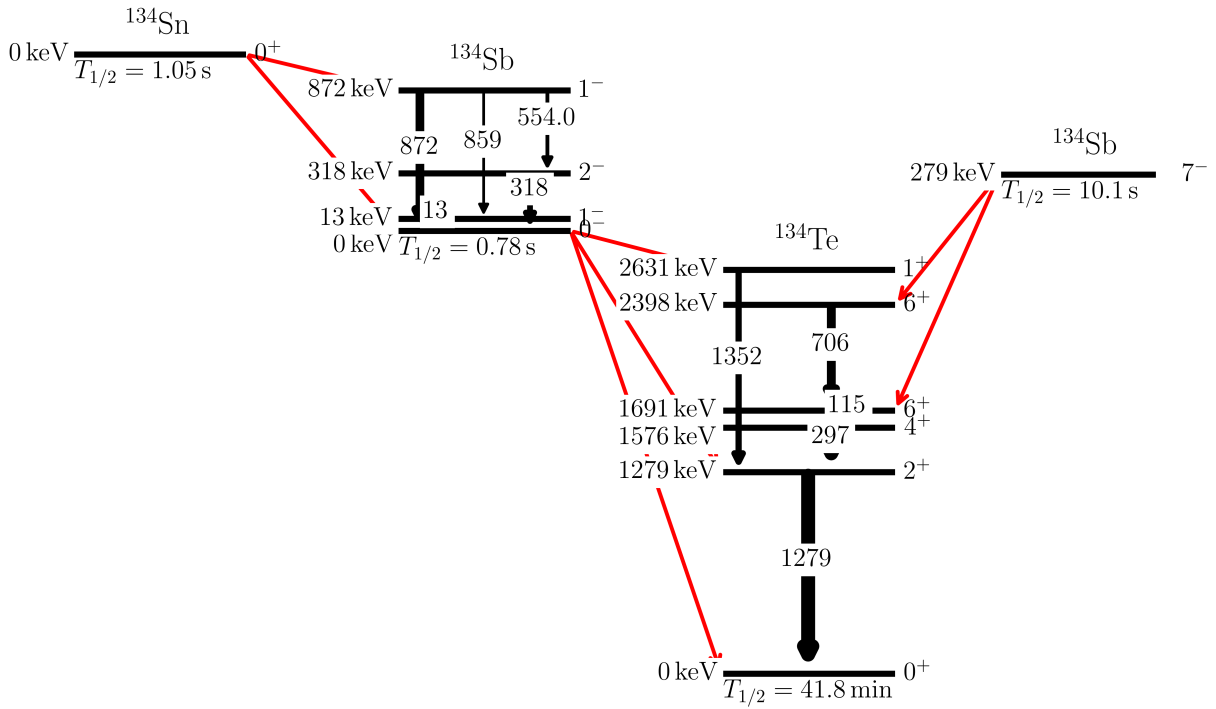


Figure 7.16. Decay scheme along the $A = 134$ isobar chain from ^{134}Sn to ^{134}Sb and ^{134}Te . The 7^- isomer in ^{134}Sb is included because a fraction of the delivered beam was observed to populate this state. The subsequent decay sequence beyond ^{134}Te is omitted, as it is not relevant for determining the $^{134}\text{Sn}/^{134}\text{Sb}$ ratio in the beam. The arrow width is proportional to the absolute intensity. For more information, see text.

7.8.2. ^{134}Sn and ^{134}Sb

Coulomb-excitation data cannot be used to determine the ^{134}Sb contribution in the beam. On the one hand, only a few counts corresponding to excitations in ^{134}Sb were observed in the experiment with the lead target (shown in Appendix B.2); on the other hand, no transition probabilities are known for ^{134}Sb to date. Therefore, the β -delay γ -ray data was used instead. The decay scheme relevant for the $^{134}\text{Sn}/^{134}\text{Sb}$ ratio in the beam is shown in

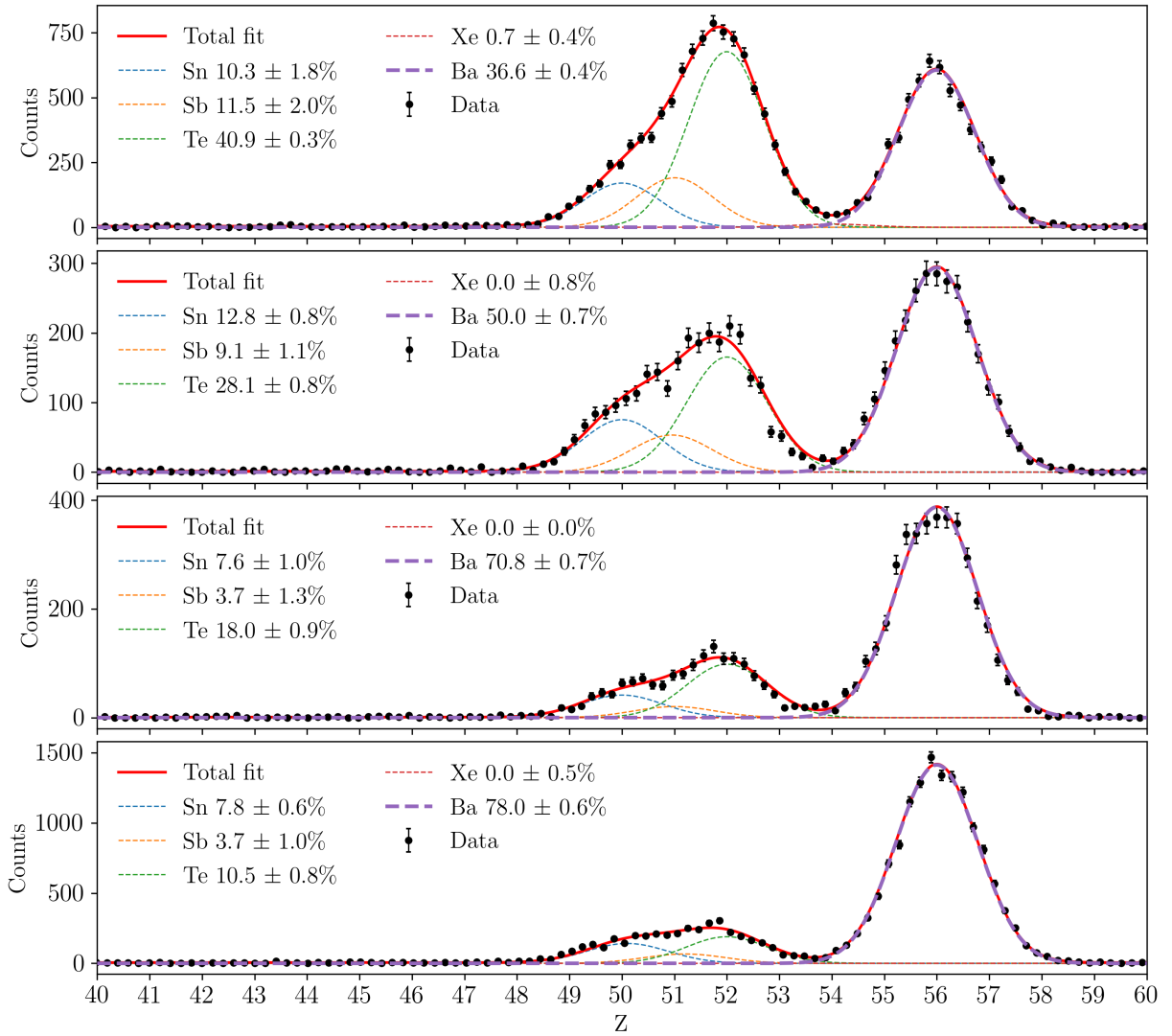


Figure 7.17. Ionization-chamber spectra recorded at four different times during the experiment (panels (a)–(d), in chronological order). The measurements were taken once per day, starting on the third day of the beam time. Gaussian peak functions were fitted to the individual Z components to extract the beam composition shown in the legend. The relative beam composition changes over the course of the experiment.

Fig. 7.16. For clarity, only the dominant decay branches are included. Since ^{134}Sb has a 0^- ground state, the β^- decay of ^{134}Sn predominantly populates the ground state of ^{134}Sb directly. The same holds for the decay of ^{134}Sb into ^{134}Te , where the ground state of ^{134}Te is populated in the dominant branch. The half-lives of ^{134}Sn and ^{134}Sb are short compared to the duration of the experiment; therefore, a constant decay equilibrium can be assumed. Under this assumption, the number of decays of ^{134}Sn and ^{134}Sb equals their respective production rates.

The isotope ^{134}Sb exhibits an isomeric 7^- state that decays via β^- decay. Figure 7.12 shows, that γ -rays attributed to the β decay of this isomeric Sb have been measured. Consequently, three contributions have to be considered in the background spectra: ^{134}Sn , $^{134}\text{Sb}_{\text{g.s.}}$, and $^{134}\text{Sb}_{\text{is.}}$. For $^{134}\text{Sb}_{\text{is.}}$, the 1279 keV transition can be used, which is populated with an absolute intensity of 100(5) %. The ground state $^{134}\text{Sb}_{\text{g.s.}}$ is identified via the β -delayed γ ray at 1352 keV with an absolute intensity of 0.93(5) %. 1.1(5) % of the decaying $^{134}\text{Sb}_{\text{g.s.}}$ isotopes emit a 1279 keV γ ray, this intensity was in the following analysis subtracted for the identification of the isomeric $^{134}\text{Sb}_{\text{is.}}$.

The limiting factor for a precise determination of the ^{134}Sn fraction is the large uncertainty of the absolute intensity of the 872 keV transition in ^{134}Sn , $I_{\text{abs.}} = 6(3)$ %. This uncertainty is largely driven by the limited knowledge of the ground-state feeding of the β^- decay; therefore, using alternative γ -ray transitions does not improve the achievable accuracy [152, 153].

To refine this value, ionization chamber (IC) measurements were performed at four different times during the experiment. The chamber was mounted downstream of the experimental setup, and the target was removed to minimize energy straggling. The IC data were calibrated using well-defined reference points. When the primary proton beam on the uranium target was switched off, no short-lived ^{134}Sn and ^{134}Sb isotopes reached the charge-breeding system, leaving only ^{134}Ba and ^{134}Te in the beam. If the beam line between the primary target and the EBIS was additionally blocked, only the noble gas xenon, remaining in the EBIS chamber, was extracted. These three well-defined beam conditions provided reliable calibration points for the IC spectra.

The resulting calibrated IC measurements are shown in Fig. 7.17. In this experiment, the ionization chamber measured the energy loss (ΔE) in the gas as a function of the residual energy deposited in the silicon detector located downstream. Owing to its superior resolution, the energy-loss signal was used for the calibration. Gaussian functions were fitted

to the peaks associated with the four expected beam components (^{134}Sn , ^{134}Sb , ^{134}Te , and ^{134}Ba), and an additional fit component was included for ^{134}Xe . The xenon contribution, however, was found to be negligible and is therefore not considered further.

A clear temporal evolution of the beam composition is observed. At the beginning of the measurement, the beam is dominated by ^{134}Te . With time, the ^{134}Te fraction decreases while the ^{134}Ba component increases. This behavior is consistent with the trend inferred from the Coulomb-excitation analysis shown in Fig. 7.15. The relative contributions of ^{134}Sn and ^{134}Sb also vary over the course of the experiment.

The combined amount of target excitation induced by the $^{134}\text{Sn} + ^{134}\text{Sb}$ beam components was determined in the previous chapter. It is thus sufficient to determine the $^{134}\text{Sn}/(^{134}\text{Sb} + ^{134}\text{Sn})$ ratio in order to constrain the overall beam composition. The approach is to combine the information from the ionization chamber—which provides the ^{134}Sn -to- ^{134}Sb ratio at four

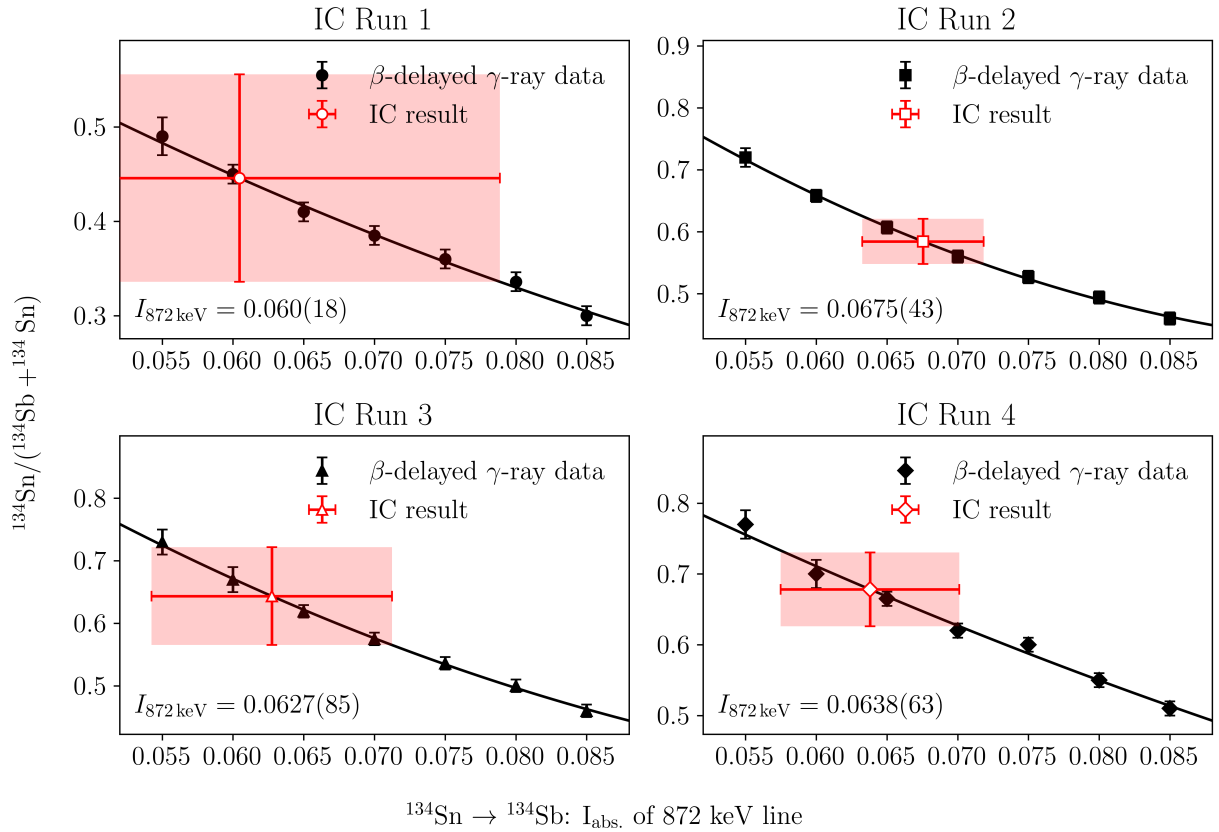


Figure 7.18. Ratio of $\text{Sn}/\text{Sn}+\text{Sb}$ for different assumed absolute intensities of the 872 keV transition in the decay of ^{134}Sn . The measured ratios from the ionization chamber are compared to the calculated values, and the corresponding absolute intensities are deduced for all four measurements.

discrete times—with β -delayed γ -ray data. From the ionization chamber measurements, ratios of $^{134}\text{Sn}/(^{134}\text{Sn}+^{134}\text{Sb}) = 0.45(11), 0.58(4), 0.64(8),$ and $0.68(5)$ were obtained for the four IC runs.

The β -delayed γ -ray spectra were analyzed in time windows spanning two hours before and two hours after each ionization chamber run. The intensities of the 872 keV, 1279 keV, and 1352 keV transitions—associated with the decay of ^{134}Sn , $^{134}\text{Sb}_{7-}$, and $^{134}\text{Sb}_{\text{g.s.}}$, respectively—were determined. The 1279 keV and 1352 keV peak intensities were corrected for the relative detection efficiency with respect to the 872 keV line, as well as for their known absolute γ -ray intensities. Using these corrected counts, the ratio $^{134}\text{Sn}/(^{134}\text{Sn}+^{134}\text{Sb})$ was evaluated as a function of the (unknown) absolute intensity of the 872 keV transition. This procedure was carried out using the peak intensities extracted for the time intervals before and after each of the four ionization chamber runs. The resulting values are shown as black points in Fig. 7.18 together with a parabolic fit function.

