Errata.

p30, figure 1.6, Y-axis is in units of keV; line 3, should read ‘$T_z=0$ nucleus [5]’
p33, line 3, ‘form’ should be ‘from’
p38, line 19, ‘or the’ should be ‘for the’; line 20, ‘for for’ should be just ‘for’; line 9, (two protons) should be $T_z=-1$ and (two neutrons) should be $T - z = +1$
p39, figure 2.2, same as above.
p46, $A_1$ and $A_2$ should be $M_T$ and $M_B$
p52, $J_i$ is defined as ‘the vector of the angular momentum of the compound nucleus.’
p53, $A_0$, $A_2$ and $A_4$ should be $a_0$, $a_2$ and $a_4$; $W(\theta)$ should be $I(\theta)$; between equation 3.5 and 3.6, ‘quadrupole radiation,’ should be ‘quadrupole and mixed quadrupole/dipole radiation,’
p56, line 6 and line 8, $m c^2$ should be $m_0 c^2$
p57, line 7, ‘70cm’ should be ‘70mm’
p58, line 1, ‘10cm’ should be ‘25cm’; line 18, ‘80%’ should be ‘60%’
p59, line 19, The $\Delta \theta$ term comes also from the change in recoil direction due to the recoil cone.
p71, figure 3.13, ‘seperation’ should be ‘separation’
p74, figure 3.15, The full scale of the X-axis is approximately 270 ns
p80, section 4.2 should be called ‘Selection of Real Events.’
p92, line 12, ‘statistical E2’ should be ‘statistical E1’
p108, line 6, ‘was’ should be ‘were’
p134, figure 5.13, caption should include ‘with background subtraction’
p136, figure 5.15, caption should include ‘with Doppler corrections’
p137, line 7, ‘froward’ should be ‘forward’; line 22, ‘si-ball’ should be ‘Si-ball’
p143, figure 5.19, The 571 gated spectra, ‘571’ peaks should be labelled ‘376’
p149, line 1, ‘if figure’ should be ‘in figure’
p151, figure 6.1, the caption should indicate that ‘the lowest lying state from each calculation is aligned with the ground state of $^{63}$Ga’. 
Spectroscopy of $N \approx Z$ Nuclei around $A=60$ Using AYEBALL and PEX.

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Abstract

This thesis is the result of a PhD study at the University of Surrey, aimed at the study of high angular momentum states in neutron deficient nuclei in the mass 60 region, with approximately equal numbers of protons and neutrons, \((N\sim Z)\). The main motivations of this work are to provide an insight into the mechanisms for generation of high angular momentum states in a limited particle valence space above the \(N=Z=28\) doubly magic core, and to investigate the role of the isospin quantum number in heavy \(N=Z\) nuclei. The decay scheme for the odd-odd \(N=Z\) nucleus \(^{62}\text{Ga}\) has been deduced for the first time, and the decay schemes for \(^{61}\text{Zn}\) and \(^{61}\text{Cu}\) have been extended. The data came from two experiments, the first using the reaction \(^{24}\text{Mg} + ^{40}\text{Ca}\) at a beam energy of 65 MeV, performed at the Argonne National Laboratory, Chicago, US, using the germanium \(\gamma\)-ray detector array ‘AYEBALL’ in conjunction with the Argonne Fragment Mass Analyzer and a gas filled ionisation chamber. The second experiment using the reaction \(^{28}\text{Si} + ^{40}\text{Ca}\) at a beam energy of 88 MeV was performed at the Niels Bohr Institute, Risø, Copenhagen, Denmark using the ‘PEX’ \(\gamma\)-ray detector array with a charged particle detector ball and an array of liquid scintillator neutron detectors. The data analysis techniques and results of the experimental analysis are presented. Gamma-ray energy spectra for different nuclei are shown according to the mass, neutron number and proton number of the nucleus. The proposed decay schemes are justified by coincidence and DCO arguments, and are compared to shell model calculations using a restricted \(pf_\frac{3}{2}g_\frac{1}{2}\) basis. In the case of \(^{62}\text{Ga}\), these are then compared with the latest \(IBM-4\) calculation, which explicitly includes \(T=0\) and \(T=1\) bosons. Suggestions are also made for future work to complement this data.
Acknowledgements

This thesis has only been made possible thanks to the EXCELLENT supervision of Dr. Paddy Regan and Prof. Bill Gelletly. Paddy has been a great supervisor, a collaborator and a mentor. Bill has been an inspiration ever since my undergraduate course, and is really the reason I started this PhD. My unofficial supervisor Prof. Dave Warner at Daresbury has also been a great influence and tutor, being able to teach me the finer points of $n - p$ pairing is no mean feat.

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The biggest thank you is saved for my family. My mum, whose unequivocal support, both financially and personally, has been immense. My late dad, whose passion for science and tireless dedication to my future inspired me from an early age, and my brother Nige for many constructive observations. Finally, my fiancee, Teri, who has been the single biggest influence in my life in the last five years. Her unwavering tolerance and constant support through my stress and poverty has been an unenviable task. I owe this thesis to her.
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6.1 Comparison of the low lying energy levels of calculations for $^{62}$Ga and
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Since the discovery of the electron by J.J. Thompson 100 years ago [1], physicists have been attempting to describe the concentration of positive charge in the atom, required to give an overall neutral electric potential. Rutherford, with an experiment performed by Geiger and Marsden in 1911 [2] proposed the notion of a nucleus, and with the advent of quantum mechanics, and interventions by Heisenberg, Planck and Chadwick, the current basic model of the nucleus was derived.

The basis of the model is that the nucleus is a quantum mechanical system of order \( fm \) \(^\dagger\) in diameter, with mass \( A \), and is composed of \( Z \) protons and \( N \) neutrons held together by a strong nuclear force. The protons and neutrons can combine in many configurations, subject to certain constraints, to produce several thousand different isotopes. The Segré chart of the nuclides (figure 1.1) is a map of all the nuclear species currently known to exist. It can be seen that for \( A \leq 40 \), the nuclei which are stable with respect to radioactive decay have approximately equal proton and neutron number. However, for \( A \geq 40 \), the Coulomb repulsion between protons necessitates an excess of neutrons in the stable nuclei.

As the neutron excess in stable nuclei increases with mass, then accordingly nuclei with equal numbers of protons and neutrons (\( N=Z \)) become increasingly exotic compared to the line of beta stable nuclei, and consequently more difficult to pro-

\(^\dagger m = \text{Fermi} = 10^{-15} \text{ m}\)
duce and investigate experimentally [3]. The neutron separation energy \((S_n)\) is the amount of energy required to remove a neutron from the nucleus, and is equal to the difference in binding of the nucleus with and without the neutron. Similarly, the proton separation energy \((S_p)\) is the energy required to remove a proton. The limits of the existence of nuclei are defined to be where the separation energy of the last nucleon is zero, and are referred to as driplines. The proton dripline is on the neutron-deficient side of the stable nuclei and is predicted to cross the \(N=Z\) line somewhere just above \(^{100}\)Sn, [4].

The low lying excited states of the even-even \(N=Z\) nuclei have been successfully studied in excited states in nuclei up to \(^{84}\)Mo. Figure 1.2 shows the energies of the first \(2^+\) state of the even-even \(N=Z\) nuclei as a function of mass from \(^{64}\)Ge to \(^{84}\)Mo [5, 6, 7, 8, 9, 10]. Also plotted is the empirically derived estimates of deformation \(\beta_2\) calculated from the energy of the \(2^+\) state using Grodzins estimate, (equation 1.1) [11] (see section 1.7).
Figure 1.2: (top) Experimentally determined values of $E(2^+)$ for $N=Z$, even-even nuclei from $^{64}$Ge to $^{94}$Mo. (bottom) Calculated values of deformation, derived using Grodzins's estimate.
However, the odd-odd \(N=Z\) nuclei between \(N=Z=28\) and 50 have remained largely unexplored, and it is this region where the phenomena of \(T=0\) \(n-p\) pairing may manifest itself most clearly [12]. This thesis is a study of nuclear structure in this region of the Segré chart using the method of gamma-ray spectroscopy.

1.1 Nuclear Spins and Parities.

The total angular momentum of a nucleon, \(j\), is defined as the coupled sum of the orbital angular momentum, \(l\), and the spin, \(s\) (denoted simply as \(j = l + s\)). Each nucleon in a nucleus can be labeled with these three quantum numbers. The vector sum of the total angular momenta of all the nucleons in a nucleus is then the total angular momentum of the nucleus, \(J\), (or sometimes \(I\)) referred to as nuclear spin. A level with total angular momentum quantum number \(j\) is allowed by the Pauli exclusion principal to accommodate \(2j + 1\) like nucleons having magnetic substates of angular momentum, \(m_j = j, j - 1, ..., -j\). The magnetic substate, \(m_j\), is equal to \(m_l + m_s\), where \(m_l\) has a \(2l + 1\) degeneracy, all being integer values, and \(m_s\) can be \(\pm \frac{1}{2}\hbar\). The magnetic substate, \(m_j\), is therefore half-integral, so an even number of nucleons will have an even number of half-integral components, resulting in an integral value for the magnetic substate (\(z\) component) of nuclear spin \(J\). Similarly, if the number of nucleons is odd, the total \(z\) component, and hence \(J\) must be half-integral. The nuclear parity operator \(P\) (or \(\pi\)), is defined as having the property of reflecting the position coordinate of every nucleon in a system through the origin and is determined by \(\pi = (-1)^l\), positive being even, and negative being odd. A nuclear state can therefore be labelled with both its spin and parity, the usual symbolism being \(J^\pi\). For nucleons in a closed subshell, the parities are all the same, and since there is always an even number of them, then the total parity must be \(\pi = +1\) (ie \(J^\pi = 0^+\)).
1.2 The Nuclear Force.

The nucleon-nucleon interaction is called the *strong force* and can be inferred as having the following properties [13]:

- It is attractive and stronger than the Coulomb force at distances up to a few fm. This enables the nucleons to overcome the repulsive effect of their electrostatic charge, and remain bound in closely packed systems.

- It is weaker than the Coulomb force at distances larger than a few fm, becoming negligible at atomic distances, thus atomic properties can be explained based only on the Coulomb force.

- It contains a repulsive component at very small distances (<fm) which prevents the nucleons collapsing into singularity, and keeps them at an average separation of around 1.8 fm. The effect of this is that to first order, the nuclear size is roughly proportional to the mass.

- It is almost *charge independent*, acting nearly equally on protons and neutrons. The effect of this can be observed in certain pairs of nuclei called *mirror nuclei*. These pairs consist of one nucleus with N neutrons and N+1 protons, and another nucleus with N+1 neutrons and N protons. For example, in the case of the nuclei $^{24}$Mn and $^{25}$Cr [14, 15], the energy level structure of the two nuclei are almost identical up to $\sim 6$ MeV, with the energies of the corresponding levels in the two nuclei similar to within 100 keV or less. The observed difference can be attributed almost entirely to the Coulomb force, inferring that the nuclear force acts essentially equally on protons and neutrons, and is therefore charge independent.

- It is dependent on the relative spins of the nucleons, which accounts for the experimental observation that the deuteron ($^2$H$_1$) has only one bound state (see section 1.3).

- It has a noncentral tensor component which does not conserve orbital angular momentum, and accounts for the deformation of the deuteron (see section 1.3).
1.3 The Deuteron.

The simplest nuclear system in which the nuclear force may be studied is the deuteron (\(^{2}\text{H}_{1}\)) [16]. It has a ground state which is bound by 2.23 MeV but has no bound excited states. The ground state is experimentally determined to have a spin of \(J=1\) and orbital angular momentum which is predominantly \(l=0\) (called an \(s\) state), and must therefore have a total intrinsic spin \(S=1\). If the nuclear force was independent of the relative spin orientations of the nucleons, then there would be a second bound state with \(S=0\) (and thus \(J=0\)) with the same energy as the \(J=1\) state. Since no such state exists, the nuclear force is deemed to be spin dependent. The deuteron has a measured magnetic moment \(\mu = 0.86\mu_{N}\) [17] which is slightly different to the value calculated for a pure \(s\) state deuteron, and an electric quadrupole moment \(Q = 2.82 \times 10^{-31} \text{m}^2\) [18], indicating that on average, it has a deformed shape (see section 1.7). This is not possible if the deuteron wavefunction is purely \(s\) state since zero orbital angular momentum implies no angular dependence and hence spherical symmetry and a zero electric quadrupole moment \((Q=0)\). The wavefunction must therefore have a higher order angular momentum component. Since nuclear states have definite parity, no mixture of parities is allowed, so the mixture must have an even parity \(\left((-1)^l\right)\), therefore \(l=1\) (called \(p\) states) are not allowed, and the only possible cases are \(l=2\) (called a \(d\) state), an \(l=4\) (called an \(g\) state), and so on with decreasing probability. It must also have the experimentally observed \(J=1\), so \(S\) must equal 1. The measured quadrupole and magnetic moments are accounted for with an admixture of \(\sim 3-5\%\) \(d\) state and the rest \(s\) state. There is a component of the nuclear force which causes this admixing, called the tensor force. The effect is that since the quadrupole moment of the deuteron is slightly positive, a configuration where the nucleon spins are oriented parallel, rather than perpendicular to the symmetry axis, must be energetically favourable.
1.4 Isospin.

The proton and the neutron differ in mass by only $\sim 0.1\%$, and because the strong nuclear force is approximately charge independent, the main difference in the way they interact is due to the much weaker electromagnetic interaction. It is therefore sometimes mathematically convenient to group them together as two states of a single particle, the nucleus. The proton and neutron are differentiated by a fictitious quantum number isospin ($T$) [13]. The two nucleon states are then said to have degenerate isospin vectors $t = \frac{1}{2}$, with orientations $m_t = +\frac{1}{2}$, or 'isospin-up' for the neutron and $m_t = -\frac{1}{2}$, or 'isospin-down' for the proton. In any nucleon system, these spin vectors can then couple in the same manner as angular momentum vectors. Therefore semiclassically, two nucleons can couple aligned to give total isospin $T=1$, or anti-aligned to give $T=0$. The $z$ axis orientation of this total isospin, $T_z$ is then the vector sum of the $z$ components of the individual nucleons, or the difference in the number of protons and neutrons.

$$T_z = \frac{1}{2}(N - Z) \quad (1.2)$$

1.5 Pairing Interaction.

There is a tendency for nucleons to couple as pairs to more stable configurations. This is due to the pairing force which can be divided into three categories: p-p, n-n, and p-n interactions. If the nuclear force is approximately charge independent, then the strength of these three interactions will be nearly equal. If we consider the low lying energy states in an isobaric triplet such as $^{26}_{12}$Mg, $^{28}_{13}$Al and $^{26}_{14}$Si (figure 1.3) [16], then the charge symmetry of the p-p and n-n force is immediately obvious in the even-even isobars where the forces lead to very similar energy states. On close inspection, there are also states in the odd-odd isobar $^{26}_{13}$Al, which have energies.

$^2$The $m_t$ values of the nucleons are sometimes defined the other way round (eg. reference [18].) This is to be consistent with the convention used by particle physicists.
close to similar states in $^{26}\text{Mg}_{14}$ and $^{26}\text{Si}_{12}$ relative to the lowest $0^+$ state, suggesting the p-n interaction is also of a similar strength. However, there are also many other states in $^{26}\text{Al}$ with no analogues in $^{26}\text{Mg}$ or $^{26}\text{Si}$, and this situation is not confined to the A=26 isobars. The difference between the $T=0$ and the $T=\pm 1$ systems, is that the p-n $T=0$ system can exist in two different configurations, namely $T=0$ and $T=1$, while the p-p and n-n $T=\pm 1$ systems are confined to a single $T=1$ state. This is discussed in section 2.3.

Figure 1.3: The lower energy levels of the A=26 isobars $^{26}\text{Mg}_{14}$, $^{26}\text{Al}_{13}$ and $^{26}\text{Si}_{12}$, with energies shown in keV. The inset shows the binding energies in keV relative to the stable isotope $^{26}\text{Mg}_{14}$ [16].

Evidence for the pairing force can be found in the fact that there are many more beta-stable nuclei which have even numbers of protons and neutrons (even-even nuclei) than either nuclei with an odd total number of nucleons, (odd-A), or
nuclei with odd numbers of both protons and neutrons, (odd-odd nuclei). Other evidence comes from the observation that all even-even nuclei have $J^\pi=0^+$ ground states, suggesting that all the nucleons are paired off into $J^\pi=0^+$ pairs. This effect must be accounted for in any theoretical calculations of the mass of the nucleus, (see equation 2.1), and also supports the assumption of the shell model that the total spin and parity of a nucleus is mainly determined by the valence (un-paired) nucleons, (see section 2.2).

1.6 Nuclear Electromagnetic Moments.

Because the electromagnetic force is much weaker than the strong force (as discussed in section 1.2), we can use electromagnetic interactions with nuclei to probe the nuclear structure without seriously affecting the motion or nature of the strongly interacting nucleus [18]. Two electromagnetic moments are often measured which reveal the way in which magnetism and charge are distributed in the nucleus, namely the magnetic dipole moment ($\mu$) and the electric quadrupole moment ($Q$).

1.6.1 Nuclear Magnetic Moments.

For an electron with charge $-e$ and mass $m$ orbiting with angular momentum $L$, the orbital magnetic moment, $\mu_L$, and the intrinsic magnetic moment $\mu_s$ associated with its spin $s$, are given by equations 1.3 and 1.4 respectively [18];

\[
\mu_L = -\frac{e}{2m}L \tag{1.3}
\]

\[
\mu_s = -\frac{g_e}{2m}s \tag{1.4}
\]

so for an electron with spin $\frac{1}{2}$,

\[
\mu_s = -\frac{1}{2}g\mu_B \tag{1.5}
\]
where $g$ is known as the $g$-factor, which is given by Dirac's relativistic wave equation as $2.00000$, and $\mu_B$ is the Bohr magneton, defined as [19]

$$\mu_B = \frac{e\hbar}{2m_e} = 9.274 \times 10^{-24} JT^{-1}.$$ (1.6)

The equivalent magnetic moments for the proton and neutron are defined as;

$$\mu_p = g_p \frac{e}{2m_p} s_p = \frac{1}{2} g_p \mu_N$$ (1.7)

$$\mu_n = g_n \frac{e}{2m_p} s_n = \frac{1}{2} g_n \mu_N$$ (1.8)

where $\mu_N$ is the nuclear magneton, defined in the same way as the Bohr magneton, except with the mass of a proton instead of an electron, and equal to $\mu_N = 5.051 \times 10^{-27}$ JT$^{-1}$ [19]. The measured values of nucleon $g$-factors are $g_p = +5.5856$ and $g_n = -3.8262$ [19]. It is then a simple step to define the total magnetic moment operator $\mu_J$ as;

$$\mu_J = g_J \frac{e}{2m_p} J$$ (1.9)

where $J$ is the spin, $g_J$ is the nuclear $g$-factor and $\mu_J = g_J \mu_N J$.

It is expected from this equation that the ground states of even-even nuclei, being $J^e=0^+$, will have no magnetic moments. For odd-$A$ nuclei, the expected value of magnetic moment can be calculated using the single-particle shell model. Experimentally measured values can then be compared with calculations to test theories and spin assignments.

1.6.2 Nuclear Electric Quadrupole Moments.

This quantity is a measure of the deviation in average shape of a nucleus from spherical symmetry. The intrinsic electric quadrupole moment is observed only in
the intrinsic nuclear frame (ie. where the nucleus is not spinning) and is defined classically as $Q_0$, by the equation [18]:

$$Q_0 = \int \rho_{ch}(r)(3z^2 - r^2) dV$$  \hspace{1cm} (1.10)

where the $z$-axis is taken to be along the axis of symmetry of the nuclear spin, and the charge density $\rho_{ch}$ is dependent on the radial vector $r$.

The experimentally measured quantity, the static quadrupole moment for a state with spin $I$, is $Q(I)$, (measured in the lab frame) and is related to the intrinsic quadrupole moment, $Q_0$ and the $K$ quantum number by equation 1.11 [20], where $K$ is the projection of the total nuclear spin onto the symmetry axis.

$$Q = \frac{3K^2 - I(I + 1)}{(I + 1)(2I + 3)} Q_0(K)$$  \hspace{1cm} (1.11)

1.7 Nuclear Deformations.

The two most common forms of nuclear deformation are *prolate* where the nucleus is stretched from two ends into a 'rugby ball shape', and *oblate* where the nucleus is squashed from two ends into a 'pancake' shape, (figure 1.4.)

