$K$-isomerism at high-spin beyond the fusion limit

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Abstract

New high-$K$ isomers have been populated in $^{181}\text{Ta}$ and $^{186}\text{W}$ by bombarding thick targets with pulsed $^{238}\text{U}$ beams at 1600 MeV. The new inelastically excited multi-quasiparticle states include $K^\pi = 21/2^-$ and $29/2^-$ 3-quasiparticle isomers in $^{181}\text{Ta}$ and $K^\pi = 7^-$ and $(16^+)$ isomers in $^{186}\text{W}$. The $K^\pi = (16^+)$ state in $^{186}\text{W}$ has $t_{1/2} \geq 3$ ms and extends the 4-quasiparticle isomer systematics beyond the limit accessible with fusion-evaporation reactions, using stable beams and targets, for the first time. In addition, a $t_{1/2} \geq 1$ ms 3-quasiparticle isomeric state feeding a strongly coupled rotational band has been populated in $^{183}\text{Ta}$ by nucleon transfer. The excitation energies of the intrinsic states are compared to predictions of blocked BCS calculations.

A different study of high-$K$ states used a radioactive $^{14}\text{C}$ beam at 67 MeV to investigate the high-spin structure of $^{180}\text{Os}$. New rotational bands built on multi-quasiparticle states with $K^\pi = 5^-, 7^-, 9^-, 10^+$ and $15^+$ are observed. The first crossing of the ground-state band at $I = 14 \hbar$, is interpreted as involving a high-$K$ $\pi$-band structure and explained using a two-band mixing model. The structure at higher angular momentum is dominated by intrinsic states, that exhibit a dramatic loss of isomerism. Potential-energy-surface calculations, with Lipkin-Nogami pairing, show these configurations to be triaxial, accounting for the breakdown in $K$ conservation. The relation between the $K$ projection and the total angular momentum is investigated for these non-axial states.
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5. Identification of a high-$K$ isomer in neutron-rich $^{185}$Ta.
“My shoes were too tight and I had forgotten how to dance.”

Londo Mollari,

by J. Michael Straczynski.
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Chapter 1

Introduction

Located at the centre of every atom is the nucleus comprised of protons and neutrons. The motion of these particles within the nucleus is complicated because they lie in a nuclear potential that arises due to the interactions between all of the nucleons, but can be approximated by an average (or mean) nuclear field. More precise descriptions for the properties of nuclei require corrections to the mean field approach to take into account, for instance, the fact that nucleons prefer to bind together in pairs. Understanding these residual effects can provide insight into the internal structure of nuclei. One phenomenon which lends itself to study and provides a great deal of nuclear structure information is the behaviour of nuclei as they rotate.

In 1968 a metastable (long-lived) state was discovered in the nucleus $^{178}\text{Hf}$ [1] with a half-life of 31 years [2] and a high excitation energy (2.4 MeV). (A metastable state is defined here to be an excited nuclear state with a half-life greater than a few nanoseconds.) Such states are now well established in the mass-180 region, although most of the known examples have considerably shorter half-lives, typically microseconds. The nuclei in this region are deformed with rugby ball (prolate) shapes. The most common mode of rotation for these nuclei is for the nucleus as a whole to turn about an axis perpendicular to the symmetry axis. This is called collective rotation. Conversely the metastable (or isomeric)
states consist of one or more unpaired particles rotating about the symmetry axis. Therefore, the angular momentum vectors for these two different types of rotation are at right angles to each other. Energy is required to break pairs of particles and form the isomeric states and so in most cases these lie higher in energy than the lowest collective rotational state of the same angular momentum. The only available decay mode for the metastable states is to lower-lying collective rotational levels which changes the angular momentum vector by 90°. This large change underlies why such states are long-lived because these isomers can only decay at all due to fluctuations in the nuclear shape (or shape orientation), weak mixing between dissimilar states or by slow transitions that carry away many units of angular momentum.

Although many isomeric states are known in the mass-180 region their behaviour in neutron-rich nuclei and at high-spins is unknown and not well predicted. In addition, although approximate calculations of the energy of the isomeric states can be performed no reliable method exists for predicting the half-lives to within orders of magnitude. A larger body of data is a necessary prerequisite for addressing these problems.

1.1 Motivation

The isomers discussed in the section above are referred to as \( K \)-isomers, where \( K \) is the projection of the total angular momentum on the symmetry axis. In well-deformed prolate nuclei found in the mass-180 region, this quantum number is approximately conserved, giving rise to long-lived states for decays involving large changes in \( K \).

In the past the most popular and successful way of studying these isomers was by using fusion-evaporation reactions. Here, two lighter nuclei are fused to form a hot compound system which then emits particles such as neutrons, protons and alpha-particles, leaving the product nucleus in a high-spin state.
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The limitation of this technique is that because only stable beams and targets have been available, the product nuclei are neutron-deficient. Predictions for $K$-isomers in neutron-rich nuclei [3] have therefore remained untested for many years.

Calculations can be performed to estimate the excitation energies of multi-quasiparticle states to within a few hundred keV, however, there is as yet no successful theory for reliably including residual nucleon-nucleon interactions, and the limited data available on these are once again largely restricted to neutron-deficient nuclei. Half-life calculations are even less reliable.

Another reason for sustained interest in $K$-isomers has been the observation of isomers with much shorter half-lives than expected. Such states are thought to occur for several different reasons. Coriolis mixing between states with the same angular momentum but different $K$ values can result in a faster than expected decay. This can be thought of as a "wobbling" effect when the nucleus rotates. The net result of this is to mix low-$K$ components into the wave function of the isomer and high-$K$ components into the wave function of the state to which it de-excites.

A weaker mixing phenomenon can accelerate the decays from isomers when they have a high excitation energy compared to the lowest lying level with the same angular momentum. This is a statistical effect because the number of states increases with energy [4]. At high energies the sheer number of levels makes it more likely that states will occur close to the isomer with the same angular momentum but different $K$, increasing the chance of mixing.

Another mechanism for unexpectedly fast isomeric decays may occur in nuclei that do not have axially symmetric shapes or that undergo axially asymmetric shape fluctuations. This results in the $K$ quantum number no longer being conserved.

It is still not clear which of the modes mentioned above are responsible for short lived states, or how $K$-isomers behave in nuclei with large neutron numbers.
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Data on neutron-rich nuclei in this mass region and at high spins will provide a rich testing ground for theoretical predictions.
Chapter 2

Nuclear structure

2.1 The Spherical Shell Model

In the first half of the 20th century a model was introduced to explain the observed shell structure in nuclei. Closed shells occur at proton and neutron “magic” numbers of 2, 8, 20, 28, 50, 82 and 126 where there are large energy gaps between successive nuclear orbitals. At these shell closures, the binding energy of the last nucleon is much larger than the corresponding value in the neighbouring nuclei, e.g. $^{208}_{82}$Pb is doubly magic and is thus very stable. The individual nucleon motions can be described as being constrained in a potential that itself arises from the combined effect of all the nucleons. The model introduced was aptly named the Spherical Shell Model [5], and it best describes spherical nuclei with a Woods-Saxon [6] potential (see Equation (2.1) below) coupled with a spin-orbit potential. This spin-orbit term arises from a coupling between the intrinsic angular momentum, $s$, and orbital angular momentum, $l$, of the individual nucleons, such that $j = l + s$. The energy levels of each $j$-shell are $(2j+1)$ degenerate, labelled by $m_j$ (the projection of $j$). The Woods-Saxon potential can be expressed as
where \( r \) is the radial distance from the centre of the potential and \( a \) is a parameter that determines how sharply the potential increases to zero. \( R \) is the radius at which \( V(R) = -V_0/2 \) where \( V_0 \) defines the depth of the potential. The nature of the spin-orbit coupling causes high-j orbitals to be pushed down in energy. When one of these orbitals appears in a shell lower than it would have been without the spin-orbit term, it is called an intruder orbital or unique parity orbital since its parity differs from the other orbitals in that shell. This model can explain many features of spherical nuclei but needs modifying to describe nuclei with many nucleons outside closed shells. The residual interactions between these many valence nucleons may be more simply described in terms of deformed potentials.

2.2 Deformation

For nuclei to rotate they must be non-spherical so that they have a preferred axis. For deformed nuclei assuming a constant nuclear volume (i.e. incompressibility) and real solutions, the nuclear radius can be described by

\[
R(\theta, \phi) = R_{av}[1 + \sum_{\lambda=2}^{\infty} \sum_{\mu=-\lambda}^{\lambda} \alpha_{\lambda\mu} Y_{\lambda\mu}(\theta, \phi)]
\]

where \( \alpha_{\lambda\mu} \) are the coefficients of the spherical harmonics \( Y_{\lambda\mu}(\theta, \phi) \) [7]. The \( \lambda = 1 \) terms are normally excluded from the sum as these correspond to a translation of the centre-of-mass providing the coefficients \( \alpha_{2\mu} \) are small. \( \lambda \) and \( \mu \) govern the deformations in \( \theta \) and \( \phi \) respectively. For instance for spheroidal nuclei,

\[
R(\theta, \phi) = R_{av}[1 + \beta_2 Y_{20}(\theta, \phi)]
\]

which is independent of \( \phi \). This means that such nuclei are axially symmetric, either oblate or prolate, see Figure 2.1. The deformation parameter \( \beta_2 (=\alpha_{20}) \),

\[
V(r) = \frac{-V_0}{1 + \exp\left(\frac{r-R}{a}\right)}
\]
can be related to the axes of the spheroid by,

$$\beta_2 = \frac{4}{3} \sqrt{\frac{\pi}{5}} \frac{\Delta R}{R_{av}}$$  \hspace{1cm} (2.4)

in which the average radius, $R_{av} = R_0 A^{1/3}$, and $\Delta R$ is the difference between the semi-major and semi-minor axes. The larger the value of $\beta_2$ the more deformed the nucleus. Positive and negative $\beta_2$ values correspond to prolate and oblate shapes respectively (see Figure 2.1).

![Diagram showing oblate, spherical and prolate shapes](image)

Figure 2.1: Diagram showing oblate, spherical and prolate shapes [8]. The arrows for the oblate and prolate shapes indicate the symmetry axis.

In some circumstances the quadrupole deformation parameters $\epsilon_2$ and $\delta$ are used. These are related to $\beta_2$ by the Equations (2.5), taken from [9].

$$\delta = \frac{\Delta R}{R_{r.m.s.}}$$

$$\epsilon_2 = \delta + \frac{1}{6} \delta^2 + \frac{5}{18} \delta^3 + \frac{37}{216} \delta^4 + \ldots$$

$$\beta_2 = \sqrt{\frac{\pi}{5}} \left[ \frac{4}{5} \epsilon_2 + \frac{4}{9} \epsilon_2^2 + \frac{4}{27} \epsilon_2^3 + \frac{4}{81} \epsilon_2^4 + \ldots \right]$$  \hspace{1cm} (2.5)

Higher order axially symmetric effects have also been observed in nuclei, such as hexadecapole deformations quantified by $\beta_4$ (or $\epsilon_4$).
The shape parameters introduced so far all describe axially symmetric nuclear shapes, but quadrupole ($\lambda = 2$) deformations can give rise to asymmetric shapes. These triaxial distortions are governed by the $\gamma$ shape degree of freedom, and this describes a stretching/squashing effect at right angles to the major nuclear axis. (This effect is shown in Figure 2.3.) From Equation (2.2) $\gamma$ is defined by,

$$\alpha_{20} = \beta_2 \cos \gamma \quad \text{and} \quad \alpha_{22} = \frac{1}{\sqrt{2}} \beta_2 \sin \gamma$$  \hspace{1cm} (2.6)

with $\alpha_{21} = \alpha_{2-1} = 0$. Gamma is measured in degrees where $\gamma = 0^\circ$ and $\gamma = 60^\circ$ correspond to prolate and oblate shapes respectively. Completely triaxial shapes have $\gamma = 30^\circ$.

The model that describes axially symmetric nuclei is called the Deformed Shell Model. In this model the Schrödinger equation is solved using a potential that describes, as closely as possible, the actual shape of the nucleus. Another result of the deformation is that the orbital angular momentum, $l$, and intrinsic spin, $s$, are no longer good quantum numbers and thus states with different $l$ values (but the same parity) can mix. The energy of the states now depends on the component of the angular momentum along the symmetry axis, $\Omega$. For each orbital with angular momentum $j$, there are $2j + 1$ values of $\Omega$ ($= m_j$ in the absence of other couplings). However, levels with $+\Omega$ and $-\Omega$ have the same energy due to the reflection symmetry of axially symmetric nuclei, so each state is now doubly degenerate, i.e. 2 particles can be placed in each state. For example the $i_{13/2}$ orbital can have $|\Omega|$ equal to $13/2$, $11/2$, $9/2$, $7/2$, $5/2$, $3/2$ and $1/2$. The ordering of these $\Omega$ levels depends on the particular shape of the nucleus, since the lowest in energy is the orbital which interacts (or overlaps) the most with the nuclear core. For prolate nuclei the states with the lowest $\Omega$ values are the most tightly bound whereas for oblate nuclei, the highest $\Omega$ orbitals occur lowest in energy. Such deformed shell model calculations were first performed by Nilsson in 1955 [10] with an anisotropic harmonic oscillator potential and the calculated states (called Nilsson orbitals) are labelled by $\Omega[N n_z A]$, where $N$ is
the total oscillator shell quantum number and determines the parity, given by \((-1)^N\). \(A\) is the projection of the particle orbital angular momentum, \(l\), on the symmetry axis (see Figure 2.2), and \(n_\lambda\) is the number of oscillator shell quanta along the direction of the symmetry axis.

2.3 Angular momentum

A nucleus can generate angular momentum in two different ways, either collectively via rotation and vibration or by few-nucleon excitations in which a small number of unpaired nucleons generate the angular momentum. In practice most states are a mixture of these two extreme modes. Both of these phenomena are illustrated in Figure 2.2, along with the notation used throughout this thesis.

The wave functions for collective states can be thought of as a coherent mixture of many single particle wave functions. A direct result of this is an increase in the E2 transition probability, between two collective members of a band, compared to the corresponding single-particle value, i.e. the wave functions of the initial and final state have a large overlap. This can be a useful way of differentiating between single-particle and collective states.

2.3.1 Rotation

Classically, the energy of a rotor is given by

\[ E = \frac{1}{2} \mathcal{J} \omega^2 \]  

(2.7)

where \( \mathcal{J} \) is the moment of inertia of the system and

\[ \omega = \frac{I \hbar}{\mathcal{J}} \]  

(2.8)

is the angular frequency and \( I \) is the angular momentum. Combining Equations (2.7) and (2.8) implies that
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10

Rotation
axis
General angular momentum coupling
\( \Omega = \Lambda + \Sigma \)
\( K = \sum_i \Omega_i \)
Symmetry axis

Figure 2.2: Notation used for angular momentum vectors in prolate deformed nuclei. \( I \) is the total angular momentum of the nucleus, \( R \) is the collective rotation and \( j \) is the single-particle component. \( K \) is the projection of \( I \) on the symmetry axis and \( i \) is the projection of \( I \) on the rotation axis. \( \Omega_i \) denotes the angular momentum projection of the \( i^{th} \) particle. In the Nilsson model, \( j \) precesses about the symmetry axis and \( \Omega \) is a conserved quantity.

\[
E = \frac{1}{2} \frac{(I\hbar)^2}{J} \quad \text{(2.9)}
\]

Transforming from a classical system to a quantum system, the angular momentum \( I^2 \) becomes \( I(I + 1) \) giving Equation (2.10).

\[
E = \frac{\hbar^2}{2J} [I(I + 1)] \quad \text{(2.10)}
\]

Here \( J \) represents the static moment of inertia (also written \( J^{(0)} \)). Thus, for an ideal rotating even-even nucleus in the ground-state band, the ratio of the energy of the first \( 4^+ \) state to the \( 2^+ \) energy is 3.33. Indeed, values close to this have been observed in many cases, e.g. for \(^{186}\text{W} \) \( E(4^+)/E(2^+) = 3.25 \).
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The spacing of rotational states is generally in agreement with Equation (2.10) until, in the $A \approx 180$ region, $I \approx 16 \hbar$ [11]. At this point there is a sharp decrease in the level spacing, known as “backbending”. This is explained by the fact that at this rotational frequency the ground-state (fully paired) band is no longer “yrast” (an yrast state is the lowest energy state for a given spin). Another band, one based on an intrinsic excitation, becomes lower in energy for the same angular momentum because the Coriolis force (Section 2.8) compensates for the reduction in binding energy due to the loss of pairing. This “$s$-band” has a higher moment of inertia $J$, and so the level spacing is less than for the ground-state band, resulting in the “backbend” in the $\gamma$-ray spectrum. An $s$-band involves particles coupled to low $K$, i.e. with a small angular momentum projection on the symmetry axis. Such a coupling is said to be rotation aligned. As $I$ increases, other bands may become yrast, for instance “$i$-bands” (tilted bands) [12]. These involve two particles excited to $i_{3/2}$ orbitals (in the mass-180 region) coupled to high $K$, with the angular momentum vector $j$, precessing around a rotation axis that is tilted with respect to the symmetry axis (Section 2.8). For a band with a non-zero $K$ value, Equation (2.10) becomes,

$$E = \frac{\hbar^2}{2J} [I(I + 1) - K^2] \quad (2.11)$$

This takes into account the non-rotational component due to the bandhead structure. The increase in the intraband transition energies, $E_\gamma$, with $I$, implied by Equations (2.10) and (2.11) arises because for any given band, the projection of $I$ ($= K$) remains constant (in the absence of increasing alignment, see Section 2.8) so as $I$ increases, more and more rotational energy is needed per unit increase in the total angular momentum.
2.3.2 Vibration

Another way in which nuclei can generate collective angular momentum is by vibrating [14]. In this case, vibrational quanta called “phonons” of multipolarity \( \lambda \), carry the energy. Dipole (\( \lambda=1 \)) vibrations correspond to translations of the centre of mass of the nucleus and therefore cannot be produced by internal forces if the protons and neutrons move together. Phonons of \( \lambda=2 \) produce low-energy quadrupole vibrations, which can take two forms. The first, \( \beta \) vibrations, are shape oscillations directed along the symmetry axis, see Figure 2.3. The angular momentum vector for such oscillations is perpendicular to the symmetry axis, therefore, such bands are based on \( I^\pi = K^\pi = 0^+ \) states. (The parity, \( \pi \), is given by \((-1)^\lambda\).) A second type of quadrupole vibration produces oscillations in the \( \gamma \) shape parameter. Gamma governs the deformation at right angles to the \( \beta_2 \) deformation axis. Thus the angular momentum vector of \( \gamma \) vibrations points along the symmetry axis (assuming a mean \( \gamma \) of zero, the nuclei will on average be axially symmetric), which gives rise to bands based on \( I^\pi = K^\pi = 2^+ \) states. Gamma and \( \beta \) phonons can couple together to produce other combinations of states, but always with positive parity.

Octupole vibrations are associated with \( \lambda=3 \) phonons. Such oscillations, represented in Figure 2.3, are more difficult to visualise but they produce bands...
based on $K^\pi = I^\pi = 0^-, 1^-, 2^-$ and $3^-$ states. Although the $I^\pi = 3^-$ level is a purely octupole vibrational state, the bandheads are able to have angular momentum of less than 3 by coupling with the back rotating nuclear core. This back rotation can partially cancel the component of angular momentum of the phonon along the rotation axis, resulting in a $K<\lambda$ bandhead. As shown in Appendix A, this produces an energy minimum.

2.4 Quasiparticle excitations

In the Liquid Drop Model the binding energy due to the pairing force ($B_p$) is accounted for by Equation (2.12).

$$B_p = \left[ \frac{(-1)^N + (-1)^Z}{A^{1/2}} \right] \delta$$

where $\delta \approx 12$ MeV empirically [15]. Such a pairing force favours even-even nuclei over odd-N and/or odd-Z nuclei, and as a result the binding energy due to the pairing is higher for even-even nuclei making them more stable. Nilsson model calculations treat this pairing effect as a perturbation of the mean field Hamiltonian of the nucleus.

The probability amplitudes for the $k^{th}$ orbital being occupied and unoccupied by a pair of particles are $v_k$ and $u_k$ respectively such that

$$v_k^2 + u_k^2 = 1$$

(2.13)

For a nucleus in its ground state pairs of nucleons occupying orbitals close to the unoccupied levels can scatter to the “empty” single-particle states. However, when a particle is excited to a previously unoccupied single-particle level the energy, $\varepsilon_k$, is changed because pairs of particles cannot scatter to singly occupied orbitals (this is called blocking). The state at the new energy is now called a single-quasiparticle state with an energy
\[ E_k = \sqrt{(\varepsilon_k - \mu)^2 + \Delta^2} \]  

(2.14)

where \( \mu \) is the Fermi energy lying in the region between the occupied and unoccupied levels. A precise definition is given by Equations (2.16) and (2.17). \( \Delta \) is the pair gap given by,

\[ \Delta = G \sum_{k \neq k_j} u_k u_k \]  

(2.15)

where \( k_j \) represents the indices of the singly occupied orbitals. \( G \) is called the monopole pairing strength and is the strength with which a pair of particles, in time reversed orbits, interact.

From the definition of \( u_k^2 \) in Equation (2.13) the total number of particles is

\[ n = 2 \sum_k u_k^2 \]  

(2.16)

and the probability of a state being occupied

\[ u_k^2 = \frac{1}{2} \left[ 1 - \frac{(\varepsilon_k - \mu)}{E_k} \right] \]  

(2.17)

Equations (2.16) and (2.17) define the Fermi surface \( \mu \). (In a macroscopic system the Fermi level is defined as the energy at which the occupation probability is equal to one half. However, for a nuclear system there is likely to be no state at such a point.)

### 2.5 \( K \)-isomers

Nuclei in the \( A \approx 180 \) region have both neutron and proton orbitals with large spin projections on the symmetry axis, \( \Omega \), occurring near the Fermi surface. This results in intrinsic multi-quasiparticle states with aligned spins, \( K \), competing with collective rotational states to generate angular momentum. Nuclei in this
mass region are strongly prolate-deformed which means that $K$ is an approximately good (conserved) quantum number. This arises from the total angular momentum vector, $I$, having a small projection on the rotation axis (i.e. $I \approx K$) which can therefore precess around the symmetry axis, conserving $K$. This gives rise to selection rules for $K$, which, for an allowed transition is as follows:

$$\Delta K \leq \lambda$$

(2.18)

where $\Delta K$ is the change in $K$ between the initial and final states and $\lambda$ is the multipolarity of the transition. The degree of $K$-forbiddenness for a transition, $\nu$, is defined in Equation (2.19).

$$\nu = \Delta K - \lambda$$

(2.19)

If a transition is not $K$-allowed the lifetime of the state is increased compared to the normal single-particle values. Empirically, each degree of $K$-forbiddenness increases the lifetime of the state by a factor of $\sim 20$ for highly forbidden transitions. This is expressed as the hindrance per degree of $K$-forbiddenness (or the reduced hindrance) defined as,

$$f_\nu = \left[\frac{t_1^{W}}{t_1^{1/2}}\right]^{1/\nu} = \left[F_W\right]^{1/\nu}$$

(2.20)

where $t_1^{1/2}$ is the partial $\gamma$-ray half-life and $t_1^{W}$ is the Weisskopf single-particle estimate. This allows transitions of different energies and $K$-forbiddenness to be compared. For low-lying high-$K$ states, $\Delta K$ can easily be as high as 8, as in the classic example of the $K^\pi = 16^+$ isomer in $^{178}$Hf [1], which decays to a $K^\pi = 8^-$ band and has a half-life of 31 years [2]. Transitions of high multipolarity can reduce $K$ by a larger amount than dipole or quadrupole transitions and at the same time lower the degree of $K$-forbiddenness. However, lifetimes increase with the multipolarity and so high-$K$ states can still be long lived even if they do not decay by $K$-forbidden transitions, especially if the transition energy is small. An
extreme case is the naturally occurring \( K^\pi = 9^- \) [16] "spin-trap" isomer in \(^{180}\text{Ta}\), which decays (if at all) by a 75 keV \( \lambda = 8 \) transition to the \( I^\pi = 1^+ \) ground-state with \( t_{1/2} > 10^{15} \) years [17].

### 2.6 Branching ratios and g-factors

States in rotational bands, of the type discussed in Section 2.3.1, can decay by either stretched E2 transitions or by mixed M1/E2 transitions (for bands with both odd and even spin members). Stretched E2 transitions are assumed here to be of pure multipolarity because any admixture of M3 (or higher multipoles) will be very small. When rotational levels decay by both \( \Delta I = 1 \) and \( \Delta I = 2 \) transitions, the \( \gamma \)-ray branching ratio \( \lambda_b \), can provide information about the underlying nuclear structure upon which the band is built.

The following expression can be derived (see Appendix B) for the intrinsic g-factor, \( g_K \), of a given rotational band [18],

\[
\frac{(g_K - g_R)}{Q_0} = 0.933 \frac{E_1}{\delta \sqrt{(I^2 - 1)}}
\]  

(2.21)

where \( g_R \) is the rotational g-factor, \( Q_0 \) is the intrinsic quadrupole moment in units of e-b and \( E_1 \) and \( E_2 \) are the transition energies (in MeV) of the \( \Delta I = 1 \) and 2 transitions respectively, from an initial state with spin \( I \). The quadrupole/dipole mixing ratio \( \delta \) is related to the quadrupole admixture \( q \) by,

\[
q = \frac{\delta^2}{1 + \delta^2} = \frac{2K^2(2I-1)}{(I-K-1)(I+K-1)(I+1)} \frac{E_2^5}{E_1^3 \lambda_b}
\]  

(2.22)

Equations (2.21) and (2.22) are particularly useful in the analysis of rotational band properties. The factor \( \frac{(g_K - g_R)}{Q_0} \) should be a constant for a given rotational band because all of the states are built on the same bandhead configuration. However, it should be noted that since the branching ratio \( \lambda_b \) is proportional to \( \delta^2 \) (Equation (2.22)), and the ratio \( \frac{(g_K - g_R)}{Q_0} \) is proportional to \( \frac{1}{\delta} \) (Equation (2.21)),
only the magnitude of the ratio \((g_K - g_R)/Q_0\) can be found from the branching ratio. The rotational g-factor is generally given by \(g_R \approx \frac{Z}{A}\). However, \(g_R\) can be significantly different from this depending on the configuration. The expressions of Belyaev [19] and Migdal [20] take into account the change in pairing energy and the deformation, yielding a configuration dependent rotational g-factor,

\[
gr = \frac{ZJ_\pi}{ZJ_\pi + N J_\nu}
\]

where \(Z\) is the proton number and \(N\) is the neutron number. \(J_\pi\) and \(J_\nu\) are the proton and neutron moments of inertia respectively, given by

\[
J_{\nu,\pi} = J_{\text{rigid}} \left(1 - \frac{\ln[x_{\nu,\pi} + (1 + x_{\nu,\pi}^2)^{1/2}]}{x_{\nu,\pi}(1 + x_{\nu,\pi}^2)^{1/2}}\right)
\]

where the classical rigid moment of inertia for a prolate spheroid is

\[
J_{\text{rigid}} = \frac{2}{5} M R_{av}^2 (1 + \frac{1}{3} \beta_2)
\]

\(M\) is the mass of the nucleus, \(R_{av} (=1.2A^{1/3}\text{ fm})\) is the average radius, \(\beta_2\) is the quadrupole deformation parameter defined in Equation (2.3) and

\[
x_{\nu,\pi} = \frac{(\hbar \omega_0 \delta)_{\nu,\pi}}{2 \Delta_{\nu,\pi}}
\]

where \(\delta\) is the quadrupole deformation parameter given by Equation (2.5), \(\Delta_{\nu,\pi}\) are the neutron and proton pairing energies which are discussed in Chapter 3 and \(\omega_0\) is the Nilsson model oscillator frequency given by,

\[
\hbar \omega_0 = 41A^{-1/3}[1 + \frac{1}{3} \left(\frac{N - Z}{A}\right)] \text{ MeV for neutrons}
\]

and

\[
\hbar \omega_0 = 41A^{-1/3}[1 - \frac{1}{3} \left(\frac{N - Z}{A}\right)] \text{ MeV for protons}
\]

from Ref. [21].