For each IC run, the independently measured value of $^{134}\text{Sn}/(^{134}\text{Sn}+^{134}\text{Sb})$, together with its 1σ uncertainty, is indicated in the corresponding panel as a vertical red line. By projecting this constraint onto the abscissa, the absolute intensity of the 872 keV transition is obtained for each of the four runs. The resulting uncertainties for the absolute 872 keV intensity and the given uncertainty for the $^{134}\text{Sn}/(^{134}\text{Sn}+^{134}\text{Sb})$ values are highlighted as red rectangle. The resulting values for $I_{872\text{ keV}}$ are given for each ionization chamber run. The weighted mean of the four determinations yields an absolute intensity of $I_{\text{abs}} = 6.56(33)\%$ for the

Table 7.2. Yields of the decay transitions observed for the ^{134}Sn and ^{134}Sb components in the beam, together with their absolute intensities and the relative detection efficiency of the setup. Using the corrected intensities, the relative contributions to the target excitation induced by the three isotopes were determined. The ratios of ^{134}Te and ^{134}Ba in the beam and for the target excitation are added as well.

Isotope	E_γ [keV]	Counts 10^3	I_{abs} [%]	Efficiency [%]	Corr. Int. [10^3]	^{194}Pt exc. [%]	Beam comp. [%]
^{134}Sn	872	256(10)	6.56(33)	8.01(2)	486(31)	6.98(60)	7.37(70)
$^{134}\text{Sb}_{\text{g.s.}}$	1352	48(2)	0.93(5)	6.02(2)	855(53)	5.42(68)	6.46(81)
$^{134}\text{Sb}_{\text{is.}}$	1279	1443(3)	100(5)	6.24(2)	231(11)	3.41(41)	4.06(49)
^{134}Te						40.0(9)	39.1(11)
^{134}Ba						44.1(9)	43.0(11)

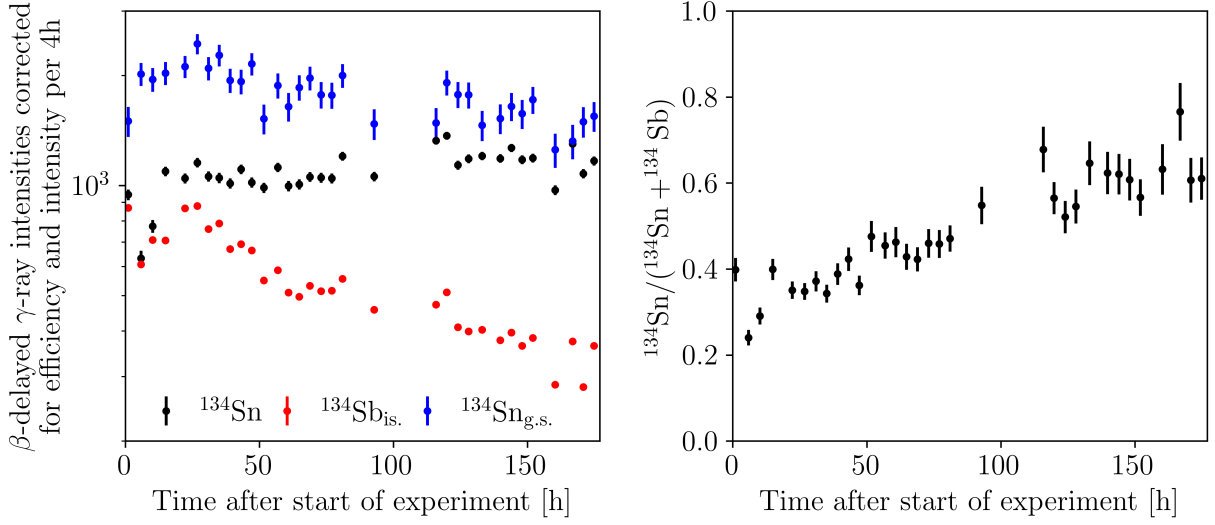


Figure 7.19. Left: β -delayed γ -ray intensities of the 872, 1279 and 1352 keV transitions of ^{134}Sn , $^{134}\text{Sb}_{\text{is.}}$ and $^{134}\text{Sb}_{\text{g.s.}}$, respectively, integrated over 4 h of beam time. The shown yields are corrected for efficiencies and absolute intensities of the transitions. Right: With these values the ratio $^{134}\text{Sn}/(^{134}\text{Sn} + ^{134}\text{Sb})$ was determined for each time.

872 keV transition in the decay of ^{134}Sn .

Using this value, together with the yields of the 872, 1352, and 1279 keV transitions over the course of the whole experiment, the fraction of target excitation induced by ^{134}Sn , $^{134}\text{Sb}_{\text{g.s.}}$, and $^{134}\text{Sb}_{\text{is.}}$ was determined. The corresponding results are summarized in Table 7.2. A total of 6.98(60) % of the observed target excitation is induced by ^{134}Sn , corresponding to a beam composition of 7.37(70) % ^{134}Sn .

Although this fraction appears small compared to the values obtained from the ionization-chamber measurements, it is consistent with the overall experimental trend. The tellurium and antimony components decrease rapidly over time: at the beginning of the experiment, Sb and Te dominate the beam composition, while only a small amount of Sn is observed. This is shown in Fig. 7.15 for the ^{134}Te component. In Fig. 7.19 the time evolution of the ratio $^{134}\text{Sn}/(^{134}\text{Sn} + ^{134}\text{Sb})$ is shown on the right, based on the time-dependent β -delayed γ -ray yields. These are shown in Fig. 7.19 on the left, corrected for efficiency and absolute intensity.

The intensity of the 1279 keV decay line from the isomeric $^{134}\text{Sb}_{\text{is.}}$ state decreases by about a factor of four during the experiment, and the rate of the Coulomb-excited $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in ^{134}Te drops by more than a factor of two even before the first ionization-

chamber measurement. Thus, although the inferred Sn fraction in the beam is lower than suggested by the ionization-chamber data, it is fully consistent with the observed temporal evolution of the beam composition.

7.9. Gosia2 Analysis of the 2_1^+ State in ^{134}Sn

The $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ and $Q_s(2_1^+)$ values were deduced similar to the values in ^{130}Sn , described in Sec.6.10. The χ^2 minimization was again performed with GOSIA2.

The data was divided into nine angular regions, corresponding to distinct ranges of scattering

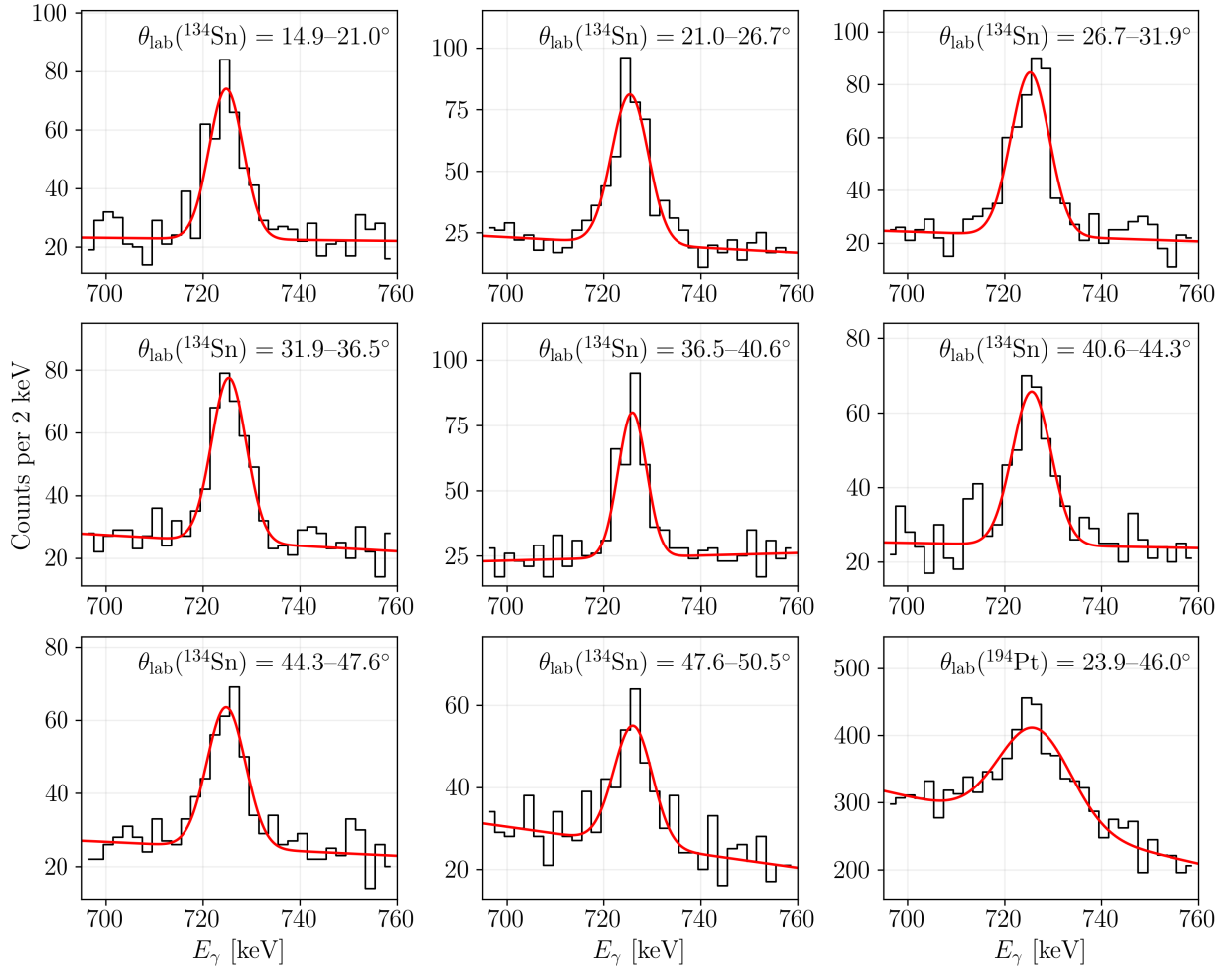


Figure 7.20. Particle-gated, Doppler-corrected γ -ray spectra for the $^{134}\text{Sn}(^{196}\text{Pt}, ^{196}\text{Pt})^{134}\text{Sn}$ reaction correct for beam-like particles. Each panel corresponds to a distinct range in the ejectile scattering angle, obtained by combining two adjacent rings of the CD detector. The red curves represent Gaussian fits superimposed on a linear background, used to extract the intensity of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition at 725 keV.

angles for the detected ejectiles. For the ejectile kinematics, two adjacent rings of the CD detector were combined in each case, while for the recoil-gated kinematics, the respective recoil-angle ranges were used. The resulting Doppler-corrected γ -ray spectra are shown in Fig. 7.20. The intensity of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in ^{134}Sn was determined by fitting a Gaussian peak on top of a linear background to the experimental data. In total, 3396(150) counts were assigned to this transition. To correct for impurities of the target material, the extracted intensities were corrected by dividing by a factor of 96.45 % target content of ^{194}Pt . The resulting yields were then used as input for the GOSIA analysis to determine the relevant electromagnetic matrix elements.

The observed yield of the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition in the ^{194}Pt target was scaled by a factor of 6.98(69) %, as determined in the previous chapter, and corrected for the relative photopeak-efficiency difference between the 328 keV transition in ^{194}Pt and the 725 keV transition observed in the ejectile channel. Since the excitation probability of the 2_1^+ state in ^{194}Pt depends on the projectile charge Z and on the scattering angle θ , an additional angular-dependent correction was applied to account for the different effective Coulomb-

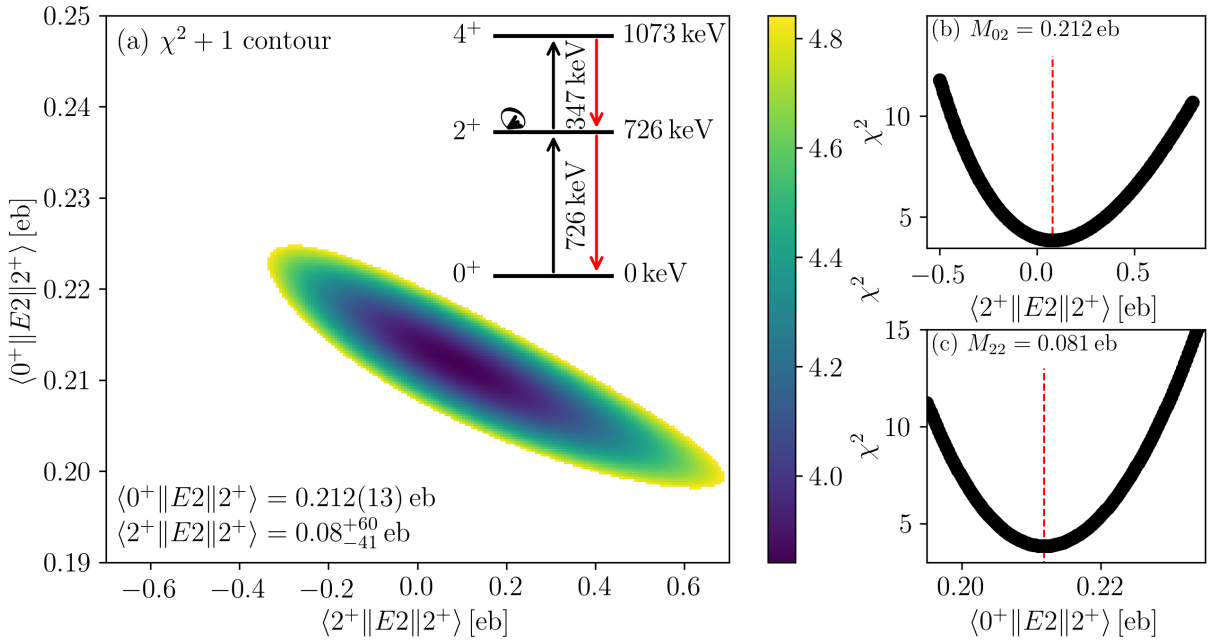


Figure 7.21. Left: The 1σ contour of the normalized total χ^2 surface as a function of the transitional matrix element and the spectroscopic quadrupole moment of the 2_1^+ state in ^{134}Sn . The χ^2 values were obtained with GOSIA2. For details of the analysis, see the text. Right: The projection onto the resulting transitional (b) and diagonal (c) matrix elements and the χ^2 distribution around the minima, which is given as a red, vertical line

excitation cross sections sampled by the individual detector rings. This correction factor was determined separately for each ring using GOSIA2.

The total energy loss in the target and the integration range of the beam energy from 510–585 MeV has been discussed in the previous section. The scattering angles have been determined based on the target-to-CD distance. Two rings were combined for ejectile-gated spectra and the scattering range is given in Fig. 7.20 for each spectra. Stopping powers and internal conversion coefficients were calculated using SRIM and BRICC. All matrix elements for ^{194}Pt listed in Table B.1 were included in the GOSIA2 analysis.

For ^{134}Sn , the first 4^+ state was included as a buffer state. Its known lifetime of $\tau(4_1^+) = 1707(57)$ ps was implemented as an additional constraint in the minimization procedure [112]. No depopulating γ -ray transitions from the 4^+ state were observed experimentally. The GOSIA2 option `OP,INTI` was again used to integrate the measured yields over the scattering angles and beam energy.

The transitional $\langle 2_1^+ \| E2 \| 0_{\text{g.s.}}^+ \rangle$ and diagonal $\langle 2_1^+ \| E2 \| 2_1^+ \rangle$ matrix elements were then varied simultaneously, and a total χ^2 was calculated with the `OP,MINI` option for both the projectile and the target excitation yields for each pair of matrix elements.

Figure 7.21 shows the $\chi^2 + 1$ contour surface as a function of the diagonal matrix element and the corresponding spectroscopic quadrupole moment, where $Q_s(2_1^+) = 0.7579 \times \langle 2_1^+ \| E2 \| 2_1^+ \rangle$, on the left. A well-defined minimum is obtained at $\langle 2_1^+ \| E2 \| 0_{\text{g.s.}}^+ \rangle = 0.212(13)_{\text{stat.}} \text{ eb}$ and $Q_s(2_1^+) = 0.06_{-0.31}^{+0.45} \text{ eb}$. This result is corroborated by the one-dimensional projections shown on the right-hand side, taken at the best-fit values $M_{02} \equiv \langle 0_{\text{g.s.}}^+ \| E2 \| 2_1^+ \rangle = 0.212 \text{ eb}$ (top) and $M_{22} \equiv \langle 2_1^+ \| E2 \| 2_1^+ \rangle = 0.081 \text{ eb}$ (bottom).

The uncertainties arising from efficiency ratios and beam composition are already included in these statistical errors. Additional systematic uncertainties were estimated by varying the scattering angles according to a 0.5 mm uncertainty in the target-to-detector distance and a 5 % uncertainty in the target thickness, which translates into a corresponding uncertainty in the beam energy of 5 mm. Including these effects leads to final adopted values of

$$\begin{aligned} \langle 2_1^+ \| E2 \| 0_{\text{g.s.}}^+ \rangle &= 0.212(13)_{\text{stat.}}(5)_{\text{syst.}} \text{ eb}, \\ Q_s(2_1^+) &= 0.06_{-0.31\text{stat.}}^{+0.45} (3)_{\text{syst.}} \text{ eb}. \end{aligned}$$

for the reduced $E2$ matrix element and the spectroscopic quadrupole moment, respectively.