The intrinsic quadrupole moment $Q_0$ is equal to 0 for spherical nuclei, is positive for prolate, and negative for oblate deformation. It is related to the deformation parameter $\beta_2$ by the expression [18];

$$Q_0 = \frac{3}{\sqrt{5\pi}} R_{av}^2 Z \beta_2 (1 + 0.16\beta_2)$$  \hspace{1cm} (1.12)

where $Z$ is the proton number and $R_{av}$ is the mean radius which is taken to be [18, 20]

$$R_{av} = R_0 A^{1/3} fm$$  \hspace{1cm} (1.13)
with $R_0$ typically around 1.2 $\text{fm}$.

Figure 1.5 shows the plot of a Woods-Saxon calculation [22] of single particle levels as a function of deformation of the nucleus, $\beta_2$, for neutrons in the mass 60 region [23]. The plot shows the spherical ‘magic numbers’ at $N=Z=28$, 40 and 50 (see section 2.4). One can also observe pronounced regions of low level density, or ‘shell gaps’, at deformation parameter $\beta_2 \sim +0.4$ for $N,Z=38$, and $\beta_2 \sim -0.35$ for $N,Z=36$. This mirrors the result of the Grodzins estimates (figure 1.2) which suggested that $^{76}\text{Sr}$ is highly deformed in its ground state. There is consequently a dramatic variation in the shape of the nucleus with the addition or removal of just a few nucleons in this mass region.

Nuclei with equal numbers of protons and neutrons ($N=Z$) inherently possess a natural symmetry with protons and neutrons simultaneously populating the same orbitals and shell gaps, thus amplifying any deformation driving effects. Calculations using both a shell-correction technique, and the pairing-self-consistent cranking method with a non-axially deformed Woods-Saxon potential [24] predict a progression in the shape of $N=Z$ nuclei from spherical at $Z=28$, through oblate/prolate...
Figure 1.5: The Wood-Saxon diagram of single particle levels for neutrons (central \(N=31\)) as a function of deformation \(\beta_2\). The plot for protons is very similar. Shell gaps are labelled, as are the nuclear orbitals. Taken from [23].

co-existence to a large \(\beta_2=0.4\) prolate deformation for \(^{76}\)Sr, then back through the same sequence to spherical at the doubly magic \(^{100}\)Sn [25, 26].

1.8 Gamma-Rays.

Most radioactive decays and nuclear reactions leave the final nucleus in an excited state, from which, if it is particle bound (and not forbidden by gamma selection rules), it de-excites by the emission of gamma-rays \((\gamma\text{-rays})\). Gamma-rays are photons of electromagnetic radiation, and as such, can be described classically in terms of multipole moments of charge distributions, namely dipoles, quadrupoles, octupoles, hexadecapoles etc. Gamma-rays have intrinsic spin \(1h\) and negative parity, and must
conserve angular momentum thus [20];

$$|J_i + J_f| \geq L \geq |J_i - J_f|$$  \hspace{1cm} (1.14)

where $L$ is the angular momentum carried by the photon, $J_i$ is the spin of the initial level and $J_f$ is the spin of the final level. Note, a $0^+ \rightarrow 0^+ \gamma$ decay is consequently forbidden. For a transition from a level with spin $2\hbar$ to a level with spin $0\hbar$, i.e. $\Delta J=2$, then $L$ must $= 2$. For a transition from a level with spin $3\hbar$ to a level with spin $2\hbar$, i.e. $\Delta J=1$, then $L$ can equal 1, 2, 3, 4 or 5.

The parity change due to the $\gamma$-ray is directly related to the angular momentum $L$ by the relations [20];

$$\Delta \pi(EL) = (-1)^L$$  \hspace{1cm} (1.15)

for electric multipole radiation, and;

$$\Delta \pi(ML) = (-1)^{L+1}$$  \hspace{1cm} (1.16)

for magnetic multipole radiation. Thus for a transition from a $3^+$ level to a $2^-$ level, the multipoles $E_1$, $M_2$, $E_3$, $M_4$ and $E_5$ are possible. In fact all of them will occur with competing probabilities and transition rates. The transition probability $T$ for radiation of electric or magnetic multipolarity $\lambda L$ is given by the expression [18, 27, 28];

$$T(\lambda L) = \frac{8\pi(L + 1)}{\hbar L((2L + 1)!!)^2} \left(\frac{E_\gamma}{\hbar c}\right)^{2L+1} B(\lambda L)$$  \hspace{1cm} (1.17)

where $B(\lambda L)$ is the reduced transition probability for a $\gamma$-ray of multipolarity $\lambda$ and angular momentum $L$, reduced meaning the spin and parity information is removed. This contains nuclear structure information about the initial and final states of a transition. The reduced transition probability for a transition from initial state
\( J_i \) to final state \( J_f \) is calculated using reduced matrix elements \( Q(\lambda L) \) and equation 1.18. These reduced matrix elements are often described in terms of Weisskopf units which are estimates of transition strengths for a single proton in a uniform charge distribution [28];

\[
B(\lambda L) = \frac{1}{2J_i + 1} | < J_f | Q(\lambda L) | J_i > |^2
\]  
(1.18)

A list of values for transition probabilities and \( B(\lambda L) \)'s are given in table 1.1 [28].

Table 1.1: Transition probabilities \( T(\text{sec}^{-1}) \) expressed by \( B(\text{EL}) \) in \((e^2 fm)^{2\lambda}\) and \( B(\text{ML}) \) in \( \left( \frac{e^2}{2mc} \right)^2 (fm)^{2\lambda-2} \). \( E \) = Gamma-ray energy, measured in MeV.

<table>
<thead>
<tr>
<th>Transition rate</th>
<th>Weisskopf estimate</th>
</tr>
</thead>
<tbody>
<tr>
<td>( T(E1) ) = 1.587 \times 10^{15} E^3 B(E1)</td>
<td>( B(E1) = 6.446 \times 10^{-2} A^{3} )</td>
</tr>
<tr>
<td>( T(E2) ) = 1.223 \times 10^{9} E^5 B(E2)</td>
<td>( B(E2) = 5.940 \times 10^{-2} A^{3} )</td>
</tr>
<tr>
<td>( T(E3) ) = 5.698 \times 10^{2} E^7 B(E3)</td>
<td>( B(E3) = 5.940 \times 10^{-2} A^{2} )</td>
</tr>
<tr>
<td>( T(E4) ) = 1.694 \times 10^{-4} E^9 B(E4)</td>
<td>( B(E4) = 6.285 \times 10^{-2} A^{3} )</td>
</tr>
<tr>
<td>( T(E5) ) = 3.451 \times 10^{-11} E^{11} B(E5)</td>
<td>( B(E5) = 6.928 \times 10^{-2} A^{1\frac{1}{3}} )</td>
</tr>
<tr>
<td>( T(M1) ) = 1.779 \times 10^{13} E^3 B(M1)</td>
<td>( B(M1) = 6.446 )</td>
</tr>
<tr>
<td>( T(M2) ) = 1.371 \times 10^{7} E^5 B(M2)</td>
<td>( B(M2) = 1.650 A^{\frac{5}{2}} )</td>
</tr>
<tr>
<td>( T(M3) ) = 6.387 \times 10^{9} E^7 B(M3)</td>
<td>( B(M3) = 1.650 A^{3} )</td>
</tr>
<tr>
<td>( T(M4) ) = 1.899 \times 10^{-6} E^9 B(M4)</td>
<td>( B(M4) = 1.746 A^{2} )</td>
</tr>
<tr>
<td>( T(M5) ) = 3.868 \times 10^{-13} E^{11} B(M5)</td>
<td>( B(M5) = 1.924 A^{1\frac{2}{3}} )</td>
</tr>
</tbody>
</table>

If the measured transition probability is much greater than the Weisskopf estimates, then the nuclear structure is favourable for such a transition, indicating similar wavefunctions for the initial and final states, and a possible collective motion involving many nucleons. Typical numbers of Weisskopf units for transitions of different multipolarity in different mass regions can be found in tables [29], and
these can be compared with experimentally determined values, possibly indicating a favourable multipolarity assignment for a previously unknown transition.

The transition probability for an electric radiation is roughly two orders of magnitude larger than for an equivalent magnetic transition, meaning electric transitions are favoured, \( \frac{A(E^1)}{A(M^1)} \ll 1 \). The probability is also inversely proportional to the multipolarity meaning higher multipoles are slower, \( \frac{A(E^1)}{A(M^1)} \ll 1 \). The effect of this is that while an E2 transition can compete with an M1, the M3 and higher orders are negligible. Also, an M2 cannot compete with an E1, which is therefore generally described as a 'pure' dipole.

### 1.8.1 Isomers.

If the decay from a nuclear state is hindered, for example due to unfavourable multipolarity, or due to dissimilarities in the wavefunctions of the initial and final state, then it is said to be isomeric. The definition of an isomer is arbitrarily decided to be a state which lives longer than a certain time. An uninhibited nuclear state typically lives for a few picoseconds. For many purposes, including those relevant to this report, 1 nanosecond is a reasonable lower limit for an isomeric lifetime.

### 1.9 Beta Decay.

Beta decay is a manifestation of the weak interaction, and is a process by which an unstable nucleus decays to a more energetically bound isobar. In figure 1.3, \(^{26}\text{Al}^*\) undergoes \(\beta^+\) decay to \(^{26}\text{Mg}^{14}\) as a proton decays to a neutron, a positron (\(e^+\)) and a neutrino (\(\nu_e\)). Similarly, \(\beta^-\) decay involves a neutron decaying to a proton, an electron and an antineutrino (\(\bar{\nu}_e\)), and electron capture sees a proton and an inner (usually K-shell) electron converting into a neutron and a neutrino. A typical beta decay energy is of the order of 1 MeV [16], and the momentum of an electron or neutrino with such an energy would be \(\sim 5 \times 10^{-22} \text{Jm}^{-1}\text{s}^{-1}\). The emitted particles carry away angular momentum \(\simeq L\hbar(L = 0, 1, 2, \ldots)\), which corresponds to an
emission at a distance \( \frac{L^2}{p} = 2L \times 10^{-13}m \) from the centre of the nucleus. This, from a nucleus of typical radius \( \sim 10^{-14}m \) (equation 1.13) limits \( L = 0 \) to be the most probable (called allowed) case, with \( L = 1 \) and \( L = 2 \) and so on, being called first forbidden\(^3\) and second forbidden, etc respectively. If \( L = 0 \), then the parity is even, and so cannot be changed by the decay. The electron, neutrino and their anti partners all have spin \( \frac{1}{2}\hbar \), thus the emitted pair can combine to have \( S = 0 \), which must keep the angular momentum unchanged (historically called a Fermi decay), or \( S = 1 \), which must remove one unit of angular momentum in the decay (called a Gamow-Teller decay). In the example of \( ^{26}\text{Al}^+ \rightarrow ^{26}\text{Mg}^+ + e^- + \nu_e \), the decay from the excited \( 0^+ \) state in \( ^{26}\text{Al} \) to the \( 0^+ \) ground state in \( ^{26}\text{Mg} \) must be pure Fermi type as no angular momentum is transferred. The decay \( ^{9}\text{He}^+ \rightarrow ^{9}\text{Li}^+ + e^- + \bar{\nu}_e \) happens between a \( 1^+ \) and a \( 0^+ \) state, \( \text{i.e. } \Delta S = 1 \) and is therefore a pure Gamow-Teller decay. There can also be mixing between Fermi and Gamow-Teller decay modes, for example in the decay of a neutron (spin \( \frac{1}{2}^+ \)) to a proton (spin \( \frac{1}{2}^+ \)), or the general case of decays between mirror nuclei, which have the same ground state spins and parities.

The lifetimes of beta decays are dependent \([18]\) on; (i) the statistical factor which accounts for the number of final states allowed for the decay; (ii) the Fermi function which accounts for the Coulomb effect; and (iii) the nuclear matrix element squared, \( |M_{fi}|^2 \), which is dependent on the initial and final nuclear states involved in the decay. If \( H \) is an operator which operates on the wave function of the initial state \( \psi_i \) to make \( H\psi_i \), then the matrix element \( |M_{fi}| \) is a measure of how much \( H\psi_i \) is like the wave function of the final state \( \psi_f \) \([30]\), \( \text{i.e.} \)

\[
M_{fi} = \int \psi_f^* H\psi_i \, d\tau \quad (1.19)
\]

Often, the comparative half-life \((ft)\) is quoted, where the \( f \) factors are tabulated against \( Q_\beta \) and \( Z \) \([31]\) and incorporate the statistical and Fermi functions, leaving

\(^3\)These decays are not actually forbidden, although they are less probable and so tend to have longer half-lives.
only the matrix element factor. In this way it is possible to make a direct comparison between decay \( ft \) values in any nuclei. It is sometimes convenient to quote \( \log_{10} ft \) values, which range from \( \sim 3 \) to \( 20 \) [18]. The form of \( ft \) for the case of a Fermi superallowed decay is given in equation 1.20 [18]

\[
f_{0^+} = 0.693 \frac{2\pi^3 h^7}{g^2 m_0^5 c^4 |M_{fi}|^2}
\]

(1.20)

where \( g \) is the \( \beta \)-decay strength constant, equal to \( 0.88 \times 10^{-4} \) MeV \( \bar{f} m^3 \) [18]. This value is experimentally determined from measurements made between two \( 0^+ \) states eg. \( ^{26}\text{Al} \beta^+ \rightarrow ^{26}\text{Mg} \), where the matrix element \( |M_{fi}| = \sqrt{2} \). This decay is also between states of identical isospin, namely \( T=1 \), and is called [32] a Fermi superallowed decay. Since these decays require no change in isospin or angular momentum, they are uninhibited, and are thus epitomised by a relatively fast decay time, having \( \log ft \) values around 3-4. The decay selection rules can be very useful in determining the spin, parity and isospin values of decaying states. If the spin, parity and isospin of the daughter state is known, and the decay is found to be of a Fermi superallowed character, then the parent must necessarily have the same spin, parity and isospin as the daughter.

1.10 Motivation For the Current Work.

In near stable nuclei up to \( N=Z=28 \), the Coulomb forces in the nucleus are small and the charge independence of the strong nuclear force is well understood. Thus, for these nuclei, the nuclear models do not differentiate between protons and neutrons, i.e. isospin is un-important in theoretical models. In the mass region above \( N=Z=50 \), the Coulomb forces are large, but the excess of neutrons required to counter this repulsive force means that the valence protons and neutrons are in different major shells, so they do not interact, and isospin remains a good quantum number. However, in the \( N,Z=28-50 \) shell, there is a non negligible Coulomb force, and the valence protons and neutrons both populate the same major shells. Thus for nuclei
with \( N \sim Z \), the protons and neutrons can interact, and consequently it is unclear how well the isospin condition holds in this region.

As the nuclei increase in mass away from the closed \( N=Z=28 \) shell, they become more deformed [6, 33, 34], so there is a transition from states which are best described by the spherical shell model, to a more collective structure which requires a deformed basis [35]. However, unlike higher mass regions where this transition has been studied, in the 28-50 shell, the valence protons and neutrons can both populate the same major shell of the oscillator potential, and this, coupled with the appreciable Coulomb force necessitates the consideration of the isospin degree of freedom in any theoretical collective framework.

In general, the lowest lying states in a nucleus have an isospin value, \( T \) which is equal to the \( T_z \) value. Thus, in \( N=Z \) nuclei, where \( T_z=0 \), one might expect the \( T=0 \) states to have the lowest energy. In the odd-odd nuclei, this is consistent with predictions of possible \( n-p \) pairing [36]. The correlations arising from the last proton and neutron can involve a \( T=1 \) or a \( T=0 \) pair. The latter involves the nucleons orbiting together in time reversed orbits and so is assumed to be of lowest energy, indeed most of the \( N=Z \) odd-odd nuclei with \( A < 40 \) have a \( T=0 \) ground state\(^4\). This is highlighted in figure 1.6, which shows the experimentally determined energy difference between the lowest lying \( T=1 \) and \( T=0 \) states in the odd-odd \( N=Z \) nuclei up to mass 80.

It has been experimentally deduced however that low lying \( T=1, I^\pi=0^+ \) states exist in the odd-odd \( N=Z \) systems \(^{67}\)Ga, \(^{66}\)As, \(^{70}\)Br, \(^{74}\)Rb, \(^{75}\)Y, \(^{87}\)Nb and \(^{85}\)Tc [37, 38, 39, 40, 41, 42] from studies of beta decays from these nuclei. The decay from \(^{67}\)Ga was measured [37] to have a half life of 116.34±0.35 ms or a \( \log_{10} T \) value of 3.49±0.1, which is indicative of Fermi superallowed beta decay (see section 1.9). In all the above cases, the decays were fast and the ground states of the daughter nuclei were \( T=1, I^\pi=0^+ \), thus it was inferred that the beta decays came from similar \( T=1, I^\pi=0^+ \) states in the parent.

\( ^4\)with the exception of \(^{34}\)Cl
Figure 1.6: Experimentally determined energy differences between the first $T=1$ state and the first $T=0$ state ($\Delta E_{T=1-T=0}$) for the odd-odd $N=Z$ nuclei up to $A=80$. The $N=Z=39$ isotope $^{78}\text{Y}$ has recently been deduced to have a $T=1$ ground state with a $T=0$ $I^\pi = 5^+$ state at no more than 500 keV [43].

Results from measurements in the $N=Z=32$ system $^{64}\text{Ge}$ [5] have suggested that the observation of E1 transitions, which are forbidden between $T=0$ states in a $T_z=0$ nucleus, is evidence of significant isospin mixing between the $T=0$ and $T=1$ states. The recent study of the high-spin yrast states of the odd-odd system $^{74}\text{Rb}$ by Rudolph et al. [44] provided evidence of a $T=1$ band, built upon the $0^+$ ground state. The evidence being that the energies of the yrast $2^+$ and $4^+$ levels were very close in energy to those of the neighbouring $T_z=1$ isobar $^{74}\text{Kr}$ [45]. This is suggested to be consistent with the presence of an odd-odd pairing gap [12] in $^{74}\text{Rb}$. This $T=1$ band appeared to be crossed by the $T=0$ configuration at about spin $5\hbar$. It is thus proposed [44] that at $N=Z=37$ at the heaviest, the $T=1$ pairs become more
attractive than the $T=0$.

It is therefore the aim of this study to investigate the phenomena of $n-p$ pairing and isospin mixing as a function of mass, and angular momentum. This should act as a test of predictions such as the presence of pair gaps [12], and isospin purity at high angular momentum due to low density of states. In addition, by studying the high spin states of exotic nuclei around the $N=Z=28$ shell gap, one can also investigate the possible sources of angular momentum in the limited valence space of the $fp$ nuclei, whether it be from the breaking of the $^{56}\text{Ni}$ core and/or excitations into the positive parity $g_{\frac{9}{2}}$ shell. It is also a chance to examine the trend in isospin dependence of the ground states in a region where there is a change from $T=0$, to $T=1$. 
Chapter 2

Theoretical Nuclear Models.

In an attempt to understand the physics of the nucleus, a full theoretical explanation would need a solution to a quantum mechanical many-body problem, with not only a set of coupled equations describing the mutual interactions of the nucleons, but also the correlations of all these nucleons as they act on others. Even for a nucleus with only a few nucleons, this quickly becomes an impossibly large dimensional problem.

It is therefore only feasible to adopt oversimplified theories, based on macroscopic phenomena and atomic shell model analogies, and adapt them to account for specific nuclear properties. The hope is then that as the model becomes more sophisticated, we can not only understand and describe these properties, but also predict trends and additional new observables. There are two main genre of nuclear model, the macroscopic property based liquid drop model, and the microscopic property based shell model.

2.1 The Liquid Drop Model.

This model provides a general overview of masses and related stabilities of nuclei, and assumes the nucleus behaves in a gross collective manner, similar to a charged drop of liquid. Intermolecular forces are relatively short ranged, thus just as the energy required to evaporate a given mass of liquid from a drop is independent of
the size of the drop, the energy required to evaporate a given number of nucleons form a nucleus is roughly independent of the size of the nucleus. Also the nucleus has a very low compressibility. With this in mind, it is then relatively straightforward to empirically determine the total mass of a nucleus as the sum of a number of different contributions. One expression for the mass energy of an atom, proposed by \textit{Weizsäcker} \cite{46} is the semi-empirical mass formula, given in equation 2.1:

\[ M(A, Z) = Zm_p + (A - Z)m_p + Zm_e - a_\sigma A + a_\delta A^{3/2} + a_{o-C} \frac{Z^2}{A^3} + a_{A_s} \frac{(A - 2Z)^2}{A} + \delta(A, Z) \] (2.1)

Where \( \delta \) (the pairing term) is generally defined as:

\[
\delta(A, Z) = -a_{o-E} A^{-\frac{1}{3}} \quad \text{for } A \text{ even, } Z \text{ even}
\]
\[
= +a_{o-E} A^{-\frac{1}{3}} \quad \text{for } A \text{ even, } Z \text{ odd}
\]
\[
= 0 \quad \text{for } A \text{ odd}
\]

The values of the parameters \( a_\sigma, a_\delta, a_{o-C}, a_{A_s} \) and \( a_{o-E} \) are determined by comparison with experimental data, and a typical set of values, suggested by Green in 1954 is given in table 2.1.

