In-beam Study of Extremely Neutron Deficient Nuclei Using the Recoil-Decay Tagging Technique

BAHARAK HADINIA

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Abstract
The low-lying structures of the extremely neutron-deficient nuclei $^{106}$Te, $^{107}$Te, $^{110}$Xe, $^{170}$Ir and $^{172}$Au have been investigated experimentally. Prompt gamma rays emitted in fusion-evaporation reactions were detected by the Jurogam HPGe array. The gamma rays were assigned to specific reaction channels using the recoil-decay tagging technique provided by the gas-filled separator RITU and the GREAT focal-plane spectrometer. The experimental set-up and the technique used to extract the information from the experimental data are described in detail. Results were interpreted in terms of the nuclear shell model and Total Routhian Surface calculations. In addition, decay studies on $^{170}$Ir, $^{172}$Au and $^{164}$Re led to the discovery of new alpha-decay branches in these nuclei.
Publications

This thesis is based on the first five publications in the list below.

1. First identification of \( \gamma \)-ray transitions in \(^{107}\)Te.

2. First identification of excited states in \(^{106}\)Te and evidence for isoscalar-enhanced vibrational collectivity.

3. In-beam \( \gamma \)-ray and \( \alpha \) Spectroscopy of \(^{170}\)Ir.

4. Identification of Excited States in the N=Z+2 Nucleus \(^{110}\)Xe : Evidence for Enhanced Collectivity Near the N=Z=50 Double Shell Closure.
   M. Sandzelius, B. Hadinia, B. Cederwall, K. Andgren, E. Ganoğlu, I. G. Darby, M. Dimmock, S. Eeckhaudt, T. Grahn, P. Greenlees, E. Ideguchi, P. Jones,
D. T. Joss, R. Julin, S. Juutinen, A. Khaplanov, M. Leino, L. Nelson, M. Ny-
man, R. D. Page, J. Pakarinen, E. S. Paul, M. Petri, P. Rahkila, J. Sarén,
C. Scholey, J. Uusitalo, R. Wadsworth and R. Wyss,

5. Identification of $\gamma$-rays from $^{172}$Au and new $\alpha$ decays in $^{172}$Au and $^{164}$Re.

B. Hadinia

B. Cederwall, K. Andgren, E. Ganioğlu, M. B Gómez Hornillos,
T. Grahn, P. Greenlees, G. Hornillos, E. Ideguchi, U. Jakobsson, A. Johnson,
P. M. Jones, R. Julin, S. Ketellhut, M. Leino, M. Niikura, I. Özgür, R. D. Page,
E. S. Paul, P. Rahkila, M. Sandzelius, C. Scholey and J. Uusitalo,
Manuscript in preparation, intended for submission to Phys. Rev. C.

Other articles the author has contributed to, which are not commented on within
this thesis.

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D. Sohler, K. Lagergren, J. Blomqvist, B. Cederwall, A. Johnson, B. Hadinia
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2. Recoil-decay tagging of $\gamma$-rays in the extremely neutron-deficient nucleus $^{162}$Os.

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3. Probing structural changes in the very neutron-deficient Os isotopes with
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D. T. Joss, N. Amzal, D. E. Appelbe, T. Bäck, C. J. Barton, M. A. Bent-
ley, B. Cederwall, J. F. C.Cocks, D. M. Cullen, S. Eeckhautd, T. Grahn,
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H. Kankaanpaa, A. Keenan, H. Kettunen, S. L.King, P. Kuusiniemi, K. Lager-
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J. Pakarinen, E. S. Paul, P. Rahkila, C. Scholey, A. Savelius, J. Simpson,
M. J. Taylor, J. Uusitalo, S. J. Williams, D. D. Warner, D. R. Wiseman and
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dt, T. Grahn, P. M. Jones, R. Julin, S. Juutinen, B. Hadinia, H. Kettunen, M. Leino, A.-P. Leppä
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6. High-spin states in the proton-unbound nucleus $^{161}$Re.
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7. First identification of excited states in $^{169}$Ir.
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glu, K. Andgren, D. E. Appelbe, C. J. Barton, T. Bück, S. Eeckhau
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M. Sandzelius, J.-N. Scheurer, G. Sletten, D. Sohler, J. Timár, M. Trotta,
9. Coulomb shifts and shape changes in the mass 70 region.


10. Nuclear levels in proton-unbound \(^{109}\)I: relative single-particle energies and octupole collectivity beyond the proton dripline.


11. Alpha decay of \(^{159}\)Re and proton emission from \(^{155}\)Ta.


12. First observation of isomeric decays in the r-process waiting-point nucleus \(^{130}\)Cd\(_{82}\).

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Chapter 1

Introduction

Atomic nuclei are complex many-body systems and exhibit an interplay between single-particle and collective degrees of freedom. In order to describe and predict nuclear structure a variety of theoretical models have been created. Experimental information is needed in order to test the validity of and/or improve the model descriptions in different regions of the nuclear chart. One important tool for studying nuclear properties is gamma-ray spectroscopy.

An impressive technical development of radiation detectors and electronics during the last few decades have provided an extensive range of results. For instance different exotic nuclear shapes have been observed. In fact, relatively few nuclei have spherical shapes in their ground states and different shapes can be found in the same nucleus. Deformed nuclei can schematically be classified as prolate (American football shaped), oblate (flattened sphere) and triaxially deformed (flattened American football). This classification depends on the relative sizes of the axes of the ellipsoidal shape approximating the nuclear mass distribution and that will be discussed in this thesis. Many nuclei also exhibit shape coexistence, i.e the possibility of assuming different shapes (such as prolate and oblate) at similar excitation energies. Depending on the shape of a nucleus, different modes of excitation are possible. For instance, the low-energy part of the level scheme in some spherical nuclei can be interpreted in terms of surface vibrations. The three lowest and most common vibrational modes of excitation are dipole, quadrupole and octupole vibrations. Collective vibrational motion in quadrupole-deformed nuclei also has been known in two forms: beta- and gamma-vibration. In deformed nuclei with collective rotational motion, the energies of the excited states obey similar rules as in a quantum mechanical rotor. A rotational band has a sequence of states with energies proportional to $I(I+1)$, where $I$ is the total angular momentum, which are linked with E2 transitions. Many nuclei also have a transitional character involving both vibrational and rotational features. Experimental evidence has indicated a tendency toward octupole deformations (pear-shaped nuclei). Octupole deformation is understood to arise due to strong coupling of adjacent orbitals with $\Delta l = 3$.
or $\Delta j = 3$. The rotational bands of an octupole deformed nucleus can display a sequence of states with alternating positive and negative parities, which are linked by strong $E1$ transitions. The superdeformed (SD) shape, corresponds to a shape with a major to minor principal axis ratio of 2:1 or 3:2 in the ellipsoid. The transitions between SD and normal deformed states have been established for some nuclei.

One of the main topics of this thesis is the study of nuclei close to the $N = Z$ line. These studies can give information on proton-neutron correlations in cases where the protons and neutrons occupy similar orbitals as well as the effects on the nuclear structure from the active single particle orbitals close to the Fermi level. The level structures of neutron-deficient nuclei, $^{106}$Te, $^{107}$Te and $^{110}$Xe, have been studied in this work. Typical cross sections for the population of these nuclei are of the order of nb ($10^{-37}m^2$). This demonstrates the level of selectivity required in these experiments, which is at the limit of what can be achieved by present-day technology.

The level structure of odd-odd nuclei is complicated due to a high level density and has been studied less extensively. Features such as signature inversion and chirality have been found in doubly-odd nuclei. Study of odd-odd nuclei can give information on the interactions between the valence particles and the role of the individual proton and neutron orbitals in driving the nuclear shape. This work describes the study of two odd-odd nuclei, $^{170}$Ir and $^{172}$Au.

The nuclei under study in this thesis are populated via fusion-evaporation reactions. In such a reaction, the target and projectile nuclei fuse together, forming a compound nucleus. The compound nucleus deexcites by evaporating particles which results in the population of a number of different final nuclei with a lower mass than the compound nucleus. The highly sensitive method of recoil-decay tagging is applied to select the reaction channels of interests.

Chapter 1 gives a flavour of nuclear structure science and the subjects of interest in this thesis. Chapter 2 contains a brief overview of the theoretical framework used to interpret the experimental results. In chapter 3 the experimental set-up and the applied technique are described in detail. In chapter 4 the analysis of the collected data is described. Chapter 5 presents summaries of the publications which this thesis is based on.

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1. The longest segment that can be obtained by joining two points on the ellipse.
2. Line segment which passes through the center perpendicular to the major axis.
Chapter 2

Theoretical Framework

2.1 Radioactive Decays

Alpha, beta and gamma decay are the three main radioactive decay modes occurring in atomic nuclei. However, beyond the proton dripline, proton decays also are observed. In the following sections some physical properties of the alpha and gamma decay modes are described.

2.1.1 Alpha Decay

Alpha decay is the emission of a helium nucleus, $^4_2\text{He}$, and is represented by

$$A^Z_X \rightarrow A^Z_X - 4 - 2\ Y_{N-2} + \alpha,$$

where $A$ and $Z$ are the mass number and the atomic number, respectively, of the decaying nucleus, which is commonly called the mother nucleus. The decay occurs spontaneously in many heavy nuclei with mass number greater than 208 and also in some neutron-deficient nuclei. The half-lives of alpha decays vary strongly as a function of the kinetic energy of the alpha particle, from $\mu$s to $10^{19}$ years.

The relation between the alpha particle energy and the half-life of the decay can be estimated by a simple semi-empirical formula developed by Taagepera and Nurminen

$$\log_{10} t_{1/2} = 1.61(ZE_{\alpha}^{-1/2} - Z^{2/3}) + 28.9,$$

where $E_{\alpha}$ is the kinetic energy of the alpha particle in the unit of MeV and $t_{1/2}$ is the partial half-life for the alpha branch in the unit of year.

It was pointed out by Hans Geiger and John Mitchell Nuttall for the first time and known as the Geiger-Nuttall law.

3
particle \(E_\alpha\):

\[ E_\alpha \approx Q_\alpha \frac{A}{A+4}, \quad (2.3) \]

where \(A\) is the mass number of the daughter nucleus. If the mass of the mother nucleus is \(m_M\), the mass of its decay daughter is \(m_D\) and the mass of the alpha particle is \(m_\alpha\), \(Q_\alpha\) is the mass difference between the initial and final system \((Q_\alpha = (m_M - m_D - m_\alpha)c^2)\).

Before the decay, the alpha particle is formed inside the mother nucleus, and how this happens has not yet been fully understood. From the simplified shell model perspective the Hamiltonian of the system before forming the alpha particle can be written as

\[
H = \sum_{i=1}^{4} \frac{p_i^2}{2m_i} + \sum_{i=1}^{4} V_{iD}(r_i) + \sum_{i \leq j=1}^{4} V_{ij}(r_i - r_j), \quad (2.4)
\]

where \(V_{iD}(r_i)\) are the interactions between the nucleons (the four nucleons which eventually become the alpha cluster) and the daughter nucleus and \(V_{ij}(r_i - r_j)\) are the interactions between the four nucleons themselves. At some point the four nucleons are bound together and form the alpha cluster.