From the extracted transition matrix element, the corresponding reduced transition strength is obtained as

$$B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.0449(57) e^2 \text{b}^2.$$

For comparison, imposing $Q_s(2_1^+) = 0$ (consistent with a spherical, non-deformed 2_1^+ state) gives $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.046(4) e^2 \text{b}^2$.

8. Discussion

The experimental results reported here provide the opportunity to test and constraint shell-model descriptions in the vicinity of the $Z = 50$ $N = 82$ shell closure. The discussion will proceed as follows. First, the new $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values are compared to available theoretical predictions, with particular emphasis on clarifying the origin of the deviations between earlier experimental results and MCSM calculations, and on placing the findings into the broader systematics along the Sn isotopic chain. Second, the newly determined quadrupole moments are discussed and compared with theory.

The matrix elements obtained from the combined Coulomb-excitation and lifetime analysis provide direct access to the reduced transition strengths. Table 8.1 summarizes the newly deduced values in $^{130,134}\text{Sn}$ and compares them to previously published results. The old $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values are taken from the conference proceedings [57–60], as no peer-review literature is published to date. The $2_1^+ \rightarrow 4_1^+$ transition strength in ^{134}Sn could not be determined from the present data set; however, it was adopted as an external constraint in the analysis and is therefore added for completeness, as well as the $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ value in doubly-magic ^{132}Sn from Ref. [23].

The lifetime of the 5_1^- state determined in this thesis, $\tau = 52(3)$ ns, is consistent with the weighted average of earlier measurements, $\tau = 54.8(22)$ ns [94, 102, 103]. Converting the deduced matrix elements into reduced transition probabilities shows that the combined approach of lifetime and Coulex substantially improves the precision for the $B(E2; 7_1^- \rightarrow 5_1^-)$ and $B(E1; 4_1^+ \rightarrow 5_1^-)$ significantly compared to the literature, while the new values remain compatible with previously reported results within uncertainties. The good agreement of the GOSIA2 analysis with the lifetime result underlines and confirms the accurate $^{130}\text{Sn}_{\text{g.s.}}$ and isomeric $^{130}\text{Sn}_{7-}$ beam composition determination of this work.

Earlier tentative determinations of $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ in ^{130}Sn and ^{134}Sn — $B(E2) = 0.023(5)$ and $0.029(5) e^2\text{b}^2$, respectively — reported in the proceeding papers Refs. [57–60] — disagree with the present values. This discrepancy can plausibly be attributed to (i) limited statistics

in both Oak Ridge data sets and (ii) insufficient γ -ray energy resolution in the ^{134}Sn case, which complicates peak separation and contaminant subtraction. In addition, for the ^{130}Sn experiment performed at Oak Ridge, the beam-composition analysis relied on absolute intensities of β -delayed γ rays that have since been revised. A refined reassessment of the beam composition indicates an increased fraction of isomeric $^{130}\text{Sn}_{7-}$ in the beam, which directly affects the normalization of Coulomb-excitation yields and, consequently, the extracted $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ strength [154]. Taken together, these considerations provide a consistent explanation for the discrepancy of the earlier values that are lower relative to the present results.

8.1. 2_1^+ States in $^{130,134}\text{Sn}$

The new $B(E2)$ information obtained for ^{130}Sn and, in particular, for ^{134}Sn , since no $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values are known for $A > 132$ Sn isotopes, serves as a sensitive probe of modern shell-model descriptions. These nuclei test both the choice of effective charges and the degree to which excitations across the $N = 82$ gap (i.e. polarization and possible breaking of the ^{132}Sn core) contribute to quadrupole collectivity.

Figure 8.1 compares the newly determined $B(E2)$ values for $^{130-134}\text{Sn}$ and ^{132}Sn with predictions from several theoretical approaches. The $B(E2)$ value for ^{132}Sn is taken from Ref. [23], while those for ^{130}Sn and ^{134}Sn are obtained in the present work. The gray bands denote the corresponding experimental uncertainties.

Table 8.1. $B(E\lambda)$ values determined in this work for $^{130,134}\text{Sn}$. The $B(E2)$ values for $0_{\text{g.s.}}^+ \rightarrow 2_1^+$ transitions for $^{130,134}\text{Sn}$ are compared to values determined in Oak Ridge from Ref. [57–60]. For completeness, the $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ value in ^{132}Sn determined in Ref. [23] is added as well as the $B(E2; 2_1^+ \rightarrow 4_1^+)$ value determined from lifetime measurements of Ref. [112].

Nucleus	Transition $J_i^\pi \rightarrow J_f^\pi$	Multipolarity $[\sigma\lambda]$	$B(\lambda L)_{\text{this work}}$ $[e^2\text{b}^\lambda]$	$B(\sigma\lambda)_{\text{lit.}}$ $[e^2\text{b}^\lambda]$
^{130}Sn	$0_{\text{g.s.}}^+ \rightarrow 2_1^+$	$E2$	0.0486(57)	0.023(5)
	$7_1^- \rightarrow 5_1^-$	$E2$	0.0042(2)	0.0040(11)
	$4_1^+ \rightarrow 5_1^-$	$E1$	$71(3) \times 10^{-6}$	$72(10) \times 10^{-6}$
^{132}Sn	$0_{\text{g.s.}}^+ \rightarrow 2_1^+$	$E2$	–	0.087(19)
^{134}Sn	$0_{\text{g.s.}}^+ \rightarrow 2_1^+$	$E2$	0.0449(57)	0.029(5)
	$2_1^+ \rightarrow 4_1^+$	$E2$	–	0.0241(9)

The MCSM calculations by Togashi *et al.* [54] are carried out in a valence space comprising the proton and neutron orbitals $1g_{9/2}$, $2d_{5/2,3/2}$, $3s_{1/2}$, $1h_{11/2}$, $2f_{7/2}$, and $3p_{3/2}$. Owing to the maximal dimension up to 7.5×10^{23} , these calculations are beyond the reach of conventional shell-model diagonalization methods. This is made tractable in the MCSM by diagonalizing the Hamiltonian in a truncated space spanned by a stochastically generated set of deformed Slater-determinant basis vectors. Consequently, the calculation can be performed across the full tin isotope chain and, in the present case, provides theoretical values beyond $^{130-134}\text{Sn}$. The underlying Hamiltonian is obtained through targeted modifications of selected two-body matrix elements (TBMEs), which are subsequently kept fixed throughout the isotope chain; importantly, these adjustments do not use information from the doubly-magic nucleus ^{132}Sn and neighboring $^{130,134}\text{Sn}$. Effective charges of $e_p = 1.25e$ and $e_n = 0.75e$ are adopted consistently for all tin isotopes.

The LSSM calculations of Naïdja, Nowacki and Sieja (hereafter denoted NNS110) employ a model space with neutron orbitals $1h_{11/2}$, $2f_{7/2}$, $1h_{9/2}$, $2f_{5/2}$, $3p_{3/2}$, and $3p_{1/2}$ and proton orbitals $1g_{9/2}$, $1g_{7/2}$, $2d_{5/2,3/2}$, and $3s_{1/2}$ outside an inert ^{110}Zr core, also spanning below and above the ^{132}Sn shell closure. As a starting point, a realistic interaction derived from the CD-Bonn potential is used, renormalized within the $V_{\text{low-}k}$ framework and mapped onto the chosen valence space using many-body perturbation theory, including all diagrams up to second order. Details of the interaction construction are discussed in Refs. [56, 155], while the results shown here are taken from Refs. [23, 24]. The calculations use effective charges of $e_p = 1.68e$ and $e_n = 0.41e$.

For ^{130}Sn , results of different realistic shell models are also shown. The Sr88MHJM interaction (in the following denoted as ^{88}Sr core) is likewise rooted in the CD-Bonn potential, with a renormalization performed via a G -matrix approach. The G matrix is derived from the underlying nucleon–nucleon interaction, including core-polarization corrections, and is subsequently adjusted to experimental data [63, 156, 157]. Single-particle energies are inferred from the one-proton and one-neutron neighbors of ^{88}Sr , and additional monopole corrections are introduced to reproduce single-neutron particle energies and single-proton hole energies in the tin isotopes. Furthermore, selected two-body matrix elements are adjusted to optimize the agreement with experimental level energies in the tin chain. The calculations were performed by A. Blazhev [158] and are taken from Ref. [86] using an effective neutron charge of $e_n = 1.0e$. The results are denoted as ^{88}Sr core in Fig. 8.1.

A third realistic approach is included, based on a G -matrix renormalization of the CD-Bonn nucleon–nucleon potential as developed by Banu *et al.* [69], but employing different valence-space choices. In this framework, either an inert ^{100}Sn core or, alternatively, a ^{90}Zr core may be adopted as the closed-shell reference. For completeness, the present comparison therefore includes the results obtained with both core assumptions, and both sets of predictions are shown in Fig. 8.1, marked with their used cores.

Beyond the doubly-magic nucleus ^{132}Sn , the $\nu(2f_{7/2})$ orbital becomes occupied and provides the dominant degrees of freedom for neutron-rich systems above $N = 82$. Various shell-model interactions have been developed for this region, based on an inert ^{132}Sn core, including empirical Hamiltonians (like SMPN) [15] as well as interactions derived from realistic free nucleon–nucleon potentials such as CD-Bonn, renormalized either via G -matrix approaches like CWG [15, 159] or through $V_{\text{low-}k}$ techniques by Covello *et al.* [160].

The SMPN calculations, developed by Sarkar and Sakar (see Ref. [15, 16, 161] and citations therein) are performed in the $z50n82$ valence space above an inert ^{132}Sn core, with protons in the $\pi(1g_{7/2}, 2d_{5/2}, 2d_{3/2}, 3s_{1/2}, 1h_{11/2})$ and neutrons in the $\nu(1h_{9/2}, 2f_{7/2}, 2f_{5/2}, 3p_{3/2},$

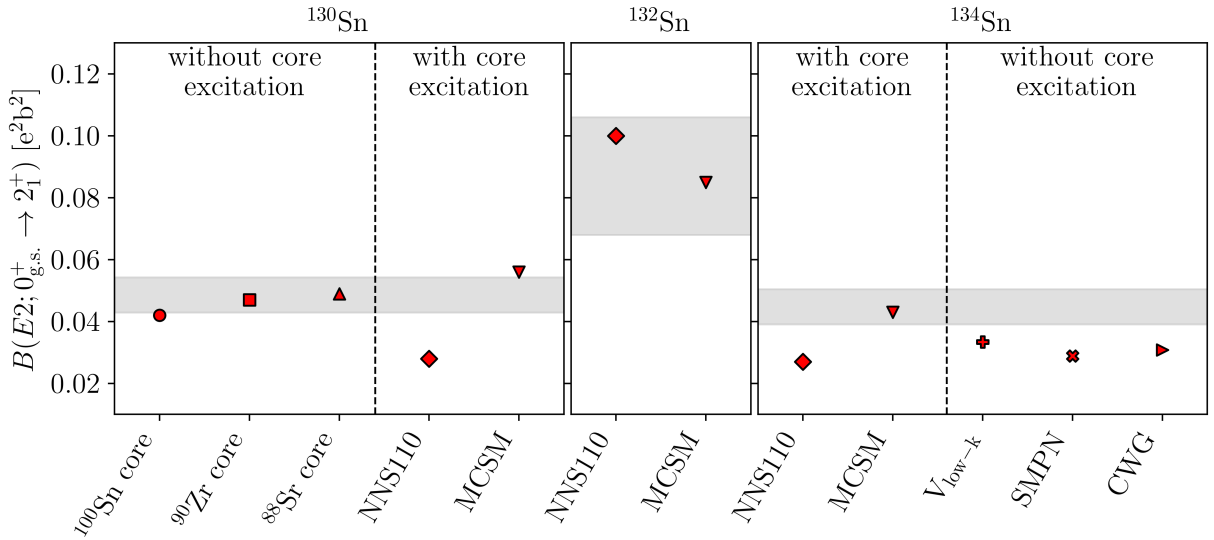


Figure 8.1. Results of various shell-model (SM) calculations (red) compared with the experimental $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values obtained in this work for $^{130,134}\text{Sn}$ and by Rosiak *et al.* [23] for ^{132}Sn (gray bands). MCSM results are from Ref. [54]; LSSM calculations based on ^{110}Zr (NNS110), ^{100}Sn , ^{90}Zr and ^{88}Sr cores from Refs. [69, 86, 155]. The latter is based on Ref. [157], values adopted from Ref. [86]. For ^{134}Sn , results from empirical SMPN [15], realistic CD-Bonn, CWG [159], and $V_{\text{low-}k}$ [160] interactions are shown, including models based solely on two-body effective interactions [161].

$3p_{1/2}, 1i_{13/2}$) orbitals. The interaction is obtained by starting from the empirical CW5082 — which was originally developed for the region above ^{208}Pb — Hamiltonian and replacing essentially all single-particle energies by experimentally determined values, except for $\pi 3s_{1/2}$ and $\nu 1i_{13/2}$, which are taken from the established systematics used in the original SMPN construction. Out of the total 2101 two-body matrix elements, 26 TBME are modified using only the binding energies and low-lying spectra of the $A = 134$ isobars (^{134}Sn , ^{134}Sb , ^{134}Te), thereby constraining in particular the ν - ν sector relevant for neutron-only valence calculations. A key feature of SMPN is a reduced neutron-neutron pairing strength: several diagonal ν - ν $J = 0$, $T = 1$ matrix elements are further weakened (about a factor of two compared to CW5082) by adjustment to reproduce the binding energy of ^{134}Sn , reflecting diminished nn pairing in the neutron-rich nuclei above $N = 82$.

The CWG calculation was also performed by Sarkar and Sarkar, and the results presented here are taken from Ref. [15]. The interaction from Brown *et al.* [159] was applied here, with ^{132}Sn as a closed core. Notably, the interaction is applicable above or below the ^{132}Sn shell closure, by using the same proton but different neutron valence spaces. For $N \geq 82$, the calculations assume an inert ^{132}Sn core, with protons in $\pi(0g_{7/2}, 1d_{5/2}, 1d_{3/2}, 2s_{1/2}, 0h_{11/2})$ and neutrons in $\nu(0h_{9/2}, 1f_{7/2}, 1f_{5/2}, 2p_{3/2}, 2p_{1/2}, 0i_{13/2})$. The residual interaction is derived from a CD-Bonn G matrix and mapped onto the valence space using the $\hat{Q}$ -box/folded-diagram approach, including all non-folded diagrams through third order in G and summing folded diagrams to infinite order. Since only neutrons are active in the present valence space, the relevant part is the nn sector; a global rescaling of the nn -renormalized G -matrix by a factor of 0.90 was noted to slightly improve the ^{130}Sn spectrum. For the latter two calculations, effective neutron charges of $e_n = 0.64e$ was used [15].

The shell-model interaction employed by Covello *et al.* [160] is derived from the realistic CD-Bonn nucleon-nucleon potential, renormalized to a low-momentum interaction $V_{\text{low-}k}$ (with cutoff $\Lambda = 2.2 \text{ fm}^{-1}$). An inert ^{132}Sn core is assumed; for neutron-rich Sn isotopes above $N = 82$ only neutrons are active in the 82–126 valence space $\nu(0h_{9/2}, 1f_{7/2}, 1f_{5/2}, 2p_{3/2}, 2p_{1/2}, 0i_{13/2})$. The effective interaction V_{eff} is obtained via the $\hat{Q}$ -box/folded-diagram expansion, including diagrams up to second order in $V_{\text{low-}k}$ and summing folded contributions following the standard procedure. The perturbative evaluation is carried out in a harmonic-oscillator basis with intermediate-state excitations taken up to five shells above the Fermi surface. For $E2$ observables, their calculations use an effective neutron charge of $e_n = 0.7e$ [160].

All models, that are excluding excitation above the ^{132}Sn shell closure, reproduce the new experimental $B(E2)$ value in ^{130}Sn reasonably well. The LSSM calculations of Banu, employing either a ^{100}Sn or a ^{90}Zr core and a valence space extending up to ^{132}Sn , show good agreement with the experimental value, while employing a ^{90}Zr core enhances the model space and the $B(E2)$ value. Similarly, the SR88MHJM interaction shows a good reproduction of the experimental result. By employing a bigger model space for the protons and neutrons, the $B(E2)$ value is described better. The largest model space of this three LSSM interaction — with the employed ^{88}Sr core — describes the collectivity best. All three calculations are in good agreement within the uncertainties of the experimental value.

The good reproduction of $B(E2)$ values in even-even nuclei next to ^{132}Sn is obtained by employing enhanced effective neutron charges (here $e_n = 1.0e$), which provides a phenomenological representation of correlations excluded from the chosen configuration space [162, 163].