\textbf{Table 2.1: Typical values of variables in the semi-empirical mass formula.}

<table>
<thead>
<tr>
<th>Atomic mass units, u MeV/c²</th>
</tr>
</thead>
<tbody>
<tr>
<td>( a_\sigma )</td>
</tr>
<tr>
<td>( a_\delta )</td>
</tr>
<tr>
<td>( a_{o-C} )</td>
</tr>
<tr>
<td>( a_{A_s} )</td>
</tr>
<tr>
<td>( a_{o-E} )</td>
</tr>
</tbody>
</table>
The binding energy \(B\) of a nucleus is the difference in mass energy between a nucleus and its constituent nucleons, i.e. [18].

\[
B(A, Z) = (Z m_p + (A - Z) m_n + Z m_e - M(A, Z)) c^2
\]

(2.2)

Where \(m_p, m_n, m_e\) and \(M(A, Z)\) are the mass of the proton, neutron, electron and the atomic mass respectively. More refined versions of the liquid drop model have been developed, for example those proposed by Nix and Swiatecki [47], and by Möller and Nix [48]. An example of one prediction of these types of model of specific interest to this thesis work is the lightest isotope of gallium which is predicted to be particle bound. Audi and Wapstra [49] predict that \(^{60}\)Ga has a binding energy of \(\sim 500\) MeV, a mass of \(\sim 59.95\) Amu, and a \(S_p = 30 \pm 118\) keV. The next lightest isotope \(^{59}\)Ga has a calculated \(S_p = -884 \pm 175\) keV i.e. proton unbound. \(^{61}\)Ga has a calculated \(S_p = 454 \pm 196\) keV, and \(^{62}\)Ga has a measured \(S_p = 2942 \pm 32\) keV. The stable isotopes of gallium are \(^{69}\)Ga and \(^{71}\)Ga.

The liquid drop model enables us to explore the systematic behaviour of nuclear properties such as mass and binding energy, but in order to predict trends in properties such as spin, parity, electromagnetic moments and energy levels, we have to use models which account for the microscopic nature of nuclear matter.

2.2 The Spherical Shell Model.

The fundamental assumption of the shell model is that the motion of a single nucleon is governed by a potential caused by all the other nucleons. In its simplest form, the shell model also assumes that the nucleus is spherical [50]. If we imagine a single neutron moving through the body of the nucleus there is, on average, no net force on it due to the other nucleons since they surround it in a fairly uniform manner. It simply experiences a roughly constant negative (attractive) potential energy due to them. As it moves towards the nuclear surface however, there will be an increasingly attractive inward force leading to a decrease in its potential energy. This reduces
to zero as the neutron moves away from the nucleus due to the short range nuclear force. The same argument applies for a proton, except for a positive potential outside the nucleus due to electrostatic repulsion. This is called the Coulomb barrier, and results in protons being more bound than neutrons in stable nuclei. Both unpaired protons and neutrons experience a centrifugal barrier due to the nuclear rotation which depends on $l(l+1)$, where $l$ is the orbital angular momentum of the final unpaired nucleon.

### 2.2.1 Shell Model Potential.

The most basic shell model potential would be the infinite square well, however this is not realistic as it does not take an infinite amount of energy to remove a nucleon from a nucleus. Another form would be the harmonic oscillator, which is finite, and falls smoothly to zero beyond the mean field radius of the nucleus, $R$. One of the most commonly used forms for this potential is the Wood-Saxon potential (figure 2.1) of the form [28]:

$$V(r) = \frac{-V_0}{1 + e^{\left(\frac{R-a}{a}\right)}}$$  \hspace{1cm} (2.3)

where $R$, $V_0$ and $a$ are the parameters shown in figure 2.1, namely;

- $R = \text{mean radius (given in equation 1.13)}$,
- $V_0 = \text{well depth (adjusted to give proper separation energies, being of order 50 MeV)}$, and
- $a = \text{skin thickness (defined by } 4a \ln 3 \text{ being the distance over which the potential changes from 0.9\% to 0.1\%, usually 0.542 fm) [28]}$.

In such a potential well, as in the atomic analogy, there will be a series of nucleon energy levels as a result of discrete solutions to the Schrödinger equation (the one-dimensional time-independent form of the Schrödinger equation is given in equation 2.4). These levels can accommodate only a certain number of nucleons according to the restrictions imposed by the Pauli exclusion principle.
Figure 2.1: A common form of the shell model potential.

\[ -\frac{\hbar^2}{2M} \frac{d^2\psi}{dx^2} + V\psi = E\psi \] (2.4)

The resulting eigenfunctions depend on three quantum numbers, \( n \), \( l \) and \( m_l \), where \( n \) is the number of levels with a particular orbital angular momentum \( l \), and \( m_l \) is the component of \( l \) on a space fixed axis. The corresponding energy levels have a degeneracy of \( 2(2l + 1) \), so can accommodate \( (2l + 1) \) nucleons with spin 'up' or 'down'.

The internucleon potential is basically attractive, apart from the repulsive core. This means that the lowest energy state of a nucleus is one in which pairs of nucleons experience maximum attraction and therefore contribute maximum negative potential energy, i.e. they are as close together as possible, namely in the same orbital. The Pauli principle however forbids like nucleons to be in identical quantum states, so to be in the same \( j \) subshell, the nucleons must be in different \( m \) states. Opposite \( m \) states will result in the maximum overlap of the nuclear wavefunctions, and thus the lowest energy state. The complete potential used in equation 2.4 must include a potential of the form \( V(r) \) (equation 2.3), and one due to the centrifugal force, of the form shown in equation 2.5.
2.3 The Two-Nucleon System.

When equation 2.4 is solved with the potential shown in equation 2.5 for each separate value of $l$, it can be seen by inspection that the simplest solution will be when $l=0$, since the centrifugal force term vanishes. It is thus noted that the total orbital angular momentum $L$ of a two-nucleon system in the lowest energy state is zero. The total angular momentum $J = L + S$ can then equal 1 or 0, for the cases when the spins are aligned to 1 and anti-aligned to 0 respectively.

If we allow for the isospin quantum number, we can treat all nuclei as identical fermions, that is, they are spin \( \frac{1}{2} \) particles. A result of the Pauli principal is the symmetrisation postulate which states [51]:

"The states of a system containing $N$ identical fermions are necessarily all anti-symmetrical with respect to permutations of the $N$ particles."

The permutations are space, spin and isospin, and this can be written for two particles, 1 and 2:

\[
\psi(1, 2) = -\psi(2, 1) \tag{2.6}
\]

\( \psi \) refers to all co-ordinates, $L, S$ and $T$, ie.:

\[
\psi(1, 2) = \sum_{L,S,T} \alpha_{L,S,T} \psi_{L,S,T}(1, 2) \tag{2.7}
\]

where $\alpha$ is a probability density coefficient. This is written:

\[
\psi(1, 2) = (1 - (-)^{L+S+T})\psi(2, 1) \tag{2.8}
\]
The result of this is that the total wavefunction of the two-nucleon system, \(\psi_L + \psi_S + \psi_T\), must be odd (or antisymmetric) otherwise it vanishes. Therefore, in the lowest energy state of a two-nucleon system, which we have shown to be \(L=0\) (or symmetric), the isospin and spin wavefunctions must add to an antisymmetric total. It was seen in section 1.3 that the bound (ground) state of the deuteron is \(J=1\), so it can then immediately be inferred that the isospin of this state is \(T'=0\).

If we now take as an example a set of two-nucleon systems, then before any Pauli considerations, there are two possible isospin values, \(T=0\) or \(T=1\), and three possible components of \(T_z\), namely; \(T_z=+1\) (two protons), \(T_z=-1\) (two neutrons), and \(T_z=0\) (one proton and one neutron). The \(T_z=0\) case can have the two spins aligned, making \(J=1\) (\(T=0\)), or opposite, making \(J=0\) (\(T=1\)), but the \(T_z=+1\) and \(T_z=-1\) case is forbidden by Pauli to have the same particles in the same energy state with their spins aligned, thus both must have \(J=0\), and therefore \(T=1\). As the nuclear force is essentially charge independent, then all three permutations of \(T=1\) have very similar energies, (except for the Coulomb contribution) while the \(T=0\) case does not. In the deuteron, the \(T=1\) case is unbound by about 60 keV [30] and the \(T=0, J=1\) state is the ground state, with a binding energy of 2.23 MeV.

Values of \(T\) can range from \(|T_z|\) to \(\frac{A}{2}\), thus in the case of the \(A=14\) isobars \(^{14}\text{C}\), \(^{14}\text{N}\) and \(^{14}\text{O}\), (figure 2.2), values of \(T\) from 0 to 7 are allowed or the \(T_z=0\) nucleus, and \(T=1\) to 7 are allowed for for the \(T_z=\pm 1\) nuclei. It can be seen from figure 2.2 that, as in the case of the deuteron, the \(J=1\) state is lower in energy than the \(J=0\) triplet. From equation 2.8, the lower lying \(J=1\) (\(L=0\)) state is assigned a value of \(T=0\). Likewise the \(J=0\) triplet is \(T=1\). If the energies of the \(T=1\) triplet are adjusted to account for the neutron-proton mass difference, and the Coulomb energy, then the resulting energies are within 13 keV of each other, consistent with the concept of the pairing force and charge independence of the nuclear force.
2.4 Magic Numbers.

In the atomic shell model, the shells are filled with electrons in order of increasing energy until they completely fill a closed shell, producing the inert core of a noble gas. These elements have highly stable properties, such as low ionic radius and high ionisation energy. As further electrons are added to shells outside the core, the atomic properties are primarily determined by these valence electrons. The shell model arises from observation of similar phenomena in nuclei, with certain numbers of nucleons being particularly stable, these numbers are called 'magic numbers'.

Nuclei with either numbers of protons or neutrons equal to Z, N = 2, 8, 20, 28, 50, 82, or 126 exhibit certain properties which are analogous to closed shell properties in atoms, including anomalously low masses, high natural abundances and high energy first excited states. The effect can be seen in a plot of separation energy versus increasing N or Z (figure 2.3) [52]. This is similar to the ionisation energy of an atom with increasing mass, i.e. there is a gradual increase with a definite sharp drop off at each of the magic numbers, corresponding to the filling of major nuclear shells.
Figure 2.3: Two proton (top) and two neutron (bottom) separation energies as a function of N and Z respectively. Taken from [18]
In order to reproduce the observed magic numbers, it was proposed in 1948 by Mayer and Jensen [50, 53] that in addition to the static potential, there is also a nuclear spin-orbit potential of the form:

\[ V(r) = V_{so}(r) = f(r) L \cdot s \quad (2.9) \]

This potential, as in the equivalent atomic case, perturbs the nuclear energy levels causing each one to split into two, characterised by the total angular momentum quantum number \( j = l \pm \frac{1}{2} \), as \( j \) is the vector coupled sum of the orbital angular momentum \( l \), and the spin \( s \). Unlike the atomic form, this effect is not due to an electromagnetic interaction, but a nucleon-nucleon interaction, and consequently has a much more dramatic effect.

The total effect of the potentials is shown in figure 2.4, and it can be seen that there are pronounced gaps at the points where the magic numbers of nucleons are filling the orbitals.

It is the assumption of the shell model that in nuclei with, for example, one additional proton or neutron outside the closed shell, the total angular momentum and parity of such a nucleus is simply that of the additional valence nucleon due to the Pauli blocking effect. eg. \(^{17}\)O has an additional neutron in the \( 1d_{\frac{5}{2}} \) state for which \( j = \frac{5}{2}, \ l = 2, \ \pi = (-1)^l = +1 \), thus we predict a \( J^\pi \) ground state of \( \frac{3}{2}^+ \), in agreement with experiment [54].
Figure 2.4: Theoretical nuclear level structure, labelled $nl_j$, including spin-orbit perturbations and showing energy gaps at the magic numbers.
2.5 Two-State Mixing.

If we have two levels with energies $E_1$ and $E_2$ and wavefunctions $\psi_1$ and $\psi_2$, then there will be an interaction, and hence a nuclear matrix element between them (see section 1.9 and equation 1.19). This perturbs the levels by an amount dependent on the spacing $\Delta E$ and the matrix element $|M_{fi}|$. The final levels then have energies $E_{1f}$ and $E_{2f}$, and a spacing $\Delta E_f$ given by [13];

$$|\Delta E_f| = \frac{\Delta E}{2} \left( \sqrt{1 + \frac{4|M_{fi}|^2}{\Delta E^2}} - 1 \right)$$

The final levels now have admixed wavefunctions $\psi_{1f}$ and $\psi_{2f}$, given by [13];

$$\psi_{1f} = \alpha \psi_1 + \beta \psi_2$$

$$\psi_{2f} = -\beta \psi_1 + \alpha \psi_2$$

where $\alpha^2 + \beta^2 = 1$, and the smaller amplitude $\beta$ is given by [13];

$$\beta = \frac{1}{\sqrt{1 + \left( \frac{\Delta E}{2|M_{fi}|} + \sqrt{1 + \frac{\Delta E^2}{4|M_{fi}|^2}} \right)^2}}$$

The consequences of this are numerous, and can easily be calculated using this method without the need for a complex diagonalisation. Some results which arise from this theory are that the final separation between two isolated states can never be less than twice the mixing matrix element. This is seen in for example the deformed Woods-Saxon plot where two admixed levels do not cross, but instead are repelled at an inflection point (figure 2.5-a), corresponding to when the mixed wavefunctions contain equal amounts of the unperturbed states. Another important effect can be seen from a series of simple mixing calculations on a set of many mixed levels, where the lowest state is pushed down in energy by an amount dependent on the initial spacing and the mixing matrix elements, while the rest of the levels are raised in
energy. Figure 2.5-b shows the case of N initial states which are degenerate and have equal matrix elements. Here the lowest state is lowered by an amount \((N-1)\) times the matrix element \(M\), while the rest are all raised by one unit in \(M\) [13]. There is also an important consequence of two state mixing related to transition rates which can allow an otherwise forbidden transition to occur. If a transition between two states \(A\) and \(B\) is forbidden, while it is not between \(A\) and a third state \(C\), then if there is mixing between \(B\) and \(C\), the forbidden transition can occur.

![Diagram](image)

Figure 2.5: Illustration of the effects of two-state mixing, (a) indicating non crossing of admixed levels with change in deformation, (b) lowering of energy of lowest state in a multistate mixing example. The initial states are degenerate and all have equal mixing matrix elements \(M\). Taken from [13].
Chapter 3

Experimental Techniques.

3.1 Heavy Ion Fusion Evaporation.

The preferred method of experimental production of neutron deficient nuclei in high spin states in the mass 50–100 region is heavy ion fusion evaporation. A beam of ionised nuclei of a specific isotope is produced by an ion source, and accelerated onto a metallic foil target of another isotope (figure 3.1). If two nuclei collide with a sufficiently small impact parameter and an energy large enough to overcome their mutual Coulomb repulsion, then their nuclear potentials overlap, and they can fuse to form a single compound nucleus. If this compound system lives as long as $\sim 10^{-20}$ seconds then it can equilibrate with respect to its energy, angular momentum and shape. The nucleus then loses all memory of the components involved in its formation process, and the kinetic energy of the collision and relative angular momentum of the projectile and target are converted into excitation energy and angular momentum. The energy is well defined by the reaction, and is given by [20, 55];

$$E^* = \frac{M_T}{M_B + M_T} E_{\text{lab}} + Q$$

(3.1)

where $M_B = $ mass of the beam nucleus, $M_T = $ mass of the target nucleus, $Q = [(M_B + M_T - M_{\text{en}})] c^2$ is the difference in binding energy of the compound nucleus and the nuclei in the entrance channel.
Figure 3.1: Schematic of a heavy ion fusion reaction. Taken from [21].

This process usually produces compound nuclei which are highly excited, and with a large amount of angular momentum. The maximum angular momentum of the compound nucleus following a reaction is given by equation 3.2 [28].

\[
I_{\text{max}}^2 = \left( \frac{2\mu R^2}{\hbar^2} \right) (E_{\text{CM}} - V_c)
\]

where, \( \mu = \frac{A_1 A_2}{A_1 + A_2} \) is the reduced mass, \( V_c \) is the Coulomb barrier \( \simeq 1.44Z_1Z_2 \) [56], and \( R \) = distance of closest approach, usually given by \( 1.36(A_1^{\frac{1}{3}} + A_2^{\frac{1}{3}}) + 0.5 \text{ fm} \) [18].

On formation, the energy of the compound nucleus is above the particle emission threshold, and consequently, the decay from hot state to cold residual nucleus is initially dominated by statistical particle emission, each nucleon taking the order of 5 MeV and each \( \alpha \)-particle taking the order of 10 MeV, but each taking only about \( 1\hbar \) of angular momentum. Linear momentum must be conserved as the evaporated nucleons and \( \alpha \)-particles leave the compound nucleus, thus the residual nucleus must recoil, shifting it from its original trajectory. These particles are evaporated approximately isotropically in the centre of mass frame of reference, but the nucleus is moving in the lab frame in the direction of the beam, resulting in a forward focussed cone of recoils in the lab frame called the recoil cone. A typical velocity for the recoiling nucleus \( v_{\text{cm}} \) from a heavy ion fusion reaction is 1–5% the speed of light, and is related to the reaction kinematics as shown in equation 3.3 [55]:

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At high energy, the density of nuclear levels (level density) can be up to $10^{12}$ MeV$^{-1}$ with level widths ($\Gamma$) typically of the order eV. This results in a large degree of overlap between levels, effectively forming a continuum. After the evaporation of a number of particles, the recoiling nucleus has cooled in terms of excitation energy, but retains the high internal angular momentum. This is ideal for the study of the high spin states of the reaction product nuclei. Below the particle separation energies, statistical $\gamma$-ray emission dominates, until, at lower excitation energy, the nature of the states changes from continuum to discrete, well defined quantum mechanical levels. The nucleus completes its decay to the ground state by $\gamma$-ray emission from these discrete states. The states which are preferentially populated in this type of decay are yrast or near yrast (for example, see reference [57]), meaning of lowest excitation energy for a given spin. Figure 3.2 depicts the sequence of the decay of the compound nucleus to the ground state, in terms of energy and spin.

For near stable nuclei, neutron evaporation is favoured over charged particle evaporation due to the Coulomb barrier in the nucleus. For very neutron deficient systems, the separation energy of the protons is increasingly less than that of the neutrons, making it energetically favourable to evaporate protons or $\alpha$-particles. Figure 3.3 shows the region of the Segré chart of interest to this study. As we increase in mass along the N=Z line, the nuclei become increasingly neutron deficient compared to the beta stable isotopes. In order to experimentally populate these neutron deficient nuclei, it is often necessary to do so via an evaporation channel which contains neutrons (eg. 2pn), but the production cross-sections for such nuclei are lower than for isobars created with all charged particles (e.g. 3p). It can be seen however in the insert of figure 3.3 that with the combinations of stable beams and targets currently experimentally available, a compound nucleus is often populated which is on the neutron deficient side of the beta stable isotopes, i.e. nearer the N=Z line.

$$v_{em} = \frac{M_B}{(M_B + M_T) \sqrt{\frac{2E_B}{M_B}}}$$ (3.3)
Figure 3.2: Schematic of the excitation energy – spin plane, showing the cooling of a nucleus formed in a heavy ion fusion evaporation reaction.

Predictions of the population probabilities for different residual nuclei from a heavy ion fusion reaction can be made by evaporation residue codes such as PACE [58].
Figure 3.3: Section of the Segré chart of interest to this study. Insert shows that addition of target and beam nuclei produces a compound nucleus on the neutron deficient side of stable nuclei.
3.2 Accelerators.

In order to overcome the Coulomb barrier, ions must be accelerated into a beam of a certain energy. There are various methods of achieving this, but the one of relevance to this study is the tandem Van de Graaff accelerator [59, 60, 61]. A negatively charged source of ions is needed, along with a method of extracting these ions from the source, to be injected into the accelerator [62]. The negative ions can be produced in a number of ways. One type of source commonly used in this type of study is a sputter source [63, 64, 65]. This method has a yield comparable with other types of negative ion source, but unlike other sources, ions can be directly generated from a solid. A spray or vapour of cesium is used to ionise a solid source of the ion to be extracted (called a ‘cone’ or a ‘pill’). Cesium is used as it has loosely bound electrons. These ions are then extracted with an electric field.