After the formation, the Hamiltonian for the two-body system of the alpha particle and the daughter nucleus, when they are far enough and the short range nuclear force is negligible, can be written as

\[
H = \frac{P_D^2}{2m_D} + \frac{P_\alpha^2}{2m_\alpha} + V_{\text{Coulomb}}, \quad (2.5)
\]

Considering the non-recoil approximation, the mass of the daughter nucleus \((m_D)\) is much greater than the mass of the alpha particle \((m_\alpha)\), the first term in equation \((2.5)\) is negligible. The Schrödinger equation

\[
H\psi = E\psi, \quad (2.6)
\]

applying the Hamiltonian from the eq.(2.5) can be written, in spherical coordinate, as

\[
\left(\frac{d^2}{dr^2} + k^2 - \frac{l(l+1)}{r^2} - \frac{2m_\alpha V_{\text{Coulomb}}}{h^2}\right)r\psi(r) = 0. \quad (2.7)
\]

where \(k = \sqrt{2mE/h^2}\) and \(l\) is the orbital angular momentum of the outgoing alpha particle. The solution of this equation is given below

\[
r\psi_{l}^{(out)}(r) = N_l[G_l(r) + iF_l(r)], \quad (2.8)
\]

where \(N_l\) is a normalization constant and \(F_l\) and \(G_l\) are the regular and irregular Coulomb functions, respectively. The probability rate per second that the alpha decay tunnels through the surface element, \(dS = r^2 \sin \theta d\theta d\phi\), is obtained
2.1. RADIOACTIVE DECAYS

by $|\psi_{\text{out}}^{(l)}(r)|^2 v dS$, where $v = \hbar k/m$ is the velocity of the alpha particle ($m$ is the reduced mass of the alpha and the daughter nucleus). For large $r$

$$\lim_{r \to \infty} |r\psi_{\text{out}}^{(l)}(r)|^2 = N_l^2 [G_l^2 + F_l^2] = N_l^2. \quad (2.9)$$

Therefore the alpha decay life-time $\tau$, which is the inverse of the decay probability per second, is obtained by $1/|N_l|^2 v$. The normalization constant $N_l$ is obtained by finding the distance $R$ where the wave function $\psi_{\text{out}}^{(l)}(R)$ and the solution $\psi_{l}^{(R)}(R)$ of the Schrödinger equation inside, which is regular at the origin and has outgoing boundary conditions, match each other. Here $\psi_{l}(R)$, $\psi_{l}^{(int)}(R)$, is the formation amplitude of the alpha cluster, which in fact is the overlap between the mother wave function and the antisymmetrized products of the daughter and alpha cluster wave functions. Thus the width and consequently the life-time can be given by

$$\Gamma_l(R) = \frac{\hbar}{\tau} = \frac{\hbar^2 k}{m} \frac{R^2 |\psi_{l}(R)|^2}{G_l^2(R) + F_l^2(R)}. \quad (2.10)$$

The evaluation of this expression (2.10), in particular the wave function $\psi_{l}(R)$, is difficult. Therefore the effective parameters frequency and penetration probability corresponding to the terms $\frac{\hbar R|\psi_{l}(R)|^2}{m}$ and $\frac{\hbar R}{G_l^2(R) + F_l^2(R)}$, respectively, are usually used.

A classical explanation of alpha emission, after the alpha cluster is formed, is that the alpha particle tunnels through the Coulomb barrier. Figure 2.1 illustrates schematically the potential of the nuclear system which alpha decays. The potential barrier at the surface of the daughter nucleus is smoothed by the superposition of a nuclear potential [2] and the Coulomb potential in a classical perspective.

The probability $P$ (the quantum mechanical quantity $\frac{\hbar R}{G_l^2(R) + F_l^2(R)}$) to penetrate the barrier can be obtained (semiclassical approximation) by integrating over the potential barrier as

$$P = \exp\left[-2 \int_{R_i}^{R_o} \frac{2m}{\hbar^2} (V(r) - Q_\alpha)^{1/2} dr\right], \quad (2.11)$$

where $m$ is the reduced mass of the alpha particle and the daughter nucleus in this two body system. $R_i$ and $R_o$ correspond to the inner and outer classical turning points (see Fig. 2.1), respectively. A common potential which was proposed by Rasmussen [3] is

$$V(r) = -1100 \times \exp\left[-\frac{r - 1.17 A^{1/3}}{0.574}\right] + \frac{2Ze^2}{4\pi\varepsilon_0 r} + \frac{\hbar^2}{2mr^2}l(l + 1), \quad (2.12)$$

where $A$ and $Z$ are the mass and atomic number, respectively, of the daughter nucleus and $l$ is the orbital angular momentum of the emitted alpha particle. As can be seen in Fig. 2.1, the width of the barrier increases as the angular momentum
carried by the alpha particle increases, which results in a smaller probability for tunnelling and a longer half-life.

A naive classical interpretation of the introduced quantum mechanical quantity of $\frac{R_0 |\psi_l(R)|^2}{m}$, is the frequency ($f$) with which the alpha particle appears at the inner edge of the Coulomb barrier. The decay constant and the half-life can then be calculated by

$$ \lambda = fP \quad t_{1/2} = \frac{ln2}{\lambda}, $$

(2.13)

where $P$ is the probability to penetrate the barrier. The frequency $f$ can be estimated as $v/R_i$, where $v$ is the velocity of alpha particles inside the nucleus. For a typical potential depth of $V_0 \approx 35$ MeV and $Q_\alpha \approx 5$ MeV $f$ is approximately\(^2\) in the order of $10^{21}$ s\(^{-1}\).

\(^2\)Assuming the released energy in the reaction, $Q_\alpha$, is transferred to the alpha particle as a kinetic energy, the frequency can be estimated from $f = \sqrt{(Q_\alpha)/(2\pi^2 m R_i^2)}$. 

Figure 2.1: The potential energy of an alpha-decaying nucleus. The solid curve displays the potential (2.12) when the angular momentum carried by the alpha particle is zero. The dashed curve illustrates schematically the potential when the alpha particle carries an orbital angular momentum.
2.1. RADIOACTIVE DECAYS

Hindrance Factor and Reduced Width

The hindrance factor is a quantity which is used to evaluate how favourable an alpha decay is in terms of the structure of the mother and daughter nuclei. It is defined as the theoretical alpha decay rate divided by the experimental one,

\[ HF = \frac{\lambda_{\text{calc}}}{\lambda_{\text{exp}}} = \frac{t_{1/2(\text{exp})}}{t_{1/2(\text{calc})}} \]  

(2.14)

Transitions between states with equal spins or similar structures give \( HF \leq 4 \), whereas transitions between states with different spin or structures are more hindered and usually have \( HF > 4 \). Another commonly used quantity to describe the properties of alpha decay is the reduced alpha decay width which is defined as

\[ \delta^2 = \frac{\lambda_{\text{exp}} h}{P}, \]  

(2.15)

where \( \lambda_{\text{exp}} \) is the experimental decay constant, \( h \) is Planck’s constant and \( P \) is the probability of barrier penetration. The reduced width for an unhindered alpha decay is \( \delta^2 \geq 40 \text{ keV} \) and for a hindered decay \( \delta^2 \) is approximately 1 keV. An alternative definition of hindrance factor in terms of alpha decay reduced width is

\[ HF = \frac{\delta^2(\text{ground state} \rightarrow \text{ground state decay of the closest even-even nucleus})}{\delta^2(\text{the decay under examination})}. \]  

(2.16)

2.1.2 Gamma Decay

When a nucleus gamma decays, it changes from a higher energy state to a lower energy state through the emission of electromagnetic radiation. The photons produced in this decay are known as gamma rays and have typical energies up to a few MeV. These photons transmit not only energy but also \( L \) units of angular momentum, where \( L = 1, 2, 3... \) is the multipolarity of the transition. The angular momentum and parity selection rules for the allowed transitions can be summarised as

\[ |I_i - I_f| \leq L \leq |I_i + I_f| \]

change in parity: even \( L \) magnetic, odd \( L \) electric,

no change in parity: even \( L \) electric, odd \( L \) magnetic,  

(2.17)

where \( I_i \) and \( I_f \) are the spins of the initial and final states. A transition is called stretched if the multipolarity \( L \) is equal to \( I_i - I_f \) i.e has the lowest value permitted by the selection rules.
CHAPTER 2. THEORETICAL FRAMEWORK

Internal Conversion

Another nuclear deexcitation process that competes with the emission of gamma rays is internal conversion. The electromagnetic field of a nucleus undergoing internal conversion interacts directly with the atomic electrons, which results in the emission of an electron usually from one of the innermost shells. The internal conversion coefficient, $\alpha$, for a gamma ray transition is defined as the ratio between the number of emitted electrons and the number of emitted gamma rays (see eq. 2.18). The conversion process is followed by the emission of low-energy photons (X-rays) when the vacancy created by the emission of an electron is filled by electrons from higher shell or by the emission of an Auger electron. Thus, the internal conversion coefficient can be measured experimentally using gamma-ray detectors by studying both the gamma rays as well as the X-rays. The conversion coefficient $\alpha$ is obtained by:

$$\alpha = \frac{I_x}{I_{\gamma}} = \frac{I_x/\epsilon_x}{I_{\gamma}/\epsilon_{\gamma} \eta},$$

(2.18)

where $I_X$ is the intensity of the X rays and $I_{\gamma}$ is the intensity of the gamma rays. $\epsilon_x$ and $\epsilon_{\gamma}$ are the efficiencies of the germanium detector for energies of the X and gamma rays, respectively. The coefficient $\eta$ is the fluorescent yield, which corrects for the Auger effect. Internal conversion coefficients depend on the atomic number, the multipolarity and the type of transition (electric or magnetic) and on the energy of the transition [4]. It is higher for high $Z$ nuclei and transitions with low energy and high multipolarity. Therefore, comparisons of the measured internal conversion coefficients with theoretical predictions can be used to identify the multipolarity and type (electric or magnetic) of the corresponding transition.

2.2 Fundamentals of Nuclear Structure

One of the major aims in nuclear physics is to gain a better understanding of the internal structure of atomic nuclei. The atomic nucleus is a quantum system with a size of $10^{-14}$ m to $10^{-15}$ m. It contains protons and neutrons which interact strongly with each other, and is an excellent environment for studying a many-body quantum system with a finite number of particles. The nuclear shell model was one of the first models that was created to describe the structure of nuclei. It has been very successful in describing nuclei near closed shells. An extension of the model using Monte Carlo techniques has made it possible to describe microscopically even heavy nuclei with many protons and neutrons outside closed shells. Today there are a number of models which have their specific domains where they can explain a limited number of experimentally observed phenomena, but no universal nuclear model exists. A goal of experimental nuclear structure physics is to provide stringent tests of nuclear models so that theory can advance towards a more fundamental understanding of nuclei. The following sections provide an introduction to a few basic nuclear models used to interpret experimental results in this thesis.
2.2.1 The Nuclear Shell Model

Maria Goeppert Mayer and a group including Otto Haxel, Hans Jensen and Hans Suess simultaneously and independently came to equivalent conclusions about the nuclear shell structure in 1949 [5]. If residual shell model interactions are neglected, each individual nucleon moves independently in an average field produced by the other nucleons. The Schrödinger equation can then be written as [6]:

\[ H_0 \psi(r) = \sum_i^A h_i \psi(r) = \sum_i^A [T_i + V(r_i)]\psi(r) = E\psi(r) \]  

(2.19)

where \( \psi \) can be the product of the eigenfunctions \( \phi_n \) of the single-particle Schrödinger equation

\[ h_i \phi_n(r_i) = \epsilon_n \phi_n(r_i). \]  

(2.20)

Solving the Schrödinger equation gives the wave functions (eigenfunctions) and the energy eigenvalues (eigenstates) of possible nuclear states that are available for a system of nucleons.

In the shell model the eigenstates are obtained by filling the single-particle energy levels while obeying the Pauli principle. Thus the wave function must be antisymmetric with respect to the exchange of spacial coordinates of any pair of neutrons or protons. This cannot be fulfilled by a simple product of single particle wave functions. An antisymmetric wave function was given by John Slater [7] and is known as the Slater determinant, which is expressed as follows:

\[ \psi(r_1, r_2, \ldots) = \frac{1}{\sqrt{N!}} \begin{vmatrix} \phi_n_1(r_1) & \phi_n_1(r_2) & \ldots \\ \phi_n_2(r_1) & \phi_n_2(r_2) & \ldots \\ \vdots & \vdots & \ddots \end{vmatrix} \]  

(2.21)

where \( N \) is the number of neutrons or protons. From this equation follows that if one tries to put two nucleons of the same kind in the same quantum state the resulting wave function \( \psi \) is zero.

The main challenge when using a shell model description is the choice of the one-body potential \( V(r) \). Since the nucleons generate the potential, it should have a radial dependence corresponding to the nuclear density. One simple form that is often used is the Wood-Saxon potential which for spherical nuclei is written as:

\[ V(r) = -V_0 [1 + \exp\left(\frac{r - R_0}{a}\right)]^{-1}, \]  

(2.22)

where \( V_0 \) is the depth of the potential well which varies with atomic and neutron numbers. \( R_0 = r_0 A^{1/3} \) is the nuclear radius (where \( r_0 \approx 1.25 \) fm) and the parameter \( a \), related to the thickness of the surface, depends upon the number of nucleons and usually is around 0.5-0.7 fm. The skin thickness over which the potential changes from \( 0.9 V_0 \) to \( 0.1 V_0 \) is given by \( 4a \ln 3 \). Solving the Schrödinger equation with this simple potential yields large gaps in the energy level distribution at the magic
particle numbers, 2, 8 and 20, in agreement with the systematics. By adding the strong nuclear spin-orbit interaction,

$$V_{so} = f(r) \vec{l} \cdot \vec{s}$$

$$f(r) = -V_{ls} \frac{\partial V(r)}{\partial r}$$

(2.23)

to the potential, eq.(2.22), the remaining magic numbers (28, 50, 82, 126, ...) can be reproduced. In eq.(2.23) $V(r)$ is the central potential and $V_{ls}$ is the strength constant of the spin-orbit interaction. However, the experimentally observed magic numbers for heavier nuclei are not reproduced by any simple central potential form.