Once the experimental g-factors have been extracted for a given band, they can be compared to the expected g-factor for a given configuration. For a band
built on a multi-quasiparticle state the value of \( g_K \) can be calculated using the following equation.

\[
K g_K = \sum (g_\Lambda \Lambda + g_\Sigma \Sigma)
\]  

(2.29)

Where \( \Lambda \) and \( \Sigma \) are the projections of the orbital angular momentum and intrinsic spin on the symmetry axis respectively for a particular particle (Figure 2.2). Their values may be calculated from the Nilsson model. \( g_\Lambda \) and \( g_\Sigma \) are the corresponding g-factors. \( g_\Lambda \) is 0 for neutrons and 1 for protons, and \( g_\Sigma^{\text{free}} \) is +5.59 for protons and −3.83 for neutrons. These values for \( g_\Sigma \) are attenuated from their “free” values by a factor commonly found to be 0.6 [22]. It should be noted that the asymptotic Nilsson wave functions and corresponding quantum numbers (Section 2.2) are valid in the limit of large deformations. As a consequence of this Equation (2.29) works well for orbitals with large \( K \) projections but is less satisfactory for low-\( K \) orbitals (e.g. the 1/2[510] neutron orbital) where there may be significant mixing in the wave function and thus contributions from orbitals with different \( l \) values or even from different shells. This mixing can be calculated and the non-asymptotic wave functions used for the calculation of the intrinsic g-factor. In this instance Equation (2.29) is usually written as

\[
K g_K = \sum g_\Omega \Omega = \sum \Omega \left[ g_\Lambda + (g_\Sigma - g_\Lambda) \frac{<s_\Sigma>}{\Omega} \right]
\]  

(2.30)

where \( <s_\Sigma> \) is the expectation value of the intrinsic spin. The values of \( <s_\Sigma> \) are tabulated in Refs.[23, 24].

### 2.7 Two-level mixing

In Section 2.3.1 the phenomena of backbending was introduced and explained in terms of band crossings [25]. When bands cross they can interact or mix. Mixing occurs when two bands with levels of the same spin and parity lie close in energy (i.e. the states have a large overlap) and the corresponding levels from each band
repel each other, thus causing an apparent change in the moment of inertia. The sections below describe quantitatively how the strength with which two bands interact can be calculated.

2.7.1 Two-state mixing

Consider two states from different bands in the region where the bands cross, as shown in Figure 2.4. Both of the states have the same spin, $I$ and parity, $\pi$. The wave functions for the two states $\psi_1$ and $\psi_2$, can be written as a linear combination of the two wave functions involved in the mixing,

$$
\psi_1 = \alpha \phi_1 - \beta \phi_2, \quad \psi_2 = \beta \phi_1 + \alpha \phi_2
$$

(2.31)

where $\phi_1$ and $\phi_2$ are the pure wave functions for bands 1 and 2 and $\alpha$ and $\beta$ are the amplitudes of the major and minor components in the wave functions respectively, i.e. $\alpha > \beta$, and $\alpha^2 + \beta^2 = 1$. From the perturbed (observed) energies ($E_{\text{obs}}$) of the two levels, it is possible to find the unperturbed energies, and vice versa, given an interaction matrix element $V$, such that

$$
\begin{pmatrix}
E_1^+ & V \\
V & E_2^-
\end{pmatrix}
\begin{pmatrix}
\phi_1 \\
\phi_2
\end{pmatrix}
= E_{\text{obs}}
\begin{pmatrix}
\phi_1 \\
\phi_2
\end{pmatrix}
$$

(2.32)

where $E_1^+$ and $E_2^-$ are the unperturbed energies of the upper and lower states respectively. The two solutions for $E_{\text{obs}}$ can be found by simply rearranging Equation (2.32) and then diagonalising the resulting matrix, leading to,

$$
E_{1,2} = \frac{1}{2} \left[ (E_1^+ + E_2^-) \pm \sqrt{(E_1^+ - E_2^-)^2 + 4V^2} \right]
$$

(2.33)

($E_1$ goes with the plus sign.) Conversely, the solutions for the unperturbed energies, expressed in terms of $E_1$ and $E_2$ are found by solving the conjugate of Equation (2.32) (i.e. $|V|^2 \rightarrow -|V|^2$), resulting in,
\[ E_{1}^{+}, E_{2}^{=} = \frac{1}{2}[(E_{1} + E_{2}) \pm \sqrt{(E_{1} - E_{2})^2 - 4V^2}] \]  

(2.34)

Expressions for the coefficients of \( \phi_{1} \) and \( \phi_{2} \), namely \( \alpha \) and \( \beta \), can be found by substituting the solution for \( E_{\text{obs}} \) into Equation (2.32). Manipulating the resulting equations leads to the following expressions.

\[
\beta = \frac{1}{\sqrt{1 + \frac{V^2}{(E_{1}^{+} - E_{1})^2}}} \\
\alpha = \sqrt{1 - \beta^2}
\]  

(2.35)

where \( |E_{1}^{+} - E_{1}| = |E_{2}^{=} - E_{2}| = \epsilon_{1,2} \) is the energy difference between the initial and final states.

Figure 2.4: Two-level mixing. The dashed lines represent the unperturbed (pure) states and the perturbed (experimental) levels are shown by solid lines. \( V \) is the mixing matrix element. For an explanation of other symbols see text.

### 2.7.2 Two-band mixing

Extending the situation of two-state mixing, as discussed in Section 2.7.1, leads to a description of two-band mixing, as presented below. Figure 2.5 shows four different situations of two-band mixing where the two bands have different \( K \) values, \( K_{1} \) and \( K_{2} \). To enable the mixing matrix element, \( V \), to be calculated...
it is useful to try and reproduce the experimental out-of-band to in-band $\gamma$-ray branching ratio. In addition to $V$, this ratio depends on the mixing amplitudes $\alpha$, $\beta$, $\gamma$, and $\delta$ (Figure 2.5), the Clebsch-Gordon coefficients between the initial and final states and the energies of the states. Equations (2.35) give expressions for $\alpha$ and $\beta$ in terms of the energies of the two highest states, $E_1$ and $E_2$. Similarly for the two lower states, $E_3$ and $E_4$,

$$\delta = \frac{1}{\sqrt{[1 + \frac{V^2}{(E_3 - E_4)^2}]}}$$  \hspace{1cm} (2.36)

where $\gamma > \delta$, $\gamma^2 + \delta^2 = 1$ and the energy difference $|E_3^+ - E_3| = |E_4^- - E_4| = \varepsilon_{3,4}$ is the perturbation caused by the mixing.

The reduced transition rate for a given $\gamma$-ray (B(L\(\lambda\)), where $\lambda$ is the multipolarity and L is either M (magnetic) or E (electric)) depends on the mixing amplitudes of the initial and final states and also on the Clebsch-Gordon coefficients. Considering a multipolarity of E2 only, a transition from a state with $I, K_1$ to another state with $I - 2, K_1$ has a Clebsch-Gordon coefficient $c_1$ [26].

$$c_1 = \sqrt{\frac{3(I - K_1)(I - K_1 - 1)(I + K_1)(I + K_1 - 1)}{(2I - 2)(2I - 1)I(2I + 1)}}$$  \hspace{1cm} (2.37)

Substituting $K_1$ for $K_2$ gives an expression for $c_2$. The ratio of these coefficients is defined as, $c = c_2/c_1$. For example, the reduced transition rate from level $E_1$ to $E_3$, as shown in Figure 2.5a is given by,

$$B(E2)_{1\rightarrow3} \propto (c_1 \alpha \delta + c_2 \beta \gamma)^2$$  \hspace{1cm} (2.38)

The quadrupole moment $Q_o$ is assumed to be the same for the initial and final states.

Using Equations (2.35), (2.36), (2.37) and (2.38), the ratio $R$, of the out-of-band to the in-band reduced transition rates, can be calculated for any given
Figure 2.5: A plot showing the four different types of two-band mixing. In-band destructive interference is represented in a and b. This is the situation where the bands are crossing. Out-of-band destructive interference is shown in c and d which occurs when the bands are not crossing. \( \alpha, \beta, \gamma, \) and \( \delta \) are the coefficients of the \( K_1 \) and \( K_2 \) components in the different levels. The \( K_1 \) component is quoted first for each level. The arrows represent in-band (vertical) and out-of-band (diagonal) \( \gamma \)-ray transitions.

The mixing matrix element \( V \). For each of the four cases shown in Figure 2.5,

\[
R_a = \left( \frac{c_1 \alpha \delta + c_2 \beta \gamma}{c_1 \alpha \gamma - c_2 \beta \delta} \right)^2 = \left( \frac{\alpha \delta + c_\beta \gamma}{\alpha \gamma - c_\beta \delta} \right)^2
\]

\[
R_b = \left( \frac{\beta \gamma + c_\alpha \delta}{\beta \delta - c_\alpha \gamma} \right)^2
\]

\[
R_c = \left( \frac{\alpha \delta - c_\beta \gamma}{\alpha \gamma + c_\beta \delta} \right)^2
\]

\[
R_d = \left( \frac{\beta \gamma - c_\alpha \delta}{\beta \delta + c_\alpha \gamma} \right)^2
\]

(2.39)

This ratio \( R_a \) is equal to the reduced out-of-band to in-band branching ratio, i.e. the \( \gamma \)-ray branching ratio with the energy dependence of \( E_{\gamma_1}^5 \), removed. Thus,
CHAPTER 2. NUCLEAR STRUCTURE

$V$ can be found for any given value of $R$, but is usually calculated between the two closest lying states in the two bands. $V$ is constrained to be less than or equal to $\frac{E_1-E_2}{2}$ with the maximum value corresponding to initially degenerate levels, and $V$ is zero when there is no mixing.

2.8 Aligned angular momentum

Bands built on configurations containing orbitals which tend to align strongly with the nuclear rotation axis were briefly discussed in Section 2.3.1.

A particle in a given orbital of angular momentum $j$ in a state with a total nuclear spin $I$ feels a Coriolis force $\propto I\cdot j$ [27] aligning it with the nuclear rotation axis. It follows therefore, that the bands exhibiting the most alignment due to the Coriolis interaction are those built on configurations containing high-$j$ particles, in particular those occupying the $i_{13/2}$ (neutron) and $h_{11/2}$ (proton) orbitals (in the mass-180 region). For these orbitals the Coriolis force can contribute over 1 MeV, compensating for the reduction in pairing energy (Chapter 3) at high spins.

The total aligned angular momentum can be obtained for a given observed transition [28] from

$$I_x = \frac{(I_i^2 + I_f^2)}{2}$$

(2.40)

where

$$I_x^{i,f} = \sqrt{I_i I_f (I_i I_f + 1) - K^2}$$

(2.41)

from Pythagoras's theorem where $I^i$ and $I^f$ are the total angular momenta of the initial and final states respectively and $K$ is the projection of the total angular momentum on the symmetry axis. This method gives the average aligned angular momentum for the transition as opposed to the standard technique [29] of calculating the aligned angular momentum for the average spin of the $\gamma$-ray transition, $I_{av}$, where

$$I_x = \sqrt{I_{av} (I_{av} + 1) - K^2}.$$ 

Although the differences between these two
methods are generally small, they can be significant at low rotational frequencies as shown by Purry et al. [28]. All the alignments extracted in the following chapters are found using Equation (2.40) and plotted against rotational frequency, $h \omega$. This quantity is calculated from the energy of a rotor, $E = \omega I_x h$. Manipulating this gives $E_{\text{initial}} - E_{\text{final}} = E_\gamma = h \omega (I^i_x - I^f_x)$, yielding the expression

$$h \omega = \frac{E_\gamma}{I^i_x - I^f_x}$$  \hspace{1cm} (2.42)

In many of the plots showing the “net” aligned angular momentum, $i_x$, for a rotational band as a function of rotational frequency, a subtraction is made of the form $[\mathcal{J}_0^{(1)}(h \omega) + \mathcal{J}_1^{(2)}(h \omega)^3]$, where $I$ is the total angular momentum and $\mathcal{J}_0^{(1)}$ and $\mathcal{J}_1^{(2)}$ are the Harris reference parameters [30] with units $h^2 \text{MeV}^{-1}$ and $h^2 \text{MeV}^{-3}$ respectively. These parameters are found by fitting the first few transitions in the ground-state band with the above function. The effect of this is to expand any differences between bands, allowing the net alignments for different bands and nuclei to be contrasted easily.

The graphs of total alignment versus frequency (with no subtraction) can be used to extract the kinematic and dynamic moments of inertia for a band, defined as

$$\mathcal{J}^{(1)} = h \frac{I_x}{\omega} \quad \text{and} \quad \mathcal{J}^{(2)} = h \frac{dI_x}{d\omega}$$  \hspace{1cm} (2.43)

respectively. For a specific rotational band in a nucleus $\mathcal{J}^{(1)}$ and $\mathcal{J}^{(2)}$ will differ (they are the same for a rigid body) revealing information about the bulk and quasiparticle motion.
Chapter 3

Calculations

In the following sections, the theoretical calculations, used for comparison with experimental data, are described.

3.1 Blocked BCS theory

For several of the nuclei discussed in the following chapters, "blocked BCS" calculations as described by Jain et al. [31], have been performed to estimate the energies of multi-quasiparticle configurations, which can then be compared to the experimental data.

BCS theory [32] (Bardeen-Cooper-Schrieffer after its developers), was first used to explain superconductivity in metals (which are macroscopic). It was suggested that the energy gap observed in the electronic excitation of superconductors is analogous to the energy gap observed in the excitation spectra of even-even nuclei [33], both of which deviate from independent particle motion (a Fermi gas). In superconductivity, electrons with equal and opposite momenta are correlated. Similarly, a pairing force between nucleons in time reversed orbits can reproduce the energy gap between the ground-state and first excited intrinsic states in nuclei.

When applied to the nucleus, a microscopic system, the sharp transitions ob-
served in the properties of superconducting metals would be expected to become gradual changes. As a superconductor is heated, at some critical temperature the superconductivity is lost. Increasing the internal energy of nuclei has the analogous effect of destroying the pairing. Similarly, rotating a nucleus is analogous to the application of a magnetic field to a superconductor. Empirically the destruction of pairing in nuclei is manifest as an increase in the effective moment of inertia, $J$ [34]. As angular momentum is added to the nucleus pairs of nucleons are broken and align with the rotation vector $I$. This results in an increase in the moment of inertia from the ground-state superfluid value which is about half that expected for a rigid rotor. The nuclear analogue of rigid rotation is all the particles acting independently (i.e. unpaired). At high angular momentum and/or high internal energy one might expect the nucleus to approximate such a case, as more and more pairs of nucleons are broken [34].

### 3.1.1 Pairing and blocking

BCS theory treats the pairing effect as a perturbation of the mean field Hamiltonian of the nucleus. To calculate the energy of a quasiparticle from the single-particle (Nilsson) energies, only the orbitals closest to the Fermi level need to be considered. This is because the probability of excitations and scattering (see Section 2.4) to an unoccupied state is highest for particles near the Fermi surface. The quasiparticle energies for a state $k$ are given by Equation (2.14) and the pair gap, $\Delta$, is given in Equation (2.15) ($\Delta = G \sum_{k \neq k_j} u_k v_k$). $\Delta$ is proportional to $G$ (the monopole pairing strength) because as the strength of the interaction between two particles increases, more energy is needed to break a pair and form multi-quasiparticle states. Thus, $G$ can be thought of as a scaling factor for the multi-quasiparticle excitation spectrum and is chosen to produce the correct excitation energy for the lowest 2-quasiparticle states. In the $A \approx 180$ region it is possible to treat neutrons and protons separately since they occupy different
CHAPTER 3. CALCULATIONS

shells. Hence, there is a neutron pair strength, $G_n$, and a proton pair strength, $G_p$, which are treated independently.

Equation (2.15), for calculating $\Delta$, involves a sum over single-particle states $k$. However, some of these states may be occupied by quasiparticles and as a result these states must be removed (or blocked) from the sum since paired particles cannot scatter to singly occupied orbitals. This leads to a reduction in the pairing energy, and is important in producing an accurate multi-quasiparticle energy spectrum. It is the main difference between blocked BCS theory and ordinary BCS theory. The latter approximates the multi-quasiparticle energies by summing over all of the single-quasiparticle states, i.e. with $\Delta = G \sum_k u_k v_k$, but this approach does not yield an accurate energy spectrum.

3.1.2 Multi-quasiparticle calculations

To perform the calculations, the deformation parameters $\epsilon_2$ and $\epsilon_4$ (Section 2.2) along with $N, Z, G_n$ and $G_p$ are required. The Nilsson single-particle energies, $\varepsilon_k$, are calculated and those near the Fermi surface can be adjusted to produce the correct single-quasiparticle energy differences $E_k$ (Equation (2.14)) in the neighbouring odd-mass nuclei. (Experimentally the Fermi level is not known, only the energies relative to the ground state.) To obtain the multi-quasiparticle energies, $E_{map}$, it is necessary to calculate $\Delta$ for each configuration of like particles (protons and neutrons), blocking the appropriate orbitals. Since $v_k$ and $u_k$ (the occupation amplitudes) are dependent on $\Delta$, this is an iterative process so as to be self consistent. Finally, the neutron and proton configurations are combined, which is a straightforward sum of the two energies, yielding

$$E_{map} = \sum_{k_n} E_{kn}^n + \sum_{k_p} E_{kp}^p$$ (3.1)
3.1.3 Residual interactions

In the mean field approach any interaction not accounted for is considered to be a residual interaction. BCS theory takes into account the pairing force and the blocking describes the non-independent particle motion, both of which are residual interactions. However, residual nucleon-nucleon interactions arising from the intrinsic spins of the particles are not included in blocked BCS calculations, although the shifts these can cause are only a few hundred keV. Such interactions favour couplings of \textit{like} particles with \textit{opposite} intrinsic spin projections and \textit{unlike} particles with the \textit{same} spin projections \cite{35}. From this empirical rule, the states which are most favoured, for a particular configuration, can be found. Jain \textit{et al.} \cite{36} have compiled splitting energies between proton-proton, proton-neutron and neutron-neutron orbitals in the mass 180 region, but the data are limited and calculations of these energies are not accurate. In the chapters that follow, 'residual interactions' refers to these nucleon-nucleon intrinsic spin coupling energies.

3.2 Potential-energy-surface calculations

To calculate the shapes and energies of specific multi-quasiparticle configurations, Xu \textit{et al.} \cite{37} have developed configuration-constrained Potential-Energy-Surface (PES) calculations that include the $\gamma$ shape degree of freedom.

The total energy for a given nucleus and configuration is calculated for a range of $\beta_2$ and $\gamma$ using,

$$E_{\text{total}}(\beta_2, \gamma, \beta_4) = E_{\text{macroscopic}} + E_{\text{shell-correction}} + E_{\text{pairing}}$$ \hspace{1cm} (3.2)

such that at every point on the $\beta_2$-$\gamma$ plane the total energy is minimised with respect to $\beta_4$. The quadrupole deformation parameters $\beta_2$ and $\gamma$ were introduced in Section 2.2, along with the hexadecapole deformation parameter $\beta_4$. $E_{\text{macroscopic}}$ is calculated using the Liquid Drop Model (see below in Section 3.2.1) with the
original parameters [38]. $E_{\text{shell-correction}}$ (also Section 3.2.1) is found using the Strutinsky shell-correction [39] (microscopic), incorporating single-particle levels from a non-axially deformed Woods-Saxon potential [40]. The shell correction term is configuration independent. The pairing energy, $E_{\text{pairing}}$, which is configuration dependent, is obtained using the Lipkin-Nogami treatment of pairing [41] discussed in Section 3.2.2.

### 3.2.1 The Liquid Drop Model and shell corrections

The Liquid Drop Model (LDM) [38] describes the bulk properties of nuclei with terms taking account of the volume, Coulomb repulsion, surface area, symmetry (neutron to proton ratio) and pairing. However, although this empirical model can reproduce the average properties of most nuclei, there are significant deviations from this near closed shells where the LDM underestimates the binding energy. It is therefore necessary to include a correction that takes account of this phenomenon in order to accurately calculate, for instance, ground-state masses as the differences can be more than 10 MeV. Near a closed shell the level density at the Fermi surface (Section 2.4) is very low and this leads to extra binding energy. Qualitatively this is because the nucleons are occupying “deeper” orbitals. The shell effects can therefore be thought of as arising from fluctuations in the level density around the Fermi surface. The method for incorporating these effects is called the Strutinsky shell correction (or renormalisation approach) [42]. The shell correction energy to the LDM energy in Equation (3.2) is

$$E_{\text{shell-correction}} = -2 \int_{-\infty}^{\mu} \tilde{g}(\varepsilon) \varepsilon d\varepsilon$$

(3.3)

where $\tilde{g}(\varepsilon)$ is the mean (or smoothed) energy density of the single-particle states at energy $\varepsilon$. The Fermi energy, $\mu$, corresponding to $\tilde{g}(\varepsilon)$ is determined from

$$n = 2 \int_{-\infty}^{\mu} \tilde{g}(\varepsilon) d\varepsilon$$

(3.4)
such that the total number of particles, \( n \), is conserved \([42]\).

### 3.2.2 Diabatic blocking and Lipkin-Nogami pairing

In order to block the correct orbitals when calculating the pairing term, they are followed by examining the expectation values of their approximate quantum numbers (diabatic blocking), \( \langle N \rangle \), \( \langle n_z \rangle \), \( \langle \Lambda \rangle \) and \( \langle |\Omega| \rangle \), introduced in Section 2.2. This differs from the adiabatic approach in which the lowest energy orbitals are blocked. As \( \beta_2 \) and \( \gamma \) are varied the single-particle orbits can cross leading to incorrect blocking. The technique of diabatic blocking, used here, identifies the orbitals by their quantum numbers, removing this ambiguity. The expression for the Lipkin-Nogami pairing energy \([37]\) that appears in Equation (3.2) is

\[
E_{\text{pairing}} = \sum_{j=1}^{S} \varepsilon_{kj} + \sum_{k \neq k_j} 2v_k^2 \varepsilon_k - \frac{\Delta^2}{G} - G \sum_{k \neq k_j} v_k^4 + G \frac{N_{\nu,\pi} - S}{2} - 4\lambda_2 \sum_{k \neq k_j} (u_k v_k)^2 
\]

where \( S \) is the seniority of the proton or neutron configuration which corresponds to the number of singly occupied (blocked) orbitals with index \( k_j \). \( N_{\nu,\pi} \) is the number of neutrons or protons in the nucleus and \( \lambda_2 \) (a Lagrange multiplier) is a function of \( u_k \), \( v_k \) and \( G \). This \( \lambda_2 \) is sometimes called the number fluctuation constant as it corrects for the fluctuation in particle number. For a definition of other terms see Section 2.4.

The quasiparticle energies for an odd nucleon in an odd-\( A \) nucleus are given by

\[
E_k = \sqrt{(\varepsilon_k - \mu)^2 + \Delta^2 + \lambda_2} 
\]

where \( \mu \) is the Fermi energy. Unlike the BCS calculations, in the actual calculations there is no adjustment of the single-particle energies, and the monopole pairing strengths are not fitted to the observed energies of the 2-quasiparticle states, but are calculated. Since the difference between the ground-state single-particle level and the Fermi surface will be small for an odd-\( A \) nucleus, \( E_k \approx \Delta + \lambda_2 \).
(For an odd-$A$ nucleus the ground-state orbital is only half occupied and so is expected to be close to the Fermi level.) In the Lipkin-Nogami model $\Delta + \lambda_2$ is the odd-even mass difference. The pairing strengths, $G$, can therefore be obtained by fitting $\Delta + \lambda_2$ to the experimental odd-even mass difference, $\Delta^{\text{exp}}$. For example, for proton pairing in an even-even nucleus,

$$\Delta_\pi^{\text{exp}} = -\frac{1}{8} [M(Z + 2, N) - 4M(Z + 1, N)$$

$$+ 6M(Z, N) - 4M(Z - 1, N) + M(Z - 2, N)] \quad (3.7)$$

This is called the average gap method [41]. Furthermore, in Ref.[37] the calculation includes blocking and deformation effects. The latter is necessary due to the differences in the shape parameters between neighbouring nuclei (which are contained in $\Delta^{\text{exp}}$). This results in a renormalisation of the monopole pairing strengths to match theoretical and experimental odd-even mass differences.

The results of these calculations are the surface potentials for different configurations the excitation energies of which can be compared to the experimental observations. As with BCS theory, for accuracies in the calculated energies better than $\sim 100$ keV, residual nucleon-nucleon interactions (Section 3.1.3) should be incorporated.
Chapter 4

Experimental techniques

4.1 Nuclear reactions

Although neutron-rich nuclei present possibilities for exploring new phenomena, such nuclei must first be produced in a quantity sufficient to allow their study. Reactions used in subsequent chapters and other reactions capable of introducing significant angular momentum to a given nucleus are discussed briefly below.

Fusion-evaporation reactions

A common method of producing highly excited nuclei is via fusion-evaporation reactions. In this technique, two stable nuclei are brought together at an energy above the Coulomb barrier; the resulting compound system subsequently evaporates nucleons before decaying by $\gamma$-ray emission. However, such a process produces nuclei on the neutron-deficient side of the valley of stability. This is because stable light nuclei such as the target and projectile have a lower $N/Z$ ratio compared to the heavy stable nuclei. In addition, the compound system preferentially emits neutrons, leaving a more neutron-deficient nucleus. The $\alpha$- and proton-emission channels are less probable for heavy compound nuclei near stability. In order to produce nuclei that are more neutron-rich with this technique, radioactive beams (or targets) can be used. One such projectile that is
Deep inelastic reactions

Another method of populating high-spin states is by the use of deep inelastic reactions. Here, at energies above the Coulomb barrier, nucleons are transferred between the target and projectile with the direction of transfer tending to equalise the N/Z ratio of the system. (Compound-nucleus formation does not compete when the Z of the target and projectile is very high.) Neutron evaporation can also occur resulting in a large spread of nuclei between the beam and target. By using the most neutron-rich target and projectile combination possible in the relevant mass region, neutron-rich nuclei can be excited by neutrons being transferred to the target or projectile. Details of the yield distributions and angular momentum transferred for these reactions are described in Refs. [43] and [44] respectively. In addition, the target and beam nuclei themselves are highly excited, possibly by the transfer and subsequent evaporation of nucleons or other inelastic processes. This provides a novel way of studying stable and neutron-rich isotopes. One drawback with this method is the difficulty in identifying isotopes produced by transfer if insufficient is known in the region of study to use γ-ray coincidences with known transitions between low-lying states. Ancillary detectors are difficult to use effectively with these reactions due to the large angular spread in the product nuclei, although there is a peak in the cross-section at the grazing angle. Often, coincidences with X-rays can be used to give the proton number. Another complementary technique uses the correlation of prompt partner nuclei. By identification of transitions in one of the products (beam-like or target-like), decays in the binary partner nucleus can be deduced [45]. Unfortunately this cannot be applied to off-beam events. An experiment employing this method of deep inelastic reactions is described in Chapter 5.
Other reactions

The Coulomb excitation of the nuclei of interest can introduce significant angular momentum, but does not populate states based on a different configuration to that of the ground state to any significant degree. Generally, only the ground state band and associated collective (i.e. vibrational) bands are populated, although if other configurations are populated in the decay of the collective structures, they might also be observed.

Fragmentation reactions involve very high-energy beams of nuclei being broken up, on a robust target, into lighter products. Such reactions produce a wide range of nuclei and so the experimental setup must include one or more ancillary detectors to select and identify the nuclei of interest. Due to the high speed of the fragments (typically $v/c \sim 60\%$) it is difficult to collect useful prompt $\gamma$-ray events. The principal advantage of these reactions is their ability to produce both neutron and proton-rich nuclei and this technique is suited to studies of isomeric decays [46] but is only sensitive to isomers with half-lives from $\sim 1$ to $\sim 100$ $\mu$s. The pursuance of this complementary approach is outside the scope of this thesis.

The fission of heavy nuclei into lighter products can lead to the population of very neutron-rich nuclei [47]. Some degree of selectivity is possible by choosing induced fission which is symmetric, or spontaneous fission which is asymmetric. Ancillary detectors are once again necessary to associate the $\gamma$-ray decays with a given nucleus. However, fission leads to nuclei with significantly lower mass than those studied in this thesis.

The techniques described above are complex often including several reaction mechanisms. A detailed discussion of these is beyond the scope of the present work. For further details see Ref.[48].
4.2 Angular distributions and DCO ratios

In fusion-evaporation reactions the compound nucleus is formed in a state with its angular momentum vector perpendicular to the axis defined by the direction of the beam. The subsequent evaporation of particles causes some smearing of the direction of the 'polarisation' but even so the nucleus retains a high degree of orientation for times of the order of nanoseconds. When a nucleus in such a state emits $\gamma$-radiation the relative intensities at different angles with respect to the beam-axis depend on the multipolarity of the transition. Figure 4.1 shows the intensity distributions for dipole and quadrupole transitions as a function of angle, $\theta$, to the beam direction.

![Figure 4.1: Angular distributions for an $I\rightarrow I - 1$ dipole transition (solid line) and an $I\rightarrow I - 2$ quadrupole transition (dashed line). Zero degrees corresponds to the positive $x$-axis.](image)

The $\gamma$-ray intensity distributions at a given angle [49] are given by,
\[ I(\theta) = \sum_{\ell = \text{even}} A_\ell P_\ell(\cos \theta) \]  

(4.1)

where \( A_\ell \) are the coefficients of the Legendre polynomials \( P_\ell(\cos \theta) \),

\[ P_0(\cos \theta) = 1 \]
\[ P_2(\cos \theta) = \frac{1}{2} (3 \cos^2 \theta - 1) \]  

(4.2)
\[ P_4(\cos \theta) = \frac{1}{8} (35 \cos^4 \theta - 30 \cos^2 \theta + 3) \]

Values for the coefficients \( A_\ell \) are tabulated in Ref. [50].