The principal of a Van de Graaff generator is that charge is continuously transferred to a terminal via a moving insulating belt. The terminal is in electrical contact with an outer shell which collects all the charge deposited by the belt, according to the laws of electrostatics. The charge which can be deposited is limited only by the insulation properties of the surrounding medium. Sulphur hexafluoride (SF₆) is used as the insulating gas because it is highly resistant to electrical breakdown. The terminal voltage of a tandem accelerator is generally around 10 million Volts, and this is used to attract the negatively charged ions with a kinetic energy equal to the charge on an electron $e$ times the voltage of the terminal, $V$. As the ions reach the terminal, they pass through a thin carbon stripper foil, which removes $n+1$ electrons from the ion, resulting in a positive charge of $ne$. This is then accelerated away from the terminal with a kinetic energy of $(n+1)eV$.

It is then possible to further boost the energy of the beam with a linear accelerator (linac) although this was not required for the experiments in this study.
3.2.1 ATLAS Facility.

The accelerator facility at the Argonne National Laboratory, Chicago, Illinois, US. is called ATLAS (Argonne Tandem Linear Accelerator System) [66] and a schematic is shown in figure 3.4. It consists of two separate ion sources coupled to a 20MV superconducting linac. The ion sources available are either a positive-ion injector and a 12MV injector linac [67, 68], or a SNICS-II negative ion inverted sputter source [62] and a 9MV tandem Van de Graaff accelerator. The latter is the one of interest to this study. The beam then passes through an ion beam buncher [69] which produces a pulsed beam incident onto a target.

![Figure 3.4: Schematic of the layout of the ATLAS facility. Taken from [70].](image)

3.2.2 Niels Bohr Institute Tandem Accelerator Laboratory.

The tandem accelerator at the Niels Bohr Institute is a 9MV Van de Graaff [71] coupled to two linac booster modules [72] (figure 3.5). As at Argonne, the linac modules were not used in the experiments performed for this study.
3.3 Gamma-Ray Angular Distributions.

Gamma-rays are emitted from a nucleus with preferred direction depending on the spin of the nucleus, and the direction of \( J \). In a heavy ion fusion reaction, a preferential direction of \( J \) is achieved from the compound nuclear reaction where, if the spin of the projectile and target are zero, \( J_{CN} = L = \ell \times P \). This produces a nucleus which is perfectly aligned perpendicular to the beam axis. There may then be some loss of this alignment due to particle emission, and by subsequent \( \gamma \)-ray emission. The general formula for the angular distribution of \( \gamma \)-ray intensity as a function of angle \( \theta \) with respect to the beam direction is given by equation 3.4 [73].

\[
W(\theta) = \sum_k A_k P_k(\cos\theta) \tag{3.4}
\]

Where \( k = 2\ell \), and \( \ell \) is the highest order multipole of the radiation. \( A_k \) is called
the angular distribution coefficient which depends on the initial and final spins, the angular momentum taken by the $\gamma$-ray, and if there is mixing of different angular momenta, then $A_k$ depends on the mixing ratio, $\delta$. $P_k$ is a standard Legendre polynomial. Both $A_k$ and $P_k$ are tabulated, eg [74].

For pure dipole radiation,

$$W(\theta) = A_0\{1 + A_2 P_2(\cos\theta)\}$$

(3.5)

where $P_2(\cos\theta) = \frac{1}{2}(3\cos^2\theta - 1)$ and $A_0$ is the 'true' intensity. For quadrupole radiation,

$$W(\theta) = A_0\{1 + A_2 P_2(\cos\theta) + A_4 P_4(\cos\theta)\}$$

(3.6)

where $P_4(\cos\theta) = \frac{1}{3}(35\cos^4\theta - 30\cos^2\theta + 3)$.

### 3.3.1 Anisotropy.

By experimentally measuring $\gamma$-ray intensities as a function of angle using detectors positioned at different angles around a target, it is possible to determine the values of the coefficients $A_2$ and $A_4$ by fitting the distribution of $\gamma$-rays to equation 3.4, and hence determine the multipolarity of the detected radiation. This technique depends on having many detectors at different angles. However, if only a few angles are accounted for, it may still be possible to distinguish between dipole and quadrupole radiation. A value of anisotropy is arbitrarily defined in equation 3.7 [75].

$$A = 2\left(\frac{W(\theta_1) - W(\theta_2)}{W(\theta_1) + W(\theta_2)}\right)$$

(3.7)

$W(\theta)$ is the intensity of a $\gamma$-ray detected at angle $\theta$, and this can be experimentally measured, and compared either with theoretical values for dipoles and quadrupoles calculated using equations 3.5 and 3.6, or with values measured for transitions of known multipolarity.
3.3.2 DCO.

With high efficiency $\gamma$-ray detector arrays [76, 77], it is possible to perform measurements of Directional Correlation from Orientated states (DCO) on pairs of $\gamma$-rays detected in coincidence with one another [73, 78, 79]. In this method, the ratio of intensities is taken for a $\gamma$-ray detected at two or more different angles around a target, provided the $\gamma$-ray is detected in coincidence with another $\gamma$-ray of known multipolarity (usually E2) [73].

$$R_{DCO} = \frac{I_n(\theta_1) \text{ gated by } \gamma_2(\theta_2)}{I_n(\theta_2) \text{ gated by } \gamma_2(\theta_1)}$$  

By requiring a coincidence, the spectra are generally much cleaner than in singles. The large amount of angular momentum transferred to the compound nucleus in a heavy ion fusion reaction (equation 3.2) is always aligned in a plane perpendicular to the beam direction, i.e. in an $m = 0$ substate relative to the beam direction. However, this alignment is lost with subsequent $\gamma$-ray emission as each transition changes the $m$-state by an amount up to the multipolarity until the nucleus is usually almost completely randomly orientated in the ground state. This method is useful as it introduces an angular correlation effect which enables differentiation of different multipoles even at the end of a $\gamma$-ray cascade. It is also possible with this method to determine the mixing ratio if two multipoles are present in one $\gamma$ transition. As with the anisotropy method discussed previously, differentiating between multipoles can be by performed theoretically [73] or more conveniently, by comparison with known transitions. For example in the AYEBALL array (section 3.4.3) the geometry is such that $R_{DCO} \simeq 1.0$ for a stretched E2 gated by an E2, and $R_{DCO} \simeq 0.6$ for a pure dipole gated by an E2, with mixed E2/M1 transitions having a ratio anywhere between 0.25 and 1.25 (see section 4.6).
3.4 Detectors.

3.4.1 Gamma-Ray Interactions.

Electromagnetic radiation can interact with matter in a number of ways, but for the purposes of a discussion on γ-ray detection, only three are of interest, namely photoelectric absorption, Compton scattering and pair production [80].

- The photoelectric effect involves an incident γ-ray being absorbed by an atom which ejects a photoelectron, usually from the K-shell, with an energy equal to the incident γ-ray energy minus the binding energy of the emitted electron. The vacancy left by the electron is filled by atomic rearrangement, producing either an Auger electron or a characteristic X-ray which is subsequently photoelectrically absorbed by a loosely bound electron. This type of interaction results in full energy deposition in the detector, and thus a photopeak in the energy spectrum.

- The Compton effect involves a γ-ray scattering from an electron, resulting in a recoiling electron and a loss of energy of the γ-ray. The energy of the γ-ray after the interaction $E'_\gamma$ is dependent on the angle of scatter and is given by equation 3.9;

$$E'_\gamma = \frac{E_\gamma}{1 + \frac{E_\gamma}{m_0c^2}(1 - \cos \theta)} \quad (3.9)$$

where $m_0c^2$ is the rest mass energy of the electron (511 keV). The kinetic energy of the recoiling electron is thus the difference in energy of the incident and scattered γ-ray, given by equation 3.10.

$$E_{e^-} = \frac{E_\gamma^2}{\frac{m_0c^2}{1 - \cos \theta} + E_\gamma} \quad (3.10)$$

If there are no further interactions then the γ-ray scatters out of the detector and the energy deposited ranges from zero to a maximum when $\theta=180^\circ$. The result is incomplete energy deposition and a spectrum with a continuum of energy called the Compton background, up to the maximum value, called the Compton edge.
The probability for a Compton scatter to occur at an angle $\theta$ can be calculated from the *Klein-Nishina formula* for the differential cross section per electron, given in equation 3.11 [18]:

$$\frac{d\sigma}{d\Omega} = r_0^2 \left[ \frac{1}{1 + \alpha(1 - \cos\theta)} \right]^3 \left[ \frac{1 + \cos\theta}{2} \right] \left[ 1 + \frac{\alpha^2(1 - \cos\theta)^2}{(1 + \cos^2\theta)[1 + \alpha(1 - \cos\theta)]} \right]$$ (3.11)

Where $\alpha$ is the photon energy in units of the electron rest mass energy ($\alpha = E_\gamma/mc^2$), and $r_0$ is a parameter called the *classical electron radius*, which although it has nothing to do with the actual radius of an electron, is a convenient parameter, equal to $e^2/4\pi\varepsilon_0mc^2=2.818 \text{ fm}$ [18].

- **Pair Production** is possible only when the $\gamma$-ray energy is at least twice the electron rest mass energy (i.e. $E_\gamma \geq 1.022 \text{ MeV}$), so that an electron-positron pair can be created at the point of disappearance of the $\gamma$-ray. This must be in the presence of an atom in order that momentum be conserved. The electron and positron then quickly lose their kinetic energy in the medium, and the positron consequently annihilates with an electron from the medium, producing two back to back $\gamma$-rays with energy equal to the rest mass energy of the two particles (511 keV each).

Figure 3.6 shows the linear attenuation coefficients versus $\gamma$-ray energy for the three types of interaction for various detector materials, indicating which processes will be important at different energies.

### 3.4.2 Gamma-Ray Detection.

In discrete $\gamma$-ray spectroscopy measurements, the $\gamma$-rays are usually detected by germanium semiconductor crystals. As the radiation interacts with the material, electrons are excited across the band gap to the conduction band, leaving behind a hole in the valence band. The electron-hole pair is then collected by an electric field created by an applied voltage, and the charge converted to an output voltage by a preamplifier. The number of electron-hole pairs, and so the output voltage, is proportional to the energy of the incident $\gamma$-ray. The band gap in germanium is
Figure 3.6: Linear attenuation with respect to $\gamma$-ray energy for (a) germanium (b) sodium iodide (c) bismuth germanate and (d) barium fluoride. Taken from [81].

$\sim 0.9$eV [82] and the $W$-value is $\sim 3$eV [82], which is small enough to produce many electron-hole pairs from the incident radiation, producing good energy resolution, but also small enough that thermal excitations can cause unwanted promotions, and hence noise. To reduce this effect, Ge crystals are operated at liquid nitrogen temperature (77 K). Typical energy resolution (FWHM/$E_\gamma$) of a germanium detector is a few tenths of a percent at 1 MeV. Crystals of hyper-pure germanium (HPGe) can be grown up to 70cm in diameter, and these are used in coaxial configurations in the modern detector arrays, each detector having a photopeak efficiency of up

\footnote{the average energy per electron hole pair production after trapping etc.}
to 80% compared to a 3" x 3" crystal of sodium iodide (NaI(Tl)) at 10cm for a 1.33 MeV γ-ray (the standard for detector relative efficiency). As can be seen in figure 3.6 for the energy range 100 keV ≤ Eγ ≤ 2 MeV, there is an appreciable chance of the interaction being a Compton scatter, rather than the photoelectric effect, resulting in incomplete energy deposition, and contributing to an unwanted Compton background in the spectrum. To reduce this effect, germanium detectors often have an outer shell of bismuth germanate (Bi₄Ge₃O₁₂), or BGO, scintillator material, coupled to photomultiplier tubes, which detect some of the events which Compton scatter out of the germanium crystal [81, 83, 84, 85]. This can be used to suppress these events from an energy spectrum, leaving only the photoelectric events plus any Compton scattered events which were missed by the suppression due to the finite detection efficiency, for example, γ-rays which scatter out of the entrance window, or through the cold finger at the back of the crystal.

Clovers and Clusters.

The current generation of germanium detectors being developed and used in the latest high efficiency γ-ray arrays, are called clovers and clusters [76, 86, 87, 88, 89]. One clover detector contains four 20% efficient high purity germanium crystals encapsulated in a single Compton suppression casing. The clusters are seven large 80% efficient crystals in one casing. The advantages of this are that there is less volume occupied by lead collimators, and the BGO shielding, so the detectors can be packed closer together to cover more of the full 4π sphere around the target. Also, γ-rays which Compton scatter from one crystal into another, can have the energies of the two individual signals summed to reproduce the full original energy, thus increasing the efficiency of the detector, while maintaining the granularity. Good granularity comes from more smaller crystals rather than one large crystal, and results in better charge collection, timing properties and position sensitivity of the γ-ray, which is important in Doppler correction techniques.
Doppler Effects.

If the geometry of the germanium detectors relative to the recoiling nucleus is known, it can be utilised in the analysis of the data in the use of Doppler correction techniques. The residual nucleus moves through the target with typical velocities between 1-5% speed of light. Consequently, when a $\gamma$-ray or particle is emitted, there is a Doppler shift associated with the detected energy, which is dependent on the angle at which it was emitted according to equation 3.12 [20];

$$E_s(\theta) = E_0 \frac{\sqrt{1-\beta}}{1-\beta\cos\theta} \approx E_0 (1 + \beta \cos\theta)$$

(3.12)

where, $\beta = \frac{v}{c}$, $v$ = velocity of the recoil, $c$ = speed of light, $E_s$ = the observed (Doppler shifted) energy, $E_0$ = the true (unshifted) energy, and $\theta$ = angle between the detector and the recoil direction.

From equation 3.12, the shift in energy is greatest at angles close to 0° or 180°, i.e. forward or backward angles. By knowing the angle at which the $\gamma$-ray is emitted, the detected energy can be adjusted in software analysis back to the true value. The granularity of the clusters and clovers is beneficial because the size of each detector crystal is kept small, so the angle $\theta$ is known more accurately, and the true energy can be more accurately calculated.

Another manifestation of Doppler effects is in 'Doppler broadening' of the peaks (equation 3.13). This effect is an artifact of the finite size of the angle subtended by the crystal face to the target ($\Delta\theta$). A $\gamma$-ray interacting at one edge of the crystal will be Doppler shifted by a slightly different amount than an event in the other side. Consequently, the two events will deposit different amounts of energy, and so there will be a spread in the corresponding peak. It cannot be accounted for in software, but can be minimised by having as high a detector granularity as possible. Approximating equation 3.12 to $E_s(\theta) \approx E_0 (1 + \beta \cos\theta)$, and differentiating, we obtain;

$$\Delta E_s = E_0 \cos\theta \Delta\beta - E_0 \beta \sin\theta \Delta\theta$$

(3.13)
The $\Delta \beta$ arises from the spread in velocity of the residual nucleus which has two main contributions. The loss in the target, which will be different if the reaction occurs at the front of the target than if it occurs at the back, and the conservation of momentum, which results in a different velocity when nucleons are evaporated in a forward direction from the compound nucleus, than if they are evaporated backwards. For studies using thin targets in the mass 130 region and higher, it is often sufficient to approximate this expression by ignoring the term in $\Delta \beta$ as the difference in energy loss in the target is small, and the mass of the evaporated nucleons and $\alpha$-particles are negligible compared with the mass of the residual nucleus. For these cases, only the $\Delta \theta$ is important, and the $\sin \theta$ dependence of the expression results in a larger Doppler broadening effect observed in detectors positioned closer to 90°. This is the reason for the positioning of smaller granulated detectors such as clovers near 90° in many large detector arrays [76], with the larger detectors such as clusters placed nearer 0°. However, as is shown in section 5.3, this assumption is not valid for $\alpha$-evaporation channels in the mass 60 region, where the first term in equation 3.13 can have a greater effect than the second term.

3.4.3 The AYEBALL Array.

Figure 3.7 shows a photograph of the AYEBALL array, (Argonne-Yale-European Ball) [90] which was constructed at the Argonne National Laboratory. AYEBALL was a $\gamma$-ray detector array consisting of 18 high purity germanium detectors, in BGO Compton suppression shields, mounted in four annular rings around the beam direction. The array comprised both 25% efficient TESSA type detectors [83, 84] and also eight 80% efficient EUROGAM type detectors [85], and a GAMMASPHERE prototype detector [91]. Table 3.1 summarises the properties of the germanium detectors used, and their position in the AYEBALL array. The array had a measured germanium photopeak efficiency of 1.1% at 1173 keV.
Figure 3.7: A photograph of the AYEBALL array, showing the liquid nitrogen filling pipes and dewars for the Eurogam detectors on the left and the smaller TESSA type detectors on the right. The beam comes from the left down the vacuum pipe. The first electric dipole of the FMA is visible on the right.
Table 3.1: Positions of the various germanium detectors in the AYEBALL array.

<table>
<thead>
<tr>
<th>Ring</th>
<th>Type</th>
<th>θ</th>
<th>φ</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>EUROGAM</td>
<td>157.6</td>
<td>36</td>
</tr>
<tr>
<td>1</td>
<td>EUROGAM</td>
<td>157.6</td>
<td>108</td>
</tr>
<tr>
<td>1</td>
<td>EUROGAM</td>
<td>157.6</td>
<td>180</td>
</tr>
<tr>
<td>1</td>
<td>EUROGAM</td>
<td>157.6</td>
<td>252</td>
</tr>
<tr>
<td>1</td>
<td>EUROGAM</td>
<td>157.6</td>
<td>324</td>
</tr>
<tr>
<td>2</td>
<td>GAMMASPHERE</td>
<td>133.6</td>
<td>54</td>
</tr>
<tr>
<td>2</td>
<td>EUROGAM</td>
<td>133.6</td>
<td>90</td>
</tr>
<tr>
<td>2</td>
<td>EUROGAM</td>
<td>133.6</td>
<td>270</td>
</tr>
<tr>
<td>2</td>
<td>TESSA</td>
<td>133.6</td>
<td>306</td>
</tr>
<tr>
<td>3</td>
<td>YALE</td>
<td>101</td>
<td>36</td>
</tr>
<tr>
<td>3</td>
<td>TESSA</td>
<td>101</td>
<td>108</td>
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<tr>
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<td>YALE</td>
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<td>216</td>
</tr>
<tr>
<td>4</td>
<td>TESSA</td>
<td>79</td>
<td>288</td>
</tr>
</tbody>
</table>

3.4.4 The PEX Array.

The Pre-EUROBALL Experiment 'PEX' [92] at the Niels Bohr Institute (NBI), Risø, near Copenhagen, Denmark comprised four cluster detectors which were being tested before use in the European collaborative γ-ray detector array, 'EUROBALL' [77] at Lengaro, Italy. Figure 3.8 shows a photograph of the array. The detectors were positioned around the target, two with the central crystal at 145.8° to the recoil direction and two with the central crystal at 104.6° to the recoil direction.
3.5 Gamma-Gamma Matrices.

The relatively high $\gamma$-ray detection efficiency of arrays such as AYEBALL and PEX can be utilised to produce two dimensional spectra or 'matrices' of correlated $\gamma$-ray energies. In each fusion-evaporation event there are typically $\sim 20\rightarrow 30$ $\gamma$-rays emitted (depending on the angular momentum transferred), but due to limited $\gamma$-ray detection efficiency, only a few of these are observed. Every time more than one is
detected from each cascade, the energies are written to the two axes of a `γ-γ' matrix. It is then possible to take a thin slice of this matrix (called a 'gate') on one axis, centred around the channel corresponding to a certain γ-ray energy, and project this slice onto the other axis. Thus, by gating on a particular γ-ray, an energy spectrum can be obtained of all the other transitions which come in co-incidence with that γ-ray.

3.6 Ancillary Detectors.

In order to study the low cross-section evaporation channels of interest in this study (\(\sigma \leq 100\mu b\)), some method of channel selection must be employed in coincidence with the γ-ray detection. For very neutron deficient nuclei, this usually takes one of two forms: (a) an array of charged particle detectors to tag the evaporated particles from the compound system [93, 94, 95, 96, 97], or (b) a recoil mass separator to tag the recoiling nuclei by their mass to charge state ratio, \(A/Q\) [70, 98, 99, 100, 101, 102, 103].