The nuclear shell model described above treats the nucleons as independent particles moving in a common potential well. A better model would take into account residual interactions between nucleons. The N-nucleon Hamiltonian limited to two-body interactions can be written:

$$H = \sum_i T_i + \sum_{i \neq j} V_{ij} = \sum_i (T_i + V(r)) + \left( \sum_{i \neq j} V_{ij} - \sum_i V(r) \right) = H_0 + H_{res} (2.24)$$

Here $H_0$ is the single-particle energy and $H_{res}$ is the sum of the two-body matrix elements due to the residual interactions. The calculation of $H_{res}$ can be simplified by choosing a nucleus with closed shells as a reference core and only study the configurations with a few particles and holes relative to it. From studies of light nuclei it has recently been found that also three-body interactions play an important role in nuclei. However, such interactions can not easily be applied to a large number of nucleons due to a rapid increase in calculational complexity.

The experimental results for $^{106}$Te that are presented in this thesis were interpreted by comparing them to shell model predictions by Jan Blomqvist [9]. The nucleus $^{106}$Te has two protons and four neutrons outside the closed $N=Z=50$ shells. In the calculations the $d_{5/2}$ and $g_{7/2}$ subshells were chosen as a basis for the valence protons and neutrons. From data on light Sn isotopes it is known that the $d_{5/2}$ and $g_{7/2}$ subshells are close in energy and separated from the higher $s_{1/2}$, $d_{3/2}$ and $h_{11/2}$ subshells by a substantial gap. Therefore the truncation of the basis to the $d_{5/2}$ and $g_{7/2}$ subshells is reasonable. From light Sn nuclei the single-particle energies for both protons and neutrons are chosen to be $\varepsilon(d_{5/2}) = 0$ and $\varepsilon(g_{7/2}) = 120$ keV. In this basis there are a total of 102 two-body interaction matrix elements ($\langle j_1j_2, JM | V_{12} | j_3j_4, JM \rangle$). This number can be reduced to 43 with the assumption of charge independence.

### 2.2.2 Deformation

The shape of a nucleus can be described by an expansion in spherical harmonics $Y_{\lambda \mu}$ and a few shape parameters, $\alpha_{\lambda \mu}$, depending on the multipolarity $\lambda$ of the
2.2. FUNDAMENTALS OF NUCLEAR STRUCTURE

vibration,

\[ R(\theta, \phi) = R_0 \left( 1 + \sum_{\lambda=1}^{\infty} \sum_{\mu=-\lambda}^{\lambda} \alpha_{\lambda \mu} Y_{\lambda \mu}(\theta, \phi) \right), \tag{2.25} \]

where \( R_0 \) is the radius of a sphere with the same volume as the nucleus. Therefore, for a spherical nucleus \( R(\theta, \phi) \) is equal to \( R_0 \). The term \( \lambda = 1 \) is called the dipole term. The effect of this term is a shift in the centre of mass. If the origin of the coordinate system is at the centre of mass the coefficients \( \alpha_{1 \mu} \) are zero. The quadrupole shapes corresponding to \( \lambda = 2 \), are described by five coefficients \( \alpha_{2 \mu} \). The quadrupole-deformed shape of a nucleus is often described by the two parameters \( \beta_2 \) and \( \gamma \), defined as follows [10]:

\[ \alpha_{20} = \beta_2 \cos \gamma; \alpha_{2\pm1} = 0; \alpha_{2\pm2} = \frac{1}{\sqrt{2}} \beta_2 \sin \gamma. \tag{2.26} \]

The range of these parameters are \( \beta_2 \geq 0 \) and \(-120^\circ < \gamma < +60^\circ\), as illustrated in Fig. 2.2. The parameter \( \beta_2 \) represents the elongation or flattening of the nuclear deformation and can be approximated by

\[ \beta_2 \simeq 1.06 \frac{\Delta R}{R_0}, \tag{2.27} \]

where \( \Delta R \) is the difference between the major and minor semi-axes of the ellipsoid. The parameter \( \gamma \) gives the degree of triaxiality in a deformed nucleus. The majority of deformed nuclei attain prolate shapes. The reason has not been fully understood but it might be explained by the fact that the nuclei with prolate shapes have larger moments of inertia in comparison with those with oblate shapes. The larger moment of inertia gives lower lying energy levels in a rotating nucleus which is more favourable (see section (2.2.7)).

2.2.3 The Liquid-Drop Model

The liquid-drop model describes the nucleus as a drop of incompressible nuclear fluid. The model is able to account for “macroscopic” properties of the nuclei that can not yet be described accurately by pure quantum models like the shell model. The liquid-drop model was first proposed by George Gamow in 1928. The Weizsäcker binding energy is an empirically refined form of the liquid-drop model and is given by:

\[ B = E_V + E_S + E_0 + E_{sym} + E_{pair} \]

\[ B = a_{vol} A - a_{surf} A^{2/3} - a_c Z (Z - 1) A^{-1/3} - a_{sym} (A - 2Z)^2 A^{-1} + \delta, \tag{2.28} \]

where the five terms give the volume energy, surface energy, Coulomb repulsion energy, symmetry energy, and the pair energy [12, 13]. The Coulomb repulsion and
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Figure 2.2: Schematic diagram of the shape parameters in rotating quadrupole-deformed nuclei. The $\gamma = 60^\circ$ axis corresponds to the non-collective oblate (rotation around the symmetry axis) shape, $\gamma = 0^\circ$ to collective prolate (rotation perpendicular to the symmetry axis) shape, $\gamma = -60^\circ$ to collective oblate and $\gamma = 120^\circ$ to non-collective prolate shape. The figure is taken from reference [11].
the surface binding energy are shape dependent; therefore, in order to apply the liquid-drop model to deformed nuclei some modification is needed. Using eq.(2.25), the Coulomb repulsion and surface energies are up to second order given by [14]

\[ E_S = E^0_S (1 + \frac{1}{8\pi} \sum_{\mu, \lambda} (\lambda - 1)(\lambda + 2)|\alpha_{\lambda\mu}|^2) \]

\[ E_c = E^0_c (1 - \frac{5}{4\pi} \sum_{\mu, \lambda} \frac{\lambda - 1}{2\lambda + 1}|\alpha_{\lambda\mu}|^2). \]  (2.29)

The energy difference between a deformed and a spherical shape is:

\[ E_S + E_c - E^0_S - E^0_c = E^0_S \left( \frac{2}{5}(1 - \frac{E^0_c}{2E^0_S})a_2^2 - 4 \frac{5}{105} \sqrt{\frac{5}{4\pi}} (1 + \frac{E^0_c}{E^0_S})a_2^2 + \ldots \right) \]  (2.30)

where \( a_2 = (5/4\pi)^{1/2} \beta_2 \). Neglecting third and higher order terms in eq.(2.30) the minimum is at \( \beta_2 = 0 \) which corresponds to a spherical shape. If \( \frac{E^0_c}{E^0_S} < 1 \) the curvature is positive, thus it is stable against small deformation. If \( \frac{E^0_c}{E^0_S} > 1 \) the repulsive Coulomb energy is greater than twice the surface energy, and the nucleus will be unstable against deformation and fission. In this model the stable shape is always found to be spherical, and it is therefore unable to explain the properties of deformed nuclei. However, the liquid-drop model has been successfully combined with the deformed shell model using a method developed by Strutinsky [15, 16]. The addition of quantal “shell effects” allow the model to reproduce stable deformations. These are abundant for nuclei far from closed spherical shell gaps.

### 2.2.4 The Nilsson Model

According to the liquid-drop model the energy minimum of a nucleus is at a spherical shape. However, many stable nuclei have been found to have a deformed shape. The stability of the deformed nuclei therefore are caused by shell effects which modify the liquid-drop energy, and which can create energy minima at finite deformation. The simple asymmetric harmonic oscillator potential gives a description of these shell effects. A more sophisticated potential is the Nilsson potential. This potential is produced by adding \( \vec{l} \cdot \vec{s} \) (spin-orbit) and \( \vec{l}^2 - < l^2 >_N \) terms to the single-particle Hamiltonian of the harmonic oscillator [17, 18]. The Nilsson Hamiltonian for an axially symmetric deformed nucleus is expressed as:

\[ H_{Ni} = -\frac{\hbar^2}{2m} \nabla_i^2 + \frac{m}{2} \omega_0^2 [x^2 + y^2] (1 + \frac{1}{3}z_2)^2 + z^2 (1 - \frac{2}{3}z_2)^2] - 2\kappa\hbar\omega_0 [\vec{l} \cdot \vec{s} - \mu (l^2 - < l^2 >_N)], \]  (2.31)

with

\[ \hbar\omega_0 = 41A^{-1/3}(1 + \frac{|N - Z|}{3A})\text{MeV}. \]  (2.32)
Figure 2.3: Single-particle energies as a function of deformation using Nilsson potential for $50 \leq Z \leq 82$ are shown (see the text for details on the labels). The deformation parameter is $\epsilon_2 = 0.95/\beta_2$ [19]. Shell gaps corresponding to the magic numbers 50 and 82 are indicated.
The $\vec{l} \cdot \vec{s}$ in eq.(2.31) is the spin-orbit force and the last term is added artificially to make the nuclear potential more flat in the centre and more attractive at the surface. The parameters $\kappa$ and $\mu$ are obtained by fitting to the experimental data. The parameter $\kappa$ is typically in the range of 0.05-0.12 and $\mu$ is between 0.0 and 0.7. An example of a Nilsson diagram in the region $50 \leq Z \leq 82$ is shown in Fig. 2.3. The labels of the orbitals are called the asymptotic Nilsson quantum numbers and are written as

$$\Omega[N_{n_z} \Lambda]$$  \hspace{1cm} (2.33)

where $N$ is the principal quantum number, $n_z$ is the number of oscillator quanta along the symmetry axis, $\Omega$ the projection of the single-particle angular momentum on the symmetry axis and $\Lambda$ is the projection of the single-particle orbital angular momentum on the symmetry axis, as illustrated in details in Fig. 2.4. In Fig. 2.3, the positive and negative values of $\epsilon_2 = 0.95\beta_2$ correspond to a prolate and oblate shape of the nucleus.

In order to interpret some of the experimental results presented in this thesis, the single-particle level energy has been calculated as a function of deformation. The calculations were performed using a different potential called the Woods-Saxon potential (see eq.(2.22)) for the nuclei under study [20]. Although the Nilsson potential gives analytical results, the Wood-Saxon potential is more realistic and better reproduces nuclear shapes.

![Asymptotic quantum numbers in the deformed Nilsson model](image)

Figure 2.4: Asymptotic quantum numbers in the deformed Nilsson model where $x$ is the axis of rotation.

### 2.2.5 Collective Motion

Collective excitations involve many of the nucleons. Nuclei having collective properties are usually those with many valence nucleons, that is, those with proton or
neutron numbers that are far from filled shells. For such nuclei the performance and interpretation of the shell model calculations are complicated. In contrast, a collective model allows calculations of spins, parities, and transition probabilities that are in good agreement with the measured properties of these nuclei. Rotation and vibration are examples of collective degrees of freedom and have been described in detail by e.g. Bohr and Mottelson [21, 22].