In practice, these enlarged effective charges are used to compensate for missing contributions from core polarization and core excitation. Core excitations correspond to cross-shell particle–hole configurations in which nucleons are excited across a major shell gap, in this case above the $N = 82$ or $Z = 50$ shell, thereby increasing the number of quadrupole configurations and, typically, the attainable $E2$ collectivity. Core polarization denotes the induced response of the core to the motion of valence nucleons; microscopically, it arises from couplings to configurations outside the truncated valence space (e.g. collective core modes) and thus renormalizes the effective $E2$ operator acting within that space, leading to larger calculated $B(E2)$ values than obtained with bare charges. Both effects therefore increase the effective quadrupole strength beyond what is accessible in a restricted valence space, and their net impact is commonly absorbed into enlarged effective neutron charges (see e.g., Ref [162, 163]).

In the LSSM descriptions of the tin isotopic chain west of ^{132}Sn , adopting $e_n = 1.0e$ effectively accounts for the missing polarization and/or cross-shell contributions and yields a very good reproduction of the experimental $B(E2)$ strengths [69, 157].

In contrast, the calculations with the NNS110 interaction employs an extended model space beyond the doubly-magic nucleus and provide a reasonable description of $B(E2)$ in ^{132}Sn (Fig. 8.1), but they underpredict the new ^{130}Sn result by nearly a factor of two. The discrepancy can be traced back to the effective charges used: $e_n = 0.41e$ and $e_p = 1.68e$ were chosen *ad hoc* to reproduce the earlier $^{130,134}\text{Sn}$ value from the Oak Ridge experiment

and the ^{132}Sn result of Rosiak *et al.* [23, 155]. As discussed in detail by Jungclaus *et al.* [108], these effective charges differ significantly from empirical values in this region, $e_n = 0.54(1)e$ and $e_p = 1.35(5)e$. The isospin dependence of the effective charges was investigated by comparing isomers based on similar proton but different neutron configurations in ^{98}Cd and ^{130}Cd , both two protons below the corresponding doubly-magic Sn isotope. It is worth noting that Jungclaus *et al.* base their discussion on calculations performed with the same NNS110 interaction, facilitating a direct comparison of the effective-charge systematics.

Since the 2_1^+ excitation in ^{130}Sn is predominantly based on neutron excitation, an enhanced neutron effective charge increases the calculated transition strength through its quadratic contribution to $B(E2)$. For ^{132}Sn , where the wave function contains a stronger proton component, the concomitantly reduced empirical proton charge tends to compensate, thereby maintaining agreement with experiment. Interestingly, the same argument also accounts for the discrepancy in ^{134}Sn , where these LSSM calculations again underpredict the new experimental $B(E2)$ value by nearly a factor of two. While the absolute $B(E2)$ strengths are not reproduced, the relative evolution between ^{130}Sn and ^{134}Sn is captured well: the experimental ratio $B(E2; ^{130}\text{Sn})/B(E2; ^{134}\text{Sn}) = 1.08(11)$ compares nicely to the theoretical value of 1.04. This further supports that the discrepancy originates from the effective-charge prescription, not from the interaction.

The MCSM calculations employ an even bigger model space, symmetrically for protons and neutrons, to describe the whole tin isotope chain between the two doubly-magic shell closures ^{100}Sn and ^{132}Sn and beyond. Good agreement between MCSM calculations and the experimental value in $^{130,134}\text{Sn}$ is obtained. Consequently, the longstanding factor-of-two discrepancy between this calculation and the old experimental values are now resolved.

In the region west of ^{132}Sn , LSSM calculations in a valence space extending up to ^{132}Sn have been shown to reproduce the available $B(E2)$ strengths well, especially for ^{130}Sn . In these calculations, excitations across the $N = 82$ shell closure are not included; the dominant configurations therefore correspond to (at most) two neutron holes below ^{132}Sn , with no particle excitations above the closed shell. By contrast, the additional LSSM calculations considered for ^{134}Sn , which employ an inner ^{132}Sn core, fail to describe the present $B(E2)$ strength (see Fig. 8.4).

All shown interactions underestimate the measured transition probabilities. Since these calculations treat ^{132}Sn as an inert core, the structure of the 2_1^+ state in ^{134}Sn is generated solely by neutron–neutron correlations, i.e. by two-neutron excitations. The persistent

underprediction of $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$, largely independent of interaction type (empirical or realistic), together with the $\sim 40\%$ deficit in $B(E2; 4_1^+ \rightarrow 2_1^+)$ reported previously [112], indicates that one or more ingredients of the standard framework are insufficient.

As discussed for ^{130}Sn , truncated LSSM calculations are only able to reproduce the experimental findings by using enhanced effective charges. For ^{134}Sn , the effective charges adopted in SMPN and CWG calculations follow the values proposed by Sarkar and Sarkar [15, 161]. Within that framework, an effective neutron charge of $e_n \approx 0.62e$ (sometimes $0.64e$) was used for $N = 84$ nuclei, motivated in part by the then-accepted small $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ strength in ^{136}Te , two-protons and two-neutrons outside of ^{132}Sn . In Ref. [164] this value was reported as $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.103(15) e^2 \text{b}^2$. This value was later revised upward in Ref. [98] to $0.122(18) e^2 \text{b}^2$ after incorrect target thicknesses had been identified in the previous analysis. Subsequently, Allmond *et al.* [165] reported a further increase to $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.181(15) e^2 \text{b}^2$ in ^{136}Te . With this revision of the ^{136}Te value, the original justification for adopting $e_n \approx 0.62e$ is correspondingly weakened.

For ^{134}Sn , where proton excitations are suppressed by the $Z = 50$ closure and the closed ^{132}Sn core in the calculation, a moderately larger neutron effective charge yields a markedly improved description using the SMPN interaction [15]. The present experimental value, $B(E2; 0^+ \rightarrow 2_1^+) = 0.0448(57) e^2 \text{b}^2$, is reproduced more closely, e.g., with $e_n \simeq 0.72e$ than with $e_n = 0.62e$ (leading to 0.0366 versus $0.0271 e^2 \text{b}^2$) [161]. The description of $B(E2; 6_1^+ \rightarrow 4_1^+)$ in ^{134}Sn is likewise improved with the enhanced neutron charge [161].

A similar observation can be made for other effective interactions by Kartamyshev *et al.* [166]. Here, reduced transition probabilities in the Sn isotope chain were calculated using an interaction closely related to CWG, with an effective neutron charge of $e_n = 0.64e$, which was scaled to the $B(E2; 6_1^+ \rightarrow 4_1^+)$ value in ^{134}Sn by Zhang *et al.* [109] and to the $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ strength determined at Oak Ridge [57–60]. In that work it was already noted that reproducing the $B(E2; 6_1^+ \rightarrow 4_1^+)$ value in ^{134}Sn would require a larger neutron effective charge. Furthermore, their calculations underestimate the isomeric $B(E2; 6_1^+ \rightarrow 4_1^+)$ strengths in $^{136,138}\text{Sn}$, yielding 0.33 and 0.27 W.u., respectively, compared to the experimental values of $0.57(9)$ and $0.36(8)$ W.u., which have been determined in Ref. [167]. The calculations of Ref. [166] are not shown in Fig. 8.4, since their $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ results in ^{134}Sn is indistinguishable from the results from the CWG calculations, as they use a similar interaction, and would not add further information in this representation.

Covello *et al.* adopt an effective neutron charge of $e_n = 0.7e$, guided by the available $E2$

transition data in ^{134}Sn , in particular $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ and $B(E2; 6_1^+ \rightarrow 4_1^+)$. Within a single, state-independent choice of e_n , however, both transitions cannot be reproduced simultaneously: the value $e_n = 0.7e$ represents a reasonable compromise, but it tends to overestimate the previously reported $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ strength while underestimating the $B(E2; 6_1^+ \rightarrow 4_1^+)$ transition.

This all supports the conclusion that the adopted effective neutron charges are too small. In particular, the new $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ value for ^{134}Sn determined in this work points to a larger effective neutron charge than previously used, providing an important constraint for theory. At the same time, the systematic underestimation of the $6_1^+ \rightarrow 4_1^+$ decays of predominantly neutron isomeric 6_1^+ states in the heavier isotopes is consistent with the need for an increased effective neutron charge. While $e_n = 0.64e$ can accommodate older results [57–60], enhanced effective neutron charge values provide a better overall description of the $E2$ transition strengths associated with the isomeric decays, e.g., the $6_1^+ \rightarrow 4_1^+$ transitions in neutron-rich Sn nuclei, as well as the present $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ data.

The systematic need for an enhanced neutron effective charge east of ^{132}Sn points to core-polarization contributions, which appear to be underestimated in the calculations employing $e_n \approx 0.62\text{--}0.70e$. This is consistent with shell-model descriptions west of ^{132}Sn , where the observed $B(E2)$ systematics can be reproduced only by adopting substantially increased neutron charges, up to $e_n = 1.0e$. In LSSM calculations restricted to the ^{132}Sn , the $E2$ strength in ^{130}Sn is dominated by two-neutron hole excitations, whereas in ^{134}Sn the collectivity can arise only from two-neutron particle excitations above $N = 82$.

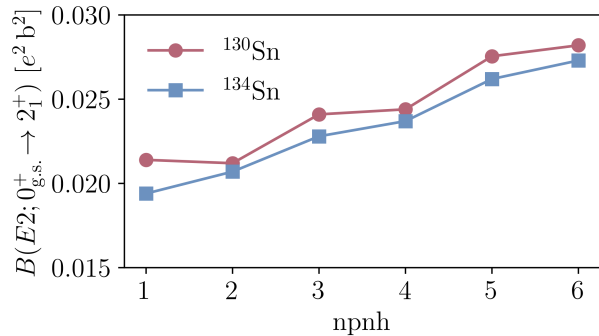


Figure 8.2. Calculated $B(E2)$ strength in ^{130}Sn (red) and ^{134}Sn (blue) employing the NNS110 interaction and effective charges of $e_p = 1.68e$ and $e_n = 0.41e$ truncated to different nph excitations. Data points extracted and taken from Ref. [24].

Taken together, the deviations from the present $B(E2)$ values suggest that the calculated quadrupole collectivity remains too weak, indicating that a larger neutron effective charge is required to reproduce the data within a fixed interaction.

Beside effective charges, the new data also elucidates the role of cross-shell excitations.

Figure 8.2 summarizes NNS110 calculations for $^{130,134}\text{Sn}$ in which the number of particle–hole excitations is varied, up to a truncation of 6p6h. Increasing the allowed excitations from 2p2h (which is the minimal number of core excitation expected when moving two neutrons away from ^{132}Sn) to 6p6h enhances the predicted $B(E2)$ values by 32 % and 40 % for ^{130}Sn and ^{134}Sn , respectively. This indicates that core excitations associated with breaking the ^{132}Sn core can account for more than one quarter of the total quadrupole strength in these nuclei. Given ^{132}Sn as a particularly strong closed-shell system, this is a large effect and provides a natural explanation for why models that treat ^{132}Sn as an inert core tend to underestimate the $E2$ collectivity of the first excited 2^+ and 4^+ states in ^{134}Sn . This also gives an explanation why only the MCSM calculation with valence spaces going below and above ^{132}Sn is able to reproduce the $B(E2)$ value in ^{134}Sn .

Notably, the MCSM calculations that show the best overall agreement with the available $B(E2)$ systematics adopt an effective neutron charge of $e_n = 0.75e$ for all their performed calculations. A. Junglaus *et al.* pointed out that the isospin dependence of the effective charges manifests itself in the present mass region [108]. In this context, the new experimental results from below and above the doubly-magic ^{132}Sn shell closure, establish a test for the MCSM calculations but also for LSSM calculations and a sensitive test of isospin-dependent effective-charge prescriptions.

Indications for missing core-polarization contributions around the doubly-magic ^{132}Sn core have been discussed not only for neutron-rich systems, but also in other nuclei like the two-proton neighbor ^{134}Te . There, the enhanced $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ strength has been attributed to proton excitations coupled to vibrational modes of the ^{132}Sn core [168], an effect that predominantly influences the properties of the 2^+ state. In contrast, the experimental results for the neutron-rich system ^{134}Sn show marked discrepancies with theory for both the 2_1^+ and 4_1^+ states. This points to a more general shortcoming in the present description of neutron-dominated quadrupole collectivity above the $N = 82$ shell closure and, as discussed above, is naturally interpreted as the consequence of missing core-polarization contributions associated with cross-shell excitations across the $N = 82$ gap.

8.1.1. Evolution of Collectivity along the Sn Isotopic Chain

Figure 8.3 places the new experimental $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values obtained in this work into the broader context of quadrupole collectivity along the tin isotopic chain. The compilation includes experimental results for even–even Sn nuclei spanning $N = 54$ to $N = 84$, together with several theoretical approaches frequently used to describe this region, namely LSSM

calculations employing empirical interactions with either an ^{100}Sn , ^{90}Zr or ^{88}Sr core (denoted “100Sn core”, “90Zr core” and “88Sr core”), and the MCSM results. The experimental values shown are taken from Refs. [23, 70–85, 169]. Where multiple measurements exist for a given nucleus, the gray band highlights the corresponding weighted mean, thereby providing a compact visual guide to the current experimental database and its spread. Only values published in this century in peer-reviewed journals are shown, together with the results of this thesis.

A first general observation is that LSSM descriptions anchored to a fixed core and limited to neutron excitation using a model space up to ^{132}Sn do not provide a global description of the full isotopic chain. While these interactions can achieve good agreement in restricted mass regions — in particular, close to the ^{132}Sn shell closure — their systematics follow

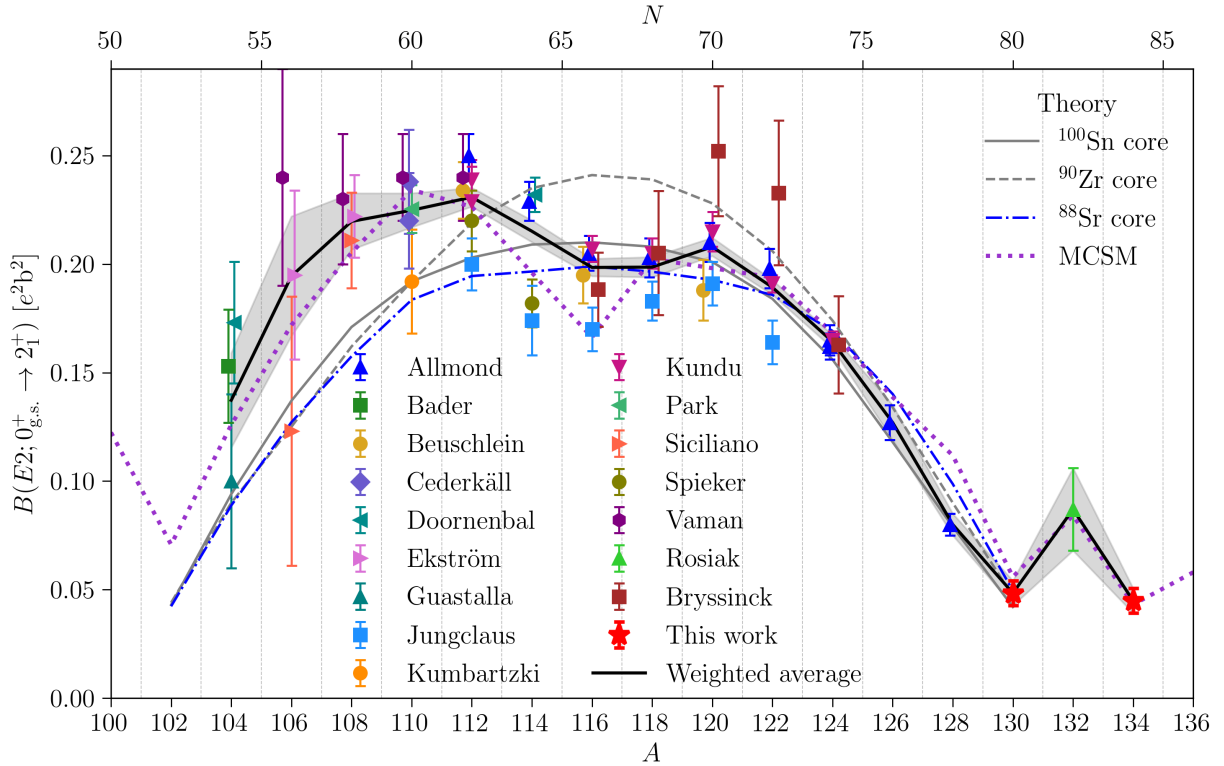


Figure 8.3. Systematics of reduced quadrupole transition strengths $B(E2; 0_{g.s.}^+ \rightarrow 2_1^+)$ along the even-even Sn isotopic chain from $N = 54$ to $N = 84$. The new experimental results from this work are shown together with selected literature data and representative theoretical descriptions, including LSSM calculations employing effective interactions based on ^{100}Sn , ^{90}Zr and ^{88}Sr cores (“100Sn core”, “90Zr core” and “88Sr core”), and MCSM results. The gray band indicates the weighted mean of the available experimental values in regions where multiple measurements exist (see text).

an approximately parabolic trend that is characteristic of a seniority-dominated evolution within a limited neutron valence space. Such a behavior is qualitatively expected for $E2$ transitions governed predominantly by the filling of a single neutron orbital (or a small set of orbitals) between major shell closures. Experimentally, however, the collectivity pattern is more complex: the pronounced enhancement of $B(E2)$ strengths in lighter Sn isotopes is not captured by these LSSM calculations, as well as the flat, or decreasing trend in the mid-shell nuclei. Furthermore, they are limited to nuclei above ^{100}Sn and below ^{132}Sn .