The different angular distributions for pure quadrupole and dipole radiation shown in Figure 4.1, allow multipoles to be distinguished by examining the relative \( \gamma \)-ray coincidence intensities at angles approximating (for example) \( 0^\circ \) and \( 90^\circ \). This method (using Directional Correlations de-exciting Oriented states) is called DCO analysis [51]. For the data obtained using the NORDBALL detector array (Chapter 7) the respective angles are \( 37^\circ, 143^\circ \) and \( 79^\circ, 101^\circ \). The DCO intensity ratio is thus defined as

\[ R_{DCO} = \frac{I_{\gamma 2}(37^\circ, 143^\circ)}{I_{\gamma 2}(79^\circ, 101^\circ)} \text{ gated by } \frac{\gamma 1(79^\circ, 101^\circ)}{\gamma 1(37^\circ, 143^\circ)} \times \varepsilon \]  

(4.3)

Where \( \gamma 1 \) is the gating transition and \( I_{\gamma 2} \) is the area of the \( \gamma \)-ray transition of interest in the projected spectrum. The efficiency correction factor \( \varepsilon \) is given by,

\[ \varepsilon = \frac{\varepsilon_{\gamma 2}^{\text{projection}}(79^\circ, 101^\circ) \times \varepsilon_{\gamma 1}^{\text{gate}}(37^\circ, 143^\circ)}{\varepsilon_{\gamma 2}^{\text{projection}}(37^\circ, 143^\circ) \times \varepsilon_{\gamma 1}^{\text{gate}}(79^\circ, 101^\circ)} \]  

(4.4)

taking into account the efficiency at both the gating energy (\( \varepsilon_{\gamma 1}^{\text{gate}} \)) and the projection energy (\( \varepsilon_{\gamma 2}^{\text{projection}} \)). Gating on a stretched quadrupole transition gives ratios of \( \approx 1.0 \) for quadrupole transitions and \( \approx 0.56 \) for pure stretched dipole radiation [52], although these numbers depend on the spectrometer. Mixed M1/E2 transitions can have a DCO ratio ranging from \( \approx 0.3 \) up to \( \approx 1.2 \). This can lead to
ambiguities in the multipolarity assignment, but one advantage is that the sign of the mixing ratio, \( \delta \) (Section 2.6), can in theory be determined for the mixed transitions [51]. A negative mixing ratio for an \( I \rightarrow I - 1 \) transition will lead to a reduction in the DCO ratio, such as for bands built on neutron configurations. A positive mixing ratio leads to an increased DCO ratio compared to that of the unmixed transition. This lends further support to the configuration assignments from the g-factor information obtained from the in-band branching ratios (Section 2.6).

### 4.3 Internal conversion

There are two possible electromagnetic decay modes for excited nuclear states, namely \( \gamma \)-ray emission and internal electron conversion. The latter of these involves processes where inner shell electrons are emitted directly from the atom. Electrons in the atomic s-shell have the highest probability of being emitted due to their large overlap with the nucleus. The energy with which the electrons emerge is the transition energy minus the electron binding energy, which is shell dependent. (Note that s-shell conversion electrons cannot be expelled when the transition energy is below the s-shell binding energy.) Characteristic X-rays accompany this process when the vacancy (left by the emitted electron) is filled by electrons from outer shells. (Aside: alternatively, Auger electrons rather than X-rays can be produced which are the atomic analogue of internal conversion electrons.)

Electron conversion competes effectively with (or dominates over) \( \gamma \)-ray emission for transitions of low-energy (\( \leq 200 \) keV) and/or high angular momentum transitions (\( \lambda > 2 \)). The probability of electron conversion is also higher (in general) for magnetic transitions than for electric transitions of the same multipolarity. For low-energy transitions the multipolarity can be determined by balancing the "missing" intensity into and out of a given state. This method has the advan-
tage of providing parity information (if there are significant differences between the possible conversion coefficients) which cannot be obtained from DCO ratios (Section 4.2). Electron conversion coefficients are tabulated in Ref.[53].
Chapter 5

$^{181}$Ta target isomers

5.1 Experimental method

Deep inelastic reaction experiments have been performed at the Argonne National Laboratory using a 1600 MeV pulsed $^{238}$U beam on the heaviest stable isotopes of ytterbium, lutetium, tantalum and tungsten. The targets used were $^{186}$W (16 mg cm$^{-2}$), $^{181}$Ta (23 mg cm$^{-2}$), $^{175}$Yb (6 mg cm$^{-2}$), and $^{177}$Lu (>20 mg cm$^{-2}$), and each target was backed by 43 mg cm$^{-2}$ of lead. The $^{238}$U beam from the ATLAS accelerator had a natural micro-pulsing period of 82.5 ns and was macro-pulsed by a beam sweeper with time ranges of 1.65 µs (4 pulses on, 16 pulses off), 16.5 µs (40 pulses on, 160 pulses off), 165 µs (400 pulses on, 1600 pulses off) and 1650 µs (4000 pulses on, 16000 pulses off). The beam energy was chosen to be 15 %$^a$ above the Coulomb barrier. Gamma-ray events were collected using the Argonne/Notre-Dame array of 12 Compton suppressed germanium detectors and a 50 element BGO inner ball. Of the 12 germanium detectors, all but the forward spectrometers had absorbers in front to reduce the number of X-rays measured; the forward emitted X-rays were absorbed by the lead backing on the target.

$^a$[54, 55, 56, 57, 58, 59] and others have used beam energies in the range 10-30 % above the Coulomb barrier successfully for deep inelastic reactions. The value of 15 % chosen here lies within this range.
The master trigger condition required at least one germanium detector firing in the beam-off interval. For coincidences between γ-ray events, there was a time window of 120 ns. The events obtained with each target were sorted into 4096 x 4096 matrices and sliced to give background-subtracted projected spectra. The first matrix was built with all γ-ray events in which two or more germanium detectors fired in the off-beam period. The analysis program XMESC [60] was used to slice these symmetrised matrices. A second matrix of γ-ray transition energy versus time was sorted to obtain half-life information.

The original aim of this experiment was to look for isomers primarily in the transfer products in the target region. However, the target nuclei were populated very strongly by inelastic excitation (which here is taken to include more complex transfer/evaporation processes, providing the final nucleus is the same as the initial target nucleus) and dominated the out-of-beam events. New high-spin isomeric states have been identified in all but $^{176}\text{Yb}$. (The 11 s two-quasiparticle isomer observed in the ytterbium target is known from previous work [61].) Three new isomers have been found in $^{181}\text{Ta}$ and two in $^{186}\text{W}$, with half-lives ranging from hundreds of nanoseconds to longer than 3 milliseconds. A new isomer was also observed in $^{175}\text{Lu}$. In addition to these new target isomers, a large range of known isomers in multi-particle transfer products have been observed, though so far these are restricted to target-like nuclei. This may be because transitions from isomers in the uranium region tend to be of low energy and consequently highly converted, below the sensitivity of this experiment. In addition, nuclei produced from transfer on the uranium beam can fission directly, and isomers resulting from this process have been identified. The results from the $^{186}\text{W}$ target are presented in Chapter 6 and those from the $^{181}\text{Ta}$ target can be found in the sections below. The $^{176}\text{Lu}$ isomer and other partial results from this work are reported in [62].
5.1.1 Cross-sections for isomeric states

For several of the new isomers observed with the tantalum target (this chapter) and the tungsten target (Chapter 6) population cross-sections are quoted. These are calculated using the formula,

\[ \sigma = \frac{A \epsilon \times 10^{24}}{I_{\text{beam}} N_{\text{target}} T_{\text{irrad}}} \]  

(5.1)

where \( A \) is the total intensity, in counts, of the de-populating transitions (including electron conversion), \( \epsilon \) is the absolute efficiency of the germanium detector array, \( I_{\text{beam}} \) is the beam intensity in particles per second, \( N_{\text{target}} \) is the number of target nuclei per square centimeter and \( T_{\text{irrad}} \) is the irradiation time in seconds. The factor of \( 10^{24} \) converts the cross-section to barns. The absolute efficiency of the Argonne/Notre-Dame array is estimated to be \( \epsilon = 0.35 \% \) at 1.3 MeV [63]. The effective target thickness for both tantalum and tungsten is taken to be \( 6 \text{ mg cm}^{-2} \). This is the thickness after which the \(^{238}\text{U}\) beam will be below the Coulomb barrier for the target nuclei. This value is approximate and therefore the final production cross-sections are quoted with a minimum uncertainty of \( \pm 10 \% \). The total transition intensities \( (A) \) are found from the peak areas in the off-beam singles spectrum. The areas are then corrected to account for electron conversion and scaled by the relative efficiency at 1.33 MeV. Using the singles spectrum avoids any further efficiency corrections due to higher fold events.

5.2 Results

The new isomeric states that have been observed in \(^{181}\text{Ta}\) are shown in Figure 5.1. They feed the previously known \( K^\pi = 9/2^- \) band [9], which has been extended here to \( 25/2^- \). A 1-quasiparticle \( K^\pi = 1/2^- \), 18 \( \mu \)s isomer [64] at 615 keV feeding the \( K^\pi = 7/2^- \) ground-state band [65], was also strongly populated. The \( K^\pi = 9/2^- \) bandhead is isomeric, decaying by a multiplicity one, 6.1 \( \mu \)s dipole
transition to the ground-state [66], but this decay was not observed due to its low energy (6 keV). The configuration of this state is $\pi\{\frac{9}{2}^-[514]\}$. More recent work by Dracoulis et al. [67] and Saitoh et al. [68] using the $^{170}\text{Yb}(^{11}B,\alpha 2n)^{181}\text{Ta}$ reaction has further extended the $^{181}\text{Ta}$ level scheme. The results from these studies are generally in good agreement with the results presented here and in Ref.[62]. Table 5.1 lists the transitions, together with the relative intensities and spin and parity assignments. The properties of the new isomeric states are discussed in Section 5.3. The half-lives for the new long-lived states are obtained from a least-squares fitting of the background subtracted time spectra. When making assignments only E1, M1 and E2 multipolarities are considered where prompt decays ($\leq 1$ ns) have been established. For states with multiple decays spin and parity assignments are chosen to give consistent relative $\gamma$-ray branching ratios. Finally, for low-energy transitions, intensity balancing is used to yield a total electron conversion coefficient (Section 4.3). Comparison with theoretical values [53] can give the multipolarity of the transition.

5.2.1 $K^\pi = 21/2^-$ isomer

The isomer at 1484 keV has been observed here with $t_{1/2} = 25\pm2$ $\mu$s, shown in Figure 5.2, populated with a cross-section of 56\pm6 mb. The half-life is confirmed by the $t_{1/2} = 23_{-2}^{+6}$ $\mu$s value published by Dracoulis et al. [67]. The spin and parity assignment has been determined from the relative transition intensities for the $\gamma$ rays de-populating the isomer. The strongest branch is via an intense 456 keV decay to the $19/2^-$ level in the $K^\pi = 9/2^-$ band. This competes with a 711 keV transition and a low energy 177 keV decay. A $21/2^+$ assignment is ruled out because a 711 keV magnetic quadrupole transition would not compete effectively with electric dipole transitions. For spin and parity assignments of $19/2^+$ and $19/2^-$ the multipolarity of the three transitions from the isomer would be E1 or M1 respectively, and the higher energy transition (711 keV) should be the most intense. A $K^\pi = 21/2^-$ assignment for the 25$\mu$s isomer is the only spin and
Table 5.1: Energies, assignments and relative $\gamma$-ray intensities for transitions observed in $^{181}$Ta. (The relative $\gamma$-ray intensities are those obtained from the 165 $\mu$s range data.)

<table>
<thead>
<tr>
<th>$E_\gamma$ (keV)</th>
<th>$E_i$</th>
<th>$E_f$</th>
<th>$I_\gamma$</th>
<th>$(I^*, K)_i$</th>
<th>$(I^*, K)_f$</th>
</tr>
</thead>
<tbody>
<tr>
<td>130.3</td>
<td>2229</td>
<td>2098</td>
<td>1.9(3)</td>
<td>29/2$^-$,29/2</td>
<td>(23/2$^+$,21/2)</td>
</tr>
<tr>
<td>151.6</td>
<td>(1955)</td>
<td>(1804)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>152.1</td>
<td>158</td>
<td>6</td>
<td>58.3(30)$^a$</td>
<td>11/2$^-$,9/2</td>
<td>9/2$^-$,9/2</td>
</tr>
<tr>
<td>176.9</td>
<td>1484</td>
<td>1307</td>
<td>7.8(5)</td>
<td>21/2$^-$,21/2</td>
<td>21/2$^-$,9/2</td>
</tr>
<tr>
<td>179.0</td>
<td>337</td>
<td>158</td>
<td>63.1(34)</td>
<td>13/2$^-$,9/2</td>
<td>11/2$^-$,9/2</td>
</tr>
<tr>
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<td>1403</td>
<td></td>
<td></td>
<td>(15/2$^-$,15/2)</td>
</tr>
<tr>
<td>205.0</td>
<td>542</td>
<td>337</td>
<td>49.4(21)</td>
<td>15/2$^-$,9/2</td>
<td>13/2$^-$,9/2</td>
</tr>
<tr>
<td>212.3</td>
<td>(1804)</td>
<td>(1591)</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>230.6</td>
<td>773</td>
<td>542</td>
<td>45.5(18)</td>
<td>17/2$^-$,9/2</td>
<td>15/2$^-$,9/2</td>
</tr>
<tr>
<td>255.1</td>
<td>1028</td>
<td>773</td>
<td>40.7(16)</td>
<td>19/2$^-$,9/2</td>
<td>17/2$^-$,9/2</td>
</tr>
<tr>
<td>279.4</td>
<td>1307</td>
<td>1028</td>
<td>13.2(7)</td>
<td>21/2$^-$,9/2</td>
<td>19/2$^-$,9/2</td>
</tr>
<tr>
<td>292.6</td>
<td>2098</td>
<td>1805</td>
<td>26.3(18)$^a$</td>
<td>(23/2$^+$,21/2)</td>
<td>(21/2$^+$,21/2)</td>
</tr>
<tr>
<td>295.4</td>
<td>2229</td>
<td>1933</td>
<td>13.3(6)</td>
<td>29/2$^-$,29/2</td>
<td>25/2$^-$,9/2</td>
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<tr>
<td>301.6</td>
<td>1608</td>
<td>1307</td>
<td>2.3(3)</td>
<td>23/2$^-$,9/2</td>
<td>21/2$^-$,9/2</td>
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<tr>
<td>321.5</td>
<td>1805</td>
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<td>31.5(18)</td>
<td>(21/2$^+$,21/2)</td>
<td>21/2$^-$,21/2</td>
</tr>
<tr>
<td>324.9</td>
<td>1933</td>
<td>1608</td>
<td>4.6(3)</td>
<td>25/2$^-$,9/2</td>
<td>13/2$^-$,9/2</td>
</tr>
<tr>
<td>331.8</td>
<td>337</td>
<td>6</td>
<td>3.6(5)</td>
<td>13/2$^-$,9/2</td>
<td>9/2$^-$,9/2</td>
</tr>
<tr>
<td>384.0</td>
<td>542</td>
<td>158</td>
<td>9.9(7)</td>
<td>15/2$^-$,9/2</td>
<td>11/2$^-$,9/2</td>
</tr>
<tr>
<td>435.7</td>
<td>773</td>
<td>337</td>
<td>14.2(9)</td>
<td>17/2$^-$,9/2</td>
<td>13/2$^-$,9/2</td>
</tr>
<tr>
<td>456.3</td>
<td>1484</td>
<td>1028</td>
<td>48.2(19)</td>
<td>21/2$^-$,21/2</td>
<td>19/2$^-$,9/2</td>
</tr>
<tr>
<td>485.7</td>
<td>1028</td>
<td>542</td>
<td>17.0(9)</td>
<td>19/2$^-$,9/2</td>
<td>15/2$^-$,9/2</td>
</tr>
<tr>
<td>534.4</td>
<td>1307</td>
<td>773</td>
<td>7.4(5)</td>
<td>21/2$^-$,9/2</td>
<td>17/2$^-$,9/2</td>
</tr>
<tr>
<td>580.8</td>
<td>1608</td>
<td>1028</td>
<td>1.8(2)</td>
<td>23/2$^-$,9/2</td>
<td>19/2$^-$,9/2</td>
</tr>
<tr>
<td>626.3</td>
<td>1933</td>
<td>1307</td>
<td>4.6(4)</td>
<td>25/2$^-$,9/2</td>
<td>21/2$^-$,9/2</td>
</tr>
<tr>
<td>711.2</td>
<td>1484</td>
<td>773</td>
<td>17.3(11)</td>
<td>21/2$^-$,21/2</td>
<td>17/2$^-$,9/2</td>
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<tr>
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<td>1403</td>
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<td>0.5(3)</td>
<td>(15/2$^-$,15/2)</td>
<td>15/2$^-$,9/2</td>
</tr>
<tr>
<td>1065.9</td>
<td>1403</td>
<td>337</td>
<td>6.9(11)</td>
<td>(15/2$^-$,15/2)</td>
<td>13/2$^-$,9/2</td>
</tr>
<tr>
<td>1245.1</td>
<td>1403</td>
<td>158</td>
<td>9.5(19)</td>
<td>(15/2$^-$,15/2)</td>
<td>11/2$^-$,9/2</td>
</tr>
</tbody>
</table>

$^a$ From intensity balancing due to contamination of peaks in the singles spectrum.
Figure 5.1: Decay scheme for the isomeric states in $^{181}$Ta discovered in the present work. The width of the arrows is proportional to the $\gamma$-ray intensities (black) and electron conversion (white), except for de-excitations from states represented by dashed lines for which the placement is tentative and intensities are generally weak.

parity that is consistent with the measured relative $\gamma$-ray intensities.

The 3-quasiproton configuration $\pi\{\frac{3}{2}^{-}[514], \frac{7}{2}^{+}[404], \frac{5}{2}^{+}[402]\}$ is proposed by comparison with blocked BCS calculations (see Section 5.3.3).

5.2.2 $K^\pi = 29/2^-$ isomer

An isomeric state with $t_{1/2} = 210\pm20$ $\mu$s, at 2229 keV has been observed that decays to the $I^\pi = 25/2^-$ member of the $K^\pi = 9/2^-$ band via a 295 keV $\gamma$-ray transition (see Figure 5.3). The observation of the 302 and 325 keV transitions and the corresponding E2 crossovers, which continue the systematics of the $K^\pi = 9/2^-$ band, allowed the band to be extended beyond the previously known $21/2^-$ state [9]. The production cross-section for this isomer is $30\pm5$ mb.

A 27/2$^-$ assignment for the isomeric state can be ruled out by the absence of
a 620 keV quadrupole transition to the $23/2^-$ state in the $K^\pi = 9/2^-$ band. A $27/2^+$ assignment is also unlikely because no competing E2 transition to the $23/2^+$ state at 1863 keV [67] in the ground-state ($K^\pi = 7/2^+$) band has been observed. This would be expected to compete with an E1 decay to the $K^\pi = 9/2^-$ band as for example in $^{175}$Hf [69] and $^{175}$Ta [70]. Intensity balancing for the 295 keV transition from the isomer yields an electron conversion coefficient of $\alpha_T(\text{exp}) = 0.02 \pm 0.16$ compared to theoretical values [53] of $\alpha_T(\text{E2}) = 0.09$, $\alpha_T(\text{E1}) = 0.02$, $\alpha_T(\text{M1}) = 0.22$ and $\alpha_T(\text{M2}) = 0.87$. This rules out an M2 assignment for the 295 keV transition that would result from a $29/2^+$ isomeric state. From this intensity analysis the favoured multipolarities are E2 and E1 although M1 decays are not discounted. Therefore, the most likely spin and parity for the 210 $\mu$s isomer is $29/2^-$ although this assignment is not firm. A low-lying $K^\pi = 29/2^-$ state with favoured residual interactions is predicted by blocked BCS calculations.
In addition to the 295 keV branch to the $K^* = 9/2^-$ band, the 210 $\mu$s state decays by a cascade of three $\gamma$ rays (130, 293 and 322 keV, shown in Figure 5.3) to the $K^* = 21/2^-$ isomer. (This results in a long half-life component in the time spectra of the 177, 456 and 711 keV transitions.) Due to the low-energy of the 130 keV transition, the intensity can be matched to that of the 293 and 322 keV $\gamma$ rays to find the implied electron conversion coefficient. This yields $\alpha_T(\text{exp}) = 18.4 \pm 3.3$, compared to theoretical values [53] of $\alpha_T(E3) = 19.3$, $\alpha_T(M2) = 15.3$. (This assumes $M1$ multipolarity for the 293 keV $\gamma$ ray, which is necessary to balance the intensity with the 322 keV supposed $E1$ transition, see below.) This favours an electric octupole assignment for the 130 keV transition, but a magnetic quadrupole transition cannot be discounted. In each case a parity change would occur, with the final state being either $23/2^+$ or $25/2^+$. The preference for $23/2^+$ depends on comparison with model calculations, as discussed in Section 5.3.3.

Considering all of the available evidence, the most likely decay route for the $K^* = 29/2^-$ isomer is via a 130 keV E3 transition to a $23/2^+$ rotational state (the M2 alternative would require an electric quadrupole transition), built on a $K^* = 21/2^+$ intrinsic level. This could decay by either a 293 or 322 keV $\gamma$ ray to the $K^* = 21/2^-$ 25 $\mu$s isomer. In Figure 5.1 the 293 keV decay has been assigned as the $23/2^+ \rightarrow 21/2^+$ transition, as this gives the closest match of the intensities in the off-beam singles spectrum, once conversion coefficients are taken into account. In addition, the energy of this transition is close to that of the corresponding decay in the $K^* = 9/2^-$ band (302 keV).

### 5.2.3 Other states

There is a $t_{1/2} = 170 \pm 10$ ns isomer (Figure 5.4) at 1403 keV, feeding the $15/2^-$, $13/2^-$ and $11/2^-$ members of the $K^* = 9/2^-$ band. This state has been observed in previous studies [71] but no spin, parity or half-life assignments were made. The 1245 keV transition (Figure 5.5) to the $11/2^-$ level is the most intense of the 3 $\gamma$
Figure 5.3: The top spectrum is gated by the 295 keV transition populating the $K^* = 9/2^-$ rotational band from the 210 $\mu$s isomer. The middle panel shows the spectrum gated by the 1066 keV transition from the (15/2$^-$) intrinsic state. The 152 and 179 keV $\gamma$-ray transitions from the first two excited states in the $K^* = 9/2^-$ band can be clearly seen. The lower spectrum is gated by the 293 keV transition lying between the 210 $\mu$s and the 25 $\mu$s isomeric states. All labelled peaks are assigned to $^{181}$Ta.
rays de-populating this isomer. This rules out a 15/2+ state because an M2 transition is unlikely to compete effectively with E1 transitions. A 13/2- assignment is unlikely because no competing 1397 keV E2 transition decaying directly to the \( K^\pi = 9/2^- \) bandhead has been observed. A spin and parity assignment of 13/2+ for the 1403 keV level would result in the three observed transitions (861, 1066 and 1245 keV, see Figures 5.3 and 5.5) having E1 character. Competing magnetic dipole transitions would be expected to populate the corresponding ground-state band levels, but these transitions have not been observed. An assignment of 15/2- gives ratios of the transition rates to the Weisskopf single-particle rates that are in close agreement with the values tabulated by Löbner [72]. Comparison with Nilsson model calculations (Section 5.3.3) yields a favoured \( K^\pi = 15/2^- \) state with the configuration \( \nu \{ \frac{1}{2}^+ [624], -\frac{1}{2}^- [510] \} \otimes \pi \{ \frac{7}{2}^+ [404] \} \).

![Figure 5.4: A time spectrum gated by the 1066 and 1245 keV γ-ray transitions de-populating the \( K^\pi = (15/2^-) \) state. The line through the data is for a 170 ns half-life.](image-url)
Figure 5.5: A $\gamma$-ray spectrum gated by 152 keV. This transition is observed in the decay of all three new isomers. This $\gamma$-ray energy is also in coincidence with itself, because it is one of the transitions above the 170 ns isomer, as well as the lowest $\Delta I = 1$ transition in the $K^* = 9/2^-$ band. All labelled peaks are assigned to $^{181}$Ta.

The level at 1403 keV has been observed by Saitoh et al. [68], with a $15/2^-$ assignment. No life-time was associated with this level, but it was reported to be fed by a $t_{1/2} = 140\pm36$ ns ($19/2^-$) isomer, that de-excites by a low-energy transition ($<50$ keV, unless obscured by higher energy $\gamma$ rays). This half-life is consistent with the measurement obtained from the current experiment. Dracoulis et al. [67] have reported two states lying at 1403 keV ($17/2$) and 1404 keV ($15/2$) respectively. Due to the similarity of the decay paths, an observed mean-life of 4.7 ns could not be uniquely assigned. A 69 keV transition was observed to decay to the 1404 keV level but no half-life measurement was obtained. It is evident that this portion of the level scheme is complex and not uniquely
established.

The states above the 1403 keV level are based on $\gamma$-$\gamma$ coincidence relations. In both the 1245 and 152 keV gates, there are $\gamma$ rays with energies of 188, 212 and 152 keV (the 152 keV transition is in coincidence with itself, see Figure 5.5). These transitions are much weaker than those in the $9/2^-$ band, but fitting the time projections gives a half-life of $\sim 100 \mu$s. This half-life is close to that of the $K^x = 29/2^-$ state with a 210 $\mu$s half-life, but due to the weak population of the levels above the 1403 keV state, it is not clear whether they are being fed by a third decay route from the $K^x = 29/2^-$ isomer or by another, different isomer. The ordering of these three levels in Figure 5.1 is tentative. Table 5.1 gives the relative intensities of the $\gamma$-rays in Figure 5.1, except those above the $(15/2^-)$ state. These intensities are not listed because of ambiguities in the coincidence spectra. The 188 and 212 keV intensities are not consistent with the single decay path shown in Figure 5.1. This problem has also been eluded to by Dracoulis et al. [67].

5.2.4 Transfer products

This experiment was initially performed with the aim of looking for new isomeric states formed from multi-nucleon transfer between the target and projectile. The isomeric states identified in the $^{181}$Ta target region are listed in Table 5.2. The sensitivity of this experiment extends only to those products containing isomers with half-lives greater than approximately a hundred nanoseconds. Many nuclei are observed in these data (listed in Table 5.2), with a typical production cross-section of $\sim 0.2$ mb, (observed for the $15^- 45 \mu$s isomer in $^{180}$Ta [73]) but no new isomers, populated by nucleon transfer, have been assigned using the $^{181}$Ta target. Nonetheless, the large range of known isomers populated, has established this technique as being suitable for populating high-spin intrinsic states in the $A\approx 180$ region. Isomers populated by the fission of $^{238}$U are given in Table 5.3. A cascade of unplaced $\gamma$ rays shown at the bottom of Table 5.3 is assumed to be
5.3 Discussion

In the sections below the new isomers are compared to those in neighbouring nuclei and the g-factor for the $K^\pi = 9/2^-$ band is calculated. The new intrinsic states are compared to predictions by blocked BCS calculations in Section 5.3.3.

5.3.1 Hindrances for isomeric decays

The 3-quasiproton configuration $\pi\{\frac{3}{2}^-[514], \frac{3}{2}^+[404], \frac{3}{2}^+[402]\}$ has been proposed for the new $K^\pi = 21/2^-$ 25 $\mu$s isomer at 1484 keV based on comparisons with Nilsson model calculations (see Section 5.3.3). $K^\pi = 21/2^-$ isomers have been observed in the isotopes $^{179}$Ta [75] and $^{177}$Ta [76] with half-lives of 320 ns (at 1252 keV) and 5 $\mu$s (at 1355 keV) respectively, both of which have the same 3-quasiproton assignment. The 711 keV stretched E2 transition from the $K^\pi = 21/2^-$ state in $^{181}$Ta to the $K^\pi = 9/2^-$ band has a reduced hindrance of $f_\nu = 41$ (Section 2.5). The corresponding isomeric state in $^{179}$Ta has $f_\nu = 7$ for the 475 keV stretched quadrupole transition, whereas in $^{177}$Ta, the 550 keV $\Delta I = 2$ transition from the $K^\pi = 21/2^-$ isomer has $f_\nu = 22$. Although there is no steady increase in the reduced hindrance with mass number, the large value for $^{181}$Ta indicates that high neutron numbers are not diminishing the conservation of $K$.

The $K^\pi = 29/2^-$ 210 $\mu$s isomer at 2229 keV has the preferred 3-quasiparticle assignment $\nu\{\frac{11}{2}^+[615], \frac{9}{2}^+[624]\} \otimes \pi\{\frac{9}{2}^-[514]\}$ (see Section 5.3.3). This involves two $i_{13/2}$ neutrons making this a $t$-bandhead (Section 2.3.1). An isomeric state based on this configuration has also been observed in the isotope $^{183}$Se [77],

...
Table 5.2: Transitions from isomers observed in the reaction $^{238}\text{U} + ^{181}\text{Ta}$ at 1600 MeV (1.65 $\mu$s range).