2.2.6 Vibration of Nuclei

A simple model for nuclear vibrations is based on small-amplitude vibrations around the spherical equilibrium of a liquid drop. The Hamiltonian is given by

\[ H = T + V = \frac{1}{2} B_\lambda \sum_\mu |\dot{\alpha}_{\lambda\mu}|^2 + \frac{1}{2} C_\lambda \sum_\mu |\alpha_{\lambda\mu}|^2. \]  

(2.34)

The coefficients \( B_\lambda \) and \( C_\lambda \) play the same role as the mass and the restoring force, respectively, in classical mechanics. Each \( \alpha_{\lambda\mu} \) oscillates with a frequency \( \omega_\lambda = \sqrt{C_\lambda / B_\lambda} \) and the vibration energy is \( \hbar \omega \). When second quantisation is applied by using operators that create \( (b_{\lambda\mu}^\dagger) \) and annihilate \( (b_{\lambda\mu}) \) vibrational quanta, the Hamiltonian can be written

\[ H = \sum_\mu \hbar \omega (b_{\lambda\mu}^\dagger b_{\lambda\mu} + \frac{3}{2}). \]  

(2.35)

The harmonic spectrum of energy levels obtained by solving the Schrödinger equation with the Hamiltonian from equation (2.35) for \( \lambda = 2 \) is shown in Fig.2.5. The ratio of 2 between the excitation energies of the first \( 4^+ \) and \( 2^+ \) states is a signature of a harmonic vibrator structure. However, it is not expected that nuclei should behave as ideal harmonic vibrators. Various effects such as static deformations and residual interactions can lead to anharmonicities and the breaking of the level degeneracies seen in Fig.2.5. In this thesis evidence for vibrational structures in tellurium and xenon isotopes is discussed in terms of the systematics of energy ratios [9, 23]. A complementary test would be to study transition rates. These show a characteristic pattern for vibrational excitations, being proportional to the number of vibrational phonons building up each excited state. The “one-phonon” rule also states that transitions can only remove one vibrational quantum at a time.

However, the nuclei studied in this work are highly unstable and situated far from the \( \beta \)-stability line. They can only be studied in reactions where they are populated with extremely small cross sections. The techniques used for measuring transition rates require high statistics and can unfortunately not be easily applied to these extremely neutron deficient isotopes.

The oscillation around a permanently deformed nuclear shape is also possible. In this case, instead of the phonon vibration quantum number, \( \lambda \), the projection of the phonon angular momentum along the symmetry axis, \( \nu \), is specified.
2.2. FUNDAMENTALS OF NUCLEAR STRUCTURE

2.2.7 Rotation of Nuclei

Collective nuclear rotation is only meaningful for deformed nuclei. Assuming that the nucleus behaves like a rotor, the energy levels $E$ are given by

$$E = \frac{h^2}{2\mathcal{J}} I(I + 1),$$

(2.36)

where $I$ is the total angular momentum of state and $\mathcal{J}$ is the static moment of inertia of the nucleus. However, the nucleus is not a rigid body, and the measured moments of inertia are typically only 30-50% of the rigid-body values at low spin [24]. A rotational band\(^3\) displays a sequence of states with energies proportional to $I(I + 1)$. For an even-even nucleus the ground-state rotational band can contain only even values of spin, $I$, since the axially symmetric deformed shape is invariant with the rotation of $180^\circ$ around the rotational axis. In an even-even nucleus the typical value for $E(4^+)/E(2^+)$ is 3.33 (rigid rotation) which can be used as a signature of a rotational band.

The kinematic moment of inertia is defined theoretically as

$$\mathcal{J}^{(1)} = \frac{h I_x}{\omega} = h^2 I_x \left( \frac{dE}{dI_x} \right)^{-1},$$

(2.37)

where $I_x$ is the projection of the total angular momentum, $I$, on the rotation axis. It is experimentally obtained by taking the differences in energy of the adjacent

---

\(^3\)Here the rotational band is considered to show only one signature. Another convention is considering two signature partners connected with M1 transitions as one rotational band.
rotational level with respect to the spin (transition from state with spin \( I \) to a
state with spin \( I - 2 \)),
\[
J^{(1)} = \frac{\hbar^2 (2I - 1)}{E_\gamma (I \rightarrow I - 2)}.
\] (2.38)

The kinematic moment of inertia requires knowledge about the absolute spin of the
levels and these are not known in many rotational bands. Therefore a more useful
parameter, the dynamic moment of inertia, is defined as below
\[
J^{(2)} = \hbar \frac{dI_x}{d\omega} = \hbar^2 \left[ \frac{d^2E_\gamma}{dI^2} \right]^{-1}.
\] (2.39)

It is experimentally obtained using the differences in energies of consecutive gamma-
ray transitions in a rotational band and is expressed as
\[
J^{(2)} = \frac{4\hbar^2}{E_\gamma (I + 2 \rightarrow I) - E_\gamma (I \rightarrow I - 2)} = \frac{4\hbar^2}{\Delta E_\gamma (I)}.
\] (2.40)

The unit of \( J \) is \( \hbar^2 \text{ MeV}^{-1} \). It can be seen from equation (2.40) that, if the
dynamical moment of inertia was a constant, the difference in transition energy
would be the same for all spin values, but in many cases this is not correct. The
value \( J^{(2)} \) is often found to change as a function of spin. For instance, an increase in
\( J^{(2)} \) can be caused by an increase in deformation of the nucleus or by rearrangement
of the nucleonic configuration.

2.3 Cranked Shell Model

The properties of a rotating nucleus can be described by the cranking model. The
single-particle cranking Hamiltonian in a coordinate system rotating with the nu-
cleus for a rotation with angular frequency \( \omega \) along an axis perpendicular to the
symmetry axis is
\[
h_\omega = h_0 - \omega j_x
\] (2.41)

where \( h_0 \) is the non-rotating Hamiltonian and \( j_x \) is the angular momentum around
the axis perpendicular to the symmetry axis. In most applications, \( h_0 \) is defined
using a deformed Nilsson or Wood-Saxon potential. The term \( -\omega j_x \) contains the
Coriolis and centrifugal terms which modify the nucleon orbitals. The Coriolis force
strives to align the angular momenta of the nucleons with the rotation axis (see
section 2.3.2). The eigenvalues of \( h_\omega \) are known as Routhians \( \epsilon'_\mu \)
\[
\epsilon'_\mu = \langle \mu | h_0 | \mu \rangle - \omega \langle \mu | j_x | \mu \rangle,
\] (2.42)

where \( \mu \) denotes the wave function of the single particle. The value of \( i_x \) is obtained
from
\[
\frac{d\epsilon'_\mu}{d\omega} = \langle \mu | j_x | \mu \rangle = i_x
\] (2.43)
A diagram of the single-particle energy as a function of the rotational frequency is known as a Routhian plot. The differential of the Routhian with respect to the rotational frequency gives the aligned angular momentum. The total Routhian is given by a sum over all occupied levels

$$E' = \sum_{\mu} \epsilon'_\mu$$  \hspace{1cm} (2.44)

The eigenfunctions of $h_0$ are not the eigenfunctions of $j_x$, and therefore the cranking causes mixing of single-particle states. Although the time-reversal symmetry is broken, there are some symmetries which divide the single-particle states into different non-mixed symmetry groups. Parity and signature (see section 2.3.1) are good quantum numbers that can be used to label the single-particle states in the cranking model. More details on the cranking model, such as consideration of pairing correlations, can be found in [25].

The deformation parameters which were introduced in section (2.2.2) can be used to describe the shape of nuclei in relation with a Routhian calculation. These calculations are called Total Routhian Surface (TRS) calculations [26]. A TRS plot shows the minimum energy for a given nucleon configuration of the nucleus as a function of $\beta_2$ and $\gamma$ at a given rotational frequency. This plot can be used to find out the model predictions for the most stable shape for a nucleus. This thesis includes TRS calculation plots for the nuclei under study. As an example, Fig. 2.6 displays a TRS plot for $^{172}$Au.

![TRS plot for $^{172}$Au](image)

Figure 2.6: Total Routhian Surfaces at a rotational frequencies of $h\omega = 0.0$ MeV, for the configuration $\pi(-, -1/2), \nu(-, -1/2)$ (left) and $\pi(+, +1/2), \nu(-, -1/2)$ (right) for $^{172}$Au are shown. Collective and non-collective energy minima are visible at a weakly deformed $\beta_2 = 0.11$, $\gamma = 0^\circ$ and $\beta_2 = 0.12$, $\gamma = -100^\circ$ for the right and left configurations, respectively.

The nucleus is considered to be “soft” with respect to a shape deformation when a range of possible shapes are allowed within a relatively small energy range.
This feature is called $\gamma$-softness if the possible range is varying in $\gamma$ values. The shape of the nucleus is sensitive to the configuration of particles outside closed shells and prolate, oblate and triaxial shapes can coexist at approximately the same excitation energy and angular momentum. The nucleus $^{106}$Te is predicted by TRS calculations to be relatively soft with respect to shape changes near a spherical shape. In particular, it is most easily deformed along the prolate axis and might therefore be susceptible to small-amplitude $\beta$-vibrations [9].

### 2.3.1 Parity and Signature

Parity, $\pi$, is defined as a symmetry of the wave function under space reflection which can either be positive or negative.

The eigenvalues of the rotation operator under 180° are,

$$ r = e^{-i\pi \alpha}, \quad (2.45) $$

where $\alpha$ ($I = \alpha \text{ (mod 2)}$) is defined as the signature of a state with spin $I$. In eq.(2.45) $\pi$ stands for the angle of rotation in the unit of radian. For even-$A$ nuclei $\alpha$ is 0 or 1 and for odd-$A$ nuclei $\alpha$ is $\pm \frac{1}{2}$. Thus the energy levels given by eq.(2.42) are divided into four groups, $(\pi = +, \alpha = 1/2), (\pi = +, \alpha = -1/2), (\pi = -, \alpha = 1/2)$, or $(\pi = -, \alpha = -1/2)$.

### 2.3.2 The Coriolis Force

The Coriolis force appears when a particle moves in a rotating frame. The classical Coriolis force is given by

$$ \vec{F}_{\text{Coriolis}} = -2m(\vec{\omega} \times \vec{v}), \quad (2.46) $$

where $m$ is the mass of the particle, $\vec{v}$ is the velocity of the particle in the rotating frame, and $\vec{\omega}$ is the angular frequency of the rotating frame. When two nucleons move in opposite directions (time-reversed) in the same orbital, the induced Coriolis forces on them act in opposite directions as shown schematically in Fig.2.7. The result might cause breaking of pairs of the nucleons coupled to $I^\pi = 0^+$. Phenomena such as backbending and signature splitting which have been observed in some nuclei are manifestations of the Coriolis force. Backbending occurs when at a certain frequency it is energetically more favourable for the nuclei to break a pair of nucleons and align their angular momenta with the rotational axis than to rotate faster to generate more collective spin. Backbending can be investigated by plotting the angular momentum or the moment of inertia as a function of rotational frequency. The plotted curve then displays an increase in the angular momentum or the moment of inertia with no changes or even a decrease in the rotational frequency.

The signature splitting can be studied by plotting the signature splitting parameter [27],

$$ S(I) = E(I) - \frac{[E(I + 1) + E(I - 1)]}{2}, \quad (2.47) $$
as a function of angular momentum. Here $E(I)$ is the energy of the state with spin $I$. The signature splitting is observed as a separation of $S(I)$ values for different signatures. In this thesis the existence of signature splitting in the observed rotational band in $^{170}$Ir is investigated and compared with the $N + 1$ isotope, $^{171}$Ir [20].

### 2.4 Angular Momenta in Atomic Nuclei

The spin-parity of the ground state in even-even nuclei is always $0^+$. In even-odd nuclei the ground state spin is equal to the angular momentum of the unpaired proton or neutron. The parity is given by $(-1)^l$, where $l$ is the orbital angular momentum of the unpaired nucleon. These rules are usually valid for spherical nuclei and also in deformed nuclei when the unpaired nucleon is decoupled (rotation aligned and has low $\Omega$). In deformed nuclei for strongly coupled nucleon (when the unpaired nucleon is symmetry aligned and has a high $\Omega$ value) the spin of the ground state is equal to the projection on the symmetry axis of the angular momentum, $\Omega$, of the unpaired nucleon (see Fig. 2.4 for more detail). Since the relative energy of the sublevels are varying with the deformation, the location of the unpaired valence nucleon and its value of $\Omega$ strongly depends on the $\beta_2$ deformation parameter.