The enhanced collectivity for the lighter isotopes is reproduced by the MCSM calculations. In the MCSM description, the increase of $B(E2)$ strengths towards neutron-deficient tin is associated with substantial proton excitations across the $Z = 50$ shell gap. In particular, proton cross-shell configurations involving the $1g_{9/2} \rightarrow 2d_{5/2}$ excitation become increasingly important for $N \approx 54\text{--}64$, with proton-excitation admixtures reaching the level of $\gtrsim 30\%$ in that region. This is discussed in detail in Ref. [54]. The resulting strengthened proton–neutron quadrupole correlations enhance deformed components of the wave functions and generate the characteristic “bump” in $B(E2)$ strengths. In terms of a simple occupancy picture, the number of proton holes in the $1g_{9/2}$ orbital reaches a maximum (of order ~ 0.8 for the 2_1^+ state) around ^{110}Sn and then decreases towards heavier isotopes, becoming effectively 0 in $^{130,134}\text{Sn}$. This naturally explains why, in the vicinity of ^{132}Sn , the MCSM systematics approach a more seniority-like behavior similar to that produced by restricted-valence LSSM calculations.

Around mid-shell, the experimental situation is less settled. Reduced $B(E2)$ values in $^{114,116}\text{Sn}$ were reported by Jungclaus *et al.* [78]; however, other measurements do not show a comparably strong reduction, as illustrated by the results of Beuschlein *et al.* [73], Kundu *et al.* [81], and Allmond *et al.* [71]. The spread of experimental values in this mass region limits the extent to which a single theoretical description can be validated unambiguously, and it remains an open issue whether the apparent deviations reflect genuine structural effects or unresolved systematic differences between measurements based on various experimental methods.

In the heavy-mass region approaching the doubly magic nucleus ^{132}Sn and beyond, the overall collectivity decreases, as expected for nuclei approaching a major shell closure and for excitations increasingly dominated by neutron configurations above the $N = 82$ gap. In this region, the new results from the present work provide an important anchor point: they extend the set of experimentally constrained quadrupole strengths to the neutron-rich side

of ^{132}Sn , where theoretical predictions begin to diverge significantly and where additional effects such as core polarization cross-shell admixtures, and the choice of effective charges can play an important role. Consequently, the measurements reported here represent a benchmark for models aiming at extrapolations further towards the neutron-rich tin isotopes relevant to shell evolution and (indirectly) to astrophysical applications.

Overall, Fig. 8.3 underlines that a quantitatively reliable, unified description of $B(E2)$ systematics along the entire Sn chain requires (i) an explicit treatment of proton cross-shell excitations in the lighter isotopes, (ii) a consistent handling of correlations and effective charges in the vicinity of shell closures, and (iii) additional experimental constraints—especially in regions where present data sets show significant scatter. The new experimental information provided by this work contributes directly to this program by tightening the constraints on collectivity on the neutron-rich side of ^{132}Sn . The result of this work highlights the excellent agreement of the MCSM calculations with the new experimental findings.

The deviations obtained with the NNS110 interaction are primarily driven by the adopted effective charges. In LSSM calculations employing an inert ^{132}Sn core, the underestimated $B(E2)$ strength in ^{134}Sn instead indicates missing core-polarization contributions associated with cross-shell excitations across the $N = 82$ gap. The new experimental value therefore provides a stringent test for theory and a valuable constraint for refining effective-charge prescriptions in this mass region.

8.1.2. The $N = 90$ Shell Closure

Towards larger neutron number in the tin isotopic chain, several calculations based on empirical shell-model interactions predict the emergence of a new subshell closure at $N = 90$, commonly associated with the filling of the $\nu 2f_{7/2}$ orbital [16]. In contrast, calculations employing realistic two-body interactions typically do not exhibit a pronounced closure at this neutron number [159, 160]. Notably, recent *ab initio* studies report signs consistent with increased shell effects around $N = 90$, such as an enhanced behavior of two-neutron separation energies S_{2n} [14]. A similar observation has been made when incorporating an effective three-body monopole contribution into a realistic interaction (CWG3M), which leads to a markedly altered shell evolution [16].

The question whether a subshell closure emerges at $N = 90$ is not only of structural interest. Shell effects in very neutron-rich nuclei influence nuclear masses and β -decay properties,

and thereby affect calculated r -process abundance patterns. In particular, the persistence or quenching of the $N = 82$ shell gap for lower mass isotones already has well-established consequences for r -process modeling [23]. An additional structural change at $N = 90$ could have a similarly strong impact on predicted abundances in the relevant mass region [170]. Direct experimental access to nuclei at and beyond $N = 90$ remains extremely challenging — a total of five ^{140}Sn nuclei were observed in the only experimental report including such exotic nuclei in 2018 [171]. Consequently, theoretical models must be constrained by complementary observables in experimentally accessible nuclei just above ^{132}Sn .

In this context, the present work provides the first experimental results of collective quadrupole properties east of the ^{132}Sn shell closure through the $E2$ strength in ^{134}Sn . While it was shown, that realistic two-body interactions reproduce the essentially single-particle character of low-lying states in ^{135}Sn without invoking explicit three-body terms [78], transition strengths probe different aspects of the effective interaction, including correlations and core polarization. The fact that the measured $B(E2)$ value in ^{134}Sn is not reproduced by the available theoretical approaches already at $N = 84$ (i.e. with only two valence neutrons outside the ^{132}Sn core) limits the possibility to discriminate between empirical and realistic interactions at present.

Plausible sources for the discrepancy include neutron effective charge, missing configuration mixing such as cross-shell excitations, and an incomplete treatment of core-polarization effects as discussed in the previous chapter. This situation emphasizes that a quantitatively reliable description of neutron-dominated collectivity above the $N = 82$ shell gap is a prerequisite for extrapolations towards $N = 90$.

Among the available large-scale approaches, MCSM calculations provide an overall good description of known $B(E2)$ systematics in the vicinity of ^{132}Sn . However, due to the massive valence space and associated computational cost, current MCSM results are restricted to nuclei with $N \leq 90$ and to a limited set of low-lying states. Moreover, despite the impressive agreement for $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ strengths, deviations have been reported for other transitions (e.g. $B(E2; 2_1^+ \rightarrow 4_1^+)$ [86]) in mid-shell tin nuclei, and seniority- $v = 2$ isomers in the $N > 82$ region cannot presently be reproduced within the published MCSM datasets. For these reasons, the existence of a putative $N = 90$ subshell closure cannot yet be stringently predicted by MCSM.

To assess the trends of standard shell-model interactions in the experimentally accessible region, Fig. 8.4 compares calculated and experimental systematics for ^{134}Sn , ^{136}Sn , and

^{138}Sn . The left panel of Fig. 8.4 shows the 2_1^+ excitation energies, the middle panel the calculated $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ values (including the experimental point for ^{134}Sn), and the right panel the predicted $\nu(2f_{7/2}^n)$ contribution to the 2_1^+ wave function for the discussed theoretical approaches. Here $n = 2, 4, 6$ corresponds to the number of valence neutrons beyond the $N = 82$ closed shell for $A = 134, 136, 138$, respectively. The values for SMPN and CWG interactions are taken from Ref. [15], the MCSM from Ref. [54] and CWG3M from Ref. [16]. Experimental values are included where available.

For the 2_1^+ excitation energies $E(2_1^+)$, the CWG interaction reproduces the well-known near constancy of excitation energies along the even–even Sn chain beyond $N = 82$: for $N \geq 84$ the calculated values, remain approximately flat, in qualitative agreement with the experiment. Quantitatively, CWG exhibits a systematic positive offset of order ~ 50 keV for the nuclei considered here. In contrast, the SMPN calculation show a pronounced decrease of $E(2_1^+)$ with increasing neutron number, resulting in a clear disagreement with the experimental trend in ^{136}Sn and ^{138}Sn . These distinct energy systematics correlate with differences in the predicted microscopic structure.

As shown in the right panel of Fig. 8.4, the $\nu(2f_{7/2}^n)$ component remains comparatively large and only weakly varying in SMPN, whereas it decreases substantially with increasing mass number in the CWG calculation. Higher-lying orbitals, thus, have enhanced influence on the 2_1^+ state in the CWG interaction compared to the SMPN, where the 2_1^+ is dominated by the $\nu f_{7/2}$ orbital.

This contrasting structural evolution is also reflected in the predicted quadrupole collectivity. In a seniority-dominated single- j scenario, enhanced $E2$ strength is expected around mid-shell. The SMPN predicts a clear maximum of $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ in ^{136}Sn compared to $^{134,138}\text{Sn}$ as shown in the mid-panel of Fig. 8.4. In contrast, CWG yields a monotonic evolution of $B(E2)$ along the isotopic chain, following the decreasing $\nu(2f_{7/2}^n)$ component implying a progressively smaller contribution of the $2f_{7/2}$ orbital to the 2_1^+ excitation. This again highlights a key difference between the interactions: in CWG the $\nu(2f_{7/2})$ and $\nu(1h_{9/2})$ single-particle orbitals lie close in energy, promoting a concurrent filling of both configurations. In contrast, SMPN predicts a larger separation between these orbitals, resulting in a predominantly sequential filling in which the $\nu(2f_{7/2})$ orbital is filled first before significant occupation of $\nu(1h_{9/2})$ sets in.

Quantitatively, the predicted strengths diverge significantly: for ^{136}Sn , CWG gives $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+) = 3.00$ W.u. and SMPN 2.16 W.u., while for ^{138}Sn the discrepancy

increases to 4.33 W.u. (CWG) versus 1.60 W.u. (SMPN). At present, experimental information on $B(E2)$ strengths in ^{136}Sn and ^{138}Sn is not available, preventing a decisive validation of either predicted trend. For ^{134}Sn , where a measured value exists, neither interaction reproduces the observed strength within the present modeling. In these calculations, an effective neutron charge of $e_n = 0.64e$ was employed; the remaining discrepancy suggests that a refined treatment of effective charges and/or core-polarization effects, and possibly additional configuration mixing, may be required for quantitative agreement.

Finally, the potential role of three-body contributions is highlighted by the work of Sarkar and Sarkar [16], who demonstrated that augmenting the CWG Hamiltonian by a three-body monopole term (CWG3M) drives the predicted systematics towards the empirical SMPN-like behavior. Since this modification affects only nuclei with at least three valence neutrons outside the ^{132}Sn core, its impact emerges at ^{136}Sn . For ^{136}Sn and ^{138}Sn , CWG3M predicts $E(2_1^+) = 639\text{ keV}$ and 633 keV , respectively. This improves the excitation-energy systematics compared to the two-body CWG interaction, albeit with a deviation of comparable magnitude and opposite sign, i.e. an underestimation rather than the $\sim 50\text{ keV}$ overestimation observed for CWG in this mass region. At the same time, the inclusion of the three-body term produces a pronounced modification further from stability: for ^{140}Sn , it predicts $E(2_1^+) \approx 1.9\text{ MeV}$, compared to $\approx 750\text{ keV}$ without three-body contributions. While spectroscopy of ^{140}Sn remains experimentally demanding, discriminating signatures

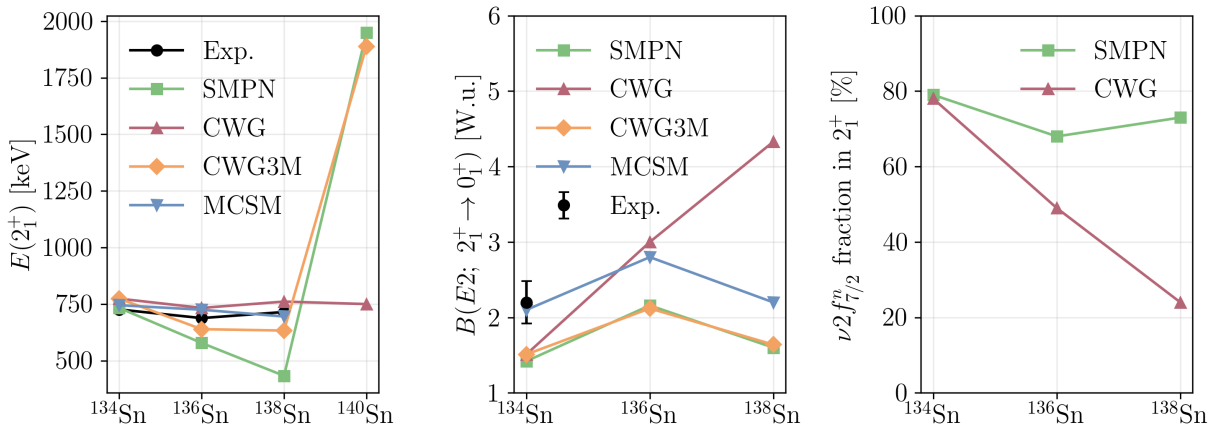


Figure 8.4. Systematics of low-lying quadrupole properties in the even-even isotopes ^{134}Sn , ^{136}Sn , and ^{138}Sn . Left: 2_1^+ excitation energies compared to experiment. Middle: calculated $B(E2; 2_1^+ \rightarrow 0_{g.s.}^+)$ transition strengths, including the experimental value for ^{134}Sn . Right: predicted $\nu(2f_{7/2}^n)$ contribution to the 2_1^+ wave function, with $n = 2, 4, 6$ for $A = 134, 136, 138$, respectively. CWG and SMPN denote the two shell-model interactions used in this work.

of three-body effects may already be accessible at ^{136}Sn through transition strengths. In particular, CWG3M yields a $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ value close to the SMPN result shown in Fig. 8.4. A stringent test therefore appears feasible by (i) constraining effective charges using the present work together with $E2$ data on long-lived states along the Sn chain (Refs. [112, 167]) and (ii) measuring $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ in ^{136}Sn . Such a program would directly test the predictive power of empirical interactions, realistic two-body Hamiltonians, and effective three-body-modified interactions in the mass region that anchors extrapolations towards the proposed $N = 90$ subshell closure. Notably, the calculated $B(E2; 2_1^+ \rightarrow 0_1^+)$ systematics obtained with the MCSM exhibit essentially the same overall trend as those predicted by the SMPN interaction and by the CWG interaction when augmented by effective three-body contributions (see Fig. 8.4).

The proposed subshell closure at $N = 90$ in the tin isotopes provides a particularly instructive analogy to the calcium isotopic chain. In Ca, the $\nu 1f_{7/2}$ orbital is filled between $N = 20$ and $N = 28$, and realistic interactions historically failed to reproduce the experimentally established $N = 28$ shell closure until three-nucleon forces were included [172]. Beside this behavior similarity, the systematics of the first 2_1^+ excitation energies show a parallel in the two regions. Following the doubly magic ^{40}Ca core, $E(2_1^+)$ energies decrease towards a minimum near ^{44}Ca , i.e. four neutrons above the closed shell and close to mid filling of the $\nu 1f_{7/2}$ orbital, before rising again towards ^{48}Ca . A closely related behavior is observed experimentally beyond ^{132}Sn : the evolution of $E(2_1^+)$ between ^{134}Sn and ^{138}Sn is nearly flat and exhibits a tendency towards lower excitation energy near mid filling of the $\nu 2f_{7/2}$ orbital. This parallelism reinforces the interpretation that the interplay of pairing, quadrupole correlations, and monopole components of the effective interaction within an $f_{7/2}$ -dominated valence space governs the low-lying quadrupole response in both cases. From this perspective, the strong sensitivity of the predicted $E(2_1^+)$ systematics in tin isotopes to effective three-body monopole terms is not unexpected. It may rather mirror the mechanism that proved essential to describe shell evolution in the calcium chain.

8.2. Quadrupole Moments

The spectroscopic quadrupole moments $Q_s(2_1^+)$ extracted in this work are compared with theoretical predictions and with available experimental information along the tin isotopic chain in Fig. 8.5. Quadrupole moments probe the static quadrupole deformation associated with the 2_1^+ state and therefore provide information complementary to transition strengths: while $B(E2)$ values quantify the collectivity of transitions, $Q_s(2_1^+)$ constrain the shape

character and the sign of the quadrupole deformation implied by the wave function of the state itself. In a spherical vibrational limit, one expects small quadrupole moments, whereas sizeable non-zero values point to static deformation or strong configuration mixing. For nuclei close to shell closures, $Q_s(2_1^+)$ is therefore a particularly sensitive observable for testing the balance between single-particle and collective components in theoretical descriptions.