3.6.1 Charged Particle Detectors.

The channel selection for the PEX experiment was obtained using the differential energy loss in a silicon charged particle detector 'Si-ball' [104] which completely surrounded the target except for the inlet and outlet holes for the beam. The ball was made from 12 hexagonal silicon wafers, each further segmented to give 31 particle detectors with a near full 4\(\pi\) coverage, (figure 3.9). The dimensions of the ball are such that it can surround a target inside an array of closely packed γ-ray detectors, it is therefore only approximately 42 mm in diameter, with the silicon wafers approximately 24 mm across. The segmentation was designed to reduce the probability of multiple hits in one detector. The particles are essentially evaporated isotropically in the centre of mass frame, but as explained in section 3.1, this produces a forward focussed cone in the laboratory frame. There is consequently an increased
probability of detection at forward angles, requiring an increased granularity as shown in figure 3.9. Even though the α-particles have higher energy on average than the protons, they are more forward focused in the laboratory frame due to their lower velocity. The particle identification technique is demonstrated in section 5.2 with the specific example of the $^{28}\text{Si} + ^{40}\text{Ca}$ reaction.

Figure 3.9: Net diagram of the PEX silicon ball charged particle detector, (a) beam inlet side, (b) beam outlet side showing segmentation of silicon wafers for increased granularity. (c) shows geometrical arrangement of complete ball, (d) indicates annular rings of segments around target. Taken from [105].
The silicon was 170\( \mu \text{m} \) thick [104], chosen to give maximum differentiation between protons and \( \alpha \)-particles. A plot of PACE predictions of evaporation probability for protons, neutrons and \( \alpha \)-particles with increasing energy for the reaction \( ^{28}\text{Si} + ^{40}\text{Ca} \) at 88 MeV is shown in figure 3.10.

![Figure 3.10: Predicted relative evaporation probabilities for protons, neutrons, \( \alpha \)-particles and \( \gamma \)-rays from the reaction \( ^{28}\text{Si} \) on \( ^{40}\text{Ca} \) at 88 MeV as a function of energy, according to PACE [58] calculations. The particle energies are relative to the centre of mass frame of the reaction.](image)

Predictions of the range of protons and \( \alpha \)-particles in silicon as a function of energy can be made using a code called 'DE-DX' [106] based on the Bethe-Bloch equation (equation 3.14) [18];
\[
\frac{dE}{dx} = \left( \frac{e^2}{4\pi\varepsilon_0} \right)^2 \frac{4\pi Z^2 N_0 Z \rho}{m c^2 \beta^2 A} \left[ \ln\left( \frac{2mc^2\beta^2}{I} \right) - \ln(1 - \beta^2) - \beta^2 \right]
\]

(3.14)

where \( \beta = v/c \) is the velocity of the particle, \( e \) is the electric charge, \( Z, A \) and \( \rho \) are the atomic number, atomic weight and density of the stopping material, \( z \) is the proton number of the projectile, \( N_0 \) is Avogadro's number, and \( m \) is the electron mass. \( I \) represents the mean excitation energy of the atomic electrons. \( I \) can be computed, but is generally regarded as an empirical constant, with a value in \( eV \) of the order of 10\( Z \).

Plots of predicted energy loss in the silicon versus particle energy are shown for protons and \( \alpha \)-particles in figure 3.11. As a particle energy increases, it deposits more energy in the silicon, until it becomes energetic enough to punch through and escape. The code predicts that a proton of energy greater than 4.3 MeV will penetrate 170\( \mu m \) of silicon. Above this, the proton will deposit less as its energy increases. There is consequently a sharp cut off point in the plots. The code also predicts that an \( \alpha \)-particle of energy 17 MeV or less will be stopped in the same material. If we look at figure 3.10 we see that this corresponds to most of the predicted range of energies of protons penetrating the silicon, while most of the range of \( \alpha \)-particles do not. The amount of energy deposited in the detector will therefore differ for protons and \( \alpha \)-particles, thus they will give characteristic signals [93, 94, 95, 96, 97], and so the number of evaporated protons and \( \alpha \)-particles associated with each detected \( \gamma \)-ray could be determined. This method of channel selection has the disadvantage that for very weakly populated channels, target contaminants can dominate the spectra, for example, even very pure \( ^{40} \)Ca targets will contain small amounts of the natural isotope \( ^{44} \)Ca.
Figure 3.11: Predicted energy loss in a 170 µm thick silicon detector for protons and α-particles, according to PACE calculations. A sharp cut off point occurs when the particle becomes energetic enough to pass through the silicon. Note the different scales for energy loss. Taken from [107].

3.6.2 Recoil Separators.

Recoil separators are spectrometers which use electric and magnetic fields and the time of flight of recoils through the separator to distinguish reaction products from beam particles, and separate them according to charge, mass and possibly total kinetic energy. (see for example [5, 70, 98, 99, 100, 101, 102, 103]). This technique can in principle give cleaner channel identification than using charged particle detectors, but is limited by the transmission efficiency of the separator for recoiling nuclei as they are often populated in a Gaussian distribution of charge states. This limitation is also dependent on the size of the recoil cone and so is reaction dependent. It will also differ for different evaporation channels as the recoil cone is larger if an α-particle is emitted rather than a proton or neutron.
The Argonne Fragment Mass Analyzer.

The Argonne National Laboratory Fragment Mass Analyzer (FMA) is an 8 metre long recoil mass spectrometer [70]. It is used to separate reaction products from primary beam particles which did not fuse with target nuclei, and disperse the residual recoils according to their Mass/Charge, \( \left( \frac{A}{Q} \right) \). Magnetic and electric fields guide the desired particles and focus them onto a position sensitive Parallel Plate Avalanche Counter (PPAC) at the focal plane. The FMA has a relatively large physical acceptance for recoils, namely a solid angle of 8 msr, relating to an \( A/Q \) acceptance of \( \pm 7\% \) around the central mass, and an energy acceptance of \( \pm 20\% \) around the central energy [70]. It can subsequently focus 2 charge states onto the PPAC. Even so, the efficiency for detection of a recoil at the focal plane for this type of reaction is of order \( \sim 5\% \).

Note that for certain masses and charge states, different \( A \) and \( Q \) values can result in the same ratio, (figure 3.12.) For example in the reaction \(^{24}\text{Mg} + \ ^{40}\text{Ca}, \ 94_{11} \approx 58_{18} \approx 5.8, \) thus \( \gamma \)-rays from \(^{64}\text{Zn} \) [108, 109, 110] (\( \alpha \)\( 2\)\( n \) channel from compound reaction \(^{24}\text{Mg} + \ ^{44}\text{Ca} \)) and \(^{58}\text{Ni} \) (\( \alpha \)\( 2\)\( p \) channel) [111, 112] are both present in the \( \left( \frac{A}{Q} \right) \) software gate and \( \gamma \)-ray spectrum for that region (see figure 4.7). However, the fact that the focal plane of the FMA was large enough to allow two charge states meant that generally such anomalies could be accounted for and removed by subtracting a normalised portion of one mass, obtained by gating on the neighbouring charge state. These anomalies can also be resolved by separating the recoils by the number of protons (explained now in section 3.6.3).

3.6.3 Z-Separation.

Information on the proton number (Z) of the recoils can be obtained using either (a) a split anode ionisation chamber [98] and/or (b) an array of neutron detectors [80, 113, 114, 115]. These are now discussed separately.

\(^{2}\)present as a natural contaminant in the \(^{40}\text{Ca} \) target
Figure 3.12: Schematic of the function of the FMA, with recoils being deflected by electric and magnetic fields. Different charge states are deflected by different amounts, leading to charge state anomalies.

**Split Anode Ionisation Chamber.**

There was an ionisation chamber placed beyond the focal plane of the FMA at Argonne (figure 3.14) which was based on the design of one previously built for the Daresbury recoil separator [98]. However, the size of the focal plane required a much larger entrance window, 150 mm wide and 50 mm high. The anode was divided into three sections, 58, 50 and 200 mm along its depth. The chamber was filled with isobutane ($C_4H_{10}$) gas to a pressure of 20 Torr, and the energy deposition was measured along the trajectory using the three anodes as shown in figure 3.13. By plotting the energy deposited in the first section against the total energy deposition, the residual nuclei are separated by the proton number ($Z$) according to the Bethe Bloch equation 3.14 (see figure 3 in reference [90] for example). This separation is also useful in resolving any charge state anomalies.
Figure 3.13: Schematic of the ion chamber showing residual nucleus' path past the three anodes. The predicted isotopic resolution according to the Bethe Bloch equation is also indicated.
Figure 3.14: A photograph of the open ion chamber at Argonne. The focal plane of the FMA is in the foreground and the mylar entrance window is visible.
Neutron Detectors.

To increase the sensitivity for identification of weakly populated neutron-evaporation channels, liquid scintillator detectors [57, 80, 113, 115] can be used in conjunction with either a charged particle detector, or a recoil separator.

As discussed in section 3.1, the reaction kinematics result in the recoils, and hence the neutrons being distributed preferentially in a forward direction in the laboratory frame. The AYEBALL array had a ring of eleven NE213 [80, 113] detectors placed in the forward direction around the entrance to the FMA. The PEX array was not used in conjunction with a recoil separator, and so could accommodate an array of 16 BC501 [115] neutron detectors, positioned in a ‘wall’ at a forward angle to the beam direction, thus maximising the detection efficiency.

The detectors are organic scintillator material and contain a large amount of hydrogen. The neutrons interact with the protons in the material via (n,p) scattering, and the protons subsequently excite electrons in the scintillator material to higher energy states. These electrons either decay directly emitting a prompt fluorescence or via another energy state with a decay time of the order of a few nanoseconds, called a delayed fluorescence. The light output is detected in a light-pipe, photo-cathode and photomultiplier tube assembly [80] which converts it into an electric charge. The two types of decay produce a fast and slow pulse from the photomultiplier tube (figure 3.15).

Gamma-rays will also deposit energy in the scintillator material via (γ,e) Compton scattering and so must be resolved from the neutron events. Gamma-rays travel at the speed of light while neutrons travel slower, so neutron-γ discrimination can be obtained using measurements of the time of detection with respect to the pulse of beam, taken from the accelerator RF. The larger the distance from target to detector, the better the time of flight (TOF) separation between neutrons and γ-rays, but this comes at the cost of a reduced overall detection efficiency. Therefore, in general, the detectors are positioned as close to the target as possible, and pulse shape discrimination techniques [115, 116] are used. The principal behind this is that the
relative amounts of prompt and delayed fluorescence depend on the interaction type, and neutron induced proton scattering events have a larger prompt component than $\gamma$-ray events. Integration of the fast and slow components of the pulse can therefore be used to distinguish the two types of event.

Figure 3.15: Fast and slow components of a scintillation light output from a neutron detector.
Chapter 4

Argonne $^{24}$Mg + $^{40}$Ca Experiment.

Two experiments were performed in September 1995 to investigate the structure of nuclei around mass 60 at the Argonne National Laboratory using the AYEBALL $\gamma$-ray detector array in conjunction with the Fragment Mass Analyzer. A beam of $^{24}$Mg at an energy of 65 MeV was incident on a target of natural $^{40}$Ca of thickness 500 $\mu$g/cm$^2$ with 300 $\mu$g/cm$^2$ Au coating (to reduce oxidation) and 60 $\mu$g/cm$^2$ Au backing. The second experiment was performed with the same beam on a target of natural $^{40}$Ca of thickness 500 $\mu$g/cm$^2$ with 300 $\mu$g/cm$^2$ Au coating and 20 mg/cm$^2$ Au backing. The aim of the first experiment was to obtain clean recoil-$\gamma$ coincidences from which identification spectra (id-spectra) could be obtained, while the second experiment used a target which was thick enough to stop the recoils in view of the AYEBALL array, thereby obtaining high resolution $\gamma$-$\gamma$ coincidence data and angular correlation data. A graph of predicted cross section for the above reaction from PACE calculations [58] is shown in figure 4.1, both as a function of beam energy (4.1-a) and at the chosen energy 65 MeV (4.1-b). Note, figure 4.1-b is a log scale and the cross section for the mass 62, pn channel $^{62}$Ga is predicted to be two orders of magnitude lower than the strongest channel $^{61}$Cu (3p).
Figure 4.1: Predicted production cross sections for the $^{24}$Mg + $^{40}$Ca reaction from PACE calculations [58]. (a) As a function of beam energy (b) at 65 MeV, (log scale) Note, the mass 62, 1 proton, 1 neutron channel is $^{62}$Ga.
The beam at the Argonne National Laboratory was provided by ATLAS as discussed in section 3.2.1. The beam was incident onto the target which was positioned in the centre of the AYEBALL array. In the thin target experiment, the reaction products continued into the Argonne FMA and were stopped beyond the focal plane in the ionisation chamber (Figure 4.2).

![AYEBALL Diagram](image)

Figure 4.2: Schematic of the AYEBALL, FMA and ion chamber set up.

Depending on how many electrons are stripped from the beam ion by the stripper foil at the anode of the tandem accelerator, the ions can enter the target chamber in a variety of charge states. An initial 'sweep' was performed to find the optimum of the distribution, and thus maximise the transmission efficiency for the recoils of interest through the FMA.

### 4.1 Argonne Electronics.

The electronics logic diagram for the AYEBALL experiment is shown in figure 4.3. Individual γ-ray events in a germanium detector were pre-amplified and amplified and passed through a discriminator with a lower level cut off to reduce the electronic noise contributions. The signal is then fed into an ‘AND’ logic gate with a ‘NOT’ logic BGO event from the same detector, producing a Compton suppressed signal. This was wired into both an *analogue to digital converter* (ADC) and a multiplicity unit. The ADC converted the size of the pulse into a digital signal which was written to tape as a sixteen bit *hex* number. The multiplicity unit counted the number of
suppressed events in a certain time window, set to correspond to one beam burst, and thus one reaction. This unit then gave a voltage with a height corresponding to the $\gamma$ multiplicity of the event. By setting the lower and upper level cut off points on a discriminator, a signal was fed into the trigger electronics as either a $\gamma$ or $\gamma-\gamma$ event. Signals were also taken from the PPAC corresponding to left, right, up and down position, as well as total energy from the PPAC anode. These were pre-amplified and amplified and fed into two CAEN ADCs. The amplified anode signal was also passed through a discriminator into the trigger electronics. The first two ion chamber anode signals were combined and sent to both a CAEN and a Silena ADC, along with the third separate anode signal. The ion chamber cathode total energy signal was also sent to both the CAEN and the Silena ADC, as well as being passed through a discriminator and into the trigger electronics.

The logic pulse from the PPAC anode was used to start a time to analogue converter (TAC), which was stopped by the logic timing signal from the ion chamber. This PPAC-IC TAC had a range 0.5 to 1 $\mu$s and was fed into the Silena ADC. The PPAC signal was also used to give a logic coincidence with the $\gamma$ and the $\gamma-\gamma$ events, thus making a recoil-\$\gamma$ and recoil-$\gamma-\gamma$ trigger. This coincidence was also used to start a recoil-\$\gamma$ TAC, which was stopped with a delayed (1 to 2 $\mu$s) $\gamma$ event. The same was done for recoil-$\gamma-\gamma$, and these were fed into the CAEN ADC. Each of the trigger conditions was passed through a rate divider to enable only a prescaled number of the different events to be used as a trigger, thus reducing the dead time. The rate division for $\gamma-\gamma$, recoil-\$\gamma$ and recoil-$\gamma-\gamma$ were set to 1, but the $\gamma$ ‘singles’ were set to 1 in 999. These rate divided events were fed into a scaler to count the number of each type of event, and a master trigger fan-in-fan-out unit, which produced the master trigger with which the ADCs were gated. Thus only events which were deemed as ‘good’ by the electronics were converted by the ADCs. These events were then written to tape by an event manager for Compton-suppressed $\gamma-\gamma$ and higher fold germanium events, or $\gamma$ and higher fold accompanied by coincident PPAC and ion chamber and/or neutron events.
Figure 4.3: Diagram of the electronics for AYEBALL, FMA and ion chamber.
4.2 Subtraction of Random Events.

The finite detection efficiency of both the γ-ray detector array and the FMA, coupled with the width of the master gate, which spans ∼15 beam bursts, unavoidably results in the writing to tape of random events, such as γ-rays which are mis-correlated with a recoil from a different reaction which occurred in either a previous or subsequent beam burst. Subtraction of these events was carried out by using timing information from the pulsed beam of ATLAS and by correlating the detection times of γ-rays in the germanium detectors with the time of flight for recoils to reach the PPAC and ion chamber at the end of the FMA. For this experiment the v/c of the recoils of interest was ∼2.4%, corresponding to a time of flight for a fusion recoil through the FMA of approximately 850 ns. The beam was pulsed by a buncher before the tandem (figure 3.4) which produced bursts separated by 82 ns, and as narrow as ∼100 ps wide [66]. In this reaction only ∼1% of these beam bursts resulted in a fusion reaction, the rest either scattered into the side of the target chamber or continued into the FMA without reacting with the target. When a reaction did occur, only ∼1% will have had a γ-ray detected, and similarly, when a γ-ray was detected, only ∼5% of the corresponding recoils will have been detected at the end of the FMA. Figure 4.4 shows how the timing could be used to software select in-beam γ-rays, thus reducing the number of random events in a spectrum. The intense region of counts in the centre of each plot represents the events of interest, having been detected at the end of the spectrometer at a time after a γ-ray was detected which is consistent with the flight time through the FMA of a fusion recoil, and the background of counts either side of the central peak are caused by random events. The separation of the beam bursts is apparent here as ‘ridges’. There is also a spread in energy of the detected nuclei, but using two dimensional software gates, only those events within a certain energy range can be selected. Those events which fall outside the software condition were ignored in the offline sort.
Figure 4.4: (a) Energy loss in the PPAC versus time of flight for recoil and (b) Energy loss in the ion chamber versus time of flight for recoil. Note the clear separation between true recoils and random events from subsequent beam bursts.

### 4.3 Mass Gating.

For a given charge state, the FMA separates the residual nuclei according to mass at the focal plane, where they are detected by a position sensitive PPAC. A two dimensional spectrum can be produced of the position of the nuclei on the X-axis of the PPAC versus the energy loss in the PPAC, (figure 4.5.)

By setting software loci around each part of this two dimensional plot, it is possible to software gate each detected $\gamma$-ray with the mass over charge state ratio of the corresponding nucleus, and hence increment $\gamma$-ray spectra for each different $A/Q$ ratio (Figure 4.6). The $\gamma$-ray energy spectra for every germanium crystal are aligned so that the channel numbers of every spectrum correspond to the same energy. This is called *gainmatching*, and is performed by incrementing a spectrum using a radioactive $\beta$-decaying source. Since the source is stationary in the target position, there is no associated Doppler shift, and because the energies of the peaks from the source are known, a quadratic function can be applied to each separate energy spectrum, transferring the dispersion to a common energy difference per
Figure 4.5: Recoil position $X$ at the PPAC versus energy deposited in the PPAC, showing the dispersion in the $X$-plane of recoils in mass/charge state ratio.

channel. This technique is illustrated in more detail in section 5.3.

As figure 4.6 shows, the mass gated spectra are contaminated by other masses due to achromatic aberrations in the FMA [70]. In particular, the strongly populated dominant 3 proton evaporation channel to $^{61}$Cu is notably visible in the $A=62$ gated spectrum. However, it is possible to subtract normalised amounts of contaminant channels from each spectrum to produce 'clean' mass spectra (figure 4.7).

The transitions labelled in figures 4.6 and 4.7 are identified as coming from specific nuclei from the following references: $^{62}$Zn [117, 118, 119, 120, 121]; $^{61}$Cu [122, 123, 124, 125]; $^{61}$Zn [90, 126, 127, 128]; $^{60}$Ni [129]; $^{60}$Cu [130]; $^{64}$Zn [108, 109, 110]; $^{58}$Ni [111, 112]; $^{57}$Ni [111, 131, 132, 133].
Figure 4.6: Gamma-ray energy spectra gated by recoil position at the PPAC, showing mass separation.
Figure 4.7: Gamma-ray energy spectra gated by mass with normalised software subtractions of contaminants. The $^{64}\text{Zn}$ is present from reactions on $^{44}\text{Ca}$ target contaminants and appears in the mass 58 spectrum due to an $A/Q$ anomaly ($^{68}\text{Ni} \sim ^{64}\text{Zn}$).
4.4 Elemental Dispersion.

4.4.1 Neutron Gating.

The evaporation channels involving one or more neutrons were resolved by way of an additional neutron detector condition on the mass gated spectra. Neutrons are distinguished from the γ-ray events in the scintillator detectors as discussed in section 3.6.3 using two dimensional spectra of (a) the time of flight versus total energy deposited in the detector, and (b) slow component of the timing signal versus total energy deposited in the detector (figure 4.8). Despite this, some γ-rays were misidentified in the neutron gates and so the spectra are partially contaminated, however, this could be resolved by de-convoluting the spectra with and without the neutron condition to leave only the neutron evaporation channel (see figure 4.9).