The coupling of the angular momenta in odd-odd nuclei has been studied by Nordheim as well as by Gallagher and Moszkowski [28]. It has been pointed out that the ground state spin of a number of odd-odd nuclei could be obtained using $j$-$j$ coupling combined with certain rules (known as strong rules):

$$I = j_p + j_n \quad \text{if} \quad j_p = l_p \pm \frac{1}{2} \quad \text{and} \quad j_n = l_n \pm \frac{1}{2},$$

$$I = |j_p - j_n| \quad \text{if} \quad j_p = l_p \pm \frac{1}{2} \quad \text{and} \quad j_n = l_n \mp \frac{1}{2},$$

(2.48)

where $j_p$ and $l_p$ (or $j_n$ and $l_n$) represent the total and orbital angular momenta of the odd proton (or odd neutron). The spin of the ground state in an odd-odd...
deformed nucleus with strongly coupled nucleons is obtained using the same rules which apply to the projections of single particle angular momenta on the symmetry axis $\Omega$ and $\Lambda$ instead of $j$ and $l$. If both the unpaired nucleons are rotational aligned (high $i_x$, low $\Omega$), it is the total spin, $j$, which contributes in the coupling instead of the projection $\Omega$. The most challenging case for a spin assignment was in this thesis for the rotational band in the $^{170}$Ir nucleus. Taking into account the theoretically obtained elongation deformation parameter, $\beta_2 = 0.15$, the proposed configuration for the odd proton and neutron was $\pi h_{11/2} \otimes \nu i_{13/2} \left((\pi h_{11/2}[505] \otimes \nu(1/2[660])) \right)$. The ground state spin resulting from the coupling of a rotational aligned neutron and a symmetry aligned proton is suggested to be $I = 9$. This value is the closest integer to 8.83 which is obtained by triangular coupling of the rotation aligned neutron with $\Omega = 0.5$, $i_x = 6.48$, and the deformation aligned proton with $\Omega = 5.5$, $i_x = 0.0$. 
Chapter 3

Experimental Details

3.1 Experimental Motivation

In neutron-deficient nuclei close to the proton dripline in the mass $A \approx 100$ region, proton-neutron correlations are believed to play an important role when valence neutrons and protons occupy identical orbitals. Of particular interest is the doubly magic nucleus $^{100}\text{Sn}$, the heaviest $N = Z$ nucleus believed to be bound. Although the excited level structure of $^{100}\text{Sn}$ is not accessible with the technology available today, important pieces of information can be obtained by studying excited states of its neighbours. Information on the structure of these nuclei will provide stringent tests of the nuclear shell model in this exotic region of the nuclear chart. The $^{106,107}\text{Te}$ nuclei which are studied in the present work have only a few protons and neutrons outside the doubly closed shell. Correlations between protons and neutrons near the Fermi level have been a topic of great interest in recent years. Theoretical models predict that isoscalar ($T = 0$) neutron-proton pair correlations may become important close to the $N = Z$ line, leading to new nuclear structure effects. The enhanced neutron-proton correlation near $N = Z$ may increase collectivity, since the neutrons and protons may form a more strongly correlated system facilitating coherent collective motion.

A strong astrophysical motivation to study the structure of $^{106,107}\text{Te}$ has emerged in recent years. The rp process is a rapid proton capture process followed by slower $\beta^+$ decay. Via the rp process a number of proton-rich isotopes are believed to be created. The process has to occur in very high temperature environments (above $10^9$ Kelvin) so that the protons can overcome the large Coulomb barrier for charged particle reactions. A hydrogen rich environment is also a prerequisite due to the large proton flux needed. For instance a binary system consisting of a compact object (neutron star) and a red giant can be a suitable environment for the rp process. The end point of the rp-process (the highest mass element it can create) is not yet well established, but Schatz et al. [29] have predicted that the rp process terminates at nuclei with $Z \leq 52$. This is due to alpha decay which lead to a closed
Sn-Sb-Te cycle, e.g., $^{105}$Sn(p,γ)$\rightarrow^{106}$Sb(p,γ)$\rightarrow^{107}$Te(γ,α)$\rightarrow^{103}$Sn. The rp process extends beyond Sn if the corresponding A+1 Sb isotope is proton bound and the (γ, p) photodisintegration cross section is small. The Sn-Sb-Te cycle occurs because the $^{106-108}$Te nuclei are not alpha bound and the (γ, α) photodisintegration cross section is large. The Sn-Sb-Te cycle includes (γ, α) reactions on $^{106}$Te and $^{107}$Te and is therefore sensitive to the level structures of these nuclei. Experimental information on the low-lying energy spectra of the extremely neutron deficient tellurium isotopes $^{106,107}$Te, which are candidates for end point elements in the rp process, is of importance for the production rates of proton-rich elements up to A=107.

The level structure of odd-odd nuclei is generally less well known than that of neighbouring even-even and odd-even nuclei. Due to existence of an odd proton and an odd neutron many possibilities of couplings are available which reflect complex structures for such nuclei. Extension of databases on odd-odd nuclei is needed in the parametrisation of residual n-p interactions in the modelling of the level structure of nuclei.

This thesis reports on the identification of the nuclei $^{106,107}$Te and $^{110}$Xe as well as two doubly odd nuclei $^{170}$Ir and $^{172}$Au.

### 3.2 Experimental Methods

The exotic nuclei studied in this work are created at high spin and high excitation energies through fusion-evaporation reactions. The ion beams were accelerated by the Jyväskylä K130 cyclotron and were used to bombard a target to create a compound nucleus. On the shortest timescales the compound nucleus primarily deexcites by evaporating particles. Different final nuclei are created corresponding to different types and number of evaporated particles. The created nuclei will then deexcite via emission of gamma rays towards their ground states. These gamma rays are detected promptly and used to investigate the excited states in the nucleus of interest. There are different ways of identifying the created nuclei. One such method is the recoil decay tagging (RDT) method which will be described later in this chapter.

#### 3.2.1 Fusion-Evaporation Reactions

The first step in these experiments, as mentioned, is population of the isotopes of interest via fusion-evaporation reactions. Depending on the impact parameter value, $b$, which is illustrated schematically in Fig. 3.1, different types of nuclear reactions can occur. Table 3.1 shows the list of beams, targets and beam energies in the experiments which are described in this thesis. The K130 cyclotron of the Jyväskylä University Accelerator Laboratory was used to accelerate the ion beams.

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1The JYFL-accelerator is a versatile, sector-focused cyclotron that can accelerate ions from hydrogen to xenon up to energies (E/A) of 130 ($q^2/A$) MeV.
3.2. EXPERIMENTAL METHODS

Figure 3.1: Heavy-ion reaction types in terms of the impact parameter $b$.

listed in Table 3.1. Thin target foils made of highly enriched materials, listed in Table 3.1, were bombarded by these ion beams to produce the nuclei of interest.

Table 3.1: Reaction specifications for the experiments performed in this work. The reaction channel is characterized by the number and type of particles which must be evaporated to populate the channel of interest.

<table>
<thead>
<tr>
<th>Beam ion</th>
<th>Beam Energy (MeV)</th>
<th>Target isotope</th>
<th>Target thickness (mg/cm²)</th>
<th>Compound nucleus</th>
<th>Reaction channel</th>
<th>Nucleus under study</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{52}$Cr</td>
<td>187</td>
<td>$^{58}$Ni</td>
<td>0.640</td>
<td>$^{110}$Te</td>
<td>3n</td>
<td>$^{107}$Te</td>
</tr>
<tr>
<td>$^{54}$Fe</td>
<td>182</td>
<td>$^{54}$Fe</td>
<td>1.1</td>
<td>$^{108}$Te</td>
<td>2n</td>
<td>$^{106}$Te</td>
</tr>
<tr>
<td>$^{54}$Fe</td>
<td>195</td>
<td>$^{58}$Ni</td>
<td>1.0</td>
<td>$^{112}$Xe</td>
<td>2n</td>
<td>$^{110}$Xe</td>
</tr>
<tr>
<td>$^{60}$Ni</td>
<td>266</td>
<td>$^{112}$Sn</td>
<td>0.8</td>
<td>$^{172}$Pt</td>
<td>pn</td>
<td>$^{170}$Ir</td>
</tr>
<tr>
<td>$^{96}$Ru</td>
<td>350</td>
<td>$^{78}$Kr</td>
<td>0.56</td>
<td>$^{174}$Hg</td>
<td>pn</td>
<td>$^{172}$Au</td>
</tr>
</tbody>
</table>

Several conditions must be fulfilled in order for a compound nucleus to be formed. First, the beam particles must have enough energy to overcome the Coulomb barrier of the beam-target system. The height of the Coulomb barrier in the center-of-mass frame is approximately given by [30]:

$$E_{CB}(\text{MeV}) = \frac{1.44Z_1Z_2}{1.36(A_1^{1/3} + A_2^{1/3}) + 0.5}$$  \hspace{1cm} (3.1)

where $Z_1$, $A_1$, $Z_2$ and $A_2$ are the atomic numbers and masses of the beam and target nuclei, respectively. Second, the transferred angular momentum should not be too large, in order to avoid overly rapid rotation which causes the centrifugal
repulsion to overcome the short-range attractive nuclear force, leading to fission. A compound nucleus has an excitation energy \( E^* \) that depends on the \( Q \) value of the reaction that produced it, on the beam energy and on the masses of the beam and the target nuclei following the equation below

\[
E^*(\text{MeV}) = \frac{A_2}{A_2 + A_1}E_{\text{lab}} + Q.
\] (3.2)

A typical excitation energy for the produced compound nuclei is 40 MeV and the maximum angular momentum is typically around 50\( \hbar \). The compound nucleus starts the deexcitation by evaporating light particles. Each particle removes a significant part of the excitation energy from the system, but only a small amount of angular momentum. The daughter nucleus is still relatively highly excited, and decays first by emission of statistical gamma rays. These are usually high energy dipole transitions which cannot be resolved in gamma ray spectroscopic studies of today. These gamma rays remove a large amount of excitation energy but small amounts of angular momentum. When the nucleus approaches the yrast line, emission of gamma rays continue to the ground state. The yrast line is defined as a line connecting the lowest possible states for every given angular momentum in a nucleus. Figure 3.2 displays schematically the stages from the creation of a compound nucleus to the ground state together with the associated time scales.

### 3.2.2 Detection of Gamma-ray Radiation with Jurogam

Photoelectric absorption, Compton scattering and pair production are the three main types of interactions which occur when gamma radiation of energies typical for nuclear deexcitations enter the detector. When a gamma-ray enters the germanium crystal of a semiconductor detector, it may interact with an atomic electron of the detector material transferring all or part of its energy to it. The resulting fast electron deposits its energy by creating electron-hole pairs through ionisation and multiple scattering. Applying an external electric field, the free electron-hole pairs are collected to create an electrical pulse which can be processed by electronics. The amplitude of this pulse is proportional to the number of released electron-hole pairs and consequently to the energy deposited by the incident photon. The semiconductor detectors used in nuclear physics function using this basic principle.

Following the population of the nuclei of interest, via evaporation of neutrons, protons or alpha particles, the prompt gamma rays are emitted in the decay from highly excited states towards the ground state. The prompt gamma rays were detected using the germanium detector array Jurogam [32] placed around the target. Jurogam consists of 43 Compton-suppressed HPGe detectors. These are distributed over six rings around the target chamber with five detectors at 158° relative to the beam direction, ten at 134°, ten at 108°, five at 94°, five at 86° and eight at 72°. The total photo-peak efficiency of the Jurogam array is 4.2% at 1.3 MeV. In Fig. 3.3 a photo of the Jurogam array is shown. Gamma rays Compton scattering out of the Ge detectors (and thus only depositing part of their energy in the detector material)
3.3 The Recoil-Decay Tagging Technique

Recoil-decay tagging (RDT) is an elegant and powerful technique which can be used for in-beam studies of excited states in heavy and proton-rich nuclei [33, 34]. For this technique to be useful, it must be possible to identify the reaction products by detecting their emitted decay particles. In an RDT experiment using a fusion-evaporation reaction the compound nuclei are produced in highly excited states. These rapidly deexcite through particle evaporation (on a time scale of \( \approx 10^{-18} \text{s} \)) and emitting statistical gamma rays. Depending on the number and

Figure 3.2: Schematic illustration of the decay of a compound nucleus. In the shown example three particles were evaporated. The dashed line corresponds to the particle evaporation threshold. The figure is adapted from reference [31].

can be identified using Compton suppression shields surrounding the germanium detectors. These are made of a scintillation material with high atomic number which leads to a high detection efficiency, in this case bismuth germanate (BGO). A gamma-ray will be disregarded if a signal is recorded simultaneously in the Ge detector and in the surrounding BGO shield. The Compton suppression shields are protected in the front by heavy metal shields to avoid direct hits by gamma rays from the reactions, which would produce false veto signals.
types of evaporated particles, different nuclei will be produced. These nuclei decay further toward their ground states by emitting gamma rays. Such prompt gamma-ray emission usually happens within a time range of the order of picoseconds up to a few nanoseconds. Taking into account the typical reaction product velocity in fusion-evaporation reactions, this means that the gamma rays are emitted close to the reaction point.