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>Particles transferred</th>
<th>$E_\gamma$ (keV)</th>
<th>$I_\gamma$</th>
<th>Half-life</th>
<th>Assignment, $I^\pi$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{181}\text{Ta}$ (target)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>133</td>
<td>19(3)</td>
<td>18 $\mu$s</td>
<td>1/2$^+$</td>
<td>5/2$^+$</td>
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<tr>
<td></td>
<td>861</td>
<td>1(2)</td>
<td>170 ns</td>
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<td>15/2$^-$</td>
</tr>
<tr>
<td></td>
<td>1066</td>
<td>6(3)</td>
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<td>(15/2$^-$)</td>
<td>13/2$^-$</td>
</tr>
<tr>
<td></td>
<td>1245</td>
<td>18(7)</td>
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<td>(15/2$^-$)</td>
<td>11/2$^-$</td>
</tr>
<tr>
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<tr>
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<td>456</td>
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<td>19/2$^-$</td>
</tr>
<tr>
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<td>($\sim$100 $\mu$s)</td>
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<td></td>
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<td></td>
</tr>
<tr>
<td>$^{182}\text{W}$ (+1p)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{180}\text{W}$ (+1p-2n)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

* From the recent work of D'Alarcao et al. [74].

b From this work [62].

* The transition is not one which directly de-populates the isomer due to very weak intensities or low transition energies for the de-populating transitions.
Table 5.3: Isomeric states in nuclei populated following the fission of $^{235}$U. The half-lives and $\gamma$-ray transition energies are shown. Relative intensities are not given because all the products listed are weakly populated.

<table>
<thead>
<tr>
<th>Nuclei</th>
<th>$\tau_{1/2}$ (s)</th>
<th>$\gamma$-rays (keV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{128}$Sn</td>
<td>11.8 $\mu$s</td>
<td>90, 197, 1023, 1171</td>
</tr>
<tr>
<td>$^{132}$Sn</td>
<td>7.2 $\mu$s</td>
<td>104, 163, 1002, 1141</td>
</tr>
<tr>
<td>$^{136}$Sn</td>
<td>6.6 $\mu$s</td>
<td>57, 112, 909, 1141</td>
</tr>
<tr>
<td>$^{137}$Sb</td>
<td>11 $\mu$s</td>
<td>806, 1114</td>
</tr>
<tr>
<td>$^{132}$Te</td>
<td>10.4 min $(^{128}$Sb $\beta$ decay)</td>
<td>314, 743, 754</td>
</tr>
<tr>
<td>$^{130}$Te</td>
<td>115 ns</td>
<td>182, 331, 794, 839</td>
</tr>
<tr>
<td>$^{132}$Te</td>
<td>28.1 $\mu$s/145 ns</td>
<td>103, 151, 697, 974</td>
</tr>
<tr>
<td>$^{134}$Te</td>
<td>165 ns</td>
<td>115, 297, 1279</td>
</tr>
<tr>
<td>$^{135}$I</td>
<td>9 s</td>
<td>74, 648, 913</td>
</tr>
<tr>
<td>$^{132}$Xe</td>
<td>8.39 ms/90ns</td>
<td>174, 538, 600, 773, 668</td>
</tr>
<tr>
<td>$^{134}$Xe</td>
<td>290 ms</td>
<td>234, 847, 884</td>
</tr>
<tr>
<td>$^{136}$Xe</td>
<td>2.95 $\mu$s</td>
<td>197, 381, 1313</td>
</tr>
<tr>
<td>$^{54}$Xe</td>
<td></td>
<td>187, 361, 503, 617</td>
</tr>
</tbody>
</table>

with a half-life of 6 ns [78]. The shorter half-life in $^{185}$Re is understood to be a consequence of the close-lying $25/2^+$ band, which provides a $K$-allowed decay path. The 295 keV stretched E2 transition in $^{181}$Ta has $f_\nu = 5$. This is in good agreement with other $\Delta K = 10$ stretched E2 decays from $t$-bandheads that have been observed in this mass region. The $K^x = 10^+ 1.4 \mu$s isomer in $^{182}$W [22] has a 1086 keV stretched quadrupole decay with $f_\nu = 5.3$. The $K^x = 25/2^+$ 1 ms isomer in $^{183}$Re [79] decays by a 194 keV stretched E2 transition to the $K^x = 5/2^+$ ground-state band with $f_\nu = 3.8$. A 20 ns $K^x = 10^+$ isomer in $^{184}$Os [80] de-excites via a 1092 keV stretched E2 transition with a reduced hindrance of $f_\nu = 3.5$. These comparisons seem to point to a general reduction in the reduced hindrance with increasing proton number. This corresponds to a fall in the number of valence protons and neutrons as particles are added moving away from the mid-shell. However, the $f_\nu$ values for all of these $t$-bands are very low, indicating that $K$ is no longer a conserved quantity. One mechanism which may be responsible for this could be related to the high alignment of the $i_{13/2}$ neutrons,
leading to increased Coriolis mixing between nearby low-\(K\) states. Although this has not been established, it seems likely that the neutrons in these configurations are having some effect because the 711 keV E2 decay from the \(K^\pi = 21/2^-\) isomer in \(^{181}\)Ta has a high reduced hindrance \((f_\nu = 41)\) suggesting that the high neutron number has a very limited effect on \(K\)-conservation.

5.3.2 Branching ratios and g-factors

For intrinsic states for which the accompanying rotational band has been observed, the in-band \(\gamma\)-ray branching ratio, \(T_2/T_1\), can be used to test configuration assignments. (The subscripts refer to the spin change associated with the transition.) The rotational model expressions (Equations (2.21) and (2.22)) are used with typical values \([22]\) of \(Q_0 = 6.5\) e-b and \(g_R = 0.3\). The calculated values in Table 5.4 are obtained using the asymptotic Nilsson quantum numbers with Equation (2.29).

Table 5.4: Branching ratios and g-factors for the \(K^\pi = 9/2^-\) band and the transitions lying above the \(K^\pi = 21/2^-\) state.

| \(K^\pi\) | \(I\) | \(E_2\) | \(E_1\) | \(T_2/T_1\) | \(|(g_K - g_R)/Q_0|\) (e-b)\(^{-1}\) |
|---|---|---|---|---|---|
| 9/2^- | 13/2 | 0.3318 | 0.1790 | 0.06±0.01 | 0.19±0.02 | +0.16 |
| 15/2 | 0.3840 | 0.2050 | 0.20±0.02 | 0.16±0.01 |
| 17/2 | 0.4357 | 0.2306 | 0.31±0.07 | 0.16±0.02 |
| 19/2 | 0.4857 | 0.2551 | 0.42±0.03 | 0.17±0.01 |
| 21/2 | 0.5344 | 0.2794 | 0.56±0.09 | 0.17±0.02 |
| 23/2 | 0.5808 | 0.3016 | 0.78±0.13 | 0.16±0.02 |
| 25/2 | 0.6263 | 0.3249 | 1.00±0.11 | 0.15±0.01 |
| 21/2^- | 25/2 | 0.6141 | 0.3215 | \(<0.026±0.003^a\) | \(>0.18±0.01\) | +0.12 |

\(a\) This assumes that the 293 and 322 keV transitions are rotational members of the \(K^\pi = 21/2^-\) band. The limit on the intensity of the E2 crossover was found by inspection of the off-beam singles spectrum. See Section 5.3.2.

The experimental \(|(g_K - g_R)/Q_0|\) values for the \(K^\pi = 9/2^-\) band are in excellent agreement with the calculated value based on the \(\pi\{9/2^-[514]\}\) configuration.
This provides additional confidence in the assignment of the $23/2^-$ and $25/2^-$ levels as rotational members of this band.

In nuclei where $K^* = 21/2^-$ isomers have been observed (see above), the first two excited states of the rotational bands built on these levels decay by transitions with energies very close to the 293 and 322 keV $\gamma$ rays observed above the 25 $\mu$s isomer in $^{181}$Ta. If the 293 and 322 keV transitions are magnetic dipole transitions in a rotational band, then a corresponding 614 keV E2 crossover transition should be present in the data. Examination of the out-of-beam singles spectrum leads to a limit on the supposed in-band $\gamma$-ray branching ratio of $\lambda \leq 0.026 \pm 0.003$. Using the rotational model Equations (2.21) and (2.22) implies $(g_K - g_R)/Q_0 \geq 0.18 \pm 0.01$ (using $Q_0 = 6.5$ e·b and $g_R = 0.3$). (This calculation is given in Table 5.4.) This is higher than the expected g-factor using both the asymptotic and non-asymptotic Nilsson wavefunctions [23] with Equation (2.29), both of which give $(g_K - g_R)/Q_0 = 0.12$. Although the discrepancy between the experimental and calculated values is small (0.06), the good agreement achieved for the $K^* = 9/2^-$ band suggests that the difference is significant. From this analysis it seems probable that the 293 and 322 keV transitions are not both members of the same rotational band. This supports the assignment of the 322 and 293 keV as transitions to and from an intrinsic state respectively. This assignment is discussed in Section 5.2.2.

The well-deformed nature of this nucleus leads to strongly coupled rotational bands for which the asymptotic Nilsson quantum numbers are a good approximation.

### 5.3.3 Multi-quasiparticle calculations

Blocked BCS calculations [31], outlined in Section 3.1, have been performed for $^{181}$Ta$_{106}$ using deformation parameters $\epsilon_2 = 0.233$ and $\epsilon_4 = 0.060$ [81]. The Nilsson single-particle energies were adjusted (consistently with the monopole pairing strengths) to reproduce the single-quasiparticle proton and neutron ener-
gies (where available [9, 82]) in $^{181}$Ta and $^{181}$Hf respectively. The proton pairing strength, $G_p = 23.0/A$ MeV, was adjusted to give approximately the correct energy of the $21/2^-$ isomer. The neutron pairing strength was chosen to be $1.0/A$ MeV lower than the corresponding proton value ($G_n = 22.0/A$ MeV), which is consistent with Ref.[31]. The resulting near-yrast multi-quasiparticle states are represented in Figure 5.6.

From Figure 5.6 there are two low-lying 3-quasiparticle configurations with $K^\pi = 21/2^-$ and $29/2^-$ respectively, in agreement with the new isomeric states found in this experiment. Above the yrast line there is a $K^\pi = 21/2^+$ state with favoured residual interactions calculated at 1872 keV (see Section 5.2.2). The $K^\pi = 15/2^-$ candidate for the 170 ns isomer does not appear in Figure 5.6 because only couplings to maximum $K$ are shown. However, the corresponding unfavoured coupling to $K^\pi = 17/2^-$ is labelled.

The only favoured intrinsic state calculated to lie between the $K^\pi = 21/2^-$ and $K^\pi = 29/2^-$ isomeric states is a $K^\pi = 21/2^+$ 3-quasiparticle level with the configuration, $\nu\{\frac{3}{2}^+[624], \frac{3}{2}^-[512]\} \otimes \pi\{\frac{3}{2}^-[514]\}$, predicted to lie at 1872 keV. From comparison with these calculations this configuration is the preferred coupling for the 1805 keV level tentatively assigned as $K^\pi = 21/2^+$ (see Section 5.2.2).

Table 5.5 summarises the information for the new isomers. The residual interactions used by Dracoulis et al. [67] shift the energies by $-77$ keV for the $K^\pi = 21/2^-$ state and by $-14$ keV for the $K^\pi = 29/2^-$ configuration. These residual interactions are not included in the values given in Table 5.5. It can be seen that the excitation energies are in agreement with the experimental energies to within $\sim 200$ keV.

It is interesting to note that the Nilsson model calculations using Lipkin-Nogami pairing, performed by Dracoulis et al. [67] predict an yrast $37/2^+$ state resulting from the $\nu\{\frac{3}{2}^+[624], \frac{7}{2}^-[503]\} \otimes \pi\{\frac{3}{2}^-[514], \frac{7}{2}^+[404], \frac{5}{2}^+[402]\}$ 5-quasiparticle coupling. However, the results of the blocked BCS calculations here
Figure 5.6: Near-yrast multi-quasiparticle states from the blocked BCS calculation. The low-lying $29/2^-$ and $21/2^-$ configurations are in agreement with the observed structure. Although there is no low-lying $15/2^-$ state, only couplings to maximum $K$ are shown. The $17/2^-$ configuration shown here, could also form a favoured $15/2^-$ state.

yield a low-lying $K^\pi = 41/2^-$ state (below the $37/2^+$ level) with a $\nu\{11/2^+[615], 9/2^+[624]\} \otimes \pi\{9/2^-[514], 7/2^+[404], 5/2^+[402]\}$ configuration. (The residual interactions would be expected to be approximately the same for both states.) Although not observed in this experiment these potentially long-lived states provide an objective for future experiments and a practical test of the two calculations.
Table 5.5: Properties of the new isomeric states in $^{181}$Ta. Only the most intense E2 ($I \rightarrow I - 2$) decays are given in this table. $t^\gamma_{1/2}$ is the partial $\gamma$-ray half-life and $f_\nu$ is the reduced hindrance (Section 2.5).

<table>
<thead>
<tr>
<th>$K^\pi$</th>
<th>$t^\gamma_{1/2}$</th>
<th>Configuration</th>
<th>Experimental Energy (keV)</th>
<th>Calculated $E_\gamma$(E2) Energy (keV)</th>
<th>$t^\gamma_{1/2}$</th>
<th>$\Delta K$</th>
<th>$f_\nu$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$21/2^-$</td>
<td>25±2 μs</td>
<td>$\pi 9/2^-[514]$</td>
<td>1484</td>
<td>1408</td>
<td>711</td>
<td>0.2 ms</td>
<td>6 41</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \pi 7/2^+[404]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \pi 5/2^+[402]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$29/2^-$</td>
<td>210±20 μs</td>
<td>$\pi 9/2^-[514]$</td>
<td>2229</td>
<td>2053</td>
<td>295</td>
<td>0.8 ms</td>
<td>10 5</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \nu 11/2^+[615]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \nu 9/2^+[624]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>(15/2$^-$)</td>
<td>170±10 ns</td>
<td>$\pi 7/2^+[404]$</td>
<td>1403</td>
<td>1546</td>
<td>(1245)</td>
<td>300 ns</td>
<td>(3)</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \nu 9/2^+[624]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\otimes \nu 1/2^-[510]$</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

5.4 Summary

New isomeric states with $K^\pi = (15/2^-)$, $21/2^-$ and $29/2^-$ have been observed in $^{181}$Ta. Configurations have been assigned by comparisons with the isotopes $^{177}$Ta and $^{179}$Ta, the isotone $^{183}$Re and predictions by blocked BCS calculations. The excitation energies are well reproduced and recent work by Dracoulis et al. [67] and Saitoh et al. [68] is in good agreement with the level scheme deduced here.
6. **Results**

A new $K^* = 7^-$ isomeric state has been observed that populates the previously known $\gamma$ [83] and octupole [85] vibrational bands. Above this, a new $K^* = (16^+)$
isomer has been identified. The level scheme resulting from the present work is shown in Figure 6.1 and the transitions de-populating the new isomers are listed in Table 6.1. Figure 6.2 shows the singles spectra for the 165 μs range. The delayed transitions can be cleanly picked out by gating between the beam pulses, where as the total singles spectrum is dominated by prompt events from Coulomb excitation.

![Diagram](image)

Figure 6.1: Decay scheme for the 7⁻ and (16⁺) isomers observed in ¹⁸⁶W. The implied 9 keV transition from the higher isomer to the state at 3535 keV (14⁺) has been omitted because it was not observed in these data. The width of the arrows is proportional to the γ-ray intensities (black) and electron conversion (white).

6.1.1 \( K^\pi = 7^- \) isomer

An isomeric state lying at 1517 keV has been observed with \( t_{1/2} = 18\pm1 \) μs (Figure 6.3) and a cross-section of 70±10 mb (see Section 5.1.1). This isomer populates the 6⁺ member of the ground-state band (g-band) directly via a 708
Table 6.1: Energies, assignments and relative intensities for transitions observed in $^{185}\text{W}$ (165 $\mu$s range data).

<table>
<thead>
<tr>
<th>$E_\gamma$ (keV)</th>
<th>$E_i$</th>
<th>$E_f$</th>
<th>$I_\gamma$</th>
<th>$(I^*, K)_i$</th>
<th>$(I^*, K)_f$</th>
</tr>
</thead>
<tbody>
<tr>
<td>92.1</td>
<td>1045</td>
<td>953</td>
<td>4.1(4)</td>
<td>3+,2</td>
<td>2+,2</td>
</tr>
<tr>
<td>118.7</td>
<td>1517</td>
<td>1398</td>
<td>65.1(45)</td>
<td>7+,7</td>
<td>6+,2</td>
</tr>
<tr>
<td>122.1</td>
<td>122</td>
<td>0</td>
<td>203(16)</td>
<td>2+,0</td>
<td>0+,0</td>
</tr>
<tr>
<td>150.1</td>
<td>1322</td>
<td>1171</td>
<td>10.4(11)</td>
<td>5-,2</td>
<td>4-,2</td>
</tr>
<tr>
<td>164.9</td>
<td>2838</td>
<td>(2673)</td>
<td>9.2(9)</td>
<td>(12,-10)</td>
<td>(11,11)</td>
</tr>
<tr>
<td>167.6</td>
<td>2286</td>
<td>2118</td>
<td>58.2(38)</td>
<td>(10,-10)</td>
<td>(9,-9)</td>
</tr>
<tr>
<td>180.0</td>
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<td>3364</td>
<td>10.2(10)</td>
<td>(16+,16)</td>
<td>(14+,13)</td>
</tr>
<tr>
<td>182.9</td>
<td>1045</td>
<td>862</td>
<td>31.9(25)</td>
<td>3-,2</td>
<td>3+,2</td>
</tr>
<tr>
<td>195.2</td>
<td>1517</td>
<td>1322</td>
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<td>7-,7</td>
<td>5-,2</td>
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<td>4-,2</td>
<td>2-,2</td>
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<td>2+,2</td>
</tr>
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<td>3-,2</td>
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<td>(12-,10)</td>
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<td>4+,2</td>
</tr>
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<td>((11,12),(11,12))</td>
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<td>(8,-7)</td>
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<td>(13+,13)</td>
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<td>4+,2</td>
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<td>809</td>
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<td>3+,2</td>
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<td>(7,-7)</td>
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<td>1007</td>
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<td>81.3(63)</td>
<td>4+,2</td>
<td>4+,0</td>
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<td>32.6(99)</td>
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<td>2+,0</td>
</tr>
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<td>621.4</td>
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<td>(13+,13)</td>
<td>((11,12),(11,12))</td>
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</tr>
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<td>925.8</td>
<td>1322</td>
<td>396</td>
<td>10.9(20)</td>
<td>5-,2</td>
<td>4+,0</td>
</tr>
<tr>
<td>1001.4</td>
<td>1398</td>
<td>396</td>
<td>17.2(32)</td>
<td>6+,2</td>
<td>4+,0</td>
</tr>
</tbody>
</table>

* From intensity balancing due to contamination of peaks in the singles spectrum.
keV γ-ray transition. The g-band has previously been observed up to 14 h in a Coulomb excitation experiment using $^{208}$Pb beams [86].

The strongest direct branch from the 18 μs isomer is via a 195 keV transition to the $K^\pi = 2^-$ octupole band (Figure 6.4), known previously up to the 4$^-$ state at 1171 keV [9]. Observation of the 150 keV decay to the 4$^-$ level and the accompanying 277 keV E2 crossover transition to the 1045 keV 3$^-$ state allows the band to be extended to incorporate the 5$^-$ rotational state at 1322 keV. From intensity balancing, the implied electron conversion coefficient for the 195 keV transition is $\alpha_T(\text{exp})=0.25\pm0.11$ (the theoretical values [53] are $\alpha_T(\text{E2})=0.34$, $\alpha_T(\text{M1})=0.73$, $\alpha_T(\text{E1})=0.07$), favouring an electric quadrupole assignment. This leads to a $K^\pi = 7^-$ assignment for the 18 μs intrinsic state.

It should be noted here that in studies of the β-decay of $^{186}$Ta to states in $^{186}$W [87, 88], there is a 277 keV transition to the 2$^+$ γ vibrational level at 738 keV. Such a transition is not observed in the coincidence data obtained from the present study, implying a second transition at this energy.

A previous ($^{185}$Ta β decay) study [88] has observed a 315 keV γ-ray decay from a state at 1322 keV to a level at 1007 keV, as found in this work. However, the other reported decays from this level are not observed here (e.g. the 1322 keV decay directly to the ground-state) and are inconsistent with an $I^\pi = 5^-$ assignment. In addition, the 150 keV and 277 keV γ-ray transitions observed here, were not seen in Ref.[88], which, given their intensity relative to the 315 keV transition, would be remarkable if the same state was being populated in each case. Therefore, it seems highly likely that a second level at 1322 keV exists.

The γ vibrational state at 738 keV has been identified in earlier studies together with the first excited state at 862 keV [83], and Günther et al. [89], tentatively assigned a level lying at 1030 keV as the 4$^+$ member of this band, following a deuteron scattering experiment. However, the authors of Ref.[88] comment that the decay of the 1030 keV state does not follow the expected systematics for a 4$^+$ γ vibrational level. In the present study, states at 1007 and 1398 keV have
Figure 6.2: The top spectrum shows the off-beam gated singles spectrum for the 165 μs range data. All of the labelled energies have been placed in $^{186}$W. The lower panel shows the total singles spectrum dominated by the prompt Coulomb excitation of $^{186}$W (energies) and $^{238}$U (Δ). The time for which each data set was recorded is given.

been assigned as the 4+ and 6+ members of the γ vibrational band respectively. The 1398 keV level was reported by Günther et al. [89] but no assignment was made. The 1007 keV level was tentatively assigned as the 2+ member of a $K = 0$ β vibrational band [89], but the observation of a strong transition to this level from the 5− member of the octupole band contradicts this assignment.

A direct 119 keV decay path from the $K^\pi = 7^-$ isomer (Figure 6.4) feeds the 6+ level of the γ band, which in turn de-excites to both the 4+ and 6+ members of the g-band and to the 4+ rotational level in the γ band. Intensity balancing for the 119 keV transition yields an electron conversion coefficient of $\alpha_T(\text{exp})=0.27 \pm 0.12$ ($\alpha_T(\text{E1})=0.24$, $\alpha_T(\text{M1})=2.99$, $\alpha_T(\text{E2})=2.03$) establishing this as an electric dipole
Figure 6.3: A time spectrum gated on γ rays below the 18 μs isomer in $^{186}$W. A long-lived component is evident. The line through the data is the fit for a two component half-life.

The preferred configuration for the $K^\pi = 7^-$ isomeric state is $\nu\{\frac{11}{2}^+|[615], \frac{7}{2}^-[512]\}$. Comparison with multi-quasiparticle Nilsson calculations is made in Section 6.2.3. A $K^\pi = 7^- 2.4$ ns isomer at 1502 keV in $^{184}$W has been assigned the same 2-quasineutron configuration [90].

### 6.1.2 $K^\pi = (16^+)$ isomer

The time spectrum gated by transitions below the $K^\pi = 7^-$ isomer, Figure 6.3, shows a two component lifetime comprising $t_{1/2} = 18 \mu$s and a much longer half-life ($\gg 100 \mu$s). Examination of the off-beam gated singles spectrum for the
1.65 ms range (Figure 6.5) shows many previously unplaced $\gamma$-rays of comparable intensity (given the long beam-pulsing period) to the transitions de-exciting the $K^* = 7^-$ isomer. The energies of these lines do not match those known in $^{185}$W or $^{187}$W, which are populated by 1-neutron transfer and are expected to be the next strongest channels after the inelastic excitation of $^{186}$W. The coincidence spectra for these transitions (e.g. Figure 6.5) show that they form a long cascade but with no simple rotational structure. The time spectra for these previously unplaced decays are compatible with the long half-life component observed for the transitions below the 18 $\mu$s level.

A comparison of the intensities for the unplaced $\gamma$-ray transitions and those below the $K^* = 7^-$ isomer reveals a good match at long times (>100 $\mu$s after the beam pulse). See Figure 6.6 for the overlaid time spectra gated by the 195 keV transition (below the 18 $\mu$s isomer) and the 221 keV (previously unplaced) transition. For times greater than 100 $\mu$s, the matching half-lives and intensities, together with the strength of the transitions in the off-beam gated singles spectrum, suggest that this previously unplaced cascade does in fact feed through the $K^* = 7^-$ 2-quasiparticle state in $^{185}$W. This new isomer is populated with a cross-section of 12±2 mb (Section 5.1.1).
CHAPTER 6. $^{186}$W TARGET ISOMERS

![Energy spectra](image)

Figure 6.5: The top spectrum shows the off-beam gated singles spectrum for the 1.65 ms range data. The strongest long-lived contaminants are denoted by $\Delta$. The lower spectrum is gated by the 552 keV $\gamma$-ray now placed in $^{186}$W. This is for the 165 $\mu$s range for greater statistics.

Adding up the $\gamma$-ray energies in the cascade leads to an excitation energy of 3544 keV for the new isomeric level.

6.1.3 Additional half-life measurements

Millisecond analysis

Fitting the background subtracted 1.65 ms range time spectra, gated by the 168, 221, 380, 399 and 552 keV transitions that feed the 1517 keV isomer, yields $t_{1/2}$=5 ms for the level at 3544 keV. However, the half-life due to the TAC* random stopping rate (estimated from an event rate of 75±13 events/sec in the

*A TAC (Time-to-Amplitude Converter) is the electronics module supplying the time information in this experiment.
beam-off period) was \( t_{1/2}^{\text{rand}} = 9 \pm 1 \text{ ms} \). To establish whether the measured half-life was real or just a random stopping effect, two \( \gamma \)-ray spectra were made from time gates, one on the first 1/3 and one on the last 1/3 of the time projection. The factor of 1/3 was used to give the optimum statistical accuracy in the spectra while retaining a large effect due to the time difference. The ratio of the \( \gamma \)-ray peak areas in the first 1/3 to the last 1/3 were calculated for both \( ^{186}\text{W} \) and known radioactivity peaks to see if there were any systematic differences between them due to the half-life. However, from Figure 6.7, it can be seen that the uncertainties in the ratios are too large, and so only a lower limit can be found for the half-life from these data, namely \( t_{1/2} \geq 3 \text{ ms} \), which is shown on the plot as the long-dashed line.

Radioactivity analysis

Radioactivity spectra were collected to look for very long-lived states once the beam was turned off. These were recorded for different time segments, from 5 minutes, up to a few hours after the beam was stopped. From the radioactive decay law, the average count rate in a time interval \( T \) ranging from an initial time \( t_1 \) up to a final time \( t_2 \) is given by,

\[
R_{av} = \frac{1}{T} \int_{t_1}^{t_2} \frac{dN}{dt} \, dt = \frac{1}{T} \int_{t_1}^{t_2} \lambda N \, dt
\]

where \( N = N_0 e^{-\lambda t} \), \( \lambda \) is the decay rate and \( N_0 \) is the initial number of nuclei. However, \( N_0 \) is not known, but the initial count rate \( \left( \frac{dN}{dt} \right)_{N=N_0} \) can be found assuming that the irradiation time is long compared to the half-life. Substituting this in place of \( \frac{dN}{dt} \) leads to the result

\[
R_{av} = \left( \frac{dN}{dt} \right)_{N=N_0} \left[ e^{-\lambda t_1} - e^{-\lambda t_2} \right]
\]

This method was used to set an upper limit on the half-life of the \( ^{186}\text{W}, 3544 \text{ keV} \) isomer. The isomer evidently has a half-life less than minutes, as no statistically
Figure 6.6: A graph of intensity versus time for the 195 keV transition (below the 7⁻ isomer) and 221 keV transition (above the 7⁻ isomer). The intensities are approximately the same at long times. Although no correction has been made for the change in efficiency and conversion coefficients, the energies are close enough for these effects to be small. The slightly reduced count level for times before \( \approx 60 \mu s \) is a dead-time (or pile-up) effect arising from the high count-rate during the immediately preceding beam-on period (ending at Time=0).

significant decay transitions could be observed in the first 5 minute spectrum.

\( R_{av} \) was found by looking at the spectra and calculating the maximum number of counts, above the background, at energies where the \(^{186}\text{W} \) \( \gamma \) rays would be expected. The decay rate, \( \lambda = \ln(2)/t_{1/2} \) was calculated to reproduce the maximum average count rate. This implies an upper limit for the half-life of \( t_{1/2} \leq 100 \mu s \).