On the experimental side, $Q_s(2_1^+)$ values are available in the stable mid-shell isotopes $^{112-124}\text{Sn}$ [71] and in the neutron-rich isotopes $^{124-128}\text{Sn}$ [70]. More recently, Park *et al.* reported the quadrupole moment of the 2_1^+ state in ^{110}Sn [86], which is the largest value determined so far along the Sn isotopic chain and differs from zero at the 2σ level. The new results presented here extend such measurements into the most exotic region yet accessed for quadrupole moments in the vicinity of a doubly magic nucleus.

Within the present uncertainties, the extracted values are consistent with zero, underscoring the pronounced spherical character of ^{132}Sn and supporting an interpretation in which the low-lying 2_1^+ excitation in this region remains dominated by weakly collective neutron

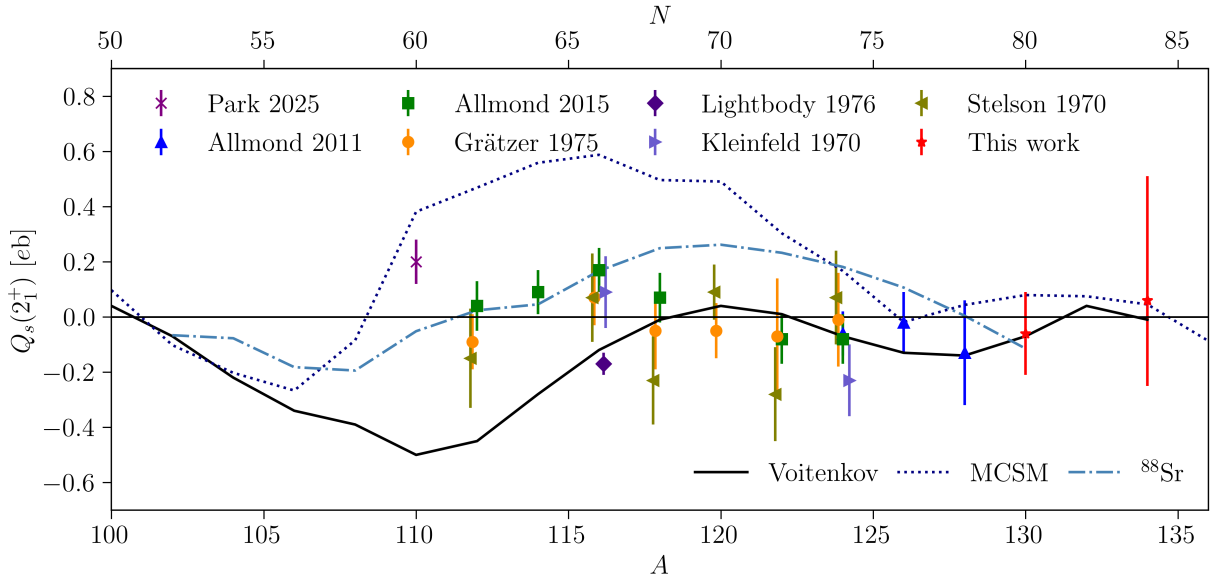


Figure 8.5. Systematics of spectroscopic quadrupole moments $Q_s(2_1^+)$ in even-even tin isotopes. The new values from this work are compared to existing experimental data for lighter and less neutron-rich Sn nuclei as well as to representative theoretical predictions, including MCSM calculations [54], QRPA results based on the Fayans energy-density functional (Voitenkov) [173, 174] and realistic LSSM calculations built on a ^{88}Sr core [86, 157]. Experimental values are taken from Refs. [70, 71, 86, 175–178]

configurations above the $N = 82$ shell gap, with limited admixture of deformed correlations. Several theoretical approaches are shown in Fig. 8.5 for comparison. The MCSM results are taken from Togashi *et al.* [54]. Self-consistent QRPA calculations based on the Fayans energy-density functional were performed by Voitenkov [173, 174]. In addition, realistic LSSM calculations built on a ^{88}Sr core are included; the interaction is described in Ref. [157], and the numerical values are taken from Ref. [86].

Overall, the comparison shows that none of the available models provides a quantitatively consistent description of $Q_s(2_1^+)$ over the isotopic chain. The MCSM predictions agree with the present data within uncertainties for ^{130}Sn and ^{134}Sn , although the sign of the quadrupole moment in ^{130}Sn is not reproduced. Moreover, in the mid-shell region, the MCSM tends to overestimate the magnitude of $Q_s(2_1^+)$, which has been attributed to a strong neutron quadrupole contribution in that mass range. The neutron-rich isotopes are described quite well by the Fayans-QRPA results of Voitenkov, but the agreement of this approach deteriorates rapidly towards lighter Sn nuclei, in some cases predicting even the wrong sign and unrealistically large values. The realistic LSSM calculation exhibits smaller deviations overall, yet it fails to reproduce the detailed systematic trend and does not capture the observed evolution consistently across the isotopic chain. Interestingly, in the light isotope ^{110}Sn , the MCSM gives a reasonable description of the measured quadrupole moment, consistent with its success in accounting for enhanced proton cross-shell excitations that cause collectivity in the neutron-deficient region. For the nuclei studied in this work, all calculations reproduce the small experimental values within the quoted uncertainties, supporting a predominantly spherical structure next to the doubly-magic system.

9. Summary

In summary, three Coulomb-excitation experiments were carried out with the Miniball spectrometer at HIE-ISOLDE between 2022 and 2025 to determine precise $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values in ^{130}Sn and ^{134}Sn . The main goal of this work was to resolve the discrepancy between theoretical approaches and previously reported experimental values published in conference proceedings, which differ strongly from state-of-the-art shell-model calculations, and to obtain first reliable $E2$ strength information of first excited states east of ^{132}Sn .

Radioactive $^{130,134}\text{Sn}$ beams were accelerated to 4.4 MeV per nucleon and impinged on ^{206}Pb and ^{194}Pt targets under safe Coulomb-excitation conditions. Recoiling nuclei were detected using a DSSSD detector at target-to-CD distances between 24.5 and 33.8 mm. γ rays detected in coincidence with the scattered particles were measured with the Miniball spectrometer. Detailed investigations of the Miniball detector angles and the recoil velocities enabled a precise Doppler correction.

The first experiment, conducted in 2022 with a ^{130}Sn beam, was the first Coulomb-excitation experiment with Miniball at ISOLDE after LS2 following the refurbishment of the spectrometer. The beam composition was analyzed using a single HPGe detector upstream of the acceleration stage, yielding a fraction of 61(7) % $^{130}\text{Sn}_{\text{g.s.}}$ in the beam. The experiment resulted in a value of $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.054_{-14}^{+10} (\text{stat.}) (7)_{\text{syst.}} e^2 \text{b}^2$ in ^{130}Sn .

To improve the precision of this value, a follow-up experiment was performed in 2023. The beam composition was determined by analyzing the β -delayed γ -ray spectra detected with Miniball throughout the experiment. A dominant component of 66(2) % $^{130}\text{Sn}_{\text{g.s.}}$ was observed, while the remaining 34(2) % corresponded to the isomeric $^{130}\text{Sn}_{7-}$ component. In total, 14713(193) counts were observed in the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition. This statistics allowed us to subdivide the data into 16 rings, providing sensitivity to the spectroscopic quadrupole moment. The experiment yielded $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.0486(57) e^2 \text{b}^2$ and $Q_s(2_1^+) = -0.06(15) \text{eb}$. In addition, Coulomb excitation of the isomeric $^{130}\text{Sn}_{7-}$ beam component to the 5_1^- state was observed. By combining lifetime information with the

Coulomb-excitation data, the transition strengths decaying transitions of the 5_1^- state could be determined with high precision.

Motivated by these results, the focus was shifted towards neutron-rich ^{134}Sn beyond the ^{132}Sn shell closure. A first attempt to measure $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ in ^{134}Sn was performed in July 2024; however, the heating device used for sulfurization proved unstable under proton irradiation. A follow-up experiment was therefore approved and conducted in October 2025. Besides ^{134}Sn , the beam contained ^{134}Te , ^{134}Ba , and ^{134}Sb . The beam composition was deduced using a combination of Coulomb-excitation yields, ionization-chamber data, and β -delayed γ -ray information, resulting in 7.37(70) % of ^{134}Sn in the beam. A total of 3396(150) counts were observed in the $2_1^+ \rightarrow 0_{\text{g.s.}}^+$ transition, allowing the data to be divided into nine angular ranges. This resulted in $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+) = 0.0449(57) e^2\text{b}^2$ and $Q_s(2_1^+) = 0.06_{-31}^{+45} (\text{stat.}) (3)_{\text{syst.}} \text{eb}$.

The $B(E2; 0_{\text{g.s.}}^+ \rightarrow 2_1^+)$ values reported here provide the first precise determination of these transition strengths in ^{130}Sn and ^{134}Sn , and the first such information for a tin isotope beyond the ^{132}Sn shell closure besides the value published in conference proceedings. The new results are not consistent with the values reported in these proceedings.

West of ^{132}Sn , the experimental $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ values are generally well reproduced by shell-model calculations in a valence space restricted to orbitals below ^{132}Sn and without explicit core excitations, provided comparatively large effective neutron charges of $e_n \approx 1.0 e$ are employed. With this choice, not only the ^{130}Sn result obtained in this work but also the systematics in lighter $A > 120$ Sn isotopes can be described.

By contrast, LSSM calculations for ^{134}Sn that assume an inert ^{132}Sn core underestimate the newly measured $B(E2)$ strength when using the smaller effective charges typically adopted in this region, $e_n = 0.62\text{--}0.7 e$. This persistent underprediction is attributed to missing contributions from core polarization and core excitations, which are not fully accounted for by the chosen effective charges.

Calculations employing enlarged valence spaces that span orbitals below and above ^{132}Sn also tend to yield $B(E2)$ values that are below the present experimental result. In particular, the NNS110 interaction uses rather small effective charges compared to values deduced empirically in this region by Jungclaus *et al.* [108]. The MCSM calculation, on the other hand, is the only approach shown here that does not rely on experimental input from $^{130\text{--}134}\text{Sn}$. It shows good agreement for $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ in ^{134}Sn while also reproducing

the $B(E2; 2_1^+ \rightarrow 0_{\text{g.s.}}^+)$ value in ^{130}Sn .

The new $^{130,134}\text{Sn}$ results thus resolve the previously reported tension between the conference-proceedings value and the MCSM description, bringing experiment and theory into much closer agreement. In contrast, the remaining discrepancy with the NNS110 interaction suggests that renewed calculations employing the effective charges determined in Ref. [108] would be valuable.

The extracted spectroscopic quadrupole moments $Q_s(2_1^+)$ of neutron-rich tin isotopes provide a benchmark for nuclear-structure theory. The present measurements extend quadrupole-moment systematics to the most neutron-rich magic nucleus studied to date and yield values consistent with zero, establishing pronounced sphericity.

The results provide direct experimental evidence for the doubly magic character of ^{132}Sn and impose tight constraints on effective interactions and model assumptions. The present study extends precise electromagnetic-structure information to tin isotopes beyond the ^{132}Sn shell closure in the vicinity of the r -process. They thus deliver nuclear-structure input for modeling neutron-rich nuclei relevant to the formation of the second r -process abundance peak and the synthesis of heavy elements in astrophysical environments.

10. Outlook

Beside establishing new experimental input for theories beyond ^{132}Sn , the revised $B(E2)$ systematics are a first step towards the possible emergence of an $N = 90$ shell closure. Since the interactions discussed in this thesis differ in their predictions around $N = 90$ while yielding rather similar (and, in several cases, still imperfect) descriptions at ^{134}Sn , additional experimental constraints in more neutron-rich isotopes are required. In this respect, ^{136}Sn is especially promising: as shown in this work, the predicted $B(E2)$ strengths in ^{136}Sn and ^{138}Sn diverge significantly between interactions, providing a clear opportunity to test their predictive power and to constrain the role of three-body contributions. Incorporating the new experimental benchmarks to re-evaluate the effective charges will, in turn, allow clean conclusions regarding the $N = 90$ shell closure.

Therefore, the next step of the experimental program with Miniball at ISOLDE after LS3 will be investigations towards the most neutron-rich Sn isotopes. This will additionally provide the opportunity to probe the microscopic origin of the structural modification predicted for ^{136}Sn by the NNS110 interaction [108, 167], in particular its impact on pairing and related correlation effects. Since the same mechanism is expected to influence the structure of the 2_1^+ state, an accurate determination of the $B(E2)$ value for the $0_{\text{g.s.}}^+ \rightarrow 2_1^+$ transition constitutes a sensitive and discriminating observable. The LS3 upgrade and the associated increase in proton-beam energy are expected to substantially enhance the production rates of very exotic nuclei, thereby rendering systems such as ^{136}Sn experimentally accessible. Complementary opportunities exist at RIKEN [179] and SPES [180] with AGATA [129] and further access will become available with the start of operations at FAIR (GSI) [181], enabling systematic studies of collectivity and shell evolution beyond ^{132}Sn .

As a short-term outlook, the new Coulomb-excitation data on ^{134}Sn should be emphasized. In the same data set Coulomb excitation of ^{134}Ba and ^{134}Te was observed as the dominant beam composition. This is particularly valuable for ^{134}Te , where experimental information on excited-state properties is still scarce: the present data allow a more detailed and quantitative analysis of its low-lying structure. Notably, the excitation of the 2_2^+ state

was observed, for which no lifetime information is currently available and for which no transition-strength values have been reported so far. The γ -ray line from this state is shown in Fig. 10.1. Coulomb excitation of the second 2^+ state in ^{134}Te provides access to a richer set of $E2$ matrix elements than the well-studied 2_1^+ state, enabling a direct test of how strongly the two valence protons are subject of ^{132}Sn core-polarization effects [168].

In combination with γ -ray branching ratios, the resulting $B(E2)$ pattern for 2_2^+ offers a test for modern shell-model descriptions of the $N = 82$ region, including the balance between proton-configuration mixing and cross-shell core excitations [164, 182]. Moreover, if complementary information (mixing ratios and/or lifetimes) reveals a significant $M1$ component in the $2_2^+ \rightarrow 2_1^+$ decay, ^{134}Te becomes a sensitive testing ground for the onset—or suppression—of mixed-symmetry quadrupole dynamics, recently identified unambiguously in neighboring ^{132}Te via a large $B(M1; 2_2^+ \rightarrow 2_1^+)$ values [183]. Taken together with ongoing

“safe” Coulomb-excitation experiments around the $Z = 50$, $N = 82$ shell closures, these matrix elements will help to chart how quadrupole collectivity develops when moving away from doubly-magic ^{132}Sn .

While state-of-the-art MCSM calculations have achieved an excellent unified description of quadrupole collectivity along the Sn chain for first excited 2^+ states as demonstrated in this work, confronting these approaches also with second-excited 2^+ states provides a deeper and more discriminating validation of the underlying microscopic structure. In this way, improved electromagnetic constraints in the $N = 82$ region will not only sharpen our understanding of shell evolution and core polarization in the vicinity of ^{132}Sn , but will also contribute to more reliable modeling of neutron-rich nuclei and the nuclear-physics inputs needed for r -process simulations [184].

Moreover, the newly determined transition strengths in the negative-parity states of ^{130}Sn call for a dedicated comparison with modern shell-model calculations. Figure 10.2 summarizes the dominant $E3$, $E1$, and $E2$ decay branches of the 5_1^- state across $^{118-130}\text{Sn}$. Any

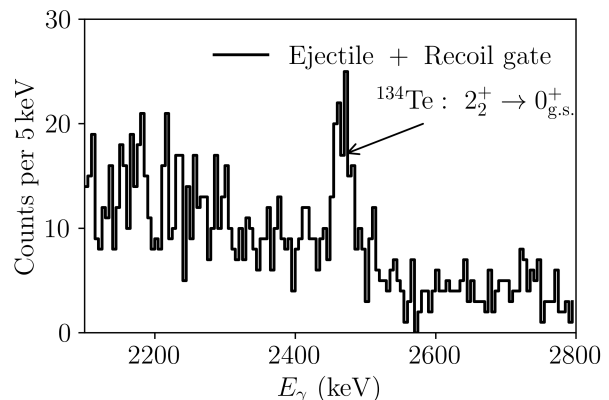


Figure 10.1. Spectra, gated on ejectiles and recoils and Doppler-corrected for ejectiles shown for high γ -ray energies above 2 MeV. The decay of the 2_2^+ into the $0_{g.s.}^+$ state is visible at an energy of 2465 keV.

interpretation of the apparent trends—in particular the weakening of the $E3$ strength between 5_1^- and 2_1^+ towards $^{128,130}\text{Sn}$ —requires a consistent theoretical description and, ideally, additional experimental constraints. A particularly promising aspect of this systematics is that the relevant states are expected to have rather clear two-quasiparticle character. The 5_1^- state is commonly associated with a $\nu(d_{3/2} \otimes h_{11/2})$ component with a possible $\nu(s_{1/2} \otimes h_{11/2})$ admixture, while the 7_1^- state is expected to be close to a pure $\nu(d_{3/2} \otimes h_{11/2})$ configuration (see, e.g., Refs. [185, 186]).