The experiments described in this work take advantage of the existence of an “island” of alpha-emitting isotopes close to the proton dripline at $A \geq 100$ as well as alpha-emitters around $A \approx 170$. The high selectivity of the RDT technique enables gamma-ray spectroscopic studies of some nuclei which would be inaccessible with conventional in-beam methods. Promptly emitted gamma rays were detected by germanium detectors which were distributed around the target chamber. A thin target foil (see Table 3.1) is used in order to avoid stopping the reaction products inside the target. The reaction products are separated from scattered beam ions by the gas-filled recoil separator RITU (Recoil Ion Transport Unit) [35, 36]. The separation is based on the difference in rigidity between the reaction products and the scattered beam when passing through a magnetic field, this will be described in section (3.3.1). A carbon charge reset foil of $40 \mu g/cm^2$ thickness was placed immediately downstream of the target to equilibrate the charge distributions of the recoiling fusion products and the beam particles. The recoiling fusion products pass through RITU and are implanted into one of the two double-sided silicon strip detector (DSSDs) of the GREAT spectrometer at the focal plane of RITU. The
unstable implanted reaction products decay inside the DSSDs with time distributions that depend on their half-lives. Signals from all detectors are recorded with a “time stamp”, and spatial and temporal correlations between the reaction products and their corresponding decays are analysed off-line. Using both energy and time correlations, each given alpha particle or proton decay energy uniquely identifies its mother nucleus. Since nuclear masses increase approximately parabolically as a function of N-Z along an isobaric chain, the alpha-decay energies increase with increasing neutron deficiency. New exotic isotopes are therefore generally characterised by larger alpha particle or proton decay energies and also by much lower production cross sections than previously known isotopes. In the present work, however, we have relied on earlier measurements of the alpha-decay energies of $^{106,107}\text{Te},^{110}\text{Xe},^{170}\text{Ir}$ and $^{172}\text{Au}$.

Since all events are available with synchronised time information from a 100 MHz clock it is possible to correlate the prompt gamma rays corresponding to each individual identified reaction product with decays occurring up to seconds later. In the set-up for an RDT experiment some additional equipment such as a multi-wire proportional avalanche counter (MWPAC), Si PIN photo-diodes, planar and Clover-type Ge detectors are placed near the RITU focal plane. Their applications will be explained in detail in the following sections. Figure 3.4 shows a schematic picture of the experimental set-up.

Other tagging methods such as tagging on electron, gamma rays following isomeric states, protons and even $\beta$-decays with high end-point energy [37] have allowed state-of-the-art research to be performed on nuclei with very low production cross sections.

### 3.3.1 The Gas-Filled Recoil Separator

The gas-filled recoil separator RITU [35, 36] consists of four strong magnets in a QDQQ configuration, where Q is a quadrupole magnet and D is a dipole magnet. Its vacuum chamber is filled with helium gas at a pressure that is typically around 1 mbar. The gas-filled region is separated from the high-vacuum region upstream of the target chamber by a powerful “differential” pumping system. The gas inside the RITU chamber decreases the angular range over which the charged ions are distributed by means of charge equilibration. This increases the transmission of the reaction products (around 35%) but a disadvantage is that there is little distinction between different masses and charge states. Figure 3.5 compares schematically vacuum and gas-filled recoil separators. In Fig. 3.6 a schematic picture of RITU in the central position of the experimental set-up is displayed.

The separation of beam particles from reaction products is based on the deflection of the moving charged particles by a magnetic field of strength $B$. This can be discussed in terms of rigidity ($R = Br$) where $r$ is the radius of the ion trajectory.
CHAPTER 3. EXPERIMENTAL DETAILS

Figure 3.4: A schematic picture of the experimental setup consisting of the germanium detector array (Jurogam), the recoil separator (RITU), and the GREAT system, which is placed at the focal plane of RITU.

Figure 3.5: Schematic illustration of two different types of recoil separators. The shaded part indicates the region filled with gas. The circle shows the region where there is magnetic field. The magnetic field is pointing perpendicular out of the page. The ions travel from left to right.
3.3. THE RECOIL-DECAY TAGGING TECHNIQUE

in the magnetic field. The rigidity is given by [39]:

\[
R = \frac{m_0 v}{e q_{av}} \approx \frac{m_0 v}{e v_0} Z^{1/3} = \frac{0.0227 A}{Z^{1/3}} [Tm] \tag{3.3}
\]

where \( v \) is the velocity of ion trajectory, \( v_0 \) (the Bohr velocity) is \( c/137 \), \( m_0 \) is the rest mass, \( q_{av} \) is the average charge state and \( Z \) is the atomic number of the ion. The maximum rigidity of RITU is 2.2 Tm. There is usually a large difference in rigidity between beam particles and fusion products since the latter generally have significantly lower velocities, but the heavier the beam particles are, the smaller is this difference. For studies of heavy neutron deficient nuclei the beam-target mass asymmetry is normally quite large and a good separation is easily achieved. However, some nuclei studied in this work can only be reached in near-symmetric fusion-evaporation reactions, leading to difficulties in separating beam and reaction products. In this case the separation is based on measurements on time of flight between the MWPAC and the DSSDs as well as the deposited energy in the MWPAC which will be described in detail in the following sections.

3.3.2 The GREAT Spectrometer

The GREAT (Gamma Recoil Electron Alpha Tagging) spectrometer is placed in conjunction with the recoil separator at the focal plan of RITU. GREAT consists of several distinct components: a multi-wire proportional avalanche counter...
(MWPAC), two double-sided silicon strip detectors (DSSD), a double-sided planar germanium strip detector, a high-efficiency segmented germanium Clover-type detector and an array of silicon PIN photo-diode detectors. In Fig. 3.7 a schematic drawing of the GREAT spectrometer is shown.

The MWPAC is placed immediately in front of the DSSDs, and the reaction products pass through it before being implanted into the DSSDs. The MWPAC signal acts as a start time for time-of-flight measurements between the MWPAC and the DSSDs and also provides information on the deposited energy. This information enables us to distinguish the reaction products from scattered beam particles and from the emitted alpha particles or protons. Fig. 3.8 shows a plot with the time-of-flight information on the x-axis and the deposited energy on the y-axis. The reaction products have larger mass than the beam ions, since the linear momentum is conserved they are slower than beam ions. They also have higher atomic number which causes them deposit to more energy in the MWPAC than the beam ions. Thus, it is possible to identify and select events that correspond to fusion-evaporation products from the distribution displayed in Fig. 3.8 and to avoid background arising from implantation of scattered beam in the DSSDs. The PIN detectors, located upstreams from the DSSD, can be used to measure conversion electron energies and to detect alpha particles which have escaped from the DSSDs. The germanium planar detector can be used to detect delayed low-energy gamma rays originating from decay of isomeric states or following alpha or proton decay of the residual nuclei. The clover germanium detector can be used to detect such gamma rays with energies up to several MeV.

![Figure 3.7: Schematic drawing of the GREAT spectrometer. The recoils travel in the positive z direction before implantation. The figure is adapted from [40].](image-url)
3.4 The TOT AL DA T A READOUT (TDR) SYSTEM

In an RDT experiment the time delay between the detection of gamma rays by the germanium detectors at the target position and the charged particle decay by the Si detectors at the focal plane plays an important role. This time delay is equal to or greater than the flight time of the fusion-evaporation products from the target to the focal plane. The total data readout (TDR) [41] system, which is a trigger-less data acquisition system, reads out all channels originating from GREAT independently with no time losses. The collected data are associated with a time stamp generated by a 100 MHz clock, which gives an accuracy of 10 ns. The data from the Jurogam array are collected within a given time interval (5µs before) relative to the events in GREAT.

The events can then be reconstructed in the software. The Grain software package [42] and the Tscan software [43] were used in order to analyse data on-line and off-line. Figure 3.9 shows a time spectrum in the trigger setup. In the software every signal in the DSSDs can be considered as a starting time for an event. It is possible to identify the time interval for processing an event. In an example shown

Figure 3.8: Plot of the deposited energy in the MWPAC versus time of flight between the MWPAC and the DSSDs. A two-dimensional gate is used to select the distribution of the reaction products. Scattered beam particles appeared below and to the right of the recoil distribution. Time increases from right to left.

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Figure 3.9: Time histogram of the events in the Jurogam with respect to the event in GREAT are plotted. The spectrum consists of peaks corresponding to events in Jurogam, MWPAC, DSSDs, planar and Clover from left to right, respectively. For instance, in this figure all signals during a time interval \(1 \mu s\) before and \(2 \mu s\) after a signal in DSSD are processed as one event by the software.

in Fig. 3.9 the collected data from \(1 \mu s\) before an event in DSSDs, until \(2 \mu s\) after are processed within an event. Since all signals are collected with a time stamp there is no limitation in choosing the trigger width and delay interval, except that the trigger delay can not be larger that \(\pm 5 \mu s\).

3.5 Cross Section Estimates

The reaction cross section, \(\sigma\), represents the likelihood for a particular reaction channel to be populated in a fusion-evaporation reaction experiment. It can be experimentally estimated using the following formula,

\[
\sigma = \frac{R}{I_b N},
\]

where \(I_b\) is the beam intensity in the unit of particles per second. \(N\) is the number of nuclei per unit of area in the target\(^2\), and \(R\) is the observed production rate of the

\(^2\)\(N = \frac{N_A d}{\rho}\). \(N_A\) is the Avogadro’s number, \(d\) is the target thickness and \(\rho\) is the intensity of the target material. Therefore \(\rho d\) is the target mass thickness which usually is the given property of a target. \(A\) is the mass number of the target material.
3.5. **CROSS SECTION ESTIMATES**

nuclei of interest. $R$ can be estimated using the rate of the detected alpha/proton decays of the corresponding nucleus. When estimating $R$ other factors such as transmission efficiency through RITU, detection efficiency in the DSSDs and the alpha decay branching ratio of the nucleus under study must be taken into account. The transmission efficiency of the RITU is around 35% and the charged particle detection efficiency of the DSSDs is about 55%. In the experiment which was aimed at studying the $^{106}$Te nucleus with 100% alpha decay branching ratio for the ground state decay, the average yield of 14 $\alpha/h$ was obtained during the experiment (without correction for efficiencies). Therefore, using 10 pnA beam intensity and a target thickness of 1.1 $mg/cm^2$ and considering the efficiency corrections, the cross section was estimated to be 25 nb. This was the lowest cross section ever reported for an in-beam gamma-ray spectroscopic study.

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31 particle nano Ampere (pnA) = \( \frac{10^{-9}}{1.6 \times 10^{-19}} \) particle/s
Chapter 4

Data Analysis

4.1 Calibration

The first step in the data analysis is to perform accurate calibrations of energy and time spectra. The energy spectra from different detectors are aligned by translating the ADC (Analog to Digital Converter) channel numbers to energies using known sources. The gamma ray sources used in the experiments described in this thesis are $^{152}$Eu and $^{133}$Ba. The DSSDs are calibrated using an alpha triple source consisting of the three isotopes $^{244}$Cm, $^{241}$Am and $^{239}$Pu.

When linear amplifiers are used the energy of the detected radiation can usually be expressed as a nearly linear polynomial function of the ADC channel number. If there is a more severe nonlinearity in the amplifier or the ADC a different function must be found. The ADCs used for the Jurogam germanium detectors have significant deviations from linearity for small signal amplitudes.

The signal corresponding to the detection of an alpha particle in the DSSDs gives the sum of the alpha particle energy and the kinetic energy of the recoiling daughter nucleus. Therefore, applying the calibration coefficients obtained using the triple source in an in-beam experiment give energy values for the alpha particles which are higher than the real alpha particle energies. One way of solving this problem is applying an internal calibration using known alpha decays in the populated region.

4.2 Doppler correction

The 43 detectors of the Jurogam Ge detector array are placed at different angles with respect to the beam direction. Since the gamma rays are emitted from recoiling nuclei moving at high velocity, the detected gamma rays are Doppler shifted. The Doppler shift depends on the angle at which the detector is placed relative to the recoil direction (which is approximately the same as the beam direction) and on the speed of the recoiling fusion product. The detected gamma-ray energy must be
corrected using the formula

$$E_\gamma = E_{\gamma D} \frac{1 - \beta \cos \theta}{\sqrt{1 - \beta^2}}$$

(4.1)

where $E_{\gamma D}$ is the measured energy of a gamma transition, $E_\gamma$ is its energy value in the rest frame, $\beta = v/c$, and $\theta$ is the angle at which the detector is placed, measured relative to the beam direction. In Fig. 4.2 gamma-ray spectra with and without applied Doppler correction for comparison have been displayed.