From the above analysis the half-life range for the 3544 keV isomeric state in \(^{186}\text{W} \) is \( 3 \text{ ms} \leq t_{1/2} \leq 100 \text{ s} \).
Figure 6.7: A graph of peak area ratios against energy for data on the 1.65 ms range. The peak areas were calculated from spectra produced from gates on the first 1/3 (beginning 120 $\mu$s after the end of the beam pulse) and last 1/3 of the total time projection. The long-dashed line is the average ratio plus one standard deviation for peaks lying between the two isomers. This gives a lower limit for the half-life. The short dashed line shows the estimated ratio due to the random stopping effect. (Larger ratios correspond to shorter halflives.)

Nanosecond analysis

The fragmentary structure of the decay scheme (Figure 6.1) above the $K^\pi = 7^-$ isomer suggests that many of these states could be intrinsic. This would lead to a stepwise reduction in the $K$ quantum number by transitions from the higher isomer to successively lower-$K$ bandheads, with the possibility of additional isomers. Since however, no additional $^{186}$W isomers were identifiable with the beam-$\gamma$ TAC, data from the second TAC were analysed to see if there were any short lifetimes (~ few ns) associated with transitions lying between the two isomers. This TAC was started by a germanium signal and stopped by 2 or more, elec-
tronically delayed, BGO signals. The events were sorted into a matrix of energy against the relative time between γ rays, and energy gates were set for many of the transitions in $^{186}$W. Figure 6.8 shows the results. (The 306 and 390 keV transition are not shown in Figure 6.8 because of contamination with γ rays below the $K^\pi = 7^-$ isomer.) The time-walk effect is demonstrated here with low energy events being delayed with respect to the higher energy lines: low energy γ rays interact closer to the surface of the germanium crystals and these signals take longer to be collected. There are statistical fluctuations in the data but there is some evidence to suggest that the 168, 221 and 380 keV transitions are delayed. These transitions are placed directly above the $K^\pi = 7^-$ state in the level scheme, based on this possibility of a short half-life ($\approx 1$ ns).

![Figure 6.8](image-url)

Figure 6.8: A graph of time centroid against energy for the Ge-BGO TAC spectra produced from background-subtracted energy gates in $^{186}$W. The 122 keV transition is labelled with # and has been corrected for the 1.0 ns half-life [91] of the $2^+$ level from which it decays. Data labelled with * correspond to transitions lying between the 2 long-lived isomers. The dashed line is a quadratic fit to the data up to $\approx 600$ keV to show the general trend.
6.1.4 Spin/parity assignments above the $K^\pi = 7^-$ isomer

The limited additional spectroscopic information leads to ambiguities in the spin and parity assignments (and even in the level ordering) for states between the two long-lived isomers in $^{186}$W. Nevertheless, the coincidence relationships and $\gamma$-ray intensities permit tentative assignments to be made, as follows.

The simplest interpretation for transitions with accompanying higher energy crossovers, is as $\Delta I = 1$ transitions since these are the fastest. Therefore, the 221 and 380 keV decays are likely to be $\Delta I = 1$ transitions. An $E2$ assignment for the corresponding 601 crossover transition is suggested. The energy of the 168 keV transition is low enough to extract the multipolarity from balancing the intensity in to and out of the 2286 keV level. The implied electron conversion coefficient is $\alpha_T(\text{exp}) = 1.6 \pm 0.24$ ($\alpha_T(M1) = 1.12$, $\alpha_T(E1) = 0.10$, $\alpha_T(E2) = 0.57$, $\alpha_T(M2) = 6.63$, $\alpha_T(E3) = 5.57$). Although outside the uncertainties, which could mean that the errors in the intensities are underestimated, this strongly suggests that the 168 keV line is an $M1$ multipole. The spin and parity of the 2286 keV level is therefore assigned ($10^-$). Feeding this state is an intense 552 keV transition (assumed to be a stretched $E2$) crossing two weaker decay paths. The intensities of the 315 and 237 keV transitions to and from the intermediate state at 2524 keV respectively, are not sensitive enough to the conversion coefficients to distinguish between $E1$ and $M1$ multipolarities. This is also the case for the 165 and 387 keV decays via an intermediate (2673 keV) level, the ordering of which could be reversed. (The intensities of the 387 keV and 165 keV transitions are closely matched if they both have $M1$ assignments, but due to their low intensities it is not possible to rule out $E1$ decays.)

Intensity balancing for the 306 keV transition leads to an electron conversion coefficient of $\alpha_T(\text{exp}) = 0.03 \pm 0.07$ ($\alpha_T(E1) = 0.02$, $\alpha_T(M1) = 0.21$, $\alpha_T(E2) = 0.08$). This was found assuming that both the 165 and 315 keV transitions are $M1$, so as to yield the maximum conversion coefficient for the 306 keV decay. Although
this is consistent with either electric dipole or electric quadrupole assignments, the simplest interpretation is with an E1 multipolarity. The presence of the 621 keV crossover transition to the 2524 keV state also favours this (13+) scenario for the 3145 keV level. It is worth noting here that the absence of a measurable half-life for the 3145 keV state does not rule out an M2 assignment for the 621 keV decay. Assuming a conservative limit of $t_{1/2} \leq 5$ ns for the 3145 keV level, gives a hindrance factor relative to the Weisskopf single-particle estimate of $F_W \leq 3$. Such a low value for an M2 transition (assuming no additional K hindrance) is not without precedent. In Ref.[72] the corresponding $F_W$ factors range from $\sim 1$ up to $\sim 10^5$.

With a definite range for the half-life of the higher isomer the Weisskopf single-particle estimates were calculated to find the transition multipolarity that is consistent with a several millisecond (or longer) half-life, shown in Table 6.2. The lowest multipolarity that can account for the observed half-life is $\lambda = 3$.

A $K$-allowed $\lambda = 3$ (0.4 MeV) transition can account for the half-life of the isomeric state, leading to a spin of (16) $h$ for the 3544 keV level. Intensity balancing supports an M3 assignment for the 399 keV transition, with $\alpha_T(\text{exp})=1.2\pm0.3$ ($\alpha_T(\text{M3})=1.05$, $\alpha_T(\text{E3})=0.13$, $\alpha_T(\text{E4})=0.53$). Although assumptions concerning the multipolarity of the 220 and 390 keV transitions were made to calculate this coefficient (both were assumed to be M1), choosing between M1 or E1 has relatively little effect giving $\alpha_T(\text{exp})=1.4\pm0.3$ for the 399 keV transition when E1 multipoles are assumed for the 220 and 390 keV transitions. On the basis of this analysis the preferred spin and parity assignment for the isomeric state is (16$^+$).

From the decay scheme in Figure 6.1 it can be seen that there is 9 keV difference between the two highest levels, and this is the primary reason for the placement of the 399 and 390 keV transitions at the top of the level scheme. Both show the same half-life, implying an unobserved 9 keV transition from the isomeric state at 3544 keV to the 3535 keV level. A transition rate for the 'missing' 9 keV supposed E2 decay can be calculated to see whether this is consistent with
Table 6.2: Weisskopf single-particle transition rate estimates [92] for $E_\gamma = 0.40$ MeV and $A = 186$. Multipolarities up to $\lambda = 4$ were considered. $T_{sp}$ is the single-particle transition rate for $E_\gamma$ in MeV and $t_{1/2}^\gamma$ is the partial $\gamma$-ray half-life.

<table>
<thead>
<tr>
<th>Multipolarity</th>
<th>Formula for $T_{sp}$</th>
<th>$T_{sp}$ (s$^{-1}$)</th>
<th>$t_{1/2}^\gamma = \frac{\ln 2}{T_{sp}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>E1</td>
<td>$1.0 \times 10^{14} A^{2/3} E_\gamma^3$</td>
<td>$2.09 \times 10^{14}$</td>
<td>3.32 fs</td>
</tr>
<tr>
<td>M1</td>
<td>$3.1 \times 10^{13} E_\gamma^3$</td>
<td>$1.98 \times 10^{12}$</td>
<td>0.35 ps</td>
</tr>
<tr>
<td>E2</td>
<td>$7.4 \times 10^7 A^{4/3} E_\gamma^5$</td>
<td>$8.05 \times 10^6$</td>
<td>0.86 ns</td>
</tr>
<tr>
<td>M2</td>
<td>$2.2 \times 10^7 A^{2/3} E_\gamma^5$</td>
<td>$7.34 \times 10^6$</td>
<td>94.4 ns</td>
</tr>
<tr>
<td>E3</td>
<td>$3.5 \times 10^3 A^2 E_\gamma^7$</td>
<td>1980</td>
<td>0.35 ms</td>
</tr>
<tr>
<td>M3</td>
<td>$1.1 \times 10^2 A^{4/3} E_\gamma^7$</td>
<td>19.14</td>
<td>36.2 ms</td>
</tr>
<tr>
<td>E4</td>
<td>$1.1 \times 10^{-5} A^{6/3} E_\gamma^9$</td>
<td>$3.25 \times 10^{-3}$</td>
<td>3.6 m</td>
</tr>
<tr>
<td>M4</td>
<td>$3.3 \times 10^{-6} A^2 E_\gamma^9$</td>
<td>$2.99 \times 10^{-5}$</td>
<td>6.4 h</td>
</tr>
</tbody>
</table>

a millisecond (or longer) half-life. The electron conversion coefficient ($\alpha_T(E2) = 60700$ [53]), which is dominant at this energy, must be included, resulting in a partial transition half-life of $t_{1/2} \approx 1.5$ $\mu$s. (This is calculated using the Weisskopf single particle estimate and the electron conversion coefficient.) Although this is several orders of magnitude shorter than the measured lower limit for the half-life, allowed E2 decays can be hindered by a factor of $\approx 1000$ [72], e.g. the 86 keV $K$-allowed E2 transition from the 192 $\mu$s isomer in $^{177}\text{Ta}$ [76] has a hindrance of $\approx 900$. With this factor the 9 keV E2 transition has an expected partial lifetime of order milliseconds in agreement with the implied experimental value. A similar hindrance factor is required for the once $K$-forbidden 180 keV (E2) transition from the $\geq 3$ ms isomer in $^{186}\text{W}$. (The order of the 220 and 180 keV transitions could be reversed, but due to their close energies this has very little effect on the calculation.) The Weisskopf single-particle estimate for this transition gives $t_{1/2}^\gamma = 47$ ns, but with a hindrance of 1000 and an additional factor of 100
increase from the one degree of $K$-forbiddenness, this partial transition lifetime is also expected to be of the order of milliseconds, consistent with the measured value.

### 6.1.5 Transfer products

As with the experiment using the $^{181}$Ta target (Chapter 5) a large number of transfer products in the target region are observed in the $^{186}$W target data. Only delayed (off-beam) events were collected which means that only nuclei containing isomers, with half-lives greater than about a hundred nanoseconds, were observed. Table 6.3 lists all the observed isomers populated using the $^{186}$W target and, where possible, gives relative $\gamma$-ray intensities. (The isomers populated in the fission of $^{238}$U are the same as for the $^{181}$Ta target, listed in Table 5.3.) The last entry in Table 6.3 is the decay of a new isomer, now tentatively placed in $^{185}$Ta, see below. Gamma-ray coincidences were used to order the transitions into a decay sequence, shown in the inset of Figure 6.9.

The strong M1/E2 and accompanying weaker stretched E2 transitions in the band, can be seen in the $\gamma$-ray spectrum of Figure 6.9. No accompanying crossover transition has been observed for the 175 keV $\gamma$ ray. Intensity balancing for this transition gives an implied electron conversion coefficient of $\alpha_T(\text{exp})=0.12 \pm 0.18$ compared to theoretical values [53] of $\alpha_T(E1)=0.08$, $\alpha_T(E2)=0.47$ and $\alpha_T(M1)=0.91$. This suggests that the 175 keV decay is an E1 transition, although an E2 assignment is not ruled out. The placement of this transition will be discussed in Section 6.2.1. The limit on the half-life for the new isomer is $t_{1/2} \geq 1$ ms. This was found by examining the counts in several different time gated spectra on the 1.65 ms TAC range as performed for the $K^{\pi}=(16^+)$ $^{186}$W isomer in Section 6.1.3. (Random stopping effects and low statistics due to a limited beam time with the 1.65 ms range prevent a higher limit being set.)

In an attempt to assign this band to a particular nucleus, the X-ray spectrum region was examined in detail. Although the use of absorbers blocked a high
Table 6.3: Transitions from isomers observed in the reaction \(^{238}\text{U} + ^{186}_{74}\text{W}\) at 1600 MeV (1.65 \(\mu\)s range). Isomers populated by the fission of \(^{238}\text{U}\) are shown in Table 5.3.

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>Particles transferred</th>
<th>(E_\gamma) (keV)</th>
<th>(L_\gamma)</th>
<th>Half-life</th>
<th>Assignment (I^\pi) Initial</th>
<th>Final</th>
</tr>
</thead>
<tbody>
<tr>
<td>(^{186}\text{W}) (target)</td>
<td></td>
<td>119</td>
<td>14.3(8)</td>
<td>18 (\mu)s</td>
<td>7(^-) 6(^+)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>195</td>
<td>56(2)</td>
<td></td>
<td>7(^-) 5(^-)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>708</td>
<td>8.9(6)</td>
<td></td>
<td>7(^-) 6(^+)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>180</td>
<td>1.3(1)</td>
<td>&gt;3 ms</td>
<td>16(^+) 14(^+)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>390*</td>
<td>5.1(3)</td>
<td></td>
<td>14(^+) 13(^+)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>399</td>
<td>3.2(2)</td>
<td></td>
<td>16(^+) 13(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{184}\text{W}) (—2n)</td>
<td></td>
<td>537</td>
<td>1.0(2)</td>
<td>8.3 (\mu)s</td>
<td>5(^-) 6(^+)</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>921</td>
<td>3.3(5)</td>
<td></td>
<td>5(^-) 4(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{182}\text{W}) (—4n)</td>
<td></td>
<td>519</td>
<td>0.3(1)</td>
<td>2 (\mu)s</td>
<td>10(^+) 10(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{180}\text{W}) (—6n)</td>
<td></td>
<td>391</td>
<td>0.4(1)</td>
<td>5.5 ms</td>
<td>8(^-) 8(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{190}\text{Os}) (+2p+2n)</td>
<td></td>
<td>503*</td>
<td>1.2(3)</td>
<td>9.9 min</td>
<td>4(^+) 2(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{183}\text{Re}) (+1p—4n)</td>
<td></td>
<td>304*</td>
<td>weak</td>
<td>1 ms</td>
<td>21/2(^+) 19/2(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{182}\text{Ta}) (—1p—3n)</td>
<td></td>
<td>172*</td>
<td>1.0(3)</td>
<td>15.8 min</td>
<td>7(^+) 6(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{181}\text{Ta}) (—1p—4n)</td>
<td></td>
<td>177</td>
<td>weak</td>
<td>25 (\mu)s</td>
<td>21/2(^+) 21/2(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{180}\text{Ta}) (—1p—5n)</td>
<td></td>
<td>456</td>
<td>0.2(1)</td>
<td></td>
<td>21/2(^-) 19/2(^-)</td>
<td></td>
</tr>
<tr>
<td>(^{184}\text{Hf}) (—2p)</td>
<td></td>
<td>368*</td>
<td>weak</td>
<td>48 s</td>
<td>6(^+) 4(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{182}\text{Hf}) (—2p—2n)</td>
<td></td>
<td>344*</td>
<td>0.4(2)</td>
<td>61.5 min</td>
<td>6(^+) 4(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{180}\text{Hf}) (—2p—4n)</td>
<td></td>
<td>332*</td>
<td>0.8(2)</td>
<td>5.5 h</td>
<td>6(^+) 4(^+)</td>
<td></td>
</tr>
<tr>
<td>(^{178}\text{Hf}) (—2p—6n)</td>
<td></td>
<td>244*</td>
<td>weak</td>
<td>10 (\mu)s</td>
<td>9(^-) 8(^-)</td>
<td></td>
</tr>
<tr>
<td>(^{185}\text{Ta})</td>
<td></td>
<td>325*</td>
<td>1.6(3)</td>
<td>4.0 s</td>
<td>6(^+) 4(^+)</td>
<td></td>
</tr>
</tbody>
</table>

\(^a\) From the recent work of D'Alarcao \textit{et al.} [74].

\(*\) The transition is not one which directly de-populates the isomer due to very weak intensities or low transition energies for the directly de-populating transitions.
CHAPTER 6. $^{186}W$ TARGET ISOMERS

Figure 6.9: The $\gamma$-ray spectrum gated by the 162, 175, 191, 246, and 266 keV transitions in the new band (see inset panel). Although the ordering of the states is not firm, a reasonable level scheme has been proposed, with a well formed rotational structure being fed by an isomeric state. The low energy isomeric transition has not been observed leading to an offset of $\Delta$ for the isomeric level. See text for details.

A proportion of the low energy events, there are significant counts in this domain in the gated spectra. Figure 6.10 shows the X-ray region for the total projections using the $^{186}W$ and $^{181}Ta$ targets, together with the spectrum gated by transitions in the new band. The X-rays from the lead backing can be seen clearly in all three spectra and though there are few counts in the lower spectrum (for the new band), small X-ray peaks do lie at the same energy as tantalum X-rays. Therefore, this new band is considered to belong to a tantalum isotope.

The relatively strong population of this isomer with respect to other transfer products suggests that a 1-particle transfer is responsible, due to the higher cross-section for such processes over multi-nucleon transfer. This, together with the X-ray information, points towards the isotope $^{185}Ta$ as being the most likely candidate, produced by 1-proton transfer from the $^{186}W$ target. The possible quasiparticle configurations of these states are discussed in the next section.
Figure 6.10: X-rays from 3 different $\gamma$-$\gamma$ spectra. The top panel shows the total out-of-beam spectrum for the $^{186}\text{W}$ target data. The middle panel is the corresponding plot for the $^{181}\text{Ta}$ target data. The lower panel shows the X-ray region of a spectrum gated by transitions in the new band. The vertical lines (left to right) are at the $K_\alpha$ and $K_\beta$ X-ray energies of tantalum and lead respectively. The X-rays in the lower spectrum match with the energies of the tantalum lines, but are offset compared with the tungsten X-rays.

6.2 Discussion

In Section 6.2.1 the possible assignments for the new band are discussed following an analysis of the band alignment and the in-band $\gamma$-ray branching ratios. In addition to this, blocked BCS calculations have been performed for $^{186}\text{W}$. The results and implications for the level scheme are discussed in Section 6.2.3. The systematics of the tungsten 4-quasiparticle isomers are examined in Section 6.2.4.
followed by a discussion of the new results in terms of the limit of access with
fusion-evaporation reactions in Section 6.2.5.

6.2.1 Properties of the new band

The observation of tantalum X-rays in coincidence with transitions in the new
band, together with the population intensity, favours an assignment to $^{185}$Ta.
This possibility is now investigated in detail.

In $^{183}$Ta (and also in $^{181}$Ta, see Chapter 5), the ground-state band is built on
a $^{7/2}^+$ [404] 1-quasiproton configuration [93]. The $^{183}$Ta ground-state is fed by an
isomeric yrast $^{9/2}^-[514]$ bandhead with $t_{1/2} = 0.11 \mu s$ [93] at 73 keV. In $^{185}$Ta (the
favoured isotope for the new band) the ground-state band has been tentatively
assigned the same $K^* = 7/2^+$ proton orbital [94]. As neutrons are added, the
trend of the $K^* = 9/2^-$ bandheads in the odd-$A$ tantalum isotopes is for in­
creasing excitation energy and shorter half-lives. The 175 keV transition is of the
right multipolarity and approximately the right energy to continue the system­
atics. The half-life limit for this transition with respect to the band members is
$t_{1/2} \leq 5$ ns, based on a time difference analysis (the Ge-BGO TAC (Section 6.1.3)
was started by the Ge signal and stopped by the electronically delayed BGO sig­
nal). This fits well with the B(E1) rates observed in the lighter odd-$A$ isotopes.
A state at approximately 163 keV has been observed with a tentative spin of
9/2 [94] from proton pick-up reactions. The uncertainty in this energy is likely
to be large due to the low energy resolution and relatively poor statistics, and
the assignment of a 175 keV transition here is not inconsistent with the previous
study.

To gauge whether the new band could be built on either a $7/2^+$ or $9/2^-$ 1-
quasiproton configuration the in-band branching ratios, $T_2/T_1$ are calculated. The
rotational model Equations (2.21) and (2.22) can be used to derive an experimen­tal
intrinsic g-factor to compare with the calculated values for each configuration.
The parameters $Q_0 = 6.5$ e-b and $g_R = 0.3$ are chosen as these reproduce the
intrinsic g-factors for the $K^\pi = 7/2^+$ and $9/2^-$ bands in $^{181}$Ta [9].

Table 6.4 shows the results based on two assumptions for the $K$ values at the bandhead, namely $K^\pi = 7/2^+$ and $9/2^-$. The g-factors fluctuate very little over the range of spins. However, the calculated quantities using the asymptotic Nilsson quantum numbers are

$$\frac{(g_K - g_R)}{Q_0} = +0.06 \text{ for } K^\pi = \frac{7}{2}^+$$
$$= +0.16 \text{ for } K^\pi = \frac{9}{2}^-$$

For $K = 9/2$ the experimentally determined g-factors agree, within the uncertainties, with the calculated value based on the $\pi_{5/2}^-$ [514] Nilsson orbital. The $|\frac{(g_K - g_R)}{Q_0}|$ values assuming $K = 7/2$ do not agree with the calculated quantity. This provides evidence for a $9/2^- [514]$ assignment, and supports the placement of this band and its associated isomer in $^{185}$Ta.

Table 6.4: Gamma-ray branching ratios and g-factors for the new band.

| $I$ (h) | $E_2$  | $E_1$  | $T_2/T_1$ | $\frac{|(g_K - g_R)|}{Q_0}$ (e-b)$^{-1}$ |
|--------|-------|-------|-----------|---------------------------------|
| $7/2^a$ | $9/2^b$ | (MeV) |           | $7/2^a$ | $9/2^b$ |
| 11/2   | 13/2  | 0.354 | 0.191     | 0.11±0.10 | 0.21±0.50 | 0.15±0.36 |
| 13/2   | 15/2  | 0.409 | 0.218     | 0.33±0.17 | 0.18±0.08 | 0.13±0.06 |
| 15/2   | 17/2  | 0.464 | 0.246     | 0.72±0.23 | 0.15±0.04 | 0.11±0.03 |
| 17/2   | 19/2  | 0.512 | 0.266     | 0.60±0.17 | 0.20±0.04 | 0.15±0.03 |

$^a$ Assuming $K = 7/2$.
$^b$ Assuming $K = 9/2$.

Adding one unit of spin for each level above an $I^\pi = 9/2^-$ bandhead would mean that the transition directly de-populating the isomer feeds the $19/2^-$ member of the band. This transition has not been observed but can be given energy limits of <100 keV (M1) and <80 keV (E1) on the basis of detection-efficiency and conversion-coefficient considerations. Comparison with blocked BCS calculations (as described in Section 3.1) favours a $K^\pi = 21/2^- \{\frac{5}{2}^+^{[402]}, \frac{7}{2}^+^{[404]}, \frac{9}{2}^- [514]\}$ 3-quasiproton configuration for the isomer, consistent with an isomeric
The aligned angular momentum, $I_a$, for the new band is plotted in Figure 6.11 along with that for the 1-quasiproton $\frac{9}{2}^-[514]$ band in $^{181}$Ta. The formulae of Purry et al. [28], Equations (2.40) and (2.41) have been used. The two sets of values lie very close together, suggesting that similar particles may be involved in both bands. The slopes of the lines are also very alike which leads, for example, to similar dynamic moments-of-inertia, $J^{(2)}$, (Section 2.8) which again might be expected if these structures are built on the same configuration. It should be noted here that when the alignment for the new band is plotted assuming $K = 7/2$ along with the $K^\pi = 7/2^+$ band in $^{181}$Ta, the relationship between the two is as for the $K = 9/2$ scenario shown in Figure 6.11. These data points have been omitted from Figure 6.11 for clarity.

Although the assignment of this band in $^{185}$Ta is not confirmed, the X-ray evidence and the relative population intensity point strongly towards this placement and this is supported by the analysis of the in-band $\gamma$-ray branching ratios and comparisons with Nilsson model calculations. Further data are necessary before a configuration for this isomer and the band that it populates can be unambiguously assigned. This new band is reported in Ref. [95].

### 6.2.2 Decays from intrinsic states

The strongest decay branch from the $K^\pi = 7^-$ isomer in $^{188}$W$_{112}$ (Figure 6.1) is the 195 keV stretched E2 transition to the $I^\pi = 5^-$ member of the octupole band. Weaker branches include decays to the $6^+$ members of the $\gamma$ vibrational and ground-state bands. The latter of these has the highest degree of $K$-forbiddenness.
Figure 6.11: A plot of aligned angular momentum, $L_z$, versus rotational frequency $\hbar \omega$ for the $\pi\{\frac{3}{2}^-[514]\}$ band in $^{181}$Ta (from Chapter 5) and the new band observed in the $^{186}$W target data. Both $\Delta I = 1$ and 2 transitions are plotted for each band.

$(\nu = 6$, Section 2.5) the hindrance of which is discussed in Section 7.3.4. $K^\pi = 7^-$ isomers are also known in the isotope $^{188}$Os$^{112}$ at 1771 keV with $t_{1/2} = 14$ ns [96], that decays by a $\Delta I = 2$ 102 keV transition to the octupole band, and in $^{186}$Os$^{110}$ [97, 98, 99], with $t_{1/2} = 8.5$ ns (Chapter 7). The latter has the same 2-quasineutron $\nu\{\frac{11}{2}^+[615], \frac{3}{2}^-[512]\}$ assignment proposed for the 18 $\mu$s isomer in $^{186}$W. In addition there is a 2.4 ns isomer in $^{194}$W [90] with the same configuration.

The $K^\pi = (16^+) t_{1/2} \geq 3$ ms isomer in $^{186}$W decays predominantly via a 399 keV M3 transition. The systematics of isomeric M3 decays are presented in Ref. [100] in terms of the Weisskopf hindrance factors, $F_W$, (Section 2.5) versus mass number. The 399 keV transition observed in $^{186}$W has $F_W \geq 0.2$ which lies on the lower limit with respect to the other examples in $A \approx 180$ nuclei. This supports the M3 assignment for the 399 keV transition. It is worth noting that all the other data points in Ref. [100] for known isomeric M3 transitions are in either odd or odd-odd nuclei. A measurement of the half-life of the $K^\pi = (16^+)$ state
in $^{186}$W would enable a quantitative comparison with isomeric M3 systematics to be made.

### 6.2.3 Multi-quasiparticle calculations

Blocked BCS calculations, as performed for $^{181}$Ta (Section 5.3.3), have been used to calculate the energies of multi-quasiparticle configurations in $^{186}$W, using deformation parameters $\epsilon_2 = 0.198$ and $\epsilon_4 = 0.061$ [81]. The Nilsson single-particle energies were adjusted to reproduce approximately the neutron and proton single-quasiparticle energies, where available [9, 82], in neighbouring odd-$A$ nuclei $^{185}$W and $^{187}$Re respectively. The neutron monopole pairing strength was optimised to give approximately the correct energy for the $7^-$ isomer in $^{186}$W resulting in $G_\nu = 22.5/A$ MeV. The proton pairing strength, $G_\pi = 23.5/A$ MeV, was chosen to be $1.0/A$ MeV higher, consistent with Ref.[31].

From Figure 6.12 it can be seen that there are two low lying $7^-$ states and a low $16^+$ state. Table 6.5 summarises the properties of the intrinsic states. The most yrast calculated 2-quasiparticle $7^-$ state is a 2-quasiproton configuration. This differs from the preferred 2-quasineutron coupling for the $18\mu$s isomer, assigned by comparison with $^{184}$W. However, the 2-quasineutron configuration is calculated to lie only $190$ keV higher in energy, and this state has favoured residual nucleon-nucleon interactions (Section 3.1.3), unlike the 2-quasiproton state. Residual interactions can alter the energies by $\sim100$ keV, therefore it is not unreasonable to assume that the 2-quasineutron $K^\sigma = 7^-$ level could lie at or below the energy of the corresponding 2-quasiproton state. A mixed configuration is also a possibility. The calculation of a low lying proton $7^-$ state could also point to an underestimated $G_\pi$ parameter. Further experimental information, such as from a rotational band based on the $7^-$ isomer, would be needed to remove these ambiguities.

The calculated $K^\pi = 16^+$ state at $3449$ keV is of the right spin and parity to account for the long-lived state observed at $3544$ keV. The residual interac-
Figure 6.12: A plot showing the calculated multi-quasiparticle states for $^{186}\text{W}$. Only configurations close to yrast with couplings to maximum $K$ are plotted.

... are favoured for this configuration which again suggests that the proton monopole pairing strength ($G_\pi$) is too low. Experimentally, this $K^\pi = (16^+)$ isomer lies below the $I^\pi = 14^+$ member of the ground-state band (3563 keV [86]), qualifying it as an “yrast trap” because it is forced to decay by transitions with multipolarities $\lambda\geq1$.