![Figure 4.8](image)

Figure 4.8: Two dimensional neutron detector spectra of (a) total energy versus time of flight, and (b) total energy versus pulse shape discrimination for the neutron detectors. The software gates used to increment events associated with neutrons or γ-rays are shown. Note the slow neutron events which are less separated from the γ-rays in (b) are the most clearly resolved from spectrum (a), and vice versa.

The neutron gating can resolve the weak evaporation channels involving neutrons, but the neutron detection efficiency is rather low (measured from figure 4.9 using the 124 keV peak as $\frac{I(G1n)}{I(G1)} = \frac{909}{89702} \sim 1\%$).
Figure 4.9: Gamma-ray energy spectra gated by mass 61, (top) only (middle) with a 1 neutron software gate, (bottom) neutron gated with normalised amount of (top) subtracted showing a pure 2pn gated ($^{61}\text{Zn}$) identification spectrum.
4.4.2 Ion Chamber Gating.

The ion chamber was also used to afford a degree of Z separation. Figure 4.10 shows the plot of total energy deposited in the ion chamber, versus the sum of the energy deposited at the first and second anodes. Figure 4.10-a shows the raw spectrum with clear separation between the true recoils and scattered beam. By setting conditions on the recoil-γ-TAC as shown in figure 4.4, this scattered beam can be removed. This pure recoil plot is then rotated in software in the direction indicated in figure 4.10-a to produce the spectrum shown in figure 4.10-b. The recoils are dispersed according to their Z across the width of this spectrum, which is then projected onto the x-axis.

![Figure 4.10](image.png)

Figure 4.10: Ion chamber spectrum used to resolve the Z of the recoil. (a) Total energy deposited verses energy deposited at first two anodes, (b) same signal gated by timing conditions to remove scattered beam and rotated for projection.

A matrix was then produced of this clean rotated ion-chamber signal, gated by different A/Q conditions, versus γ-ray energy. By setting slices on the ion-chamber axis, and projecting onto the energy axis, Z gated γ-ray spectra can be produced. Figure 4.11 shows the projections of the rotated plot for the two mass 61 isobars $^{61}$Cu and $^{61}$Zn, produced by gating on γ-rays from transitions in each
nucleus in turn. Figure 4.12 shows the $\gamma$-ray spectra gated on either side of this plot, and demonstrates the $Z$ separation afforded in the $^{24}$Mg on $^{40}$Ca reaction. By de-convoluting these and the neutron gated spectra, it is then possible to produce spectra which are isotopically clean. Such spectra have been isolated for the nuclei $^{61}$Cu, (3p channel) [122, 123, 124, 125], $^{61}$Zn (2pn) [90, 126, 127, 128] and $^{62}$Zn (pn) [108, 110, 117, 118, 119, 120], and these are shown in figure 4.13.

![Figure 4.11: Projections of ion chamber signal gated by $\gamma$-rays from two mass 61 isobars $^{61}$Cu (1310 keV) and $^{61}$Zn (124 keV).]
Figure 4.12: Gamma-ray energy spectra gated by ion chamber, (a) mass A=61 only, (b) lower Z side, enhancing the 3p channel $^{61}$Cu and (c) higher Z side, enhancing the 2pn channel $^{61}$Zn.
Figure 4.13: Isotopically pure identification spectra for (top) $^{61}$Cu, (middle) $^{61}$Zn, and (bottom) $^{62}$Zn.
4.5 Gamma-Gamma Gating.

A two dimensional γ-ray energy coincidence matrix was sorted in off-line analysis. By gating this matrix by the transitions identified in the clean mass and Z gated spectra, it was possible to produce spectra which identify all the detected energy transitions in particular nuclei (see section 3.5). Gamma-ray coincidence spectra are shown for the nuclei $^{61}$Cu in figures 4.14 and 4.15, and for $^{62}$Zn in figures 4.16 and 4.17. By noting which transitions were present and which were absent in each gate, it was possible to build up the decay scheme of a nucleus. Ordering of transitions in a cascade may also be inferred from such spectra because the intensity of transitions in a cascade generally increases at lower excitation energy due to side feeding from other non-yrast states. Thus the intensity of transitions above a gate will generally decrease with increasing excitation energy in the spectrum. However, transitions below a gate will generally have the same intensity (if there is no decay out of the cascade) as the feeding decay path has been defined by the gate. By measuring intensities in different gates, it is therefore possible to ascertain the ordering of the decay.

4.6 Spin and Parity Assignments.

In order to use the DCO technique discussed in section 3.8, the detectors of the AYEBALL array were divided into two sets. The detectors at angles 79°, 101° and 134° were summed to give roughly equal overall detection efficiency to the sum of the detectors at 158°. By gating on a transition of known multipolarity, detected at one angle and projecting other transitions which are detected in coincidence at the other angle, a DCO ratio was calculated from equation 3.8.

$$R_{DCO} = \frac{I(158^\circ) \text{ gated at (79°, 101°, 134°)}}{I(79^\circ, 101^\circ, 134^\circ) \text{ gated at (158^\circ)}} \times e$$  \hspace{1cm} (4.1)$$

where $I$ is the measured intensity of a peak, and $e$ is an efficiency multiplication.
factor with which the experimental value is corrected for the detection efficiencies of both the gate and the projected transition. This factor is equal to;

\[ \varepsilon = \frac{\text{Efficiency of gate at (158°) \times Efficiency of projection at (79°, 101°, 134°)}}{\text{Efficiency of gate at (79°, 101°, 134°) \times Efficiency of projection at (158°)}} \]

The difference in intensity of quadrupoles and dipoles, gated by an E2 transition in \(^{65}\text{Cu}\) is clearly illustrated in figure 4.18. This is gated by the known 1310 keV \(\frac{7}{2}^- \rightarrow \frac{5}{2}^-\) stretched E2 transition in \(^{61}\text{Cu}\) [123]. The 1317 keV \(\frac{11}{2}^- \rightarrow \frac{7}{2}^+\) and 1361 keV \(\frac{13}{2}^+ \rightarrow \frac{9}{2}^+\) E2 transitions [123] shown in the figure are of roughly equal intensity in both the angle projections, while the 1410 keV \(\frac{9}{2}^+ \rightarrow \frac{7}{2}^-\) E1 transition [123] is clearly more intense in the spectra of detectors nearer 90°.

Care was taken when determining the statistics in a peak, as for higher spin transitions, there was a Doppler shifted component to the peaks. This is associated with the fast feeding of these levels by statistical E2 transitions, as discussed in section 3.1. The decay from these rapidly populated high spin states can therefore occur while the nucleus is still slowing down in the target, and thus there may be a lineshape associated with the Doppler shift [134]. The velocity of the residual nucleus at formation can be calculated from equation 3.3 and so the maximum amount of shift can be deduced for each transition of known energy at each angle from equation 3.12. Figure 4.19 shows this effect for the 1705 keV, \(\frac{21}{2}^+ \rightarrow \frac{15}{2}^+\) transition in \(^{61}\text{Cu}\) for the two angular projections, with the maximum shift calculated and labelled at 1694 keV for 101°, 1671 keV for 134° and 1660 keV for 158°. The 1733 keV transition from a lower lying state is shown for comparison and has no shifted component.
Figure 4.14: Gamma-gamma coincidence spectra for $^{61}$Cu from the backed target experiment. The clean mass and ion-chamber gated id-spectrum for this nucleus is also shown for comparison (a). Gated by (b) 1410 keV, (c) 1317 keV and (d) 970 keV.
Figure 4.15: Gamma-gamma coincidence spectra for $^{61}$Cu from backed target experiment. Gated by (a) 1705 keV, (b) 1871 keV, (c) 1952 keV and (d) 850 keV.
Figure 4.16: Gamma-gamma coincidence spectra for $^{\text{61}}\text{Zn}$ from the backed target experiment. The clean mass and ion-chamber gated id-spectrum for this nucleus is also shown for comparison (a). Gated by (b) 418 keV, (c) 1403 keV and (d) 1273 keV.
Figure 4.17: Gamma-gamma coincidence spectra for $^{61}$Zn from backed target experiment. Gated by (a) 124 keV, (b) 1079 keV, (c) 1289 keV and (d) 1539 keV.
Figure 4.18: DCO $^{61}$Cu gated spectra (gated by 1310 keV $\frac{7}{2}^{-} \rightarrow \frac{5}{2}^{-}$ pure E2) highlighting the difference in the angle gated spectra for $\Delta I = 2$ and $\Delta I = 1$ transitions.
Figure 4.19: DCO $^{61}$Cu gated spectra (gated by 529 keV $\frac{17^+}{2^-} \rightarrow \frac{13^+}{2^-}$ pure E2) showing 1705 keV, ($\frac{21^+}{2^-}$) → $\frac{17^+}{2^-}$ transition with shifted components labelled at 1694 keV for 101°, 1671 keV for 134° and 1660 keV for 158°. The 1733 keV $\frac{7^-}{2^-} \rightarrow \frac{3^-}{2^-}$ transition in comparison has no shifted component.
The measured DCO ratio, gated by both dipoles and stretched quadrupoles is plotted for the strongly populated levels in $^{61}$Cu and $^{61}$Zn in figure 4.20. The gating dipoles are chosen to be as pure as possible, being either an $E1$, or an $M1$ with very little $E2$ admixture. Clear separation is apparent between transitions of different multipolarities. Ratios for all levels observed with sufficient statistics are given in tables 4.1 and 4.2 respectively.

Figure 4.20: DCO ratios for $^{61}$Cu gated by (a) an $E2$ and (b) a pure dipole; for $^{61}$Zn gated by (c) an $E2$ and (d) a pure dipole. Weighted averages of previously known transitions are indicated by the lines (see section 5.4.2).
4.7 Decay Schemes.

The decay schemes for $^{61}$Cu and $^{61}$Zn derived from this work are shown in figures 4.21 and 4.22 respectively. Spin and parity assignments, taken from DCO ratios, along with intensities, measured in both mass gated singles data and $\gamma$-$\gamma$ data, are tabulated in tables 4.1 and 4.2 respectively. The intensities taken from the backed target $\gamma$-$\gamma$ data are indicated in figures 4.21 and 4.22 by the arrow widths. Spin and/or parity labels in brackets indicate tentative assignments.

4.7.1 $^{61}$Cu.

Prior to this study, states had been observed in the nucleus $^{61}$Cu up to 4081 keV, (spin $^{13/2}^+$) using the following reactions: Light ion; $^{58}$Ni(α,p)$^{61}$Cu [123, 124] and $^{50}$Ni(p,γ)$^{61}$Cu [135, 136]; heavy ion fusion evaporation; $^{40}$Ca($^{24}$Mg,3p)$^{61}$Cu [125]; and stripping; $^{60}$Ni($^{3}$He,d)$^{61}$Cu [137]. Theoretical studies of $^{61}$Cu have also been performed with calculations using the interacting boson model [138] (see section 6.3) and the shell model [139].

All the yrast or near yrast energy level assignments made in the previous work are in agreement with the decay scheme derived from this data. As examples, three intense transitions in this data set are the 1310, 1410, and 970 keV, which have a measured DCO ratio when gated by an E2 of 1.10±0.04, 0.56±0.05 and 0.45±0.03 respectively. The multipolarity assignments are thus made as a 1310 keV $^{7/2}^-$ → $^{9/2}^-$ stretched E2, a 1410 keV $^{9/2}^+ → ^{7/2}^-$ E1, and a 970 keV $^{5/2}^- → ^{3/2}^-$ M1/E2, all in agreement with assignments by Sarantites et al. [124], Sziklai et al. [135] and Tingwell et al. [136]. Similarly, when gated by the 1410 keV E1, the DCO ratios for the 340 keV and the 1310 keV transitions are 0.94±0.14 and 1.64±0.04 respectively, consistent with the M1/E2 and stretched E2 assignments previously reported [124, 135, 136].

However many of the non-yrast states previously identified are not observed in this data. For example, the Sziklai experiment [135] identified most of the states observed in this data up to ~4 MeV, but also a number of non yrast states which
are not populated in this heavy ion fusion evaporation experiment, including a $\frac{7}{2}^-$ state at 2399 keV and a $\frac{9}{2}^+$ state at 4132 keV. The light ion experiment reported by Tingwell et al. also identifies transitions which are not observed in this data, namely a low lying $\frac{1}{2}^-$ state at 475 keV, and another at 2089 keV, two $\frac{3}{2}^-$ states at 1660 keV and 1933 keV, and two non-yrast $\frac{5}{2}^-$ states above the yrast 970 keV level, at 1394 keV and 1904 keV. The fact that there is little observation of non-yrast states supports the idea of near yrast population in heavy ion fusion evaporation reactions, and hence the assumption that when making spin assignments from DCO data, the spins increase with increasing excitation energy.

During this analysis, a parallel study was performed by Hatsukawa et al. [122] using the reaction $^{40}\text{Ca}(^{28}\text{Si},\alpha3p)^{61}\text{Cu}$. The study produced a decay scheme which is in general, consistent with the current work, identifying all but the 210, 326, 353, 566, 909 and 1975 keV transitions (see figure 4.21). Reference [122] also proposes a number of transitions not observed in this data set, linking the band based on the 2336 keV $\frac{3}{2}^-$ state into a band built on the ground state, which is not observed with sufficient intensity in this data. It also places the 937 keV and the 1042 keV transitions in the opposite order between the 4590 keV $\frac{13}{2}^+$ and 2612 keV $\frac{9}{2}^-$ states, although this is not in agreement with intensity arguments from the current work.

The spins and parity assignments in the current study are made where possible from DCO ratios, assuming an $I^r=\frac{3}{2}^-$ ground state. This assumption is made on the basis of the beta decay study by Singh et al. [140]. Spin and parity cannot be assigned to some of the levels on the basis of DCO ratios alone due to either poor statistics or contaminant lines in the gates. However, in some cases, other arguments can be made which can be used to infer the most likely assignment. For example, the 909 keV transition feeding the 4081 keV yrast $I^r=\frac{13}{2}^+$ state cannot be determined with certainty from DCO arguments, but it is consistent with an M1/E2, and this is favoured over an E2 assignment due to the weak intensity, which suggests a non-yrast state. If the decaying level had a spin of $\frac{17}{2}^+$, it would be yrast, and so expected to be more strongly populated. If the state had a negative parity, one
might expect to observe an E2 transition feeding the 3016 keV \( I^\pi = \frac{11}{2}^- \) level. The DCO ratio for the 1705 keV transition feeding the 5120 keV \( I^\pi = \frac{17}{2}^+ \) state has good statistics, but has a value which is one standard deviation away from the value for an E2, and so cannot be said to be either E2 or dipole with certainty. The levels above it are consequently labelled tentatively with spins and parities, however when compared to shell model calculations these levels are consistent with having negative parity (see section 4.8). The spin and parity assignments in the band built on the 2336 keV \( I^\pi = \frac{9}{2}^- \) state can not be made with confidence based on DCO ratios alone above the 4288 keV level, leaving the 353, 565 and the 850 keV cascade un-assigned. However, the value of 0.63±0.1 for the 850 keV transition suggests a pure \( \Delta I = 1^+ \), while the 565 keV transition has a ratio between 0.75 and 1.20, depending on which gate is used, and could conceivably be a \( J \rightarrow J \) transition. The 353 keV transition has a DCO value of 0.49±0.11, suggesting a dipole, which is consistent with the 6056 keV level having an assignment of \( \frac{17}{2}^+ \). If the 1975 keV transition to the 4081 keV yrast \( I^\pi = \frac{13}{2}^+ \) state is an E2 transition, then this would make the decaying state also a 6056 keV \( \frac{17}{2}^+ \) state. The energies of the \( \gamma \)-ray transitions below each state add up to 6056.1±1.3 for the level decaying via the 1975 keV \( \gamma \)-ray, and 6055.6±1.6 for the level decaying via the 353 keV \( \gamma \)-ray, and it is therefore suggested that the 1975 and the 353 keV transitions are decays from the same 6056 keV level. Spin and parity assignments for the three transitions directly below this level are thus tentatively made on this basis.
Figure 4.21: Partial decay scheme for $^{61}\text{Cu}$ observed in the current work.
Table 4.1: Transitions identified in $^{61}$Cu in the present work with AYEBALL. The number of the DCO gate refers to the transitions labelled with a gate number in bold type underneath the energy. Transitions with their energies labelled in bold type were known prior to this study [123, 124]. DCO information in italic refers to information from an E1 gate. Intensities are relative to the 1310 keV $\frac{7}{2}^+ \rightarrow \frac{3}{2}^+$ transition. [123]

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<th>$E_f$ (keV)</th>
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<th>DCO gate</th>
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Table 4.1 continued: Transitions identified in $^{61}$Cu.

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105
Table 4.1 continued: Transitions identified in $^{61}$Cu.

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Table 4.1 continued: Transitions identified in \(^{61}\text{Cu}\).

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4.7.2 $^{61}$Zn.

Prior to this study, states had been observed in $^{61}$Zn up to an excitation energy of 4415 keV, and a spin and parity of $\frac{17}{2}^-$, using the following reactions: Light ion; $^{58}$Ni($\alpha$,n)$^{61}$Zn [126, 127, 128]; heavy ion fusion evaporation; $^{58}$Ni($^6$Li,p2n)$^{61}$Zn, $^{54}$Fe($^{10}$B,p2n)$^{61}$Zn, $^{40}$Ca($^{24}$Mg,2pn)$^{61}$Zn [126]; and particle transfer; $^{58}$Ni($^{12}$C,$^3$Be)$^{61}$Zn [141], $^{58}$Ni($^6$Li,t)$^{61}$Zn [142].

The ground state spin and parity of $^{61}$Zn was theoretically suggested to be $\frac{3}{2}^-$ by Webber et al. [141] and Sandhu [143], and confirmed experimentally with $\beta$ decay studies by Dulfer et al. [144] and Hoffman and Sarantites [145]. All the yrast or near yrast energy level assignments made in the previous work are in agreement with the decay scheme derived from this data. For example, in the study by Schubank et al. [126], the low lying yrast and non-yrast levels were observed, including all the states identified in this work up to 3336 keV. However, the current work indicates a spin/parity assignment for the 2400 keV level of $\frac{9}{2}^-$ rather than $\frac{11}{2}^-$ as published by Schubank. The current analysis also confirms the tentative $\frac{11}{2}^-$ assignment for the 2270 keV level. As with the $^{61}$Cu states, many of the non-yrast states previously identified are not observed in this work. For example, Schubank observes a $\frac{1}{2}^-$ state at 88 keV, a $\frac{5}{2}^-$ state at 755 keV, another $\frac{1}{2}^-$ state at 938 keV, a $\frac{3}{2}^-/\frac{5}{2}^-$ state at 1361 keV, and a $\frac{7}{2}^-$ state at 1402 keV, all of which are absent in the current study. It should be noted that a 755 keV transition is identified in the $^{61}$Zn id-spectrum (figure 4.13), in agreement with one seen by Schubank to decay from a 755 keV $\frac{5}{2}^-$ level to the $\frac{3}{2}^-$ ground state. However the transition could not be linked with the rest of the decay scheme, and thus has not been assigned in this study.

The decay scheme for $^{61}$Zn derived from this work is shown in figure 4.22. The 4264 keV level has a tentatively assigned spin of $\frac{15}{2}$ from the DCO ratio for the two transitions depopulating it. If the level had a spin/parity of $\frac{13}{2}$, one might expect a stretched E2 transition to feed the 2400 keV $\frac{9}{2}^-$ state. The non observation of this transition in the data is thus consistent with the spin assignment shown.
Figure 4.22: Partial decay scheme for $^{61}$Zn observed in the current work. Previous studies [126] had identified transitions up to 4413 keV, spin $\frac{17}{2}^{-}$. 
Table 4.2: Transitions identified in $^{61}$Zn in the present work with AYEBALL. The number of the DCO gate refers to the transitions labelled with a gate number in bold type underneath the energy. Transitions with their energies labelled in bold type were known prior to this study [126], DCO information in italic refers to information from an El gate. [126]

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4.8 Shell Model Comparison and Discussion.