Figure 4.1: In the upper panel the total gamma ray spectrum (sum of the data from detectors placed at six different rings), from the reaction $^{96}\text{Ru}(^{78}\text{Kr},pn)^{172}\text{Au}^*$, without Doppler correction is shown. The gamma rays in the spectrum in the bottom plane have been Doppler corrected using the average velocity of the reaction products and the angle between the beam direction and the direction in which the gamma ray was detected. The value of $\beta = v/c$ for this reaction was measured to be 0.0396.

The Doppler effects also lead to uncertainties in the measured gamma-ray energies. One contribution to this “Doppler broadening” comes from the opening angles of the Ge detectors. Differentiation of eq.(4.1) with respect to $\theta$ shows that the Doppler broadening is approximately proportional to $\beta \sin \theta$, i.e., it is worst at $\theta = 90^\circ$. Another contribution comes from the straggling of the reaction products in the target, leading to a spread in the reaction product directions. A kick from the evaporated particles can also make a contribution to this effect which is most severe when evaporating alpha particles.
4.3 Reaction Products Identification Using Correlations

As discussed briefly in section (3.3.2) the signature for implantation of a reaction product in the DSSDs is the presence of a signal in both the MWPAC and the DSSD within a time interval corresponding to the flight time of a reaction product between the MWPAC and the DSSD. When an alpha particle is detected, a signal will only be produced by the DSSD and not in the MWPAC. Figure 4.2 illustrates

![Diagram of spectra](image)

Figure 4.2: The upper panel shows a raw spectrum obtained from the DSSD using the reaction $^{96}$Ru+$^{78}$Kr $\rightarrow ^{174}$Hg. The middle panel shows the alpha-like spectrum from the DSSD. It corresponds to the spectrum in the upper panel except events which are vetoed by the MWPAC. The bottom panel shows the corresponding alpha-like spectrum in the middle panel which is obtained by adding an additional condition that the detection of a reaction product must be followed by an alpha-like particle event in the same DSSD pixel within a time difference up to 7 ms. The alpha peak corresponding to alpha particles which are emitted from states with half-lives of the order of several ms are enhanced due to the condition of the 7 ms correlation time.

the spectra obtained by the DSSD. The upper panel shows the “raw” spectrum (everything detected by the DSSD), while the middle panel shows “alpha-like”
events that give rise to no signal in the MWPAC. The punched through particles\(^1\) and the primary beam as well as reaction products are significantly suppressed after the process of veto by the MWPAC which is shown in the middle part of Fig. 4.2. The bottom panel shows the alpha-like events which were preceded by a recoil implantation in the same pixel of the DSSD. In this spectrum the time differences between detection of the reaction product and its corresponding alpha decay was required to be within 7 ms. After applying an MWPAC veto and a correlation condition as shown in the bottom panel it is only the correlated protons and alphas and escaped particles that are visible in this spectrum.

The detected gamma rays are assigned to specific reaction channels by associating them with the corresponding implanted recoil nuclei and their subsequent decay. In an RDT experiment the identification of the recoil is, as discussed in section (3.3), done by observing its decay particle. It is done by requiring time and position correlations between an implanted recoil and its particle decay. The spacial and temporal correlations are found by searching for a recoil implantation followed within a suitable search time by an alpha decay with a characteristic energy detected in the same DSSD pixel. If the half-life of the nucleus of interest is known the search time is usually taken to be around three half-lives. However, care must be taken to avoid random background events if the recoil rate is large and the alpha decay half-life of the channel of interest is long.

In fusion evaporation reactions aimed at studying nuclei far from stability the reaction channels of interest are often hidden in a large background arising from strong fusion reaction channels, incomplete fusion, fission and Coulomb excitation. Sensitive selection criteria like in the recoil decay tagging technique therefore are important in order to separate the gamma rays of a reaction channel with low production cross section from the background. Figure 4.3 shows gamma-ray spectra from the \(^{106}\)Te experiment in order to illustrate the high selectivity that has been achieved here.

### 4.4 Half-Life Measurements

The time information of the recorded data can be used to measure the half-lives of nuclear states that decay by emitting alpha particles or protons. The half-lives of the alpha-emitting states in the nuclei under study are measured. This is done by creating histograms of the time differences between the detection of a recoil implantation and its subsequent alpha decay in the same pixel of the DSSDs. These histograms reflect the Poisson distribution of the decay of the implanted recoils.

One way of extracting the half-life is by fitting two exponential functions to the data points:

\[
N = A \exp(-\lambda t) + B \exp(-rt) .
\]

\(^1\)Promptly evaporated particles from the compound nucleus which are created at the target position passing through RITU and MWPAC.
4.4. HALF-LIFE MEASUREMENTS

In the bottom panel is a spectrum of gamma rays associated with a recoil implantation which was followed by an alpha-like decay event in the same pixel. The upper panel requires the alpha decay to have the characteristic decay energy of $^{106}$Te. The cross section for population of $^{106}$Te, which in this reaction was estimated to be 25 nb, compared to the total fusion cross section of about 250 mb, gives a selectivity of the order of $10^{-7}$. This selectivity is at the limit of what is achievable with in-beam gamma-ray spectroscopy using today’s technology.

![Gamma-ray spectrum](image)

Figure 4.3: The top spectrum is tagged on a recoil followed by a $^{106}$Te alpha particle within 200$\mu$s of search time. The bottom panel has a spectrum that is recoil alpha tagged (any alpha particle).

The parameter $r$ in eq.(4.2) is related to the rate of random background events and can be estimated from the recoil rate per pixel [44] during the experiment or determined from the result of a fit to the data. $\lambda'$ can also be obtained from the fitting. These parameters can be used to calculate the half-life of the decay by

$$t_{1/2} = \frac{\ln 2}{\lambda' - r}.$$  (4.3)

If there is no random background in the time histogram a simpler function of

$$N = A \exp(-\lambda t),$$  (4.4)

can be used to obtain constant decay, $\lambda$, directly. The half-life then is obtained by

$$t_{1/2} = \frac{\ln 2}{\lambda}.$$  (4.5)
Another method is the so-called Maximum Likelihood method, which is useful in cases of low statistics and for time spectra with no random background contribution [45]. Considering \( n \) decays at times \( t_1, t_2, ..., t_n \leq T \), the probability of observing a decay between time \( t \) and \( t + dt \) is given by

\[
P(\lambda, t)dt = \lambda e^{-\lambda t} dt (1 - e^{-\lambda T})^{-1}
\]  

(4.6)

So, the probability of observing the decays at time \( t_1, t_2, ..., t_n \) is

\[
P(\lambda, t_1, ..., t_n = \prod_{i=1}^{n} P(\lambda, t_i)
\]  

(4.7)

The decay constant can be obtained by finding the value of \( \lambda \) which has a maximum probability via the following procedure:

\[
\frac{\partial \log P(\lambda, t_1, ..., t_n)}{\partial \lambda} = 0
\]  

(4.8)

which gives

\[
\frac{1}{\lambda} = \frac{1}{n} \sum_{i=1}^{n} (t_i) + \frac{T}{e^{\lambda T} - 1}
\]  

(4.9)

where \( \lambda \) is the decay constant and \( \frac{1}{\lambda} = \tau \), which is the life time. The life time is obtained by iteration until the eq.(4.9) converges. The upper and lower uncertainty limits for the life time, corresponding to the standard errors, can be approximated as [46]

\[
\tau_u \approx \frac{\tau}{1 - \frac{1}{\sqrt{n}}}, \quad \tau_l \approx \frac{\tau}{1 + \frac{1}{\sqrt{n}}}
\]  

(4.10)

For a large value of \( n \) the error bars become more symmetric and can be approximated as

\[
\tau_u \approx \tau (1 + \frac{1}{\sqrt{n}}), \quad \tau_l \approx \tau (1 - \frac{1}{\sqrt{n}}).
\]  

(4.11)

The half-life can then been obtained from

\[
t_{1/2} = \tau \ln 2.
\]  

(4.12)

The half-lives measured in this thesis have mainly been obtained using the Maximum Likelihood method. Figure 4.4 shows the time spectrum for the decay of the ground state in \(^{172}\text{Au}\). The half-life is deduced to be 23±2 ms.
4.5 Gamma-ray Coincidence Measurements

In a gamma-ray decay cascade from a highly excited state to the ground state there are a large number of gamma rays which are emitted close in time (within ps → ns). Gamma rays detected in different detectors within a short time interval are said to be in coincidence and are interpreted to belong to the same gamma-ray cascade. These coincident events can be recorded and analysed off-line in order to reconstruct the level structure of the nucleus under study. In this work the gamma-ray coincidence analysis has mainly been used to study excited states in \(^{170}\text{Ir}\) \[20\]. In Fig. 4.5 are shown spectra of gamma rays detected in coincidence with the 267 keV and 341 KeV gamma rays. In the right part of Fig. 4.5 a partial level scheme of \(^{170}\text{Ir}\) is shown for comparison.

4.6 B(M1)/B(E2) Ratios

Using the experimental data it is possible to extract the ratio between the reduced transition probabilities of magnetic dipole and electric quadrupole transitions originating from the same state \[30\]. This is done by measuring the energies and intensities of the transitions together with the E2/M1 mixing ratio for the dipole transition. The ratio of reduced transition probabilities can be expressed as

\[
\frac{B(M1; I \rightarrow I - 1)}{B(E2; I \rightarrow I - 2)} = \frac{0.697 [E_{\gamma}(E2; I \rightarrow I - 2)]^{0.5} I_{\gamma}(M1)}{[E_{\gamma}(M1; I \rightarrow I - 1)]^{1.5} I_{\gamma}(E2) (1 + \delta_{E2/M1})}
\]

(4.13)

where \(E_{\gamma}\) is the energy of a given gamma-ray transition, \(I_{\gamma}\) is the intensity and \(\delta\) is the multipole mixing ratio of the \(I \rightarrow I - 1\) transitions (in our calculation the mixing ratio was assumed to be zero).

When \(E_{\gamma}\) is measured in MeV the unit of the B(M1)/B(E2) ratio is \((\mu N/eb)^2\). Using the semi-classical formalism of Dönau-Frauendorf \[47\] the B(M1)/B(E2) ratio

Figure 4.4: The time spectrum for the alpha decay of the ground state of \(^{172}\text{Au}\). A half-life 23(2) ms was deduced for the alpha particle emission.
Figure 4.5: In the left-hand part of the figure are shown the background subtracted spectra of gamma rays detected in coincidence with the 267 and 314 keV gamma rays which both have been assigned to $^{170}\text{Ir}$. The right-hand part of the figure shows a partial level scheme of $^{170}\text{Ir}$ deduced in the present work.

can be calculated. The full theoretical expression is written as

$$
\frac{B(M1; I \rightarrow I - 1)}{B(E2; I \rightarrow I - 2)} = \frac{12}{5Q_0^2 \cos^2(\gamma + 30^\circ)} \left[ 1 - \frac{K^2}{(I - 2)^2} \right]^{-2} \times
$$

$$
\left[ (1 - \frac{K^2}{I^2})^{1/2} [k_1 (g_1 - g_R) (1 \pm \frac{\Delta e}{\hbar \omega}) + \sum_n k_n (g_n - g_R)] - \frac{K}{T} [ (g_1 - g_R) i_1 + \sum_n (g_n - g_R) i_n ] \right]^2
$$

where the quadrupole moment is

$$Q_0 = \frac{3}{\sqrt{5\pi}} R_{av}^2 Z \beta (1 + 0.16\beta).$$

$k_i$ are projections on the nuclear symmetry axis of the individual angular momenta of quasiparticles (section (2.2.4)) and $K$ is the total projection ($\sum k_i$). $g_i$ are the g-factors, and the $i_i$ are the alignments which are the projections of quasiparticle angular momenta on the axis perpendicular to the symmetry axis (section (2.2.4)). $g_R$ is the rotational g-factor which is set equal to $Z/A$. If there is signature splitting, the parameters with subscript 1 refer to the nucleon causing the signature splitting. The term $\frac{\Delta e}{\hbar \omega}$ is a signature dependent term which can be ignored if there is no signature splitting in the rotational band. Since the $B(M1)/B(E2)$ ratios
4.7 Gamma-ray Transition Probabilities

The B(E2) transition probabilities are related to quantities such as the life time of states, deformation parameters and quadrupole moments. According to nuclear global systematics, there is a strong correlation between the energy of the first excited 2\(^+\) state, \(E(2^+)\) in an even-even nucleus, and the corresponding B(E2;2\(^+_1 \rightarrow 0^+_1\)) values. The empirical relationship given by Raman [48] is

\[
B(E2; 2^+_1 \rightarrow 0^+_1) \approx 0.66E(2^+_1)^{-1}Z^2A^{-0.69},
\]

(4.16)

where \(E(2^+)\) is measured in keV and B(E2) in \(e^2b^2\). Information about the deformation parameter, \(\beta\), and the intrinsic quadrupole moment can be used to obtain the B(E2) value and vice versa [48].