There are several configurations predicted to lie in the energy range between the two isomers which could be the intrinsic states populated in the decay of the $(16^+)$ level. When a preliminary consideration of the residual interactions is made, the lowest of these levels would be $9^-$, $10^-$, $11^-$, $12^-$, $13^+$ and $14^+$ (see Table 6.5). Figure 6.13 shows a comparison of the experimental and calculated...
Table 6.5: The energies for calculated and observed levels are compared. Only the lowest calculated energy for each spin is shown, except for the $7^-$ states for which there are two candidates. In the first column (f) indicates the most favoured by residual interactions for the given orbitals, though the energies (“Calculation”) are quoted without residual interactions. Brackets around the experimental energy imply that the spin and parity assignments for the levels are tentative.

<table>
<thead>
<tr>
<th>Assignment $K^\pi$</th>
<th>Multi-quasiparticle Energy (keV)</th>
<th>Configuration</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Experimental</td>
<td>Calculation</td>
</tr>
<tr>
<td>$7^-$</td>
<td>1517</td>
<td>1539</td>
</tr>
<tr>
<td>$7^-$ (f)</td>
<td>1517</td>
<td>1729</td>
</tr>
<tr>
<td>9$^-$</td>
<td>(2118)</td>
<td>1910</td>
</tr>
<tr>
<td>10$^-$ (f)</td>
<td>(2286)</td>
<td>2549</td>
</tr>
<tr>
<td>11$^-$</td>
<td>(2673)</td>
<td>2682</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>12$^-$</td>
<td>(2838)</td>
<td>3271</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>13$^+$ (f)</td>
<td>(3145)</td>
<td>3389</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>14$^+$</td>
<td>(3535)</td>
<td>3267</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>16$^+$ (f)</td>
<td>(3544)</td>
<td>3449</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

intrinsic states. Generally, there is good agreement between the two, with most of the states favoured by residual interactions calculated higher in energy than the corresponding observed level. The converse is also true. This demonstrates that the structure above the $K^\pi = 7^-$ isomer can be reproduced with reasonable agreement to the data, but additional information would be needed for firm assignments.

The shapes of high-$K$ intrinsic states are investigated in detail in Chapter 7, but configuration-constrained potential-energy-surface calculations (Section 3.2 and Ref.[37]) have been performed for the ground state, $7^-$ states (both proton and neutron) and $16^+$ state. These show that all of these isomers are axially symmetric with a triaxiality parameter (Section 2.2) of $\gamma<2^\circ$ in each case. The 2-quasineutron $K^\pi = 7^-$ state is predicted to lie below the corresponding 2-
Figure 6.13: Experimental and calculated intrinsic states in $^{186}$W. The favoured labels for the calculated states indicate the most favoured by residual nucleon-nucleon interactions for the given orbitals. The calculated 12$^-$ state at 3271 keV has been omitted for clarity although it should be noted that this level is calculated higher in energy than the observed (12$^-$) state but does not have the most favoured residual interactions.

quasiproton state by these calculations and unlike the blocked BCS calculations there is no adjustment of the single-particle energies or the monopole pairing strengths.

6.2.4 Tungsten 4-quasiparticle systematics

Four-quasiparticle isomers have been previously observed in all the even-even tungsten isotopes from $^{176}$W to $^{182}$W inclusive [28, 101, 102, 103]. The new $K^* = (16^+)$ isomer in $^{186}$W discovered in the present work extends the systematics to N=112. However, there is a notable gap at N=110 corresponding to $^{184}$W, where no structures with seniority greater than two have yet been observed. In
order to try and predict where such a state might lie, blocked BCS calculations have been performed for the even-even tungsten isotopes $^{176}$W to $^{186}$W, using deformation parameters from Ref.[81]. There was no adjustment of the single-particle energies and the monopole pairing strengths were kept as $G_\nu = 22.5/A$ MeV and $G_\pi = 23.5/A$ MeV, as these gave reasonable results for $^{186}$W. In each case the most yrast calculated 4-quasiparticle state is compared in Figure 6.14 to the most yrast seniority-4 isomer that has been observed. The "yrastness" of the levels was taken with respect to the corresponding ground-state band member. Where this band was not known to sufficiently high angular momentum ($^{184}$W and $^{186}$W), an extrapolation was performed from the known rotational states.

This approach makes as few assumptions as possible and treats all the isotopes consistently.

Figure 6.14 shows the excitation energies of both the calculated and experimental levels with the ground-state band reference energy subtracted for each isotope. The experimental states show an almost linear trend, becoming more yrast with increasing neutron number. The calculated energies are in good agreement with the experimental levels for $^{180}$W, $^{182}$W and $^{186}$W. (The agreement for $^{186}$W is expected based on the calculations performed above in Section 6.2.3.) The deviations for the lighter isotopes are likely to be due to an over estimation of the monopole pairing strengths since there is no a priori reason for these to remain constant. Even so, the general behaviour of decreasing relative excitation energies as neutrons are added, is reproduced. These systematics suggest that $^{184}$W may yield a long-lived isomeric state. The calculations predict a favoured $K^\pi = 17^-$ configuration at 4113 keV. Experimentally, this could be populated with the same experimental setup used here, but bombarding a target isotopically enriched in $^{184}$W. The half-life of the supposed isomer would need to be greater than about 100 ns, and not longer than many hours, for the decays to be recorded in the off-beam counting time with good efficiency.
CHAPTER 6. $^{186}$W TARGET ISOMERS

Figure 6.14: A graph showing the lowest experimental (filled circles) and calculated (solid line) 4-quasiparticle states as a function of neutron number, in the even-even tungsten isotopes from $^{176}$W to $^{186}$W. The energy of the corresponding ground-state band rotational member has been subtracted from the excitation energies for each case, in order to illustrate the degree to which the isomers become yrast at high neutron numbers.

### 6.2.5 Beyond the fusion limit

The new $K^\pi = (16^+)$ isomer observed in $^{186}$W extends the 4-quasiparticle $K$-isomer systematics beyond the limit of access using fusion-evaporation reactions, for the first time. This limit can be defined by the ($^4\text{He},2n$) and ($^3\text{Li},3n$) reactions, which bring in sufficient angular momentum to populate seniority-4 states. Figure 6.15 shows the new result together with the other known high-$K$ isomers.

Two other data points lying to the neutron-rich side of the fusion limit are for $^{180}\text{Hf}$ and $^{182}\text{Hf}$, discovered in a follow-up experiment [74] using a $^{180}\text{Hf}$ target with the setup as described in Section 5.1. Note that $^{184}_{74}\text{W}_{110}$, where no 4-quasiparticle isomer is known, is also out of reach of fusion-evaporation reactions.
It is notable that the new results establish long-lived isomers in a previously inaccessible region of the nuclear chart. Many long-lived isomers should now be identifiable with the new generation of gamma-ray detector arrays.

![Nuclear chart](image)

Figure 6.15: The mass 180, high-$K$-isomer region of the nuclear chart. Circles represent observed isomers with seniority $\geq 4$ and half-lives more than a few nanoseconds (large circles for $t_{1/2} > 1$ ms). The bold line dividing the region gives the limit of access with fusion-evaporation reactions (see text for details). The new result for $^{125}_{74}$W$_{112}$ is well to the neutron-rich side of this limit.

### 6.3 Summary

New isomeric states have been populated in $^{186}$W by inelastic reactions using $^{238}$U beams. A 2-quasiparticle $K^* = 7^-$ isomer with $t_{1/2} = 18$ $\mu$s has been assigned to $^{186}$W, decaying to the ground-state band and the $\gamma$ and octupole vibrational bands. A new 4-quasiparticle state with a probable $16^+$ configuration has been placed above the 18 $\mu$s level. The half-life could not be measured in this work but lies in the range $3$ ms $\leq t_{1/2} \leq 100$ s. The non-rotational structure of the states
populated in the decay of the seniority-4 isomer suggests that many of the intermediate levels are intrinsic. Although the level ordering remains ambiguous, lifetime information and intensity balancing have been used to construct a reasonable level scheme. The excitation energies of the proposed intrinsic states are in good agreement with predictions by blocked BCS calculations. The new results extend the known 4-quasiparticle isomers in to neutron-rich nuclei beyond the limit of access with fusion-evaporation reactions. The systematics of the known seniority-4 states in the even-even tungsten isotopes are well reproduced by blocked BCS calculations and a corresponding yrast intrinsic state in $^{184}$W is predicted.

In addition, a new $t_{1/2}\geq1$ ms isomer, populated by nucleon transfer, has been tentatively assigned to $^{185}$Ta. A $K^\pi = 21/2^-$ 3-quasiproton configuration is proposed. This extends the systematics of the 3-quasiparticle tantalum isomers to the neutron-rich side of $\beta$ stability for the first time.

The observation of long-lived states, the excitation energies of which are in good agreement with Nilsson model predictions, suggests that these calculations are reliable at high neutron numbers. The advent of suitable radioactive beams may soon allow predictions [3] of very high-spin very long-lived states in neutron-rich nuclei to be tested.
Chapter 7

\(^{186}\text{Os}\)

As protons and neutrons are added, beyond \(Z=72\) and \(N=106\) the nuclear shape becomes more susceptible to non-axial fluctuations, induced by multi-quasiparticle excitations. The observation of the anomalously fast decay of the \(K^\pi = 25^+\), \(t_{1/2} = 130\) ns isomer directly to the ground-state band in \(^{182}\text{Os}\) [80], is a possible indicator for the loss of axial symmetry. Recent potential-energy-surface calculations by Xu et al. [37] yield a modest value for the triaxiality parameter, \(\gamma = 10^\circ\) for this isomeric state. However, the non-yrast status of this isomer means that statistical \(K\)-mixing due to the high density of states, could be playing a role in the breakdown of the \(K\)-selection rules [4]. At higher neutron numbers Nilsson model calculations of the type described by Jain et al. [31] suggest that such intrinsic states may become yrast, which would help in pinpointing the mechanism responsible for the collapse of \(K\)-conservation. However, the experimental limit of using stable beam and target combinations prohibits access to neutron-rich nuclei.

In this chapter, new data are presented on the stable isotope \(^{186}\text{Os}\), which has been populated to high-spin by using a neutron-rich radioactive \(^{14}\text{C}\) beam. A recent experiment to study this nucleus [104], with the \(^{186}\text{W}(^4\text{He},4n)^{186}\text{Os}\) reaction observed a breakdown in the collective structure at 18 \(\hbar\), but due to the light-ion beam, the behaviour above 22 \(\hbar\) remained out of reach. The new data have consid-
erably extended the level scheme, and a series of 2- and 4-quasiparticle bands have been identified. The structure at high-spin has been clarified and new states have been observed. The results are discussed in conjunction with potential-energy-surface calculations [37] (Section 3.2). In addition, the first backbend in the yrast states [97] has been elucidated and the nature of the crossing band established. This provides the clearest example of a high-K $\iota$-band crossing in an even-even nucleus. The results of this work are reported in Refs. [98, 99].

7.1 Experimental method

DC beams of radioactive $^{14}$C have been used to bombard a 5 mg cm$^{-2}$, self-supporting target of ytterbium enriched to 95.7 % in $^{176}$Yb. High-spin states in the residual nucleus $^{186}$Os were populated in the fusion-evaporation reaction $^{176}_{70}$Yb($^{14}$C,4n)$^{186}_{70}$Os using a 67 MeV beam provided by the NBI tandem accelerator. The reaction products were stopped in the target, at the centre of the Nordball array, comprising 18 coaxial and 2 planar Compton suppressed germanium detectors. The germanium detectors were arranged in rings at angles of ±37°, ±79°, ±101° and ±143° to the beam direction. A high efficiency, 50 element ball of BaF$_2$ detectors allowed $\gamma$-ray multiplicity to be measured, and a 30 element silicon-detector inner ball was included to enable the charged-particle evaporation channels to be studied [105]. A master trigger condition required at least 2 germanium detectors and either 3 BaF$_2$ or 1 silicon element(s) to fire in coincidence. There was a time window of 550 ns for coincidences, started by the BaF$_2$ or silicon ball signal. Efficiency and energy calibrations were obtained using $^{152}$Eu and $^{133}$Ba sources.

For the above reaction the expected angular momentum input in to the fusion products is 29 $\hbar$, which is in good agreement with the highest observed state in $^{186}$Os at (30) $\hbar$. 
7.1.1 γ—γ analysis

The level scheme was constructed by analysing γ-ray coincidences with previously known, low-spin states [97]. The events were sorted into $4096 \times 4096$ matrices which were sliced to obtain background-subtracted projected spectra. The first matrix was built with all γ-ray events in which two or more coaxial germanium detectors fired. A second matrix was generated with the additional constraint that the coaxial detector signals were delayed (by between 120 and 550 ns), to search for decays from long-lived states. The timing was taken with respect to the signal from either the BaF$_2$ or silicon ball. The analysis program XMESC [60] was used to slice these symmetrised matrices. To examine coincidences between low-energy γ rays and X-rays, a third matrix (γ-X) was constructed. This contained events in which at least one planar and one coaxial detector fired.

7.1.2 DCO ratios

The method of Directional Correlations de-exciting Oriented states (DCO) [51] was used to distinguish between quadrupole and dipole radiation (see Section 4.2). A matrix was constructed by incrementing energy signals from coaxial detectors at angles ±79° and ±101° on one axis, and angles ±37° and ±143° on the other axis. This matrix was then fitted using matrix peak search and fitting programs. The ratios of γ-ray intensities projected onto the two different axes were then examined using Equation (4.3). Gating on a stretched quadrupole transition, for example, gave efficiency corrected ratios of ≈0.56 for pure $\Delta I = 1$ dipole γ-rays and ≈1.00 for stretched quadrupole transitions [52].

7.1.3 Half-lives

The half-lives measured in $^{186}$Os are all less than 10 ns. These were obtained using two methods. The first involved sorting events into a γ-γ-time cube, with γ-ray energy on two axes and the time difference between the two energy signals
on the third axis. The projected time spectra were fitted with a prompt Gaussian peak convoluted with an exponential decay. A fitting routine [106] of the type described in Ref.[107] was used. A second method, capable of greater statistical accuracy, employed a $\gamma$-time matrix. Plotting time centroids as a function of $\gamma$-ray energy enabled small shifts to be identified.

7.2 Results

Figure 7.1 shows the level scheme deduced in the current work. Spin and parity assignments have been made in several different ways. DCO ratios have been used to distinguish between dipole and stretched quadrupole radiation (Section 4.2). Quadrupole transitions are assumed to have electric character, except where half-lives have been measured. For low-energy transitions, a total electron conversion coefficient could be calculated, such that the intensity into and out of a given level is balanced (Section 4.3). These quantities can then be compared to theoretical values [53] to give a multipolarity. Thirdly, important consideration is given to the relative $\gamma$-ray intensities for states with several decay branches.

Quasiparticle configuration assignments to intrinsic states have been made, taking into account several factors. The first, is that for a rotational sequence, $K$ is equal to the spin of the bandhead. The validity of this assumption will be investigated in Section 7.3.6. The in-band E2/M1 branching ratios can be used in conjunction with the rotational model [13] to obtain $|(g_K - g_R)/Q_0|$ which can be compared with theoretical values to distinguish between different configurations (see Section 7.3.3). Plotting the rotation-aligned component of the angular momentum against rotational frequency for a given band can also provide evidence for configurations. This is discussed in Section 7.3.2. For instance, high-$j$ orbitals (e.g. $i_{13/2}$ neutrons) are generally more rotationally aligned than low-$j$ orbitals, due to larger Coriolis forces (Section 2.8). In Section 7.3.5, Nilsson model calculations [31] are discussed, and compared to the experimental data.
For a given spin and parity, the lowest calculated configuration is taken.

Table 7.1 lists the transitions, together with the relative intensities and spin and parity assignments. For most transitions an example DCO ratio is also listed, though typically there are many other values that corroborate each multipolarity assignment.

The following sections describe the main features of the $^{186}$Os level scheme and give details of how assignments were made.

![Diagram of the $^{186}$Os level scheme](image)

**Figure 7.1:** (a) Partial decay scheme for the $K^\pi = 0^+, 2^+, 4^+, 5^-, 7^-, 9^-$ and $10^+$ bands. The high-spin negative parity states are also shown; (b) Partial level scheme for the decay of the high-spin structures in $^{186}$Os. The most intense decay path through the low-$K$ bandheads is shown in addition to the $K^\pi = 10^+$ and ground-state bands. The width of the arrows is proportional to the $\gamma$-ray intensities (black) and the electron conversion (white). Vertical arrows indicate assigned in-band transitions. Energies are in keV. Half-lives (or half-life limits) are given for several states.
7.2.1 The ground-state band (band I)

No new levels have been firmly assigned to the ground-state band (g-band) beyond the previously known $14^+$ state [97, 108], although a new band (Ia) has been identified that feeds directly into this level. This is a candidate for a rotational aligned low-$K$ s-band (see Sections 2.3.1 and 7.3.1). From Section 7.3.2, the alignment for this band is much higher than that of the g-band. A new level (Ib) at 4352 keV has been observed but the spin and parity assignments remain tentative because a reliable DCO ratio cannot be obtained for the low intensity 793 keV transition. However, this state is a candidate for the $16^+$ member of the
Table 7.1: Energies, assignments, relative intensities, and DCO ratios for transitions observed in $^{186}$Os.

<table>
<thead>
<tr>
<th>$E_i$ (keV)</th>
<th>$E_i$</th>
<th>$E_f$</th>
<th>Assignment$^a$</th>
<th>$I_i$</th>
<th>$R_{DCO}$ [gate (keV)]</th>
</tr>
</thead>
<tbody>
<tr>
<td>118.9</td>
<td>3936</td>
<td>3817</td>
<td>$16^+\text{VIII} \rightarrow 15^+\text{VIII}$</td>
<td>4.56(16)</td>
<td>—</td>
</tr>
<tr>
<td>120.9</td>
<td>6154</td>
<td>6033</td>
<td>(24$^+$) + $\pi^c \rightarrow 22^+$ + $\pi$</td>
<td>9.8(3)</td>
<td>1.1(2) [530]</td>
</tr>
<tr>
<td>132.3</td>
<td>2564</td>
<td>2432</td>
<td>$10^+\text{VIII} \rightarrow 10^+\text{VII}$</td>
<td>3.1(2)</td>
<td>—</td>
</tr>
<tr>
<td>137.3</td>
<td>137</td>
<td>0</td>
<td>$2^+\text{I} \rightarrow 0^+\text{I}$</td>
<td>799(26)</td>
<td>1.00(1) [648]</td>
</tr>
<tr>
<td>140.0</td>
<td>911</td>
<td>768</td>
<td>$3^+\text{III} \rightarrow 2^+\text{III}$</td>
<td>4.27(20)</td>
<td>—</td>
</tr>
<tr>
<td>143.3</td>
<td>1772</td>
<td>1629</td>
<td>$6\rightarrow 5$ + $\pi$</td>
<td>45.1(14)</td>
<td>0.47(2) [585]</td>
</tr>
<tr>
<td>146.1</td>
<td>1775</td>
<td>1629</td>
<td>$7^+\text{VI} \rightarrow 5^+$</td>
<td>244(7)</td>
<td>0.96(2) [585]</td>
</tr>
<tr>
<td>(148.5)</td>
<td>3411</td>
<td>3294</td>
<td>$14^+\text{VIII} \rightarrow 13^+\text{VIII}$</td>
<td>3.0(1)</td>
<td>—</td>
</tr>
<tr>
<td>153.8</td>
<td>7145</td>
<td>6991</td>
<td>$28^+ + \pi^c \rightarrow 26^+ + \pi$</td>
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<td>1.00(6) [560]</td>
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<td>158.8</td>
<td>1071</td>
<td>911</td>
<td>$4^+\text{III} \rightarrow 3^+\text{III}$</td>
<td>4.70(23)</td>
<td>—</td>
</tr>
<tr>
<td>167.3</td>
<td>1940</td>
<td>1772</td>
<td>$7^+\rightarrow 6^-$</td>
<td>19.7(7)</td>
<td>0.53(7) [585]</td>
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<tr>
<td>193.8</td>
<td>1969</td>
<td>1775</td>
<td>$8^-\text{VI} \rightarrow 7^-\text{VI}$</td>
<td>264(8)</td>
<td>0.60(1) [146]</td>
</tr>
<tr>
<td>195.0</td>
<td>2134</td>
<td>1940</td>
<td>$8^-\rightarrow 7^-$</td>
<td>12.1(5)</td>
<td>—</td>
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<tr>
<td>197.3</td>
<td>2166</td>
<td>1969</td>
<td>$9^-\text{VII} \rightarrow 8^-\text{V}$</td>
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<td>1352</td>
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</tr>
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<td>216.5</td>
<td>2350</td>
<td>2134</td>
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<td>7.46(28)</td>
<td>0.58(5) [362]</td>
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<tr>
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<td>5246</td>
<td>5027</td>
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<td>13.9(5)</td>
<td>0.67(5) [560]</td>
</tr>
<tr>
<td>219.2</td>
<td>1492</td>
<td>1276</td>
<td>$6^+\text{III} \rightarrow 5^+\text{III}$</td>
<td>1.9(2)</td>
<td>—</td>
</tr>
<tr>
<td>219.7</td>
<td>2189</td>
<td>1969</td>
<td>$9^-\text{VI} \rightarrow 8^-\text{VI}$</td>
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<td>0.56(2) [146]</td>
</tr>
<tr>
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<td>2166</td>
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<td>$9^-\text{VII} \rightarrow 7^-\text{V}$</td>
<td>4.1(2)</td>
<td>—</td>
</tr>
<tr>
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<td>5562</td>
<td>5333</td>
<td>$(20^-) - \pi^d \rightarrow 19^- - \pi$</td>
<td>14.4(5)</td>
<td>0.85(7) [560]</td>
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<tr>
<td>232.9</td>
<td>3040</td>
<td>2807</td>
<td>$12^+\text{VIII} \rightarrow 11^+\text{VIII}$</td>
<td>29.6(9)</td>
<td>0.54(3) [523]</td>
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<tr>
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<td>$10^-\rightarrow 9^-\text{V}$</td>
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<td>0.65(4) [413]</td>
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<td>$13^+\text{IX} \rightarrow 12^+\text{IX}$</td>
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<td>$11^-\rightarrow 10^-\text{VII}$</td>
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<tr>
<td>$E_\gamma$ (keV)</td>
<td>$E_x$</td>
<td>$E_f$</td>
<td>Assignment a</td>
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<td>$R_DCO$ [gate (keV)]</td>
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<td>------</td>
<td>--------------</td>
<td>--------</td>
<td>------------------</td>
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<td>2166</td>
<td>10−VI→9−VII</td>
<td>46.9(14)</td>
<td>0.67(3) [146]</td>
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<td>332(10)</td>
<td>0.79(3) [146]</td>
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<td>4×IV→4×III</td>
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<td>0.92(6) [934]</td>
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<td>13−VII→12−VII</td>
<td>22.3(7)</td>
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<td>302.5</td>
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<td>768</td>
<td>4×III→2×III</td>
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<td>(22+) + πc→21(+) + π</td>
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Table 7.1 — continued

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<th>$E_f$</th>
<th>Assignment*</th>
<th>$I_\gamma$</th>
<th>$R_{DC0}$ [gate (keV)]</th>
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<td>1775</td>
<td>$9^+VI \rightarrow 7^+VI$</td>
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<td>420.8</td>
<td>1492</td>
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<td>$6^+III \rightarrow 4^+III$</td>
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Table 7.1 — continued

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<th>$I_\gamma$</th>
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<td>$4^+<em>{\text{II}} \rightarrow 2^+</em>{\text{I}}$</td>
<td>1.5(3)</td>
<td>—</td>
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</table>

$^a$ the spin and parity, $I^\pi$ are given along with the band label from the level scheme, for both the initial and final states.

$^b$ The gating transition is not E2.

$^c$ These transitions above the $K^\pi = 18^+$ intrinsic state are associated with positive parity levels (some tentative). See Figure 7.1b.

$^d$ These transitions above the $K^\pi = 18^+$ intrinsic state are associated with negative parity levels (some tentative). See Figure 7.1a.

7.2.2 The vibrational bands (bands II, III and IV)

The previously identified $K^\pi = 0^+ \beta$ vibrational band (band II) [97] has been observed in this work from the first excited state at 1207 keV, up to $I^\pi = (10^+)$, representing an additional two (tentatively assigned) states. The $0^+$ bandhead, lying at 1061 keV, has not been observed in this study. DCO information was not available for any of the new transitions due to the extremely weak population of this band. This reflects the high excitation energy of this band with respect
to the yrast line. The energies of the new transitions are consistent with the low alignment of this band (see Section 7.3.2).

The $K^\pi = 2^+ \gamma$ vibrational band was previously known up to $10^+$ [108] and $7^+$ [109] in the even and odd spins respectively. Three new states have been added, but above spin 10 the intensity of both the in-band and out-of-band decays is too low to allow firm spin and parity assignments to be made.

In the present work, band IV has been observed up to $5^+$ [97], although the $6^+$ member of this band was identified in Ref.[108]. In Ref.[97], the authors interpreted band IV in terms of a two $\gamma$ phonon excitation. However, a $K^\pi = 4^+$, hexadecapole phonon excitation is an alternative possibility.

7.2.3 The $K^\pi = 5^-$ band (band V)

No excited states in this band had been reported before this experiment, though the $K^\pi = 5^-$ bandhead, at 1629 keV, was well established [97], decaying by four intense E1 transitions to bands IV, III and I. The DCO ratios for the stretched E1 transitions lie in the range 0.7—>0.9, which is higher than the expected value for pure dipole transitions. However, nearly all the intensity feeding the $K^\pi = 5^-$ state is delayed by the $K^\pi = 7^-$ and $9^-$ isomers (see Sections 7.2.4 and 7.2.5 below). This may give the nucleus time to de-orientate [51], and thus the DCO ratio tends towards 1.0. Small admixtures of higher order multipoles could also be responsible for the increase in the DCO ratio.

By inspection of $\gamma$-ray time centroids, the half-life limit for the $K^\pi = 5^-$ state is found to be $t_{1/2} < 1$ ns. This band now extends to $I^\pi = (17^-)$, and in-band $\Delta I = 1$ transitions have been observed for the levels up to $11^-$. The $\gamma$-ray spectrum gated by the 411 keV $9^- \to 7^-$ transition in band V is shown in Figure 7.2. The alignment curve for this band (shown in Section 7.3.2), is very close to that for the $11/2^+$ band [110] in the odd-A neighbour, $^{185}$Os, suggesting that the same $i_{13/2}$ neutron is involved. The $\nu\{\frac{11}{2}^+[615], -\frac{1}{2}^-[510]\}$ coupling to non-maximum-$K$ is therefore proposed for the configuration of the $K^\pi = 5^-$
band. It should also be noted that there is a $t_{1/2} = 8.33 \mu s$, 5$^-$ isomer [90] at 1285 keV in the isotope $^{184}_{74}$W$_{110}$. This has been assigned the same 2-quasineutron configuration. The difference in half-life between these two $K^\pi = 5^-$ intrinsic states may be partly due to the absence of a low-lying $K^\pi = 4^+$ band in $^{184}$W, which in $^{186}$Os, provides a $K$-allowed decay route for the $K^\pi = 5^-$ level.

It is worth noting that the 1195 keV transition from the $K^\pi = 5^-$ state to the 4$^+$ state in the g-band has not been observed. This is surprising when the lower energy transition (760 keV) to the corresponding 6$^+$ state is quite intense. Yates et al. [111] and Spanhoff et al. [97] have also commented on the “remarkable” absence of this $\gamma$-ray transition. The intensity limit for the 1195 keV transition from Ref. [97], is <15 % of the 760 keV $\gamma$-ray intensity. From the present work the limit is now <1 % i.e. $I_\gamma^{1195} < 1$ in Table 7.1.