A shell model calculation was performed for $^{64}$Cu and $^{64}$Zn by M. Jensen [146] using as a model space, the $f_{\frac{5}{2}}, p_{\frac{3}{2}}, p_{\frac{1}{2}}$ and $g_{\frac{9}{2}}$ orbitals and the charge-dependent Bonn potential [147] as a nucleon-nucleon potential. The model assumes a closed $^{56}$Ni$_{28}$ core and does not allow for core breaking. The energies of these orbitals are calculated for the protons, relative to the lowest $p_{\frac{3}{2}}$, to be 1.04 MeV for the two degenerate $f_{\frac{5}{2}}$ and $p_{\frac{1}{2}}$ orbitals, and 3.51 MeV for the $g_{\frac{9}{2}}$. For the neutrons, the levels are calculated relative to the lowest $p_{\frac{3}{2}}$ to be 0.77 MeV for the $f_{\frac{5}{2}}$, 1.11 MeV for the $p_{\frac{1}{2}}$, and 3.70 MeV for the $g_{\frac{9}{2}}$ orbital.

4.8.1 $^{64}$Cu Shell Model.

This model space leaves 4 valence neutrons and 1 valence proton in the four active orbitals, and so the maximum angular momentum this basis can generate is $(\nu g_{\frac{9}{2}})_{\frac{9}{2}^+} \otimes (\nu f_{\frac{5}{2}})_{\frac{1}{2}^-} \otimes (\pi g_{\frac{9}{2}})_{\frac{9}{2}^+} = \frac{33}{2}^{-} h$. The results of this calculation are compared with the experimental levels in figure 4.23 for positive and negative parity states separately. The open circles are the calculated levels and the filled circles are the experimental data. The crosses are experimental data which have tentative spin/parity assignments, and so each of these is shown in both the positive and negative parity plots. Those crosses plotted joining two or more spin values could be either value, while those placed on the spin value have only tentative parity.

The first $\frac{5}{2}^-$ state at 970 keV is well reproduced, but the $\frac{7}{2}^-$ at 1310 keV is too high by $\sim$400 keV. This may suggest that excitation from the $f_{\frac{5}{2}}$ orbital in the $^{56}$Ni core is important. The yrast $\frac{9}{2}^+$ state is also too high by $\sim$800 keV, and since the $g_{\frac{9}{2}}$ is positive parity, it may be that this orbital is not correctly modelled. The calculated high spin negative parity states are in general also too high, although they agree with the tentative data above $\sim 6.5$ MeV much more closely than the positive parity calculations at this spin. This suggests, albeit cautiously, that these experimentally unassigned states may have negative parity. Although there is no theoretical preference as to whether the 1705 keV transition feeding the $I^\pi=\frac{17}{2}^+$
state is a dipole or a quadrupole, if it changes parity, the DCO ratio of \(>1\) would favour an M2 over an E1. A 1705 keV M2 transition in \(^{61}\text{Cu}\) would have a \(T(M2)\) value, calculated using table 1.1, of \(5.03 \times 10^9\), corresponding to a lifetime of 0.2 ns. The 6826 keV level is therefore consistent with an \(I^\pi = \frac{3}{2}^-\), with two M1/E2 transitions above it.

The 1942 keV level certainly agrees much more closely with the negative parity calculations, and there is a closer match at \(\frac{5}{2}\) rather than \(\frac{7}{2}\). The 3260 and 3780 keV levels have experimentally undetermined parity, but comparison with the shell model calculations suggest a negative parity may be preferred for these states. The 4990, 5138, 5465, 5703, and the 6056 keV levels have better agreement with the positive parity rather than the negative parity calculations. The 4990 keV level matches closely with a \(\frac{15}{2}^+\) prediction, consistent with the argument made earlier regarding such an assignment from its intensity (see section 4.7.1). The 5138 keV level also better fits the calculation as a \(\frac{15}{2}^+\) state, in agreement with the suggestion that the 850 keV transition is an E1. Indeed all the states which decay into this level agree well with calculated values which concur with the suggestion that the 6056 keV level at the top of this cascade is a \(\frac{11}{2}^-\), and that this level also depopulates via the 1975 keV transition (see section 4.7.1).

### 4.8.2 \(^{61}\text{Zn}\) Shell Model.

This model space leaves 3 valence neutrons and 2 valence protons in the four active orbitals and so the maximum angular momentum this basis can generate is \(\left(\nu g_{9/2}\right)_{15^-} \otimes \left(\nu f_{7/2}\right)_{3^-} \otimes \left(\pi g_{9/2}\right)_{15^+} = \frac{37}{2}^- \hbar\). The results of this calculation are compared with the experimental levels in figure 4.24. Only negative parity states are produced by the calculation. The open circles are the calculated levels, the filled circles are the experimental data, and the crosses are experimental data which have only tentative spin assignments. There is in general reasonable agreement, although it is suspected that the degrees of freedom outside the chosen model space are important in the low lying part of the spectrum.
Figure 4.23: Comparison of experimental and theoretical data for $^{61}$Cu for (a) negative and (b) positive parity. The open circles are calculations, the filled circles are experimental data. The crosses are experimental data with only tentative spin/parity assignments. Where they span two or more spin values, they could be either value.
Figure 4.24: Comparison of experimental and theoretical data for $^{61}\text{Zn}$ for negative parity. The open circles are calculations, the filled circles are experimental data. The crosses are experimental data with only tentative spin/parity assignments. Where they span two or more spin values, they could be either value.

4.9 $^{62}\text{Ga}$ at Argonne.

The 1 proton 1 neutron evaporation channel from the AYEBALL experiment produced $^{62}\text{Ga}$. Figure 4.1 shows how relatively small the cross section was. Consequently the isotopically pure identification spectrum (figure 4.25) has very few counts and assignments could not be made for this nucleus with certainty from this data set. The 375 keV transition is notably absent from other mass gated spectra, it is therefore consistent with a mass 62 isobar.
Figure 4.25: Channel selected $\gamma$-ray spectra, (top) with no channel selectivity conditions, (middle) mass 62 gated, showing transition in dominant channel $^{62}\text{Zn}$, (bottom) ion-chamber gated with candidate transitions in $^{62}\text{Ga}$ marked with energies.
An ion chamber projection gated by the tentatively assigned 376 keV transition is shown in figure 4.26 versus the projection gated by the strong $4^+ \rightarrow 4^+$ 557 keV transition in $^{62}$Zn. While the statistics are low, there is an apparent shift in the centroid of the two peaks, suggesting that the 376 keV transition is from a nucleus with a higher proton number i.e. Gallium.

Figure 4.26: Projections of ion chamber signal gated by all masses, and $\gamma$-rays from two mass 62 isobars $^{62}$Zn (557 keV) and $^{62}$Ga (376 keV).
Chapter 5

PEX $^{28}\text{Si} + ^{40}\text{Ca}$ Experiment.

An experiment was conducted at the Niels Bohr Institute, Risø, Copenhagen, Denmark to investigate the near yrast states of nuclei around the mass 60 region, with particular interest in the $N=Z=31$ nucleus $^{62}\text{Ga}$, using the reaction $^{40}\text{Ca}(^{28}\text{Si},\alpha\text{pn})^{62}\text{Ga}$. The 88 MeV $^{28}\text{Si}$ beam was provided by a tandem accelerator and bombarded a 1 mg/cm$^2$ self supporting target of enriched (99.96%) $^{40}\text{Ca}$. The beam energy was chosen by comparison with the evaporation residue code PACE [58], which suggested a cross-section of approximately 7 mb for the $\alpha\text{pn}$ channel, $^{62}\text{Ga}$ (figure 5.1). This is an order of magnitude increase in the production cross section for the $^{62}\text{Ga}$ channel in the Argonne experiment, however it is also apparent that there are more reaction channels open and the total cross section is approximately 1.5 times greater than for the Argonne experiment. The apparatus used was 'PEX', which comprised four cluster germanium detectors, a silicon ball charged particle detector and a neutron wall at forward angle to the target. Figure 5.2 shows a schematic of the array.
Figure 5.1: Predicted production cross sections for the $^{28}$Si + $^{40}$Ca reaction from PACE calculations [58], (a) as a function of beam energy (b) at 88 MeV, (log scale)

Note, $^{62}$Ga is the αpn mass 62 channel.
Figure 5.2: Schematic of the PEX experimental array at the Niels Bohr Institute.

5.1 PEX Electronics.

The electronics logic diagram for the PEX apparatus is shown in figure 5.3 from the pre-amplifiers to the ADCs. Note, only 31 of the Si-ball amplifiers were needed. Data were written to tape for Compton-suppressed $\gamma-\gamma$ and higher fold germanium events, accompanied by coincident silicon ball and/or neutron events.
Figure 5.3: Diagram of the electronics for the PEX array.
5.2 Channel Selection.

Isotopic identification was performed using the differential energy loss measurements from each element of the silicon ball as discussed in section 3.6.1, and pulse shape analysis and total energy signals from the neutron detectors as discussed in section 3.6.3. The neutron detector total energy signals and slow component of the scintillator pulses are plotted on two axes of two dimensional plots, figure 5.4. The neutrons are separated from $\gamma$-rays by their having slower scintillator pulses and less total energy. Neutron events are identified using software loci to create a gate as shown in the figure (see for examples [115, 116]).

![Figure 5.4: Total energy versus slow component of the scintillator pulse for a neutron detector showing neutron/$\gamma$ separation.](image)

The energy signal from silicon detectors at different angles to the beam direction is shown in figure 5.5 and they highlight the separation between proton and $\alpha$-particle signals. The majority of the protons punch through the detector and so they deposit less energy than the $\alpha$-particles which are stopped in the detector.
Figure 5.5: Total energy deposited in the silicon detector (a) nearest 0° (b) at 70° (c) at 110° (d) nearest to 180° to the beam direction. The separation between proton and α-particle characteristic signals is illustrated.
5.3 Data Analysis.

As in the AYEBALL experiment discussed in the previous chapter, the germanium energy spectra are gainmatched using a radioactive source and a quadratic correction function. Figure 5.6 shows the energy spectra from a $^{152}$Eu source for the central crystals of the four cluster detectors in PEX with this gainmatch function applied. For the reaction data, a Doppler shift is applied to every gainmatched spectrum which is calculated from the angle that each crystal makes with the beam direction, and the $\beta$, $\left(\frac{v}{c}\right)$ of the reaction recoils, according to equation 3.12. Figure 5.7 shows spectra for the reaction $^{28}$Si + $^{40}$Ca for the central crystals of the four PEX clusters with a $\beta$ value of zero. The effect of the Doppler shift is shown, the shift being of negative energy as the angles are greater than 90°, and larger for the detectors nearer to 180°. Figure 5.8 shows the same spectra with the correct $\beta$ value for the reaction, the peaks consequently being correctly gainmatched.

The effect of Doppler broadening is also apparent in these figures. It is worth noting that the effect is greater in the detectors nearer to 180°, and not those closer to 90° (figure 5.9) as may be expected simply from the sine dependence generated from the differential of the Doppler shift equation with respect to $\Delta \theta$. The measured value of full width at half maximum height (FWHM) is 14 keV at 1.096 MeV for clusters centred at 146° and reduces to 8 keV for clusters centred at 105° (see figure 5.9). This compares with the FWHM of a stationary source line form $^{152}$Eu of 3 keV at 0.972 MeV (see figure 5.6). This effect arises because the full differential expression for the Doppler shift equation must also include a term which is differentiated with respect to velocity, $\Delta \beta$ (equation 3.13). While this term is generally small for nuclei with $A \geq 100$, it must be considered in this lower mass region, particularly for channels involving $\alpha$-particle evaporation. The reason is that the mass of the evaporated particles (particularly an $\alpha$-particle) is comparable to the mass of the residual nucleus, which will therefore have a different velocity depending on whether the $\alpha$-particle is evaporated in a forward or backward direction. This $\Delta \beta$ term has a cosine dependence in equation 3.13, and so is largest at angles approaching 180°.
Figure 5.6: Gamma-ray spectra for $^{152}$Eu source data with gainmatch (a) for cluster centred at $105^\circ$, (b) for cluster centred at $146^\circ$, (c) for cluster centred at $146^\circ$, (d) for cluster centred at $105^\circ$. Clusters 1 and 4 and clusters 2 and 3 are at different azimuthal angles $\phi$. 
Figure 5.7: Total $\gamma$-ray spectra gainmatched to the source calibration spectrum coefficients, but with no Doppler correction applied; (a) for cluster centred at 105°, (b) for cluster centred at 146°, (c) for cluster centred at 146°, (d) for cluster centred at 105°. The effect of the particle emission on the Doppler broadening is clearly visible.
Figure 5.8: Gamma-ray spectra with no channel selectivity with gainmatch and Doppler shift correction, using a value of beta=0.018 (a) for cluster centred at 105°, (b) for cluster centred at 146°, (c) for cluster centred at 146°, (d) for cluster centred at 105°.
Figure 5.9: Gamma-ray spectra with no channel selectivity (top) for clusters centred at 146°, (bottom) for clusters centred at 105°.
Approximately $5 \times 10^8$ unfolded $\gamma-\gamma$ coincidence events were obtained and sorted in off-line analysis into the form of a two dimensional $\gamma-\gamma$ coincident energy matrix. This matrix was then software gated by different multiplicities of detected protons, neutrons and $\alpha$-particles. Gamma-ray spectra could then be incremented for each evaporation channel by demanding that specific numbers of protons, $\alpha$-particles and neutrons were detected in coincidence with the $\gamma$-ray. Examples of such spectra are shown in figure 5.10. Note the contaminant nuclei in each spectra due to the finite detection efficiency of the detectors. For example, the 3 proton evaporation channel is present in the 1 and the 2 proton gated spectrum whenever protons are not detected (figure 5.10-b), but the 2 proton evaporation channel should not be present in the 3 proton gated spectrum. There are also pure charged particle evaporation channels in neutron gated spectra as the neutron identification process cannot resolve neutrons from $\gamma$-rays with 100% efficiency.

The particle detection efficiency was calculated for the 2p, 3p, 4p, 1$\alpha$ and 2$\alpha$ evaporation channels by integrating the number of counts in a peak corresponding to a known transition in each of these channels. The ratio of counts was taken for each peak in the 1p, 2p, 3p, 4p, 1$\alpha$ and 2$\alpha$ gated $\gamma$-ray spectra. It was hence found [107] that the detection probability was $\sim 0.5$ for each proton from either the 4p or 3p channel, but was $\sim 0.7$ for each proton from the 2p channel. This difference was due to the recoil cone being forward focused and the finite granularity of the silicon detector.

The transitions labelled in figure 5.10 are identified as coming from specific nuclei from the following references. $^{66}$Ge [148], $^{65}$Ga [149, 150], $^{65}$Ge [120, 151], $^{65}$Cu [152], $^{64}$Ga [153], $^{64}$Ge [154], $^{62}$Zn [117, 118, 119, 120, 121], $^{61}$Cu [122, 123, 124, 125], $^{61}$Zn [90, 126, 127, 128], $^{60}$Ni [129], $^{60}$Cu [130], $^{64}$Zn [108, 109, 110], $^{58}$Ni [111, 112], $^{57}$Ni [111, 131, 132, 133], $^{56}$Ni [155], $^{55}$Co [156].
Figure 5.10: Gamma-ray channel selected spectra with (a) no channel selectivity, (b) 1 proton condition, (c) 1 neutron condition, (d) 1 \( \alpha \) and 2 protons required, and (e) 2 \( \alpha \)-particles and 2 protons required.
5.4 $^{62}$Ga at NBI.

Figure 5.11 shows projections from matrices with no channel selection condition, and also gated by the detection of $1\alpha$, $0$ or $1p$, $1n$, $1$ (henceforth called $\alpha pn$) $2\alpha 1p$, and $1\alpha 2p$ respectively. The $1p1\alpha 1n$ channel leading to $^{62}$Ga is highlighted in figure 5.11-b and clearly identifies the lines at 246, 376, 571 and 1179 keV, which we assign to $^{62}$Ga. The 246, 376 and 571 keV are also evident in the Argonne $^{62}$Ga gated spectrum (figure 4.25). Contaminant lines from the $1\alpha 2p$ channel ($^{62}$Zn [117]), and the $2\alpha 1p$ channel ($^{65}$Ge [151]), which appear from a combination of the finite detection efficiency of the silicon ball and the misidentification of $\gamma$-ray events in the neutron detectors are also indicated. The transitions associated with $^{62}$Ga are present in the $\alpha pn$ gated matrix but notably absent in both the $2\alpha p$ and $2p\alpha$ gated $\gamma-\gamma$ projections, which are, as expected, dominated by lines from $^{60}$Cu [126, 157] and $^{62}$Zn [108, 117] respectively.

5.4.1 Gamma-Gamma Coincidences.

By measuring the intensities of transitions in different gates of a $\gamma-\gamma$ matrix, and by observing which lines are absent from certain gates, it is possible to order the transitions into a decay scheme for that nucleus. The lines identified in the $\alpha pn$ spectrum were used as gates for a $\gamma-\gamma$ matrix. Figure 5.12 shows the spectra produced by gating on the three strongest lines in the $\alpha pn$ spectrum (figure 5.11). The same gates were then projected from the Argonne backed target $\gamma-\gamma$ matrix, and the coincidence spectra are shown in figure 5.14. It is noticeable that the 621, 1024, 1179, 1487 and 1748 keV transitions are absent from the Argonne data set. This is possibly due to the different population pattern associated with the $pn$ evaporation channel in the Argonne experiment, compared to the $\alpha pn$ channel in the PEX experiment.

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1This condition is used to detect the $\alpha pn$ channel with maximum statistics by accounting for protons missed by the Si-ball. This is reasonable since the $\alpha n$ evaporation channel is much less intense than the $\alpha pn$ channel in this data.
Figure 5.11: Gamma-ray channel selected spectra (a) with no channel condition, (b) αpn gated, (c) 2αp gated and (d) α2p gated. The lines labelled with energies in (b) are assigned to $^{62}$Ga, the channels (c) and (d) are indicating where they are contaminating the αpn spectrum.
Figure 5.12: Alpha,p,n gated $\gamma$-$\gamma$ coincidence spectra from the PEX data, gated by the three strongest lines identified in the $\alpha$pn gated PEX id-spectrum. (a) The sum of the three gates with all the lines observed in coincidence marked with energies, (b) 246 keV gate, (c) 376 keV gate and (d) 571 keV gate.
Figure 5.13: Alpha,p,n gated $\gamma$-$\gamma$ coincidence spectra for $^{62}$Ga from the PEX data, gated by (a) 1241 keV, (b) 621 keV and (c) 946 keV.
Figure 5.14: Gamma-gamma coincidence spectra from the AYEBALL backed target data, gated by the three strongest lines identified in the apn gated PEX id-spectrum. (a) The sum of the three gates with all the lines observed in coincidence marked with energies, (b) 246 keV gate, (c) 376 keV gate and (d) 571 keV gate.
5.4.2 Angular Distributions.

In order to ascertain the spin and parity of the levels in the decay scheme, the multipolarities of the transitions must be determined using angular distribution techniques. Gamma-ray energy spectra were incremented for the two clusters centred at 146° and for the two clusters centred at 105° for the PEX data, gated by 1α, 1 neutron and either 0 or 1 protons, and corrected for detection efficiency at that energy (figure 5.15).

![Figure 5.15: Expanded efficiency corrected spectra of the angular projections from the PEX data gated by 1 α, 1 neutron and either 0 or 1 protons showing lines in 62Ga. (top) From the clusters centred at 105° and (bottom) is from the clusters centred at 146°. The lines marked ‘c’ are contaminants from other channels, the line marked ‘e+e−’ is due to pair production in the target.](image)
Note, the pair production lines in figure 5.15 originate from an annihilation which occurs stationary in the laboratory frame of reference, so there is no Doppler shift associated with this line. It therefore appears in a different channel at different angles. Inspection of the spectrum for the 146° cluster indicates the 571 keV peak to have both a shifted and an unshifted component in its lineshape, suggesting a decay from a short lived (~ns) isomeric state. The unshifted γ-decays come from $^{62}$Ga recoils which have implanted in stopper foils placed in front of the forward direction silicon detectors in the charged particle detector [104] (see figure 5.16). The measured centroids of the peaks are separated by 6.4±0.9 keV. The target to silicon distance was 11.5±1.4 mm, and the average recoil velocity was experimentally determined by Doppler shift measurements for the detectors at the various angles to be 1.4±0.2% of the speed of light. This corresponds to a flight time for the recoils between exiting the target and being stopped in the silicon ball of 2.7±0.6 ns. The measured ratio of counts in the moving and stopped peaks for the 571 keV transition centered at 146° was 1.31±0.26, however, there was a 3 mm diameter beam exit hole in the silicon ball from which the stopped component is calculated to lose 34±10% of its counts. Correcting for this, a lifetime for the isomeric state of 4.6±1.6 ns was obtained. A check was made to test this method using the known isomeric state from the 2pn channel to $^{65}$Ge, where the $\frac{9}{2}^+$ state decays via a 326 keV E1 [151], with a lifetime of 7(1) ns. The peak in the 146° gated spectrum had a shoulder on the high energy side, indicating a stopped decay component, and the measured lifetime of this state, without accounting for any loss through the beam exit hole in the si-ball, was found to be 3.0±0.8 ns. The area of the two components of the peak could only be determined with large errors, and consequently, this is not a valid proof of the amount of recoil cone which was estimated to be lost through the beam exit hole, however even after accounting for the loss of the stopped decays, the measured value is consistent with the quoted lifetime, and is therefore a valid confirmation of the method.