As described in section (2.2.5), a collective excitation in a nucleus involves many nucleons. With many nucleons participating they can couple in many different ways in order to form e.g. the 2\(^+\) state. This consequently results in a higher transition probability and a large B(E2; 2\(^+_1 \rightarrow 0^+_1\)) value. Since the B(E2) values are a measure of nuclear collectivity, a plot of B(E2) values over an isotopic chain can be used to study the collectivity as a function of neutron number. The trend of the B(E2) values in the lightest Xe isotopes are calculated empirically and compared with the theoretical values [23]. In these calculations the values of deformation parameter and the quadrupole moment are taken from TRS calculations. In these plots the comparison is more relevant for trends rather than absolute B(E2) values.

4.8 Angular Distribution and Correlation Measurements

The angular distribution of gamma rays emitted from the reaction products can provide some information about the multipolarity of the transitions under study. If there is no mixing of different multipoles and a reasonable spin alignment of the reaction products is present, the angular distribution can to lowest order be written

\[
P(\theta) = 1 + \alpha_2 A_2 P_2(\cos \theta) + \alpha_4 A_4 P_4(\cos \theta),
\]

(4.17)

where \(P_2(\cos \theta)\) and \(P_4(\cos \theta)\) are the Legendre-polynomials. The values of the parameters \(A_2\) and \(A_4\) and the attenuation coefficients \(\alpha_2\) and \(\alpha_4\) are tabulated in [49]. The attenuation coefficients \(\alpha\) can also be extracted from a transition with known multipolarity. When using an array of detectors, it is possible to study how the
intensity of a particular transition is distributed over different angles. Comparing the experimentally obtained coefficients for the distribution function (4.17) with the theoretically calculated values [49] allows us to identify the multipolarity of the transition.

A rough approximation which can be made in case of low statistics is to measure and compare the intensity of the transitions in the “backward” or “forward” detectors (relative to the beam direction) with those at ninety degrees:

$$\frac{I_\gamma(180° \text{ or } 0°)}{I_\gamma(90°)} \quad (4.18)$$

In Fig. 4.6, angular distributions are shown for a dipole and quadrupole transition from an in-beam experiment. A gamma ray corresponding to a dipole transition is more likely to be emitted at 90° than 180° or 0°, while the converse is true for a quadrupole distribution. For typical values of $\alpha_2 A_2 = 0.3$ and $\alpha_4 A_4 = -0.1$ for a stretched quadrupole transition and $\alpha_2 A_2 = -0.2$ and $\alpha_4 A_4 = 0.0$ for a stretched dipole transition as discussed in reference [50] the ratio (4.18) using eq.(4.17) is calculated to be 1.47 for a quadrupole transition and 0.7 for a dipole transition.

Figure 4.6: Typical two dimensional angular distributions for quadrupole (left) and dipole (right) transitions. The beam direction is from the left to the right. The coefficients have been taken from reference [50].

If there is enough statistics, a traditional directional correlation of oriented state (DCO) ratio, $R_{DCO}$, measurement can be performed. It is defined as following:

$$\frac{L_1(\theta_1) \text{gated by } L_2(\theta_2)}{L_1(\theta_2) \text{gated by } L_2(\theta_1)} \quad (4.19)$$

If $\theta_1 = 157°$, $\theta_2 = 94°$, $\gamma_2$ is a quadrupole transition and, assuming the transitions are stretched (zero mixing ratio), complete alignment and the spin of initial state is 4, the calculated values are 1.0 if the $\gamma_1$ is a quadrupole transition and about 0.3 if the $\gamma_1$ is a dipole transition [52]. These angles are typical for the Jurogam detector array described in section (3.2.2). The calculated values assigned to a
quadrupole or dipole transition applying the described methods are extremely sen-
sitive to the assumptions and experimental conditions, therefore they can not be
taken as universal values. A more accurate technique to evaluate the DCO ratio
values is to examine them for several well known strong gamma-ray transitions in
each individual experiment.

The knowledge on the multipolarities of the transitions can be a great help for
spin and parity assignments of excited states.

4.9 Comparison Between Experimental Data and
Theoretical Total Routhian Surface Calculation

The cranked shell model can be used to make theoretical predictions for the level
energies in rotating nuclei. In order to compare the experimental energy levels with
the theoretical ones, the experimental data must be transferred from the laboratory
system to a rotating frame of reference. The total experimental Routhian (see 2.3)
for a transition from $I+1 \rightarrow I-1$ is defined as:

$$E'(I) = \frac{1}{2}[E(I+1) + E(I-1)] - \omega(I)I_x(I),$$

where the experimental rotational frequency, $\omega$, at the discrete intermediate spin
value $I$ is obtained from two adjacent states ($\Delta I = 2$) of the rotational band by
using the relationship below

$$\omega(I) = \frac{dE}{dI} \approx \frac{E(I+1) - E(I-1)}{I_x(I+1) - I_x(I-1)},$$

where the alignment, $I_x$, is defined as

$$I_x = \sqrt{(I + 1/2)^2 - K^2},$$

$K$ is the projection of the angular momentum on the symmetry axis. For a strongly
coupled band $K$ can be approximated by the spin of the band head in the rotational
band.

The experimental quasi-particle Routhian and the alignment are defined as:

$$i(\omega) = I_x(\omega) - I_{xp}(\omega)$$
$$e'(\omega) = E'(\omega) - E_{pr}'(\omega).$$

The $I_{xp}(\omega)$ and $E_{pr}'(\omega)$ are the values of a reference configuration. Details about
this subject and how to choose the reference function can be found in reference [53].
For instance the reference function corresponding to the ground state band can be
parameterised using the Harris polynomial:

$$I_{xp}(\omega) = \omega J_0 + \omega^3 J_1$$
$$E_{pr}'(\omega) = -\frac{\omega^2}{2} J_0 - \frac{\omega^4}{4} J_1 + \frac{\hbar^2}{8 J_0}$$

(4.24)
where the parameters, $J_0$ and $J_1$, can be obtained by plotting $J^{(1)}$, 

$$J^{(1)}(\omega) = \frac{I_{x \text{ yrast}}}{\omega},$$  

(4.25)  

as a function of $\omega^2$ and then fitted to a polynomial function, 

$$J^{(1)}(\omega) = J_0 + \omega^2 J_1.$$  

(4.26)  

In this thesis the alignment for the observed rotational band in $^{170}$Ir has been plotted and compared with the ground state band of $^{171}$Ir [20]. Lack of statistics in the present experiment allowed the rotational band in $^{170}$Ir to be observable only to relatively limited spin values. There was not enough data in order to deduce the Harris parameters for this rotational band. Thus, the Harris parameters used by Bark et al. [54] to extract the alignment in $^{171}$Ir were used for the $^{170}$Ir nucleus.
Chapter 5

Summary of Papers

This thesis describes experiments based on the recoil-decay tagging technique which were performed at the Accelerator Laboratory of the University of Jyväskylä, Finland. The extremely neutron deficient nuclei $^{107}\text{Te}$, $^{106}\text{Te}$ and $^{110}\text{Xe}$ were produced in fusion-evaporation reactions induced by $^{52}\text{Cr}$ and $^{54}\text{Fe}$ ion beams from the K130 cyclotron. The nuclei $^{170}\text{Ir}$, $^{172}\text{Au}$ were populated using $^{60}\text{Ni}$ and $^{78}\text{Kr}$ ion beams. Prompt gamma rays produced in the reactions were detected by the Jurogam gamma-ray spectrometer. The gamma rays belonging to the reaction channel of interest were selected based on the recoil identification provided by the RITU gas-filled recoil separator and the GREAT focal plane spectrometer. The recoil-decay tagging technique is suitable to probe the inner structure of these nuclei due to the emission of alpha particles with half-lives in the about 100 microsecond to second range.

5.1 Paper I

Gamma-ray transitions in the nucleus $^{107}\text{Te}$ have been identified for the first time. The reaction $^{58}\text{Ni}(^{54}\text{Cr},3\text{n})^{107}\text{Te}^*$ was used to populate excited states in $^{107}\text{Te}$. The production cross section was estimated to 1µb. For the first excited state at 90 keV, tentatively assigned to have spin and parity $7/2^+$, a $g_{7/2}$ character is proposed. This assignment is supported by the observation by Schardt et al. of two alpha energies at 3480(30) keV and 3580(30) keV associated with the decay of $^{111}\text{Xe}$ [55]. We suggest that a 631 keV line deexcites the 9/2$^+$ state to the first excited 7/2$^+$ state at 90 keV, whereas a 721 keV line deexcites the same state to the 5/2$^+$ ground state. The low excitation energy of the first excited state is interesting in relation to astrophysical models of the rp process, which is predicted to end at the nucleus $^{107}\text{Te}$ due to $(\gamma,\alpha)$ reactions. By comparing the excitation energy of this state with those of the corresponding states in the heavier Te isotopes and in the Sn isotopic chain we found different trends as N approaches Z. This points to increased influence from n-p correlations near N=Z.
5.2 Paper II

Paper II describes the identification of gamma-ray transitions in $^{106}$Te for the first time. The reaction $^{54}$Fe($^{54}$Fe,2n)$^{106}$Te$^*$ was used with an estimated production cross section of 25 nb, a new limit for in-beam gamma-ray spectroscopy. A ground-state band, tentatively extending up to $I^\pi = 10^+$, is proposed. The systematics of low-lying yrast states in the Te isotopes is discussed within the context of vibrational excitations and residual nucleon-nucleon interactions. Evidence for enhanced collectivity near N=Z due to neutron-proton correlations is discussed.

5.3 Paper III

The neutron-deficient nucleus $^{110}$Xe with a cross section of 50 nb was produced using the reaction $^{54}$Fe($^{58}$Ni,2n)$^{110}$Xe$^*$. Gamma ray transitions have been identified and the three lowest energy levels in the ground state band are proposed. The systematic trend of the first and second excited states ($2^+$ and $4^+$) for even-even Xe isotopes in the mass region of $110 \leq A \leq 136$ was investigated. It shows a decrease in the $2^+$ and $4^+$ energies levels of $^{110}$Xe, relative to the expected trend. An increase in the empirically obtained reduced transition probability, ($B(E2)$), for the $2^+$ state has been found which is opposite to the trend predicted by TRS calculations. These observation can point to an increase in collectivity which can be the result of enhanced proton-neutron interactions (proton-neutron pairing).

5.4 Paper IV

Excited states have been investigated in the highly neutron-deficient odd-odd nucleus $^{170}$Ir. The experiment was performed using the $^{112}$Sn($^{60}$Ni,pn)$^{170}$Ir$^*$ reaction. A level scheme of $^{170}$Ir is presented which partly includes a rotational band. The influence on the nuclear shape by the odd proton and neutron (and their coupling) is discussed. A significantly smaller signature splitting in the observed rotational band compared to the odd-even neighbour ($^{171}$Ir) was found, indicating a smaller degree of triaxility in the nuclear shape. The $B(M1)/B(E2)$ values for several possible configurations of the observed rotational band have been calculated and compared with the experimental values. New alpha decay branches are assigned to $^{170}$Ir using the “inverse tag” method. The alpha-decay fine structure is confirmed by the observation of $\alpha - \gamma$ coincidences at the focal plane of the recoil separator. The study of the alpha decay fine structure allowed us to deduce several low-lying excited states in $^{166}$Re, the alpha decay daughter of $^{170}$Ir.

5.5 Paper V

The fusion-evaporation reaction $^{96}$Ru($^{78}$Kr,pn)$^{172}$Au$^*$ was used to populate excited states in $^{172}$Au. Gamma-ray transitions in $^{172}$Au were identified for the first time.
The alpha decay from the ground state of $^{172}$Au has been observed. By requiring mother-daughter-granddaughter correlation, a new alpha decay branch from the isomeric state in $^{164}$Re, the granddaughter of $^{172}$Au, has been discovered. Possible configuration assignments for the $\alpha$-decaying states in $^{172}$Au are discussed in terms of systematics of the region and the results of TRS calculations.
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