7.2.4 The $K^\pi = 7^-$ band (band VI)

The $K^\pi = 7^-$ isomer lying at 1775 keV, was measured [97] to have $t_{1/2} = 8.1\pm 0.4$ ns, although the spin and parity assignments were tentative. Here, the time-difference analysis leads to $t_{1/2} = 8.5\pm 0.3$ ns (Figure 7.3a) in agreement with the previous value [97]. This state de-populates via an intense 146 keV transition, which, from both intensity flow measurements and DCO analysis, is determined to be an E2 transition. A much weaker decay branch has now also been identified, through a 906 keV E1 transition to the g-band. Although the 8$^-$ and 9$^-$ states were observed in Ref. [97], they were not assigned as rotational members of the $K^\pi = 7^-$ band. This band has now been identified up to (22$^-$) in the even spins and (17$^-$) in the odd spin states. The preferred assignment for this band is $\nu \{ \frac{11}{2}^+ [615], \frac{5}{2}^- [512] \}$, which is consistent with the alignment for the other bands containing an $i_{13/2}$ neutron (see Section 7.3.2). As with band V, comparison with the tungsten isotone, yields a $K^\pi = 7^-$, 2.4 ns isomer with the same configuration [90]. This state lies at 1502 keV and decays by a strong E2 transition to the corresponding $K^\pi = 5^-$ bandhead, as observed here in $^{186}$Os.
7.2.5 The $K^{\pi} = 9^-$ band (band VII)

Spanhoff et al. [97] measured $t_{1/2} = 5.3\pm0.2$ ns for the level lying at 2166 keV, and observed the first excited state at 2432 keV, though no spin and parity assignments were made. The band has now been extended to $17^-$ and using time differences across the isomer, $t_{1/2} = 6.1\pm0.2$ ns has been obtained (Figure 7.3b). The isomer de-excites via an intense $K$-forbidden (197 keV) M1 transition to the $8^-$ level of the $K^{\pi} = 7^-$ band. The multipolarity was obtained from intensity balancing ($\alpha_T(\text{exp})=0.94\pm0.08$, $\alpha_T(\text{M1})=0.85$, $\alpha_T(\text{E2})=0.35$). There is also a second decay path, via a very weak transition, to the $7^-$ level of band V. The spin and parity assignment for this intrinsic state is $K^{\pi} = 9^-$, resulting from the
Figure 7.3: Summed time difference spectra across: (a) the 7⁻ isomer, at 1775 keV using 277 and 194 keV gates; (b) the 9⁻ isomer at 2166 keV gated by the 194, 197 and 397 keV transitions and; (c) the 434 keV level in the g-band using the 297 and 435 keV gates (prompt). The solid lines are the fits of a Gaussian peak (the FWHM's of which are (a) 33 ns; (b) 26 ns; and (c) 23 ns) convoluted with an exponential decay with the half-lives shown. A larger width for the Gaussian component corresponds to lower gating energies.
\( \nu \{ \frac{11}{2}^+ [615], \frac{7}{2}^- [503] \} \), 2-quasineutron coupling. Once again, the alignment curve closely matches that of bands V and VI and the \( 11/2^+ \) band in \(^{186}\text{Os} \) \([110]\), all containing the same high-\( j \) neutron.

Strong out-of-band transitions between bands VII and VI arise due to the near energy degeneracy of the states and the crossing of the two bands. For instance, the \( 10^- \) levels lie only 4 keV apart. These transitions provide additional confidence in the spin and parity assignments given to the levels of both bands. Towards the top of the bands there is another crossing and this leads, potentially, to an ambiguity in the interpretation, since it is not obvious to which of the two bands, each \( 17^- \) state belongs. This is resolved by the observation of a weak decay branch from a high-\( K \) structure with \( K^\pi = 19^- \) at 5333 keV, which is assumed to preferentially proceed to the band with the higher \( K \) value, namely the \( K^\pi = 9^- \) band.

### 7.2.6 The \( K^\pi = 10^+ \) band (band VIII)

The level lying at 2564 keV and the first two excited states of band VIII at 2807 and 3040 keV respectively were observed in Ref.\([97]\), but no spins or parities were given. In the present study the spin and parity assignments for band VIII have been determined from the competing 397 keV dipole transition to the \( K^\pi = 9^- \) isomer and the quadrupole transitions that feed the \( 8^+, 10^+, 12^+ \) and \( 14^+ \) members of the \( g \)-band. The most intense out-of-band transition is a 659 keV \( \gamma \) ray from the \( 14^+ \) state at the band crossing.

The bandhead has a half-life limit of <1 ns and a \( K^\pi = 10^+ \) assignment. Once again, the DCO ratio for the 397 keV E1 decay to the \( K^\pi = 9^- \) band is greater than that for a pure dipole decay, namely \( R_{DCO} = 0.72(5) \). This might be attributed to a slight de-orientation or to a possible (small) M2 admixture. \( \Delta I = 1 \) transitions have been observed between the excited states of band VIII, up to spin 17, and the corresponding \( \Delta I = 2 \) E2 crossover transitions have been observed up to (21) and (26) \( \hbar \) respectively. This is well beyond the band
crossing that takes place between the g-band and the \( K^\pi = 10^+ \) band. The observation of the yrare as well as the yrast states is critical in understanding the interaction at the crossing. Band mixing arguments can be used to support the high-\( K \) assignment at the crossing and this is discussed in detail in Section 7.3.1. Figure 7.4 shows the g-band and 10^+ band (even spin) transitions. The transitions between the odd spin states of band VIII are shown in Figure 7.5. In \(^{182}\text{W}\) \(^{[103]}\) and \(^{184}\text{Os}\) \(^{[80]}\) similar \( K^\pi = 10^+ \) bandheads have been observed with half-lives of 1.4 \( \mu \)s (at 2230 keV) and 20 ns (at 2236 keV) respectively. In both cases an M1/E2 transition to the 10^+ member of the ground-state band competes with the higher energy stretched E2 decay to the 8^+ level in the ground-state band. In \(^{186}\text{Os}\) no such transition has been observed and the intensity limit from the present study for the unobserved 10^+\( \rightarrow \)10^+ 495 keV decay is \( I_{495}^{\text{ upper limit}} < 2 \) in the units given in Table 7.1.

The \( K^\pi = 10^+ \) band is interpreted as being built on a \( \nu\{\frac{11}{2}^+[615], \frac{9}{2}^+[624]\} \) intrinsic state. These two high-\( j \) orbitals feel a strong Coriolis force, resulting in a large alignment, approximately twice that of the other 2-quasineutron bands (see Section 7.3.2). The \( 2\cdot\frac{13}{2} \) neutron structure, coupled to high \( K \), leads to the designation of the \( K^\pi = 10^+ \) band as a \( \pi \)-band \(^{[12]}\). The crossing with the g-band at \( I = 14 \) is discussed in Section 7.3.1.

### 7.2.7 The 4-quasiparticle structures

Two new levels (IX) have been observed at 3187 keV and 3433 keV that decay directly to the \( K^\pi = 10^+ \) band and the g-band by a series of dipole and quadrupole transitions. The strongest of these \( \gamma \) rays, with energies 624, 626 and 1118 keV, are \( \Delta I = 2 \) transitions from their DCO ratios, allowing a firm \( K^\pi = 12^+ \) assignment to be made. The spectrum gated by the 1118 keV transition is shown in Figure 7.5. The \( K^\pi = 12^+ \) state is tentatively assigned the \( \nu\{\frac{11}{2}^+[615], \frac{9}{2}^+[624], \frac{3}{2}^-[512], \frac{1}{2}^-[510]\} \) configuration, based on comparisons with BCS calculations (see Section 7.3.5).
Figure 7.4: A spectrum gated by the 663 keV transition in the $K^\pi = 10^+$ band. The in-band even-spin transitions in band VIII can be seen, together with the ground-state band transitions.

The first excited state of this structure has a weak branch to level IXa, at 3222 keV with a tentative $(12^+)$ label. The weak nature of the 211 keV and 1153 keV transitions to and from the level respectively, implies that this state is non-yrast. This does not rule out an $11^+$ assignment, but the ordering of these two transitions is obtained from the relative intensities and the higher energy $\gamma$ ray is preferred as the quadrupole decay, thus favouring $I^\pi = (12^+)$. Level IXb represents a level de-exciting via a 320 keV $\gamma$ ray to band IX. The DCO analysis favours a $\Delta I = 1$ transition, though due to the low intensity this remains unconfirmed.

The $I^\pi = 13^+$ level at 3433 keV is fed by a strong 299 keV E2 transition ($R_{DCO} = 0.94 \pm 0.07$) from an yrast $15^+$ state. There is a well formed rotational band built on this level, extending up to 20 $\hbar$, providing evidence that the 3732 keV $15^+$ state is a bandhead. The first excited level in the band is fed by a weak E2 branch from the yrast intrinsic $18^+$ state (discussed below). The configuration $\nu\{11^+[615], 9^+[624], 7^-[503], 5^-[512]\}$ is tentatively assigned to this band (see Section 7.3.5).

At 4496 keV there is an yrast $18^+$ state that de-excites by an intense 560 keV E2 $\gamma$ ray to the $K^\pi = 10^+$ band. There is also a weak decay branch from this level to the $K^\pi = 15^+$ band, which suggests that the $18^+$ state is non-
collective. It is improbable that this is a rotational excitation decaying out-of-band, because it is yrast and the $18^+$ members of the other low-lying rotational structures have all been identified. The simplest interpretation of this state is an intrinsic configuration. In addition, blocked BCS calculations, discussed later, predict a low-lying intrinsic state with a 4-quasineutron $K^\pi = 18^+$ configuration: $\nu\{\frac{11}{2}^+ [615], \frac{9}{2}^+ [624], \frac{9}{2}^- [505], \frac{7}{2}^- [503]\}$. The half-life limit for this state from the centroid-shift analysis is $t_{1/2} < 0.5$ ns. This half-life is exceptionally short for a level that predominantly decays via a $\Delta K = 8$ transition, which ought to be highly $K$-forbidden. This anomaly is investigated in detail in Section 7.3.4.

![Figure 7.5](image-url)

Figure 7.5: The top spectrum was produced by gating on the 1118 keV transition that feeds the $I^\pi = 10^+$ state in band I. The link between the $K^\pi = 18^+$ yrast state and band X can be seen at 395 keV. The bottom spectrum is gated by the 971 keV $\gamma$-ray transition from the $I^\pi = 12^+$ state in band VIII. The 523, 598 and 693 keV transitions linking the odd-spin states of the $K^\pi = 10^+$ band are indicated.
7.2.8 The positive parity, high-spin states

Above the $18^+$ yrast level, the intensity fragments between many different states, as reported in Ref. [104]. However, most of the positive parity states are populated in the decay of an $I^\pi = 28^+$ state at 7145 keV with $t_{1/2} < 2$ ns. A 154 keV transition de-populates this yrast level, and both DCO ratios and intensity balancing are consistent with an electric quadrupole decay. Figure 7.6 shows the $\gamma$-ray spectrum gated by this transition and Figure 7.7 shows the DCO ratios for coincident transitions gated by the 154 keV decay. The absence of a regular rotational structure is apparent in this decay, but even near the limit of angular momentum seen using this reaction, the intensity of the transitions remains high. Further down in the cascade, beyond the $I^\pi = 24^+$ state, the fragmentation through multiple decay branches is manifest. This, in conjunction with DCO ratios, allows the spin and parity assignments for all the positive parity states, populated in the decay of the $28^+$ state, to be determined. The negative parity states are discussed in Section 7.2.9.

Multi-quasiparticle calculations predict a low-lying $K^\pi = 28^+$ state with the configuration: $\nu \{ \frac{1}{2}^+[615], \frac{3}{2}^+[624], \frac{5}{2}^-[505], \frac{7}{2}^-[503] \} \otimes \pi \{ \frac{1}{2}^-[505], \frac{3}{2}^-[514] \}$. In addition to this state, there are many low-lying intrinsic states calculated to lie in the spin range between the 18$^+$ and 28$^+$ levels. This could account for the division of intensity between many, seemingly non-rotational states, causing the breakdown in rotational structure observed here and in Ref. [104]. Configuration assignments cannot be made unambiguously for these high-spin states because no accompanying rotational bands have been firmly established. However, potential rotational members have been identified above the $I^\pi = 20^+$ states at 5498 and 5503 keV and above an $I = 21$ level at 5704 keV. The weak population of these states arising from their non-yrast status makes definite spin and parity assignments difficult. Other weakly populated structures include $(24^+)$ and $(25^+)$ states at 6154 and 6995 keV respectively. No configuration assignments
have been made for these levels, because the nature of the states that they feed has not been determined.

Figure 7.6: Gamma-ray spectrum gated by the 154 keV, $28^+ \rightarrow 26^+$ transition in $^{186}$Os. A contaminant at 600 keV is denoted by c. Even in this high-spin region there are sufficient counts in the spectrum for useful DCO ratios to be obtained.

7.2.9 The negative parity, high-spin states

As mentioned in Section 7.2.5, the $K^\pi = 9^-$ band is fed by a $K^\pi = 19^-$ state at 5333 keV. This level was observed by Balabanski et al. [104], but was near the limit of the angular momentum input and some of the associated $\gamma$-ray coincidences are inconsistent with the new data. The discrepancy is most probably due to the presence of 3 transitions within 1.5 keV (530.1, 530.7, 531.5) in the same region of the level scheme. Apart from this, the two decay schemes (from Ref.[104] and the present study) are in excellent agreement, though the present study extends the level scheme considerably.

The level at 5333 keV decays to an (18$^-$) state and also to the $K^\pi = 9^-$ band.
CHAPTER 7. $^{185}$Os

Figure 7.7: DCO ratios for transitions gated by the 154 keV assigned E2 decay feeding the positive parity high-spin states. The open circles are dipole transitions and the filled circles are quadrupole transitions.

From DCO ratios, the 708 keV γ ray, to the 2-quasiparticle band, is consistent with a quadrupole transition which leads to a probable spin and parity of 19$^-$, in the absence of a measurable half-life ($t_{1/2} < 1$ ns). The 708 keV transition can be seen in the γ-ray spectrum of Figure 7.2 gated by the 229 keV transition above the $K^\pi = 19^-$ state. The assignment of the (18$^-$) level at 5027 keV as an intrinsic state is based on the observation of a strong decay to the yrast $K^\pi = 18^+$ state at 4496 keV. The DCO ratio for the 531.5 keV transition to this 18$^+$ state is 1.04±0.03, which is consistent with an $I \rightarrow I$ dipole transition. A level at 5246 keV is a possible candidate for a rotational excitation built on this state. The suggested configurations for these two negative parity states are, 18$^-$, $\nu\{\frac{11}{2}^+[615], \frac{9}{2}^+[624]\} \otimes \pi\{\frac{5}{2}^+[402], \frac{11}{2}^-[505]\}$ and 19$^-$, $\nu\{\frac{7}{2}^-[503], \frac{11}{2}^+[615]\} \otimes \pi\{\frac{9}{2}^-[514], \frac{11}{2}^-[505]\}$ by comparison with the results of blocked BCS calculations, described later in Section 7.3.5.
CHAPTER 7. \(^{186}\)OS

7.3 Discussion

In the following sections some aspects of the level scheme will be described in more detail. In Section 7.3.1 the \(K^\pi = 10^+\) band is discussed in conjunction with a two-band mixing scenario. In Sections 7.3.2 and 7.3.3 the determination of band alignments and g-factors for the rotational structures are explained. Hindered transitions are discussed in Section 7.3.4. In Sections 7.3.5 and 7.3.6 multi-quasiparticle and potential-energy-surface calculations are compared to the experimental observations and the shapes of the intrinsic states are discussed along with the consequences for the \(K\) quantum number.

7.3.1 Two-band mixing

In order to explain the intensities of the interband transitions between the \(K^\pi = 10^+\) and the g-band band, a two-band mixing calculation was performed as described in Section 2.7.

Table 7.2 shows the calculated \((R)\) and experimental \(B(E2)_{\text{out}}/B(E2)_{\text{in}}\) ratios between the \(K^\pi = 10^+\) t-band and the g-band for two different mixing scenarios. (Note, values of \(K(\text{band I})=0\) and \(K(\text{band VIII})=10\) have been used to obtain the branching ratios in Table 7.2.) The interaction strengths are calculated using Equations (2.39). Firstly, a calculation using a constant mixing strength, \(V_c\), was performed, so that at the crossing \((I = 14)\) the observed out-of-band to in-band branching ratio was reproduced. This matrix element was then fixed and used to calculate the branching ratio away from the crossing. It can be seen from Table 7.2 that the mixing between the 12\(^+\) and 10\(^+\) states gives a \(B(E2)\) ratio from the 12\(^+\) state that is an order of magnitude too high. A spin dependence was included in the mixing matrix strength of the form

\[ V_{SD} = A[I(I + 1) - K_{\text{max}}^2] \]

(7.1)

where \(A (=0.47\ \text{keV here})\) is a constant, \(I\) is the spin of the two mixing levels
Table 7.2: Two-level mixing strengths for a constant interaction, $V_c$, and a spin-dependent interaction, $V_{sd}$ (see text for details). The experimental out-of-band to in-band $B(E2)$ ratios are given in the third column. $R$ represents the corresponding calculated ratios for the mixing strengths shown. The initial state is in band VIII. $T^{\gamma}_{\text{out}}/T^{\gamma}_{\text{in}}$ is the out-of-band to in-band $\gamma$-ray intensity ratio.

<table>
<thead>
<tr>
<th>$I^+_\text{initial}$</th>
<th>$T^{\gamma}<em>{\text{out}}/T^{\gamma}</em>{\text{in}}$</th>
<th>$B(E2)<em>{\text{out}}/B(E2)</em>{\text{in}}$</th>
<th>$V_c$ (keV)</th>
<th>$R_c$</th>
<th>$V_{sd}$ (keV)</th>
<th>$R_{sd}$</th>
<th>$R_{\text{calc}}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>10$^+$</td>
<td>—</td>
<td>—</td>
<td>46.9</td>
<td>4.70</td>
<td>—</td>
<td>—</td>
<td></td>
</tr>
<tr>
<td>12$^+$</td>
<td>1.43</td>
<td>0.04</td>
<td>46.9</td>
<td>0.45</td>
<td>26.3</td>
<td>0.19</td>
<td></td>
</tr>
<tr>
<td>14$^+$</td>
<td>32.10</td>
<td>2.68</td>
<td>46.9</td>
<td>2.67</td>
<td>51.5</td>
<td>2.68</td>
<td></td>
</tr>
<tr>
<td>16$^+$</td>
<td>0.01</td>
<td>0.04</td>
<td>46.9</td>
<td>0.05</td>
<td>80.8</td>
<td>0.04</td>
<td></td>
</tr>
</tbody>
</table>

and $K_{\text{max}}$ is the higher $K$ value of the two bands (=10). The spin dependence given in Equation (7.1) was proposed by Walker et al. [112] to reproduce the observed $t$-band crossing in $^{180}$W [112] and $^{170}$W [22]. However, this is an empirically determined relation and other forms have been successfully used, for example, $V = A[I(I + 1)]$ [113]. The mixing for $I < 14$ is overestimated by both the constant and spin-dependent interactions. However, including the spin dependence relation lowers the calculated reduced branching ratio from the $I = 12$ state by more than a factor of 2, giving a ratio closer to the observed value. This may indicate that the interaction between bands I and VIII has a stronger spin dependence than that given in Equation (7.1).

If the $K$ values for bands I and VIII are set equal, the observed $B(E2)$ ratios at the crossing cannot be reproduced with either of the above spin dependence relations, supporting the high-$K$ assignment of band VIII. Further theoretical work is needed to provide insight into the role of $I$ and $K$ in the mixing matrix elements.
The unperturbed energies (with the appropriate mixing interaction, $V_{SD}$, removed) are plotted against the spin $(I(I+1))$ of the levels in Figure 7.8. The "pure" band comprises levels with energies that vary smoothly with spin.

Figure 7.8: A plot of level energy minus an arbitrary constant-rotor reference, versus $I(I + 1)$. The even-spin members of the g-band and $K^\pi=10^+$ band (solid lines) are shown, together with the unperturbed energies (dashed lines) for $K$(band I)=$0$ and $K$(band VIII)=$10$. The mixing strengths ($V_{SD}$) are given in Table 7.2. The scale has been expanded to enlarge the band crossing region. (The 16$^+$ level plotted for the g-band is the 4352 keV state (level Ib) discussed in Section 7.2.1.)

The high-$K$ crossing band is interpreted as arising from the aligned $t$-configuration [12] of two $i_{13/2}$ neutrons, which has been observed in other even-even nuclei in this mass region (Section 2.3.1). However, in $^{182}$W [114] the $t$-band appears not to be responsible for the first backbend in the yrast states, and in $^{180}$W [112] a strong ($\approx 140$ keV) interaction leads to ambiguities in the interpretation. Recently, Shizuma et al. [115] have reported a $t$-band crossing in $^{184}$Os, but there is a third band involved, complicating the behaviour. In $^{186}$Os, the weak mixing
results in a sharp crossing, that is explained by a two-band mixing model.

The signature splitting that is observed in the \( t \)-band, leading to a suppression of the even-spin energies with respect to the odd-spin states, probably arises due to a combination of deformation and Coriolis-mixing effects. A large hexadecapole deformation can move the low-\( \Omega \), high-\( j \) orbitals close to the Fermi surface [116]. A possible consequence of this effect is the occurrence of a close-lying \( s \)-band, comprising only even spins, lowering the corresponding \( t \)-band states. Band Ia is a candidate for such a structure due to its high alignment (Figure 7.9) and the observation of only one signature.

7.3.2 Alignments

The alignments shown in Figure 7.9, and used in Section 7.2 to support the configuration assignments, are calculated using the formulae of Purry et al. [28] given in Section 2.8 (Equations (2.40) and (2.41)).

The alignments for the \( K^{\pi} = 5^- \), \( 7^- \) and \( 9^- \) 2-quasiparticle bands are equal to that of the \( 11/2^+ \) band in \( ^{185}\text{Os} \) [110], which provides evidence that these configurations contain the same aligning \( i_{13/2} \) neutron. The \( K^{\pi} = 10^+ \) band has an alignment approximately twice that of the other seniority-2 bands, suggesting that it is built on two \( i_{13/2} \) neutrons. The alternative \( K^{\pi} = 10^+ \) assignment for band VIII is an \( (h_{11/2})^2 \) 2-proton configuration. These particles have a much lower alignment than the \( i_{13/2} \) neutrons. This can be seen in Figure 7.9, which shows the alignment of the \( \pi \{g_{3/2}^{-}[514]\} \) band in \( ^{185}\text{Re} \) [117].

The only 4-quasiparticle state upon which a well formed rotational band has been observed, is the \( K^{\pi} = 15^+ \) band (band X). The alignment for this band is also plotted in Figure 7.9. Although the assigned configuration contains the same two aligning neutrons as the \( K^{\pi} = 10^+ \) band, the alignment is significantly lower. This phenomenon can be explained in terms of a reduction in neutron pairing with increased neutron seniority. Although a decrease in the pairing interaction leads to an increase in the collective moment-of-inertia (equivalent to an increase
Figure 7.9: A plot of alignment versus rotational frequency. Harris parameters $\mathcal{J}_0^{(1)} = 22.1 \hbar^2 \text{MeV}^{-1}$ and $\mathcal{J}_1^{(2)} = 67.0 \hbar^2 \text{MeV}^{-3}$ have been used: (a) the 0- and 2-quasiparticle positive-parity bands. The $K^\pi = 11/2^+$ band in $^{185}\text{Os}$ and the $9/2^-$ band in $^{185}\text{Re}$ are plotted for comparison; (b) the negative-parity 2-quasiparticle bands and band X.
in alignment) the lowering of the $\Delta K = \pm 1$ mixing due to the loss of pairing has a greater affect, resulting in an overall quenching of the alignment [118]. This picture is consistent with the observations in the present work, and supports the configuration assignment for the $K^* = 15^+$ intrinsic state.

### 7.3.3 Branching ratios and g-factors

From the in-band $\gamma$-ray branching ratios the rotational model equations [18] can be used to extract g-factors, providing an observable that can be compared to predicted values, to support configuration assignments. The relevant equations are (2.21) and (2.22). The quadrupole moment $Q_0$ was taken from Ref.[9] and the potential-energy-surface calculations (Section 7.3.6) [37], both of which give 5.7 e·b. The rotational g-factors are calculated for each configuration, using the expressions of Belyaev [19] and Migdal [20] with the BCS pairing energies. Section 2.6 describes this approach in detail.

Table 7.3 summarises the information obtained from the in-band branching ratios for all the 2- and 4-quasiparticle bands for which $\Delta I = 1$ and 2 transitions have been observed. The calculated $g_R$ factors are also given. Only the magnitude of the value $(g_K - g_R)/Q_0$, can be obtained from the data because the branching ratio is related to the square of the mixing ratio. In principle the sign of the mixing ratio can be obtained from the DCO ratio [51]. However, due to a combination of large uncertainties and delayed feeding, the only in-band $\Delta I = 1$ transition for which the DCO ratio is statistically different from $\approx 0.56$ (for a pure dipole) is the 143 keV transition in band V. The DCO ratio is 0.47(2), implying a negative $\delta$, consistent with the calculated $(g_K - g_R)/Q_0$.

For a given orbital with angular momentum projection $\Omega$ on the symmetry axis, the Nilsson quantum numbers can be used to calculate the intrinsic g-factor. The calculated g-factors shown in Table 7.3 are found using Equation (2.30) with the appropriate Nilsson wave functions [24, 23]. In addition to this, the empirical g-factor from the $11/2^+$ band [110] in $^{186}$Os was used for the $(i_{13/2})$ neutrons, thus
Table 7.3: Branching ratios and g-factors for the 2- and 4-quasiparticle bands in $^{186}$Os.

| $K^\pi$ | I (h) | $E_2$ (MeV) | $E_1$ (MeV) | $T_2/T_1$ | $g_R$ | $|(g_K - g_R)/Q_0| (\text{e} \cdot \text{b})^{-1}$ |
|---------|-------|-------------|-------------|----------|------|-----------------------------------------------|
|         | Expt. | Calc.       |             |          |      |                                               |
| 5$^-$   | 7     | 0.3104      | 0.1673      | 1.22±0.06| 0.20 | 0.027±0.001                                  |
|         |       |             |             |          |      | -0.016                                       |
|         | 8     | 0.3619      | 0.1950      | 2.94±0.15| 0.20 | 0.025±0.001                                  |
|         | 9     | 0.4102      | 0.2165      | 4.01±0.20| 0.20 | 0.030±0.001                                  |
|         |       |             |             |          |      | 0.035±0.001                                  |
| 7$^-$   | 9     | 0.4134      | 0.2197      | 0.87±0.04| 0.21 | 0.028±0.001                                  |
|         |       |             |             |          |      | -0.024                                       |
|         | 10    | 0.4668      | 0.2472      | 1.36±0.06| 0.21 | 0.035±0.001                                  |
|         | 11    | 0.5106      | 0.2636      | 12.90±0.8| 0.21 | —                                             |
| 9$^-$   | 11    | 0.5492      | 0.2830      | 0.42±0.02| 0.21 | 0.044±0.002                                  |
|         |       |             |             |          |      | -0.062                                       |
|         | 12    | 0.5758      | 0.2924      | 0.30±0.02| 0.21 | 0.082±0.003                                  |
| 10$^+$  | 12    | 0.4761      | 0.2329      | 0.46±0.02| 0.26 | 0.035±0.001                                  |
|         |       |             |             |          |      | -0.053                                       |
|         | 13    | 0.4880      | 0.2547      | 1.08±0.05| 0.26 | 0.026±0.001                                  |
|         | 14    | 0.4011      | 0.1485      | 3.72±0.18| 0.26 | 0.025±0.001                                  |
| 15$^+$  | 17    | 0.7524      | 0.3831      | 0.39±0.02| 0.13 | 0.028±0.001                                  |
|         |       |             |             |          |      | -0.032                                       |
|         | 18    | 0.7698      | 0.3860      | 0.85±0.05| 0.13 | 0.027±0.002                                  |

Taking into account the rotational alignment of these particles. The uncertainties in the calculated values of $(g_K - g_R)/Q_0$ are expected to be approximately ±0.01 (e·b)$^{-1}$.

The experimental $(g_K - g_R)/Q_0$ values for the $K^\pi = 5^-$ band show a general increase with spin, which may be due to configuration mixing at higher angular momenta. However, the g-factors at $I = 7$ and 8 agree to within the errors. These values are reproduced well by the calculation, within about 0.01 of the experimental value. The sign of these factors is also in agreement with the negative
mixing ratio, determined from the DCO ratio for the $6^\rightarrow 5^-$ transition. The situation is similar for the $K^\pi = 7^-$ band, with the calculated value lying close to the experimental quantities. The two experimental values are significantly different from each other, which may be due to the mixing with the $K^\pi = 9^-$ band. The $K^\pi = 9^-$ band shows even larger deviations with spin, induced by the mixing with the $K^\pi = 7^-$ band. The calculation lies in between the two experimental values.

The $K^\pi = 10^+$ and $K^\pi = 15^+$ configurations both contain two $i_{13/2}$ neutrons. The pronounced alignment effects could be related to the poorer agreement with the calculation for the $K^\pi = 10^+$ band.

From the discussion above it is evident that the alignment of the $i_{13/2}$ neutrons in the 2- and 4-quasiparticle configurations can be approximately taken into account by using the empirical value for the single-quasiparticle band in $^{185}$Os. This results in generally good agreement between the calculated g-factors and those extracted from the experimental branching ratios in conjunction with the rotational model.