This makes the chain an attractive benchmark for SM calculations with interactions such as the Sr88MHJM interaction. If the configurations are indeed as clean as anticipated, the evolution of the $B(E3)$, $B(E2)$, and $B(E1)$ strengths should be describable within a coherent SM framework and may provide better insight into the structure and mixing of the 5_1^- and 7_1^- states than experimental trends alone. An important missing observable is the spectroscopic quadrupole moment $Q_s(5_1^-)$, which would provide a direct constraint on configuration mixing and could not be extracted reliably in the present work. In this context, an on-line perturbed angular correlation (PAC) program at ISOLDE (CERN) has been proposed to determine electromagnetic moments of short-lived excited states in neutron-rich tin isotopes [187]. The campaign aims at the chain of 5^- states up to ^{130}Sn and includes selected 7^- and 10^+ states, with the goal of testing additivity for two-quasiparticle configurations and quantifying mixing. Since a $\nu(s_{1/2})$ component does not contribute to the quadrupole moment, measured Q_s values—compared to single-hole reference states in neighboring odd-mass Sn isotopes—can be used to assess the role of a $\nu(d_{3/2})$ admixture. Complementary off-line PAC measurements have been suggested to improve precision and control systematic effects [188]. These ongoing efforts highlight the particular relevance of the 5_1^- state as a sensitive probe of two-quasiparticle structure and configuration mixing in the neutron-rich Sn isotopes.

Progress on the structure of the 5_1^- and 7_1^- states in the neutron-rich Sn isotopes will thus

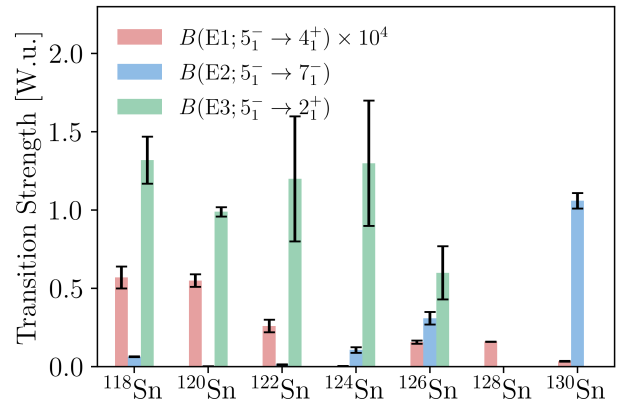


Figure 10.2. Transition strengths of the 5_1^- state to the 2_1^+ , 4_1^+ , and 7_1^- states via $E3$, $E1$, and $E2$ transitions along the Sn isotopic chain from ^{118}Sn to ^{130}Sn . The $E1$ strengths are scaled by a factor of 10^4 .

require both (i) targeted experimental information on electromagnetic moments and key transition branches and (ii) systematic SM calculations for the full isotopic chain. The combination of these two ingredients is expected to provide a consistent interpretation of the observed $E3/E2/E1$ strength evolution and to clarify the underlying configuration content towards ^{130}Sn .

Furthermore, the large uncertainties of the spectroscopic quadrupole moments along the Sn isotopic chain motivate new high-precision measurements. For ^{114}Sn , only a single measurement is available, indicating a weak oblate character within its 1σ uncertainty interval. For $^{116-122}\text{Sn}$, the evaluated experimental quadrupole moments even exhibit inconsistent signs. As an example, Allmond *et al.* report $Q_s(2_1^+) = 0.17(8) \text{ eb}$ for ^{116}Sn [71], whereas Lightbody *et al.* obtain $Q_s(2_1^+) = -0.17(4) \text{ eb}$ [176]. These discrepancies underline the need for dedicated experiments that substantially reduce the present uncertainties and establish a consistent systematics.

For the stable Sn isotopes, Coulomb-excitation measurements employing isotopically enriched Sn targets provide direct sensitivity to the spectroscopic quadrupole moments. Such experiments can be carried out at existing facilities, for example using the EAGLE setup in Warsaw [189] in combination with PIN-diode detectors or a DSSSD, as demonstrated recently in Coulomb-excitation studies of ^{110}Cd with ^{32}S and ^{14}N beams [190]. Similar measurements are also feasible at the Cologne tandem accelerator, for instance at the HORUS spectrometer in conjunction with the existing silicon particle-detector setup SONIC [191], or by combining the γ -ray array with DSSSD detectors.

The SONIC setup is currently limited to proton and α -particle detection at three fixed angles, but its energy resolution enables an event-by-event identification of target excitation energies. The addition of a DSSSD detector would extend the accessible range of scattering angles and facilitate the use of heavier projectiles. This, in turn, would allow the use of beams with precisely known transition matrix elements of low-lying states, such as ^{24}Mg , where the lifetime of the 2_1^+ state is well-known with $\tau = 1.96(5) \text{ ps}$ [68]. In this case, the target excitation could be normalized to the beam excitation, providing an additional constraint for the extraction of electromagnetic matrix elements.

An alternative approach is to rely on information from the scattered particles. By employing the LYCCA (Lund–York–Cologne CALorimeter), which combines large angular coverage and granularity with an energy resolution of $\Delta E/E < 1\%$, scattered α particles can be used to determine the angle-dependent excitation cross sections. In contrast to conventional

γ -ray based Coulomb-excitation measurements, the excitation path is directly observed, providing direct access to both transitional and diagonal matrix elements. A related strategy has recently been applied at the Q3D spectrometer by Garrett [192] using three different scattering angles. Owing to its large angular acceptance and good energy resolution, LYCCA could enable a detailed investigation of $E2$ strength distributions and spectroscopic quadrupole moments in stable nuclei, including the Sn isotopes.

Lastly, the $4_1^+ \rightarrow 2_1^+$ transition is considered in more detail. While the MCSM calculations reproduce the energies and collectivity of the first excited 2^+ states convincingly, the description of higher-lying states is less accurate. As shown on the left-hand side of Fig. 10.3, the experimental systematics of $B(E2; 4_1^+ \rightarrow 2_1^+)$ are not reproduced by the calculations. At the same time, the available experimental values carry sizable uncertainties, in particular for the stable isotopes $^{116,118}\text{Sn}$, and more exotic nuclei $^{106,108,128}\text{Sn}$. Notably, no $B(E2; 4_1^+ \rightarrow 2_1^+)$ values are reported for $^{126,130}\text{Sn}$, and the first value for ^{128}Sn became available only recently.

To address this gap in ^{126}Sn , a test experiment was already performed at the FN tandem accelerator of Cologne. A two-neutron transfer reaction was performed using an ^{18}O

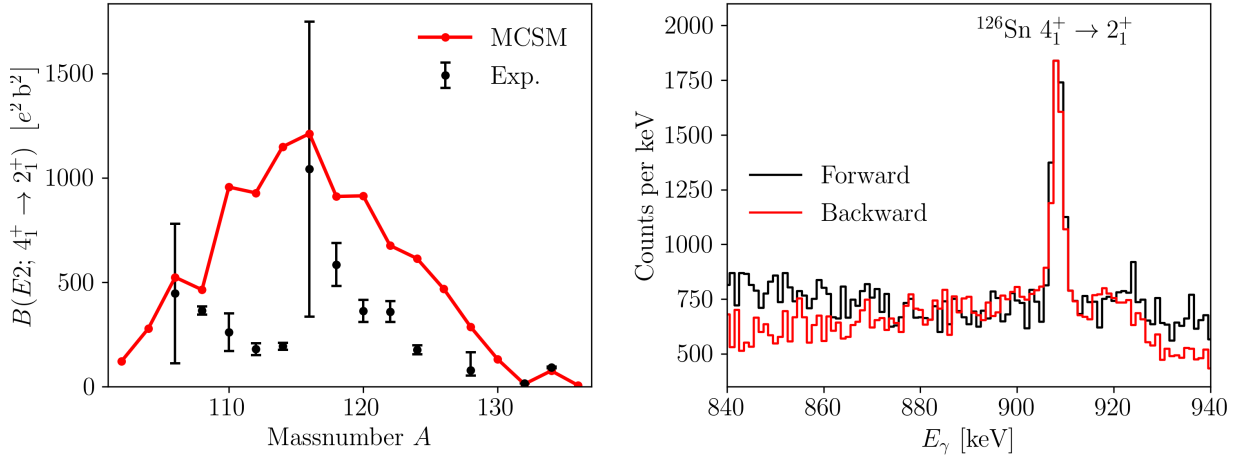


Figure 10.3. Left: Experimental $B(E2; 4_1^+ \rightarrow 2_1^+)$ values along the Sn isotopic chain compared to results from MCSM calculations. Experimental values are taken from NNDC [68], except for $^{106,108}\text{Sn}$ [82], ^{110}Sn [86], ^{128}Sn [193], and ^{134}Sn [112]. Right: γ -ray spectra measured following a two-neutron transfer reaction in the CATHEDRAL spectrometer with a gate on particles detected in the solar-cell array. Spectra are summed over six HPGe detectors at forward angles (black) and five at backward angles (red) and normalized to the 908 keV $4_1^+ \rightarrow 2_1^+$ transition observed at backward angles.

beam on a ^{124}Sn target to populate excited states in ^{126}Sn . The Cologne CATHEDRAL (Coincidence Array at the Tandem accelerator for High-Efficiency Doppler Recoil And LaBr₃(Ce) fast-timing measurements) spectrometer was operated with six HPGe detectors at forward angles and five at backward angles. Scattered oxygen isotopes were identified by solar-cell detectors at backward angles. Using the Cologne plunger device, a stopper foil was positioned at a distance of 250 μm behind the target to stop recoiling nuclei. The measurement was carried out for 8 h.

Assuming a recoil velocity of $v \simeq 0.02c$ and an upper limit of $\tau(4_1^+) \leq 25$ ps (motivated by the decreasing trend of $B(E2; 4_1^+ \rightarrow 2_1^+)$ towards heavier Sn isotopes and the value $\tau = 26(13)$ ps in ^{128}Sn), 80 % of the directly populated 4_1^+ states would decay in flight and thus emit a Doppler-shifted γ ray.

The measured spectra gated on the solar-cell detectors are shown on the right-hand side of Fig. 10.3 for forward (black) and backward (red) detector angles. No Doppler-shift-related difference is observed in the vicinity of the 908 keV $4_1^+ \rightarrow 2_1^+$ transition. This indicates that the 4_1^+ state is not populated directly but predominantly via feeding from the 5_1^- state ($T_{1/2} = 10.8(7)$ ns), i.e. after the recoil has been stopped. Consequently, Doppler-shift-based approaches are not suitable to constrain the lifetime of the 4_1^+ state with this reaction and setup.

Nevertheless, CATHEDRAL provides a perspective for lifetime measurements via fast-timing by complementing the HPGe array with LaBr₃(Ce) detectors. Recent work (e.g. Ref. [194]) demonstrates that the current experimental setup and data acquisition system can resolve lifetimes on the order of 20 ps. Therefore, employing the full configuration of CATHEDRAL (24 HPGe and eight LaBr₃(Ce) detectors) together with increased beam time should enable a determination of the presently unknown lifetime of the 4_1^+ state in ^{126}Sn .

Finally, while the results of this thesis underline the overall quality of the MCSM description along the Sn isotopic chain, key questions in the vicinity of ^{132}Sn remain open. In particular, quadrupole moments and $B(E2)$ values involving higher-lying states are not yet described consistently, and the existence of an $N = 90$ subshell closure is still under debate. At the same time, the new $B(E2)$ results presented here provide a direct experimental prove for the doubly magic character of ^{132}Sn and impose stringent constraints on effective interactions and model assumptions. The present work therefore serves as a reference point for future experimental campaigns and for further theoretical developments aimed at a quantitative understanding of quadrupole collectivity on the neutron-rich side of ^{132}Sn .

A. Derivation of Formulas for Beamcomposition

In this appendix, formulas for the determination of ^{130}Sn , ^{130}Sb and their daughter nuclei will be derived. The rate of isotopes per second in the beam is denoted by $R_{\text{isotope, state}}$. It is assumed that at $t = 0$, none of the species have accumulated in the target, i.e.,

$$N(^{130}\text{Sn}_{\text{g.s.}})(0) = N(^{130}\text{Sn}_{\text{is.}})(0) = N(^{130}\text{Sb}_{\text{g.s.}})(0) = N(^{130}\text{Sb}_{\text{is.}})(0) = 0,$$

and that the incoming rates R of all four components remain constant over time. Figure 6.18 shows a simplified decay scheme of $^{130}\text{Sn}_{\text{g.s.}}$ and $^{130}\text{Sn}_{\text{is.}}$, along with their daughter nuclei. Only γ -ray transitions with relative intensities exceeding 10 % are included. Individual β -decay branches of $^{130}\text{Sb}_{\text{g.s.}}$ and $^{130}\text{Sb}_{\text{is.}}$ are indicated, highlighting excited states that are exclusively populated by either the ground-state or isomeric decay, respectively. The number of $^{130}\text{Sn}_{\text{g.s.}}$ nuclei present at time t is then given by:

$$N(^{130}\text{Sn}_{\text{g.s.}}) = \frac{R_{^{130}\text{Sn}_{\text{g.s.}}}}{\lambda_{^{130}\text{Sn}_{\text{g.s.}}}} \left(1 - e^{-\lambda_{^{130}\text{Sn}_{\text{g.s.}}} t} \right)$$

where t is the time after the experiment started and λ is the decay constant. The same is true for $^{130}\text{Sn}_{\text{is.}}$. Since $^{130}\text{Sb}_{\text{is.}}$ is populated by $^{130}\text{Sb}_{\text{is.}}$, $^{130}\text{Sb}_{\text{is.}}$ and $R_{^{130}\text{Sb}_{\text{is.}}}$, the change of $N(^{130}\text{Sb}_{\text{is.}})$ is given as

$$\begin{aligned} \frac{dN_{^{130}\text{Sb}_{\text{is.}}}(t)}{dt} = & R_{^{130}\text{Sb}_{\text{is.}}} + R_{^{130}\text{Sn}_{\text{g.s.}}} \left(1 - e^{-\lambda_{^{130}\text{Sn}_{\text{g.s.}}} t} \right) \\ & + y R_{^{130}\text{Sn}_{\text{is.}}} \left(1 - e^{-\lambda_{^{130}\text{Sn}_{\text{is.}}} t} \right) - \lambda_{^{130}\text{Sb}_{\text{is.}}} N_{^{130}\text{Sb}_{\text{is.}}}(t). \end{aligned}$$

Here, y is the fraction of $^{130}\text{Sn}_{\text{is.}}$ which is decaying into the isomeric state of $^{130}\text{Sb}_{\text{is.}}$ and highlighted in blue in Figure 6.18. This equation is a first-order linear inhomogeneous differential equation of the general form:

$$\frac{dN(t)}{dt} + \lambda N(t) = f(t),$$

where $\lambda = \lambda_{130\text{Sb}_{\text{is.}}}$ and $f(t) = R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} (1 - e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t}) + yR_{130\text{Sn}_{\text{is.}}} (1 - e^{-\lambda_{130\text{Sn}_{\text{is.}}} t})$. First, the homogeneous part of the equation is analyzed. The homogeneous equation is:

$$\frac{dN_{130\text{Sb}_{\text{is.}}}^{(h)}(t)}{dt} + \lambda_{130\text{Sb}_{\text{is.}}} N_{130\text{Sb}_{\text{is.}}}^{(h)}(t) = 0.$$

And the solution is given as:

$$N_{130\text{Sb}_{\text{is.}}}^{(h)}(t) = C e^{-\lambda_{130\text{Sb}_{\text{is.}}} t},$$

where C is an integration constant. To find a particular solution, the form:

$$N_{130\text{Sb}_{\text{is.}}}^{(p)}(t) = A + B_1 e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t} + B_2 e^{-\lambda_{130\text{Sn}_{\text{is.}}} t}$$

is assumed. Substituting this into the original equation and differentiating:

$$\frac{dN_{130\text{Sb}_{\text{is.}}}^{(p)}(t)}{dt} = -B_1 \lambda_{130\text{Sn}_{\text{g.s.}}} e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t} - B_2 \lambda_{130\text{Sn}_{\text{is.}}} e^{-\lambda_{130\text{Sn}_{\text{is.}}} t}.$$

Substituting $N_{130\text{Sb}_{\text{is.}}}^{(p)}(t)$ and its derivative into the original equation:

$$\begin{aligned} & -B_1 \lambda_{130\text{Sn}_{\text{g.s.}}} e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t} - B_2 \lambda_{130\text{Sn}_{\text{is.}}} e^{-\lambda_{130\text{Sn}_{\text{is.}}} t} \\ & + \lambda_{130\text{Sb}_{\text{is.}}} (A + B_1 e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t} + B_2 e^{-\lambda_{130\text{Sn}_{\text{is.}}} t}) \\ & = R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} (1 - e^{-\lambda_{130\text{Sn}_{\text{g.s.}}} t}) + yR_{130\text{Sn}_{\text{is.}}} (1 - e^{-\lambda_{130\text{Sn}_{\text{is.}}} t}). \end{aligned}$$

This can be grouped into constant terms:

$$\lambda_{130\text{Sb}_{\text{is.}}} A = R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}}.$$

which can be solved for A :

$$A = \frac{R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sb}_{\text{is.}}}}.$$

And terms with coefficients of $e^{-\lambda_{130_{\text{Sn}_{\text{g.s.}}}} t}$:

$$-B_1 \lambda_{130_{\text{Sn}_{\text{g.s.}}}} + \lambda_{130_{\text{Sb}_{\text{is.}}}} B_1 = -R_{130_{\text{Sn}_{\text{g.s.}}}} .$$

leading to a value for B_1 :

$$B_1 = \frac{-R_{130_{\text{Sn}_{\text{g.s.}}}}}{\lambda_{130_{\text{Sn}_{\text{g.s.}}}} - \lambda_{130_{\text{Sb}_{\text{is.}}}}} .$$

Lastly terms with coefficients of $e^{-\lambda_{130_{\text{Sn}_{\text{is.}}}} t}$ are solved

$$-B_2 \lambda_{130_{\text{Sn}_{\text{is.}}}} + \lambda_{130_{\text{Sb}_{\text{is.}}}} B_2 = -y R_{130_{\text{Sn}_{\text{is.}}}} .$$

which gives the value for B_2 :

$$B_2 = \frac{-y R_{130_{\text{Sn}_{\text{is.}}}}}{\lambda_{130_{\text{Sn}_{\text{is.}}}} - \lambda_{130_{\text{Sb}_{\text{is.}}}}} .$$

The general solution is then the sum of the homogeneous and particular solutions:

$$\begin{aligned} N_{130_{\text{Sb}_{\text{is.}}}}(t) = & C e^{-\lambda_{130_{\text{Sb}_{\text{is.}}}} t} \\ & + \frac{R_{130_{\text{Sb}_{\text{is.}}}} + R_{130_{\text{Sn}_{\text{g.s.}}}} + y R_{130_{\text{Sn}_{\text{is.}}}}}{\lambda_{130_{\text{Sb}_{\text{is.}}}}} \\ & - \frac{R_{130_{\text{Sn}_{\text{g.s.}}}}}{\lambda_{130_{\text{Sn}_{\text{g.s.}}}} - \lambda_{130_{\text{Sb}_{\text{is.}}}}} e^{-\lambda_{130_{\text{Sn}_{\text{g.s.}}}} t} \\ & - \frac{y R_{130_{\text{Sn}_{\text{is.}}}}}{\lambda_{130_{\text{Sn}_{\text{is.}}}} - \lambda_{130_{\text{Sb}_{\text{is.}}}}} e^{-\lambda_{130_{\text{Sn}_{\text{is.}}}} t} . \end{aligned}$$

The initial condition $N_{130_{\text{Sb}_{\text{is.}}}}(0) = 0$ is applied by substituting $t = 0$ into the general

solution:

$$\begin{aligned}
 N_{130\text{Sb}_{\text{is.}}}(0) &= Ce^0 \\
 &+ \frac{R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sb}_{\text{is.}}}} \\
 &- \frac{R_{130\text{Sn}_{\text{g.s.}}}}{\lambda_{130\text{Sn}_{\text{g.s.}}} - \lambda_{130\text{Sb}_{\text{is.}}}} e^0 \\
 &- \frac{yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sn}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{is.}}}} e^0.
 \end{aligned}$$

Since $e^0 = 1$, this simplifies to:

$$\begin{aligned}
 0 &= C \\
 &+ \frac{R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sb}_{\text{is.}}}} \\
 &- \frac{R_{130\text{Sn}_{\text{g.s.}}}}{\lambda_{130\text{Sn}_{\text{g.s.}}} - \lambda_{130\text{Sb}_{\text{is.}}}} \\
 &- \frac{yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sn}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{is.}}}}.
 \end{aligned}$$

Rearranging for C , it is found that:

$$\begin{aligned}
 C &= - \frac{R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sb}_{\text{is.}}}} \\
 &+ \frac{R_{130\text{Sn}_{\text{g.s.}}}}{\lambda_{130\text{Sn}_{\text{g.s.}}} - \lambda_{130\text{Sb}_{\text{is.}}}} \\
 &+ \frac{yR_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sn}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{is.}}}}.
 \end{aligned}$$

$$\begin{aligned}
N_{^{130}\text{Sb}_{\text{is.}}}(t) = & \left(-\frac{R_{^{130}\text{Sb}_{\text{is.}}} + R_{^{130}\text{Sn}_{\text{g.s.}}} + yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sb}_{\text{is.}}}} \right. \\
& + \frac{R_{^{130}\text{Sn}_{\text{g.s.}}}}{\lambda_{^{130}\text{Sn}_{\text{g.s.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} + \frac{yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sn}_{\text{is.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} \left. \right) e^{-\lambda_{^{130}\text{Sb}_{\text{is.}}} t} \\
& + \frac{R_{^{130}\text{Sb}_{\text{is.}}} + R_{^{130}\text{Sn}_{\text{g.s.}}} + yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sb}_{\text{is.}}}} \\
& - \frac{R_{^{130}\text{Sn}_{\text{g.s.}}}}{\lambda_{^{130}\text{Sn}_{\text{g.s.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} e^{-\lambda_{^{130}\text{Sn}_{\text{g.s.}}} t} \\
& - \frac{yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sn}_{\text{is.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} e^{-\lambda_{^{130}\text{Sn}_{\text{is.}}} t} .
\end{aligned}$$

The same is of course true for the number of $^{130}\text{Sb}_{\text{g.s.}}$ nuclei, which only depends on $^{130}\text{Sn}_{\text{is}}$ and $R_{^{130}\text{Sb}_{\text{gs}}}$.

The background spectra do not allow the number of $^{130}\text{Sb}_{\text{is.}}$ nuclei to be measured, but the respective activities instead. These are defined as $A(t) = \lambda N(t)$ and for the activities of the four beam components these equations apply:

$$\begin{aligned}
A_{^{130}\text{Sn}_{\text{g.s.}}}(t) &= R_{^{130}\text{Sn}_{\text{g.s.}}} \left(1 - e^{-\lambda_{^{130}\text{Sn}_{\text{g.s.}}} t} \right) \\
A_{^{130}\text{Sn}_{\text{is}}}(t) &= R_{^{130}\text{Sn}_{\text{is}}} \left(1 - e^{-\lambda_{^{130}\text{Sn}_{\text{is}}} t} \right) \\
A_{^{130}\text{Sb}_{\text{is.}}}(t) &= \lambda_{^{130}\text{Sb}_{\text{is.}}} \left(-\frac{R_{^{130}\text{Sb}_{\text{is.}}} + R_{^{130}\text{Sn}_{\text{g.s.}}} + yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sb}_{\text{is.}}}} \right. \\
& + \frac{R_{^{130}\text{Sn}_{\text{g.s.}}}}{\lambda_{^{130}\text{Sn}_{\text{g.s.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} + \frac{yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sn}_{\text{is.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} \left. \right) e^{-\lambda_{^{130}\text{Sb}_{\text{is.}}} t} \\
& + R_{^{130}\text{Sb}_{\text{is.}}} + R_{^{130}\text{Sn}_{\text{g.s.}}} + yR_{^{130}\text{Sn}_{\text{is.}}} \\
& - \lambda_{^{130}\text{Sb}_{\text{is.}}} \frac{R_{^{130}\text{Sn}_{\text{g.s.}}}}{\lambda_{^{130}\text{Sn}_{\text{g.s.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} e^{-\lambda_{^{130}\text{Sn}_{\text{g.s.}}} t} \\
& - \lambda_{^{130}\text{Sb}_{\text{is.}}} \frac{yR_{^{130}\text{Sn}_{\text{is.}}}}{\lambda_{^{130}\text{Sn}_{\text{is.}}} - \lambda_{^{130}\text{Sb}_{\text{is.}}}} e^{-\lambda_{^{130}\text{Sn}_{\text{is.}}} t} .
\end{aligned}$$

$$\begin{aligned}
 A_{130\text{Sb}_{\text{g.s.}}}(t) = & \lambda_{130\text{Sb}_{\text{g.s.}}} \left(-\frac{R_{130\text{Sb}_{\text{g.s.}}} + (1-y)R_{130\text{Sn}_{\text{is}}}}{\lambda_{130\text{Sb}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{g.s.}}}} + \frac{(1-y)R_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sn}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{is.}}}} \right) e^{-\lambda_{130\text{Sb}_{\text{g.s.}}} t} \\
 & + R_{130\text{Sb}_{\text{g.s.}}} + (1-y)R_{130\text{Sn}_{\text{is.}}} \\
 & - \lambda_{130\text{Sb}_{\text{is.}}} \frac{(1-y)R_{130\text{Sn}_{\text{is.}}}}{\lambda_{130\text{Sn}_{\text{is.}}} - \lambda_{130\text{Sb}_{\text{g.s.}}}} e^{-\lambda_{130\text{Sn}_{\text{is.}}} t} .
 \end{aligned}$$

For $t \rightarrow \infty$ this simplifies to

$$\begin{aligned}
 A_{130\text{Sn}_{\text{g.s.}}}(t) &= R_{130\text{Sn}_{\text{g.s.}}} \\
 A_{130\text{Sn}_{\text{is}}}(t) &= R_{130\text{Sn}_{\text{is}}} \\
 A_{130\text{Sb}_{\text{is.}}}(t) &= R_{130\text{Sb}_{\text{is.}}} + R_{130\text{Sn}_{\text{g.s.}}} + yR_{130\text{Sn}_{\text{is.}}} \\
 A_{130\text{Sb}_{\text{g.s.}}}(t) &= R_{130\text{Sb}_{\text{g.s.}}} + (1-y)R_{130\text{Sn}_{\text{is.}}}
 \end{aligned}$$

B. ^{134}Sn Beam Composition

B.1. Low-energy Coulomb-excitation with a ^{206}Pb target

Figure B.1 presents γ -ray spectra obtained with kinematic gates on both target-like recoils and beam-like ejectiles. The spectrum on the left is shown without Doppler correction, whereas the spectrum on the right is Doppler-corrected for the ejectiles. In the uncorrected spectrum, a peak at 162 keV is visible, which may indicate excitation associated with the $^{134}\text{Sb}_{7-}$ component of the beam. In the Doppler-corrected spectrum, a peak at 318 keV is observed, corresponding to population of the 2_1^- state via the $0_{\text{g.s.}}^- \rightarrow 2_1^-$ excitation and its subsequent $2_1^- \rightarrow 1_1^-$ decay.

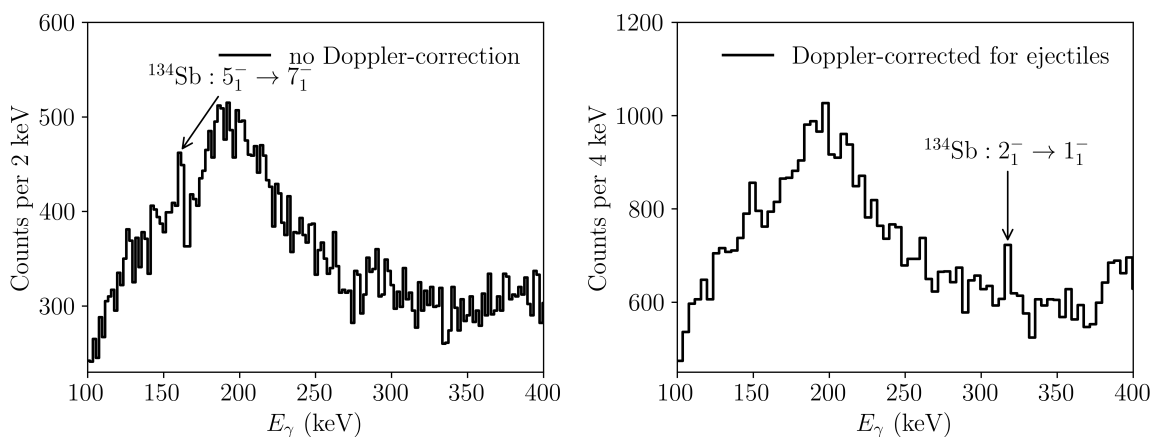


Figure B.1. γ -ray spectra for the 2025 ^{134}Sn Coulomb-excitation experiment using a ^{206}Pb target, obtained with kinematic gates on recoils and ejectiles. Left: spectrum without Doppler correction. Right: the same spectrum Doppler-corrected for the ejectiles.

B.2. Matrix Elements in $A = 134$ Isobars and ^{194}Pt

Table B.1 lists the matrix elements adopted for the GOSIA2 analysis for the $A = 134$ beam composition. In addition to observed excited states, a buffer state was introduced to effectively account for possible feeding from unobserved higher-lying states.

Table B.1. Transition data for selected states in ^{134}Ba , ^{134}Te and ^{194}Pt from NNDC [195] if not other stated. Values used as input values for the GOSIA analysis.

I_i^π	E_x [keV]	τ [ps]	I_f^π	E_γ [keV]	BR [%]	Mult. λL	δ	$B(\lambda L) \downarrow$ [W.u.]	$\langle \lambda L \rangle$ [eb $^{L/2}$]
^{134}Ba									
2_1^+	604.7	7.39(13)	0^+	604.7	100	E2	—	33.6(6)	0.827(7)
2_2^+	1168	3.90(1.15)	2_1^+	563.3	21.47(6)	M1	-7.4(9)	$7(3) \times 10^{-3}$	0.079(17)
						E2	-7.4(9)	73(22)	1.22(18)
			0^+	1168.0	100	E2	—	0.42(13)	0.093(14)
4_1^+	1400.6	1.20(13)	2_1^+	795.9	100.00	E2	—	52(6)	1.38(8)
			2_2^+	232.6	0.01(1)	E2	—	0.16(16)	0.077(77)
6_1^+	2299.7	—	4_1^+	899.1	100	E2	—	—	—
8_1^+	2835.9	—	6_1^+	536.0	100	E2	—	—	—
^{134}Te									
2_1^+	1279.1	1.14(5)	0^+	1279.0	100	E2	—	$5.1(2)^i$	0.323(6)
2_2^+	2464.9	< 1440	0^+	2465.3	100^{ii}	E2	—	—	—
			2_1^+	1185.9	9^{ii}	M1/E2	—	—	—
4_1^+	1576.1	1960(160)	2_1^+	297.0	100	E2	—	4.3(4)	0.40(2)
^{194}Pt									
2_1^+	328.5	60.1(6)	0^+	328.5	100	E2	—	49.5(5)	1.285(7)
2_2^+	622.0	50.5(5.8)	0^+	622.0	13.7(2)	E2	—	0.286_{-15}^{+44}	0.0977_{-26}^{+72}
			2_1^+	293.6	100	E2	0.22(8)	89_{-10}^{+12}	1.72_{-10}^{+11}
						M1	0.22(8)	$8.8_{-27}^{+45} \times 10^{-5}$	0.028_{-5}^{+6}
4_1^+	811.4	5.34(29)	2_1^+	482.9	100	E2	—	85.1_{-44}^{+48}	2.26(6)
4_2^+	1229.5	5.48(87)	2_2^+	607.5	100	E2	—	21.5_{-34}^{+46}	1.14_{-9}^{+12}
			4_1^+	418.1	14.6(14)	E2	>3	18_{-4}^{+8}	1.0_{-1}^{+2}
						M1	>3	$9.2_{-92}^{+29} \times 10^{-4}$	0.122_{-122}^{+18}
			2_1^+	901.0	8.6(11)	E2	—	0.26_{-7}^{+9}	0.12(2)
6_1^+	1412.0	2.31(72)	4_1^+	600.6	100	E2	—	67_{-16}^{+30}	2.4_{-3}^{+5}
6_2^+	1925.9	1.88(29)	4_2^+	696.4	100	—	—	—	—
			6_1^+	514.0	6(4)	—	—	—	—
			4_1^+	1114.5	12(4)	—	—	—	—
8_1^+	2099.7	1.59(43)	6_1^+	687.7	100	E2	—	50_{-11}^{+18}	2.38_{-28}^{+40}
8_2^+	2689.3	0.88_{-16}^{+13}	6_2^+	763.4	100	E2	—	53_{-7}^{+12}	2.45_{-17}^{+26}
10_1^+	2848.7	1.52_{-32}^{+43}	8_1^+	749.0	100	E2	—	34_{-8}^{+9}	2.18(27)

ⁱ Values for 2^+ state are taken from Ref. [168].ⁱⁱ NNDC values are given without uncertainty.

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