### 7.3.4 Hindrances of intrinsic-state decays

The $K^\pi = 7^-$ bandhead de-excites via both a $K$-allowed E2 transition to the $K^\pi = 5^-$ state and a $K$-forbidden E1 transition to the $6^+$ member of the g-band. In order to compare the rate of this decay with similar transitions in nearby nuclei, the hindrance per degree of $K$-forbiddenness (or reduced hindrance) can be calculated, given by Equation (2.20) ($f_\nu = \left[F_W\right]^{1/\nu}$). The 906 keV transition from the 8.5 ns isomer in $^{186}$Os has $f_\nu = 37$. Although no such transition has been observed from the 2 ns $7^-$ isomer in the isotone $^{184}$W, there is an 18 $\mu$s $7^-$ isomer in $^{184}$W [62] (Chapter 6). This state at 1517 keV decays, in part, through a 708 keV E1 $\gamma$ ray, directly to the $6^+$ level in the ground-state band, with $f_\nu = 83$. This is a factor of three higher than for the 906 keV transition in $^{186}$Os. There is also a $7^-$ isomer in the isotone $^{183}$Pt. The decay of this 0.2 ns level, at 1768
keV [9, 119], proceeds via a 584 keV electric dipole transition to the 6+ level in
the ground band, with $f_\nu = 9$. This is a factor of three smaller than the $^{186}$Os
transition, suggesting a decline in the reduced hindrance with increasing proton
number. The reduced hindrance for the 906 keV transition in $^{186}$Os, $f_\nu = 37$,
suggests that at low spins, the $K$ quantum number is at least partially conserved.

The $K^\pi = 10^+$ $t$-bandhead in $^{186}$Os has a half-life limit of <1 ns leading to
a reduced hindrance of $f_\nu < 3.1$ for the 1142 keV stretched E2 transition to the
ground-state band. The $f_\nu$ systematics for $t$-band decays are discussed in detail in
Section 5.3.1, but this upper limit of 3.1 observed for $^{186}$Os is in good agreement
with the low reduced hindrances for $\Delta I = 2$, E2 transitions de-exciting these
states.

For the following discussion of high-seniority states it is more transparent to
compare $F_W$ values rather than reduced hindrances. The yrast $K^\pi = 18^+$ state
at 4496 keV that feeds the $t$-band, decays by a $\Delta K = 8$ transition and yet the
half-life of this level is less than 0.5 ns, corresponding to $F_W \leq 3$ (or $f_\nu \leq 1.2$).
However, mixing with the 18+ member of the $K^\pi = 10^+$ band, which lies only 10
keV higher, could be playing an important role in the decay rate. To investigate
the possible effects of this, a comparison with $^{179}$W is made, where there is
a similar near superposition of intrinsic and collective states. A $K^\pi = 35/2^-$,
750 ns isomer in $^{179}$W [22], lies 22 keV above the $35/2^-$ member of a $K=23/2$
rotational band. In Ref.[22] the authors can explain the strong E2 decay from
the isomer to the $K=23/2$ band and the 750 ns half-life, with a mixing strength
of 24±3 eV between the two states. In view of the possibility of other mixing
mechanisms being significant, this direct mixing strength could be considered
to be an upper limit. This is a very weak interaction and it is only the close
proximity of the states that gives rise to any observable consequences. If the
same small mixing strength is used between the 18+ states in $^{186}$Os, a hindrance
factor of $F_W = 1.7 \times 10^3$ is calculated for the 560 keV transition. The observed
hindrance of $F_W \leq 3$ is more than 500 times smaller. It is therefore unlikely that
direct mixing with the $\ell$-band is responsible for the fast decay from the yrast $K^\pi = 18^+$ state. A similar comparison of hindrances can be made with the 6 ns isomer in $^{182}$Os \cite{80}. This $K^\pi = 16^+$ state decays via a $K$-forbidden 1530 keV transition, with $F_W \approx 3 \times 10^4$ (with a substantial uncertainty arising from the low intensity of the 1530 keV transition). Although this decay is extremely fast for a $K$-forbidden transition \cite{4}, the hindrance is still a factor of $\geq 10^4$ higher than for the 560 keV transition in $^{186}$Os.

A similar situation exists for the 154 keV E2 transition from the $K^\pi = 28^+$ state in $^{186}$Os, which has a hindrance of $F_W \leq 0.04$. Although this could be a $K$-allowed decay, such a fast transition rate is even collective. Statistical $K$-mixing with higher energy (unobserved) states is unlikely to explain this, due to the yrast nature of the $K^\pi = 28^+$ level. An additional mechanism must be facilitating the fast decay of these $^{186}$Os $K^\pi = 18^+$ and $28^+$ yrast states. (The energies of the intrinsic states, as a function of spin ($I$), relative to the yrast band are shown in Figure 7.10.) A candidate for such an effect is the onset of triaxiality. Once an intrinsic state is axially asymmetric, $K$-conservation would be absent, allowing it to de-excite by transitions with increased collectivity. The shapes of the multi-quasiparticle nuclear potentials are investigated in Section 7.3.6.

### 7.3.5 Multi-quasiparticle calculations

Multi-quasiparticle Nilsson model calculations \cite{31}, have been performed using blocked BCS pairing \textit{without} residual interactions (Section 3.1). The experimental energies are compared to the lowest calculated energies for a given configuration in Table 7.4. The single-particle proton and neutron energies were adjusted (consistently with the monopole pairing strength; see below) to reproduce the single-quasiparticle energies \cite{9, 82}, where available, in the odd-$A$ neighbours, $^{185}$Re and $^{185}$Os. The deformation parameters, $\varepsilon_2 = 0.198$ and $\varepsilon_4 = 0.056$ \cite{81} were used with axial symmetry ($\gamma = 0^\circ$). The neutron monopole pairing strength, $G_\nu = 21.5/A$ MeV, was chosen to reproduce (within $\sim 200$ keV), the observed en-
energies of the 2-quasineutron bandheads. Due to the absence of any 2-quasiproton states in this study, the proton pairing strength was chosen to be $1.5/A$ MeV higher than the corresponding neutron value ($G_p = 23.0/A$ MeV), which is consistent with Ref. [31].

Residual nucleon-nucleon interactions (Section 3.1.3) typically shift the energies of the states by $\sim 100$ keV, with favoured configurations (like particles with intrinsic spins antiparallel and unlike particles with intrinsic spins parallel) being moved to lower energies and unfavoured couplings being shifted towards higher energies. When a qualitative allowance is made for these interactions the agreement between calculation and observation is generally good. The two unfavoured 2-quasiparticle states ($K^\pi = 9^-$ and $K^\pi = 10^+$) are predicted (in the absence of residual interactions) to occur lower in energy than observed, and the favoured
Table 7.4: Multi-quasiparticle states in $^{186}$Os.

<table>
<thead>
<tr>
<th>$K^*$</th>
<th>Configuration$^a$</th>
<th>Energy (keV)</th>
<th>$\Delta E^c$ (keV)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Neutrons</td>
<td>Protons</td>
<td>Expt.</td>
</tr>
<tr>
<td>$5^-$</td>
<td>$\nu{11/2^+, -1/2^-}$</td>
<td>1629</td>
<td>1707</td>
</tr>
<tr>
<td>$7^-$</td>
<td>$\nu{11/2^+, 3/2^-}$</td>
<td>1775</td>
<td>1855</td>
</tr>
<tr>
<td>$9^-$</td>
<td>$\nu{11/2^+, 7/2^-}$</td>
<td>2166</td>
<td>1828</td>
</tr>
<tr>
<td>$10^+$</td>
<td>$\nu{11/2^+, 9/2^+}$</td>
<td>2564</td>
<td>2338</td>
</tr>
<tr>
<td>$12^+$</td>
<td>$\nu{11/2^+, 9/2^+, 3/2^-, 1/2^-}$</td>
<td>3187</td>
<td>3252</td>
</tr>
<tr>
<td>$15^+$</td>
<td>$\nu{11/2^+, 9/2^+, 7/2^-, 3/2^-}$</td>
<td>3732</td>
<td>3390</td>
</tr>
<tr>
<td>$18^+$</td>
<td>$\nu{11/2^+, 9/2^+, 9/2^-, 7/2^-}$</td>
<td>4496</td>
<td>4286</td>
</tr>
<tr>
<td>$(18^-)$</td>
<td>$\nu{11/2^+, 9/2^+}$</td>
<td>$\pi{11/2^-, 5/2^+}$</td>
<td>5027</td>
</tr>
<tr>
<td>$19^-$</td>
<td>$\nu{11/2^+, 7/2^-}$</td>
<td>$\pi{11/2^-, 9/2^-}$</td>
<td>5333</td>
</tr>
<tr>
<td>$28^+$</td>
<td>$\nu{11/2^+, 9/2^+, 9/2^-, 7/2^-}$</td>
<td>$\pi{11/2^-, 9/2^-}$</td>
<td>7145</td>
</tr>
</tbody>
</table>

$^a$ Neutrons: $11/2^[615], 9/2^[624], 9/2^-[505], 7/2^-[503], 3/2^-[512], 1/2^-[510]$

$^b$ Protons: $11/2^-[505], 9/2^-[514], 5/2^+[402].$

$^b$ Blocked BCS calculation without residual interactions [31].

$^c$ Calculated energy minus experimental energy.
2-quasiparticle states are calculated to lie higher than the experimental value. This is further evidence in support of the assignments made to these states and is supported by the alignments and E2/M1 branching ratios (Section 7.3.3).

The energies calculated for the seniority-4 and -6 states show a similar relation with the experimentally observed levels. However, for the negative parity configurations, there is a greater difference between the two energies. For instance the 19\(^{-}\) state occurs 649 keV higher than predicted. This is probably due to the choice of the proton pairing strength. A larger value for G\(_p\) would increase the calculated energies for all the states involving protons. Although this would increase the energy difference for the K\(^\pi\) = 28\(^{+}\) state, residual interactions would be expected to lower the energy for this configuration.

These axially symmetric calculations successfully predict the energies of the intrinsic states to within a few hundred keV. However, the half-lives of the K\(^\pi\) = 18\(^{+}\) and K\(^\pi\) = 28\(^{+}\) states are anomalously short. As discussed in the previous section, the 18\(^{+}\) yrast state at 4496 keV, seemingly decays by a \(\Delta K = 8\) transition to the K\(^\pi\) = 10\(^{+}\) band, and yet the half-life is less than 0.5 ns. To understand the mechanism responsible for this behaviour, potential-energy-surface calculations have been performed (see Section 7.3.6).

### 7.3.6 Potential-energy-surface calculations

Calculations describing the shape evolution for \(^{186}\)Os with increasing rotational frequency have been presented by Balabanski \textit{et al.} [104]. These Total Routhian Surface (TRS) calculations include particle alignment to give the yrast minimum in the \(\beta-\gamma\) plane as the nucleus is cranked. A substantial change in structure is predicted to occur at \(I \approx 18\hbar\), from collective to non-collective (high-\(K\)) prolate rotation. However, although this adiabatic approach provides a general understanding of the nuclear potential, the intrinsic states using this technique are always mixed [120]. To identify the specific orbitals responsible for the dramatic change in behaviour, a new method [37] of diabatic blocking has been employed.
Configuration-constrained, Potential-Energy-Surface (PES) calculations [37], as described in Section 3.2, have been performed for the intrinsic states observed in $^{186}$Os. The results are summarised in Table 7.5. The excitation energies calculated using this method are compared to the experimental energies and those found using the BCS method. The expectation value of $K$ ($\tilde{K}$ in Table 7.5) is calculated for the intrinsic states by projecting the single-particle wave functions on to the long nuclear axis and summing over all the occupied orbitals. Figure 7.11 shows the PES's for some of the states. States are still referred to by their axially symmetric $K$ value for convenience.

Table 7.5: Properties of intrinsic states in $^{186}$Os, observed and calculated in the present study.

| $K^\pi$ | $t_{1/2}$ $(\text{ns})$ | Energy (keV) | | Deformation $^b$ | | $\tilde{K}^b$ |
|---------|-----------------|--------------|----------------|--------------------|----------------------------------|
|         |                 | Expt. | Calc.I $^c$ | Calc.II $^b$ | $\beta_2$ | $\beta_4$ | $|\gamma|$
| gs      | —               | 0    | 0            | 0                | 0.189    | -0.054    | 1°      | 0       |
| 5$^-$   | <1              | 1629 | 1707         | 1899             | 0.193    | -0.051    | 0°      | 5.0     |
| 7$^-$   | 8.5(3)          | 1775 | 1855         | 1970             | 0.196    | -0.061    | 0°      | 7.0     |
| 9$^-$   | 6.1(2)          | 2166 | 1828         | 2079             | 0.186    | -0.051    | 6°      | 9.0     |
| 10$^+$  | <1              | 2564 | 2338         | 2070             | 0.192    | -0.055    | 1°      | 10.0    |
| 18$^+$  | <0.5            | 4496 | 4286         | 4000             | 0.186    | -0.039    | 23°     | 16.9    |
| 28$^+$  | <2              | 7145 | 7141         | 6020             | 0.179    | -0.032    | 26°     | 26.5    |

$^a$ No residual interactions are included in the calculated energies.

$^b$ Calculated using the method of Xu et al. [37].

$^c$ Calculated using the method of Jain et al. [31].

For the negative parity 2-quasineutron states the energies from both calculations are in good agreement with the observations, but for the higher lying states
the PES calculations underestimate the excitation energies by an increasing margin. The energy of the $K^\pi = 28^+$ level is calculated to be more than 1 MeV lower than observed. The reason for this is most likely to do with the choice of (configuration independent) monopole pairing strengths \cite{37} but a critical evaluation of this aspect is beyond the scope of the present study. Nevertheless, adjustments to better reproduce the energies would not affect the calculated shapes.

Inspection of the calculated shape parameters (Table 7.5) shows that, as the nucleus gains angular momentum, the shape initially remains axially symmetric, which is consistent with the observation of low-lying isomeric states and well-formed rotational bands. Indeed, the potential minimum for the $7^-$ isomer (see Figure 7.11) is more well defined than for the ground-state. This persistence of axial symmetry, up to and including the $K^\pi = 10^+$ configuration, means that for these configurations $K$ is equal to the total angular momentum of the intrinsic state.

However, the $K^\pi = 18^+$ and $28^+$ states are no longer axially symmetric, see Figure 7.11. The yrast $18 \hbar$ configuration has a triaxial parameter, $\gamma = 23^\circ$. This dramatic shape change results in $K$ being reduced to $16.9 \hbar$. This is due to the occupation of the two $9/2^- [505]$ and $7/2^- [503]$ neutron orbitals which drive the nucleus towards large $\gamma$ deformations. The triaxiality is even more pronounced for the $K^\pi = 28^+$ state ($\gamma = 26^\circ$), because of the excitation of a proton from the axially symmetric $9/2^- [514]$ orbital to the $\gamma$-driving $11/2^- [505]$ single-particle state. Again, the calculated shape of the potential leads to a reduction in $K$, to $26.5 \hbar$. These two high-spin states are almost completely triaxial ($\gamma = 30^\circ$), which can explain why $K$ is not conserved in their decay. The very low hindrance of the 560 keV and 154 keV E2 transitions (Section 7.3.1), from the $K^\pi = 18^+$ and $28^+$ yrast levels respectively, supports the results of these configuration-constrained PES calculations. Indeed, given these results it would be anomalous if these states were long-lived.
7.4 Summary

The nucleus $^{186}$Os has been studied up to high-spins following fusion-evaporation reactions with a radioactive neutron-rich $^{14}$C beam. New 2- and 4-quasiparticle bands built on the $K^\pi = 7^-$ and $9^-$ isomers, together with $K^\pi = 5^-, 10^+$ and $15^+$ intrinsic states, have been observed, and the $\beta$ and $\gamma$ vibrational bands have been extended. Configuration assignments have been made by analysis of alignments, $g$-factors, energy systematics and comparison with blocked BCS calculations. The first crossing of the $g$-band has been attributed to a high-$K$ $\pi$-band and the relative transition strengths are explained by a two-band mixing calculation.

At higher excitation energies there is a wealth of intrinsic states on or near the yrast line with an absence of associated half-lives. Analysis of transition hindrance factors reveals that, while $K$ is at least partially conserved in the low-spin domain, it is completely eroded for $I \geq 18$ h. These discoveries are in excellent agreement with predictions by configuration-constrained PES calculations, of triaxial nuclear potentials. The observed $K^\pi = 18^+$ and $28^+$ yrast intrinsic levels are calculated to have $\gamma = 23^\circ$ and $26^\circ$ respectively, corresponding to almost complete triaxiality.
Figure 7.11: Diabatic potential-energy-surface calculations. Top left: the ground state, top right: the 5\(^{-}\) state, middle left: the 7\(^{-}\) state, middle right: the 10\(^{+}\) state, bottom left: the 18\(^{+}\) state, bottom right: the 28\(^{+}\) state. There is reflection symmetry in the \(\gamma = 0^\circ\) prolate axis.
Chapter 8

Conclusion

High-spin $K$-isomers have been populated in neutron-rich nuclei in the mass-180 region. New 2- and 4-quasiparticle isomers have been observed in $^{186}$W with half-lives of 18 $\mu$s ($K^\pi = 7^-$) and $\geq 3$ ms ($K^\pi = (16^+)$) respectively. The latter of these is a 4-quasiparticle yrast trap decaying through a multitude of high-$K$ structures which feed the 2-quasiparticle $K^\pi = 7^-$ state. In $^{181}$Ta new 3-quasiparticle isomers with half-lives of 170 ns, 25 $\mu$s and 210 $\mu$s have been observed de-exciting to the $K^\pi = 9/2^-$ band. Predictions of blocked BCS calculations are in good agreement with the excitation energies of the observed intrinsic states.

A new 1-quasiparticle band populated in the decay of a $t_1/2 \geq 1$ ms 3-quasiparticle isomer has been tentatively placed in $^{185}$Ta. A $K^\pi = 21/2^-$ assignment has been proposed for this isomer by comparison with the lighter tantalum isotopes and blocked BCS calculations. An in-band $\gamma$-ray branching ratio analysis yields an intrinsic $g$-factor that is consistent with a $\pi\{\frac{3}{2}^-[514]\}$ assignment for the new band. Additional measurements with longer beam pulsing are necessary in order to measure the long half-lives of the $^{186}$W and $^{185}$Ta isomers.

The strong production of the $^{181}$Ta and $^{186}$W target isomers ($\sigma \approx 10 - 100$ mb) by inelastic and transfer mechanisms using a $^{238}$U beam, establishes this technique as suitable for populating $K$-isomers that are inaccessible with fusion-evaporation reactions using stable beams and targets. This method is likely to
yield further results at higher spins and higher neutron numbers when coupled to more sensitive detector arrays. In addition, other isotopes and target materials in this region, such as rhenium and osmium, offer access to a wider range of product nuclei that are as yet unstudied.

An alternative method of populating neutron-rich nuclei is to use fusion-evaporation reactions with radioactive beams. One of the few radioactive beams available with sufficient intensity is $^{14}$C and this has been used successfully to explore the high-spin structure of $^{186}$Os. New rotational bands have been observed built on the $K^\pi = 7^-$ ($t_{1/2} = 8.5$ ns) and $9^-$ ($t_{1/2} = 6.1$ ns) isomers as well as on the $K^\pi = 5^-$, $10^+$ and $15^+$ multi-quasiparticle states. Previously known $\beta$ and $\gamma$ vibrational bands have been extended to higher spins, and a candidate for a rotational aligned low-$K$ $s$-band has been observed. Configuration assignments have been made following an analysis of alignments, g-factors, energy systematics and comparisons with blocked BCS calculations. The first backbend in the ground-state band at $14 \hbar$ is caused by the crossing of a high-$K$ $t$-band. The out-of-band to in-band transition strengths can be explained by a two-band mixing calculation.

The high-spin yrast structure of $^{186}$Os is dominated by intrinsic states that have no measurable lifetimes ($\leq 1$ ns). Comparisons with configuration-constrained potential-energy-surface calculations suggest that the onset of axially asymmetric nuclear shapes is responsible, resulting in $K^\pi = 18^+$ and $28^+$ yrast states with triaxiality parameters of $23^\circ$ and $26^\circ$ respectively. This leads to a breakdown in $K$-conservation manifest as a loss of isomerism at high angular momentum.

The measurement of short half-lives ($<1$ ns) and the observation of rotational bands built on the triaxial states may in the future allow testing of the proposed configuration assignments above $I = 18$.

The work discussed in this thesis has enabled the behaviour of neutron-rich nuclei to be explored and the reliability of the calculations to be proven at high neutron numbers. The elucidation of the intrinsic structures will be important for
future studies trying to reach the very high spin domain (>30 ħ). Observation of high seniority states are of interest in order to better understand the persistence of the pairing interaction, although heavier radioactive projectiles will be needed for such experiments.
Appendix A

Back rotation

The bands observed in deformed nuclei have bandheads with $I = K$, i.e. all the angular momentum is aligned along the symmetry axis. This is because of a phenomenon called "back rotation". For a particular single-particle orbital, the angular momentum $j$, can have a projection $\Omega (=j, j-1, j-2...1/2)$ on the symmetry axis. Each of these energy states are doubly degenerate since the projections $-\Omega$ and $+\Omega$ have the same energy. Under these circumstances, the total angular momentum of the given state would not, in general, be equal to $K$ ($=\sum \Omega_i$, for $i$ single particles). However, the nucleons in the core of the nucleus can rotate collectively, independent of the single-particle rotations, so that the total angular momentum is equal to the sum of the single-particle and collective components.

Consider a prolate deformed nucleus with a nucleon rotating at an angle $\theta$ to the rotation axis (which is perpendicular to the symmetry axis) about which the core is rotating. See Figure A.1.

Let the angular frequency of the nuclear core in the laboratory frame be $\omega_0$, and the angular frequency of the single-particle in the same frame be $\omega_p$ at an angle $\theta$ to that of the core. The energy of this system in the laboratory frame can be written as,
Figure A.1: Diagram showing single-particle and collective rotations in a prolate deformed nucleus.

\[ E = \frac{1}{2} J_0 \omega_0^2 + \frac{1}{2} J_p \omega_p^2 \]  \hspace{1cm} (A.1)

where \( J_0 \) and \( J_p \) are the moments of inertia for the core and single nucleon respectively. The angular frequency of the nucleon can be split into two components; parallel to the symmetry axis \( \omega_p^\parallel \) (\( = \omega_p \sin \theta \)) and perpendicular to the symmetry axis \( \omega_p^\perp \) (\( = \omega_p \cos \theta \)). The corresponding components in terms of the angular momentum of the particle are

\[ \omega_p^\parallel = \frac{i\tau_p}{J_p} \]  \hspace{1cm} (A.2)

and

\[ \omega_p^\perp = \frac{i\tau_p}{J_p} + \omega_0 \]  \hspace{1cm} (A.3)

where \( \frac{i\tau_p}{J_p} \) is a constant equal to the particle angular frequency, along the rotation axis, in the rotating frame of the core. Substituting Equations (A.2) and (A.3)
in to (A.1) gives,

\[ E = \frac{1}{2} J_0 \omega_0^2 + \frac{1}{2} J_p \left[ \left( \frac{i_p}{J_p} + \omega_0 \right)^2 + \left( \frac{i_p^r}{J_p} \right)^2 \right] \]  

(A.4)

To find the rotational frequency of the core, at which the energy of the system is a minimum, Equation (A.4) must be differentiated with respect to \( \omega_0 \) and set equal to zero (see below).

\[ \frac{dE}{d\omega_0} = J_0 \omega_0 + J_p \omega_0 + i_p^r = 0 \]  

(A.5)

This implies that \( J_0 \omega_0 = -J_p \omega_0 \), which means the angular momentum of the core is equal and opposite to the rotational aligned component of the particle angular momentum \( (i_p^r) \). This means that the lowest energy state in a given rotational band is one with the total angular momentum pointing along the symmetry axis, i.e. \( I = K \). The rotational component of the particle angular momentum is cancelled out by the collective back rotation of the core.
Appendix B

Branching ratios

In Section 2.6 Equations (2.21) and (2.22) were introduced that relate the in-band γ-ray branching ratio, \( \lambda_b \), to the intrinsic and rotational g-factors (gyromagnetic factors) \( g_K \) and \( g_R \) respectively. These expressions are derived below. The definition of \( \lambda_b \) is

\[
\lambda_b = \frac{T_2(E2)}{T_1(E2) + T_1(M1)} \quad \text{(B.1)}
\]

where \( T_1 \) and \( T_2 \) are the γ-ray transition rates (\( \propto \) intensity) for \( \Delta I = 1 \) and 2 transitions respectively. Stretched E2 transitions are assumed here to be of pure multipolarity, because any admixture of M3 (or higher multipoles) will be very small. Note that the ratio \( \frac{T_1(M1)}{T_1(E2)} = \frac{1}{2} \), where \( \delta \) is the quadrupole/dipole mixing ratio. The transition rates [72] are related to the reduced transition probabilities \( B(XL) \) by Equations (B.2).

\[
B(E2; I_i \rightarrow I_f) = \frac{1}{1.225 \times 10^9 E_\gamma^5} T(E2; I_i \rightarrow I_f)
\]
\[
B(M1; I_i \rightarrow I_f) = \frac{1}{1.758 \times 10^{13} E_\gamma^3} T(M1; I_i \rightarrow I_f) \quad \text{(B.2)}
\]

Where \( E_\gamma \) is the γ-ray transition energy in MeV. Substituting for the transition probabilities in Equation (B.1) using Equations (B.2) yields,
The reduced in-band transition probabilities [72] are given by,

\[ B(E2; I_i K \rightarrow I_f K) = \frac{5}{16\pi} e^2 Q_0^2 | < I_i 2K0|I_f K > |^2 \]
\[ B(E2; I_i K \rightarrow I_f K) = \frac{5}{16\pi} e^2 Q_0^2 | < I_i 1K0|I_f K > |^2 \] \hspace{1cm} (B.4)
\[ B(M1; I_i K \rightarrow I_f K) = \frac{3}{4\pi} e^2 | < I_i 1K0|I_f K > |^2 (g_K - g_R)^2 K^2 \]

where \( Q_0 \) is the intrinsic quadrupole moment in units of e-fm\(^2\) and \( g_K \) and \( g_R \) are the intrinsic and rotational gyromagnetic ratios respectively. The relevant Clebsch-Gordon coefficients [26] are,

\[ E2 : | < I_i 2K0|I_f K > | = \left[ \frac{3(I - K)(I - K - 1)(I + K)(I + K - 1)}{(2I - 2)(2I - 1)I(2I + 1)} \right]^{1/2} \]
\[ E2 : | < I_i 1K0|I_f K > | = -K \left[ \frac{3(I - K)(I + K)}{(I - 1)I(2I + 1)(I + 1)} \right]^{1/2} \] \hspace{1cm} (B.5)
\[ M1 : | < I_i 1K0|I_f K > | = - \left[ \frac{(I - K)(I + K)}{I(2I + 1)} \right]^{1/2} \]

where \( I \) is the angular momentum of the initial state. \( K \) is assumed to be a good quantum number throughout. Substituting Equations (B.4) and (B.5) into Equation (B.3) results in,

\[ \lambda_b = \left( \frac{E_2}{E_1} \right)^5 \frac{B_2(E2) \times 1.225 \times 10^9}{B_1(E2) \times 1.225 \times 10^9 + B_1(M1)E_1^{-2} \times 1.758 \times 10^{13}} \] \hspace{1cm} (B.3)

Rearranging these terms gives Equation (B.7).

\[ \lambda_b = \left( \frac{E_2}{E_1} \right)^5 \left( \frac{5}{16} \left[ \frac{3}{(I - 1)(I + 1)} \right] + \frac{3}{4} \left[ \frac{(g_K - g_R)^2 K^2}{Q_0^2 E_1^2} \left( \frac{1.758 \times 10^{13}}{1.225 \times 10^9} \right) \right] \right) \] \hspace{1cm} (B.7)
This leads to the following expression for the branching ratio,

\[ \lambda_b = \left( \frac{E_2}{E_1} \right)^5 \left( \frac{\left[ (I-K-1)(I+K-1)(I+1) \right]}{2I(I+1)(2I^2 K^2)} \right) \]

\[ \left( 1 + \frac{4}{3} \left[ \frac{(g_K - g_R)^2}{Q_0 E_1^2} \right] \right) \]

\[ (B.8) \]

This equation is identical to that quoted by Alexander et al. [18]. To convert \( Q_0 \) from units of e-fm\(^2\) to e-b, there is an additional factor required, namely, \( Q_0 \rightarrow \frac{Q_0}{100} \). Making this substitution yields Equation (B.9)

\[ \frac{(g_K - g_R)}{Q_0} = \sqrt{\frac{5}{4}}(0.8347) \frac{E_1}{\sqrt{(I^2 - 1)}} \left( \frac{1}{\lambda_b} \left( \frac{E_2}{E_1} \right)^5 \left[ \frac{(I-K-1)(I+K-1)(I+1)}{2K^2(2I-1)} \right] - 1 \right)^{1/2} \]

\[ (B.9) \]

This result can also be expressed in terms of the quadrupole/dipole mixing ratio \( \delta \),

\[ \frac{(g_K - g_R)}{Q_0} = 0.933 \frac{E_1}{\delta \sqrt{(I^2 - 1)}} \]

\[ (B.10) \]

where \( \delta \) is related to the quadrupole admixture \( (q) \) by,

\[ q = \frac{\delta^2}{1 + \delta^2} = \frac{2K^2(2I-1)}{(I-K-1)(I+K-1)(I+1)} \frac{E_2^3}{E_1^3} \lambda_b \]

\[ (B.11) \]

Equations (B.10) and (B.11) here correspond to Equations (2.21) and (2.22) in Section 2.6.
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Chapter 6


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