An analysis of the peak shape of the 246 keV transition in the 146° gated PEX
spectrum shows that this line is also of a two-component nature with both a shifted and unshifted component, with a measured overall FWHM of 5.3±0.3 keV. This compares with a value for the 376 keV line in ⁶²Ga of 4.4±0.3 keV. Assuming a linear decrease in Doppler broadening with decreasing γ-ray energy, the expected FWHM for the 246 keV line at this angle is 3.7±0.3 keV. If the 571 and 246 keV lines decay from states with the same apparent lifetime, the energy difference between the shifted and unshifted components will have a simple linear dependence on γ-ray energy. Under this assumption the expected value for the energy difference between
the centroids of the shifted and unshifted components of the 246 keV line in the 146° spectrum is 2.8±0.3 keV. The expected peak widths of the stopped and shifted components of the 246 keV line are approximately 3.1 and 3.7 keV respectively, which assuming a similar shifted to unshifted intensity ratio as measured for the 571 keV peak, results in a single line, with a FWHM consistent with the measured value. In summary, the data suggest that the 246 and 571 keV transitions lie below an isomeric state.

By comparing the statistics in peaks in the two spectra shown in figure 5.15, a ratio of anisotropy can be obtained for the intense lines in ⁶²Ga as discussed in chapter 3. This ratio can then be compared with values obtained for known transitions in neighbouring nuclei to determine the multipolarity of the transition. Anisotropy values were also obtained for previously known transitions in ⁶²Zn [108, 117, 118, 119, 120] and ⁶⁵Ga [149, 150] from the appropriate matrices. Clear separation was shown between γ-rays of a stretched E2 and pure dipole nature with the weighted average for the measured anisotropies for stretched E2 transitions being 0.20±0.01, compared to -0.16±0.02 obtained for the ΔI=1 dipole decays. Since there were few E1 transitions, M1 transitions were chosen which had as small an E2 admixture as possible. These data are shown in figure 5.17 along with values obtained from the data for the three strongest lines identified in ⁶²Ga. The anisotropy obtained for the 376 keV line in ⁶²Ga was found to be 0.14±0.06, consistent with a stretched E2 decay. The values obtained for the 571 and 246 keV transition, which are composed of both stopped and shifted components were both consistent with isotropic decays, as expected for transitions being fed by an isomer.

A comparison of id-spectra from the AYEBALL thin target and the PEX data (figure 5.18) reveals an appreciable difference in intensity of both the 571 keV and the 246 keV lines relative to the 376 keV transition. The important difference in the experimental arrangements is that approximately 66(10)% of the recoils in PEX experiment remain inside the Si-ball, and thus in view of the γ-ray detectors, meaning any isomeric decays will be detected in this data set. In the AYEBALL experiment,
the recoils continue into the FMA and out of view of the $\gamma$-ray detectors, so an isomeric decay with a lifetime longer than the flight time through the target chamber will result in a loss of intensity for that detected transition. This difference in the spectra is consistent with the 571 keV and 246 keV transitions being fed by an isomeric state.

Figure 5.17: Anisotropy values obtained from efficiency corrected ratios of the intensities of transitions in $^{62}$Ga, $^{62}$Zn and $^{65}$Ga. These data are taken from the two pairs of cluster detectors centred at 105° and 146° to the beam line in the PEX array. The weighted average of stretched E2 transitions in $^{62}$Zn (954 and 1232 keV) and $^{65}$Ga (311 and 1096 keV), along with the weighted average of $\Delta I=1$ dipole decays at $\gamma$-ray energies of 191 and 885 keV in $^{65}$Ga and at 1197 keV in $^{62}$Zn are shown.
Figure 5.18: Comparison of id-spectra from (top) PEX data, gated by 1 α, 1 neutron and either 0 or 1 proton, and (bottom) AYEBALL thin target data, mass 62 and ion chamber gated with normalised subtractions of contaminants.
From the PEX experiment, it was established that the 376 keV transition had a value of anisotropy consistent with a stretched E2. Multipolarities for other transitions in $^{62}$Ga were then derived from a DCO analysis using the AYEBALL backed target data [73]. This involved constructing a $\gamma-\gamma$ coincidence matrix of events measured in detectors at 158° versus events in any of the 79°, 101° or 134° detectors. As discussed in the previous chapter, coincidences between previously assigned [123] stretched E2 (1039 keV, 1310 keV, 1317 and 1361 keV), $\Delta I = 1$ E1 (1066 keV and 1410 keV) and $\Delta I = 1$ M1 (340 keV) transitions from the strongly populated channel of $^{61}$Cu were used to verify this method. The weighted average $R_{dco}$ values for E2 transitions gated by E1 and E2 gates were found to be 1.60(3) and 0.95(3) respectively. Similarly, E1 transitions corresponding to essentially pure, $\Delta I = 1$, dipole decays were found to have weighted average $R_{dco}$ values of 0.92(6) and 0.62(2) when gated by E1 and stretched E2 transitions respectively. Figure 5.19 shows the DCO projection spectra gated by the 376 keV and the 571 keV transitions in turn. It is clear that whilst the 246 keV and the 1241 keV lines are the same intensity in the 376 keV gate, the 571 keV lines are different. Also, the 246, 376 and 1241 keV transitions gated by the 571 keV transition have the same relative intensities in each spectra, although they have different absolute intensities in the two projections.

Figure 5.20 shows the measured ratios for known transitions in $^{61}$Cu gated by a known E2 and also gated by a known E1. The same is done for the transitions in $^{62}$Ga, assuming an E2 multipolarity for 376 keV transition (from its measured anisotropy in the PEX data), and a $\Delta I = 1$ decay for the 571 keV transition from the 376 keV gated ratio. Table 5.1 shows the $R_{dco}$ values for gates set on the 571, 376 and 246 keV lines in $^{62}$Ga from the backed target AYEBALL data. The ratios obtained for the 246 keV and the 1241 keV transitions are both consistent with stretched quadrupole decays.
Figure 5.19: DCO projections from the AYEBALL data, efficiency corrected for gate and projection energy, projected from the backed target DCO matrix; (left) gated by the 376 keV transition and (right) the 571 keV transition. The top spectra are gated on the 79°+101°+134° axis and projected onto the 158° axis, while the bottom spectra are gated on the 158° axis and projected onto the 79°+101°+134° axis.
Figure 5.20: DCO ratio values for transitions in the backed target AYEBALL data, efficiency corrected for gate and projection energy. Open diamonds are quadrupoles and filled squares are dipoles. (a) Gated by the $^{13}_{2}^{+} \rightarrow ^{9}_{2}^{+}$ E2 1310 keV transition in $^{61}$Cu [122], with weighted averages for E2's and E1's in $^{61}$Cu, (b) gated by 376 keV transition in $^{62}$Ga identified as E2 from PEX anisotropy data, and point (x) gated by 246 keV transition, with same weighted averages as (a). (c) Gated by the $^{9}_{2}^{+} \rightarrow ^{7}_{2}^{-}$ E1 1410 keV transition in $^{61}$Cu [122], with weighted averages for E2's and E1's in $^{61}$Cu, (d) gated by 571 keV transition in $^{62}$Ga with same weighted averages as (c).
5.4.3 Decay Scheme for $^{62}$Ga.

The partial decay scheme deduced for $^{62}$Ga from this work is shown in figure 5.21. Beta decay studies of $^{62}$Ga by Davids et al. [37] have found that the ground state in $^{62}$Ga decays to the $T=1, I^e=0^+$ ground state in $^{62}$Zn with a lifetime of $116.34\pm0.34$ ms, corresponding to an $ft$ value of $3081\pm47$ s, or a log$ft$ of 3.49. This is indicative of a Fermi superallowed decay, suggesting that a $T=1, I^e=0^+$ ground state exists in $^{62}$Ga. The decaying state was deduced to be the ground state since all the decay intensity was accounted for in the experiment [37]. Spin and parity assignments in the current work are made from anisotropy and DCO measurements, and assuming the lowest level is the $I^e=0^+$ ground state. Tentative transitions are shown as dashed lines, and a list of transition energies, intensities and DCO ratio’s are given in table 5.1.

Table 5.1: Transitions identified in $^{62}$Ga in the current work. The transition energies and $R_{dco}$ values are taken from the AYEBALL backed target data.

<table>
<thead>
<tr>
<th>$E_\gamma$ (keV)</th>
<th>Intensity</th>
<th>$E_i \rightarrow E_f$ (keV)</th>
<th>$I_i \rightarrow I_f$</th>
<th>$R_{dco}$ (376 gate)</th>
<th>$R_{dco}$ (571 gate)</th>
<th>$R_{dco}$ (246 gate)</th>
</tr>
</thead>
<tbody>
<tr>
<td>246.3</td>
<td>213(20)</td>
<td>818$\rightarrow$571</td>
<td>$(3^+) \rightarrow (1^+)$</td>
<td>0.98(10)</td>
<td>1.52(11)</td>
<td></td>
</tr>
<tr>
<td>376.4</td>
<td>180(20)</td>
<td>1194$\rightarrow$818</td>
<td>$(5^+) \rightarrow (3^+)$</td>
<td>1.87(33)</td>
<td>0.97(16)</td>
<td></td>
</tr>
<tr>
<td>571.2</td>
<td>225(20)</td>
<td>571$\rightarrow$0</td>
<td>$(1^+) \rightarrow (0^+)$</td>
<td>0.66(8)</td>
<td>0.63(5)</td>
<td></td>
</tr>
<tr>
<td>621.4</td>
<td>10(5)</td>
<td>1439$\rightarrow$818</td>
<td>$(4^+) \rightarrow (3^+)$</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>946.3</td>
<td>58(10)</td>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>1108.3</td>
<td>21(3)</td>
<td>6846$\rightarrow$5738</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>1178.5</td>
<td>12(3)</td>
<td>2372$\rightarrow$1194</td>
<td>$(6^+) \rightarrow (5^+)$</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>1241.3</td>
<td>136(10)</td>
<td>2435$\rightarrow$1194</td>
<td>$(7^+) \rightarrow (5^+)$</td>
<td>1.06(15)</td>
<td>1.24(33)</td>
<td>1.07(13)</td>
</tr>
<tr>
<td>2356.3</td>
<td>61(8)</td>
<td>(4792$\rightarrow$2435)</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>
Figure 5.21: The partial decay scheme for $^{62}$Ga derived from this work. Tentative transitions are shown as dashed lines, spin parity assignments made assuming a $0^+$ ground state.
The 571 keV and 246 keV transitions are the most intense, and are in coincidence with all the other observed transitions. They are consequently placed as the lowest two transitions below the 4.6 ns isomeric level. The 621 keV is notably present in the 246 keV gate (see figure 5.12-b), thus preventing them from being positioned parallel to one another, with the 376 keV transition above the 246 keV transition to form two branches between the same levels. In fact the 621 keV transition is not in coincidence with the 376 keV line (figure 5.12-c and 5.13-b) so are both placed feeding into the same 818 keV level above the 246 keV transition. Note that while the intensities of the 571 keV dipole and the 246 keV stretched quadrupole are the same within experimental errors, there are a number of arguments which strongly favour the placement of the 571 keV line below the 246 keV transition. On the basis of transition rate arguments, the measured lifetime of the excited state at 818 keV of 4.6±1.6 ns, corresponds to a transition strength of 7.3±2.5×10⁻⁷ Wu or 1.0±0.4×10⁻⁸ Wu for a 571 keV E1 or a pure M1 respectively. While both are possible in principle, these values are very small compared with other measured transition strengths for similar multipolarity decays in this mass region [29]. However, the corresponding value of 13.5±4.7 Wu for a 246 keV stretched electric quadrupole decay is very typical for the region [29]. Also, as discussed in the next chapter, shell model calculations performed on this nucleus reproduce the B(E2) value for the isomeric decay, supporting the argument that it is the 246 keV, and not the 571 keV transition which depopulates the 818 keV isomeric level.

The ordering of the γ-rays above the isomer comes from their measured intensities in backed target AYEBALL data, normalised to the measured intensities of the 246, 376 and 571 keV lines in the PEX data, open gated spectrum. The 1241 keV is the next most intense line in the spectrum (figure 5.11) and is in coincidence with the 571, 246 and 376 keV lines (figure 5.13-a), seeing them with equal intensity. It is therefore assumed to be feeding directly from above the 1194 keV level. The 1179 keV transition is not present in the 1241 keV gate, (figure 5.13-a) but is in the 376 keV gate (figure 5.12-c). It is therefore also placed feeding the 1194 keV level. In
the 571, 246, 376 and 1241 keV gates, there are also lines of decreasing intensity at energies of 2356, 946 (figure 5.13-c) and 1108 keV respectively. These three transitions are all in mutual coincidence and are thus placed in a cascade above the 1241 keV transition with the more intense 2356 keV transition feeding the 2435 keV level. The ordering of the 2356 and 946 keV transitions is deduced from intensity measurements from both sets of data, but with the poor statistics at this high spin, the ordering of these two transitions is tentative.

There are also three other lines present in the γ-γ spectra, 1024.1(1), 1386.5(10) and 1747.8(10) keV, but they are weakly populated, and so while they can be definitely identified as transitions in $^{62}$Ga, they are not positioned in the decay scheme. All three are present in the 571 and 246 keV gates (figures 5.12-d and b respectively).
Chapter 6

Conclusions.

6.1 Discussion of $^{62}$Ga Decay Scheme.

The decay scheme of $^{62}$Ga derived from this work is shown if figure 5.21. The 571 keV transition is consistent with a pure dipole character from angular correlation measurements using both the AYEBALL backed target, and PEX data. The next three transitions above the 571 keV are assigned stretched E2 multipolarity, and an isomeric state with a mean-lifetime of $4.6 \pm 1.6$ ns was identified at an excitation energy of 818 keV.

There are many isomers identified in nuclei in this mass region [158] so checks were made to see if any decays from previously unobserved isomeric states could be found. The beam bursts of ATLAS were separated 82 ns, and the electronic coincidence acceptance time (master gate) was $1 \rightarrow 2 \mu s$. The recoils were stopped in the backed target, but passed through the thin target in a matter of picoseconds. Consequently, any isomeric decay with a lifetime of a few ns or more would be in the field of view of the germanium detectors in the backed target AYEBALL data set, but not in the thin target data set. The intensity differences observed in the $\gamma-\gamma$ coincidence spectra from both AYEBALL experiments, and the PEX experiment, are all consistent with the 571 keV and the 246 keV transitions being fed by an isomer, but no other differences were noted. It was therefore concluded that there
were no other isomers of lifetime up to \( \sim 1 \mu s \). A study of isomeric states in this mass region by Grzywacz et al. [158] at the GANIL laboratory, Caen, France using the LISE3 spectrometer [159] did not observe any isomeric states in \(^{63}\text{Ga}_{31}\) in the time range 100 ns \( \rightarrow \sim 500 \mu s \).

The beta decay study by Davids et al. [37] suggest a \( T=1, I^e = 0^+ \) ground state for \(^{63}\text{Ga}\). This previous study was also populated with a heavy ion fusion evaporation reaction, suggesting a population decay via yrast or near yrast states. Since there are no other likely candidates for isomers, we assume the lowest lying state that we observe to be this \( 0^+ \) ground state. Note that if this were the case, the 571 keV dipole transition would correspond to a \( 1 \rightarrow 0^+ \) decay. Given the difficulty of generating low-lying negative parity states from the available single particle orbitals, this would be a pure magnetic dipole, with no E2 admixture. The experimentally derived DCO ratio for this transition when gated by an E2 transition is 0.66\( \pm 0.08 \), which is consistent with this picture. The spin and parity assignments of the 571, 818, 1194 and 2435 keV levels are then made from DCO data with this assumption.

Both experiments were sensitive to \( \gamma \)-rays in the energy range \( \sim 80 \text{ keV} \rightarrow 8 \text{ MeV} \). While it is not expected that there would be transitions above 8 MeV, it is possible for there to be some below 80 keV. It is therefore feasible that there is an unobserved low energy (< 80 keV) transition below the 571 keV transition which feeds the \( 0^+ \) ground state.

Following the logic described by Rudolph et al. [44] regarding the isospin assignment of the band he observed in \(^{74}\text{Rb}\), the decay scheme of \(^{63}\text{Ga}\) deduced from this work bears little resemblance to that of the \( T_z =-1 \) isobar \(^{62}\text{Zn}\), as shown in figure 6.1. This is not surprising since due to its low deformation, the first excited state in \(^{62}\text{Zn}\) lies at 954 keV so that, if the assumption of the direct feeding of a \( 0^+ \) ground state for \(^{63}\text{Ga}\) is correct, the first \( 2^+ \), \( T=1 \) state is likely to be non-yrast. The decay scheme is therefore interpreted as a cascade of stretched E2's feeding a \( T=0 I^e=1^+ \) bandhead, which decays directly to the \( T=1 0^+ \) ground state via a 571 keV pure magnetic dipole.
Figure 6.1: Comparison of the low lying energy levels of shell model calculations [160], and IBM4 calculations [161] for $^{62}$Ga and the levels derived for $^{62}$Ga from the data. The low lying $T=1$ levels of the $T_z=1$ isobar are shown for comparison. Spin and parity of the levels are labelled on the right, energy on the left. For the IBM-4 levels, the number of $T=0$ bosons as a percentage of the total is indicated.
6.2 Parallel Studies.

This region of the Segré chart is the focus of much experimental and theoretical interest at present. Two other experimental data sets which include this nucleus are known to be currently under analysis, with results presented in conference proceedings. The work by de Angelis et al. [162] populated states in $^{62}$Ga using the reaction $^{40}$Ca($^{32}$S,2αpn)$^{62}$Ga at 140 MeV at the Legnaro National Laboratory. Reaction channels were selected using the Si-ball ISIS [163] and an 80 element inner BGO multiplicity filter. Coincident $\gamma$-rays were detected using the GASP array [164] of 40 Compton suppressed, high efficiency HPGe detectors. Transitions in $^{62}$Ga were identified using the Si-ball to gate on charged particle emission, and the BGO multiplicity filter to determine $\gamma$-ray fold, and hence neutron evaporation. Gates were then set on the 246 and 376 keV transitions on a 2α gated $\gamma$-$\gamma$ matrix, and the resulting identification spectrum is presented in reference [162]. It is believed that there was no way of identifying the isomer, and consequently no decay scheme is presented. However, the spectrum presented identifies the 246, 376, 571, 946, 1108, 1241 and 2355 keV transitions, and is consistent with figure 5.12-a.

A further result has been submitted as an abstract to a forthcoming conference by Skoda et al. [165]. The data come from an experiment performed using the GAMMASPHERE array [76] in conjunction with the MICROBALL [166] and an array of 15 neutron detectors. The reaction used was $^{40}$Ca($^{38}$Si,αpn)$^{62}$Ga at 125 MeV, and a decay scheme was derived from triple coincidence events, which agrees with the yrast cascade of the scheme produced in this work (figure 5.21). The 621, 1024 and 1179 keV transitions are not observed, but the 1487 keV transition is placed feeding the 2435 keV level, with an 867 keV transition above it, forming a parallel branch to the 4792 keV level. Also, the 1748 keV transition is placed on top of the 1108 keV, with a 1388 keV transition above it in a cascade. Angular distributions are quoted as being used to assign multipolarities, but some of the results differ from those obtained in this study. The 571 keV transition is assigned as an M1, and the 1241 as an E2 in agreement with this work, but the 246 and 376 keV transitions are
interpreted as M1/E2 decays, rather than stretched electric quadrupoles. The whole cascade is then placed feeding a 2 state at an excitation energy of not more than 10 keV above the ground state. This assignment is based on extrapolation of level energies in odd-odd Ga isotopes, and the proposed level structure is quoted as being in good agreement with systematics. In the abstract submitted to the conference [165], there is no space to show any data to support these spin assignments, and so the following interpretation of the decay scheme of 62Ga is based on the evidence presented in this thesis.

6.3 Theoretical Interpretation.

The observed levels in 62Ga and the lowest T=1 states in 62Zn have been interpreted in a shell model type analysis [160]. There are two natural shell-model spaces that can be adopted for 62Ga, one consists of the entire pf shell with a 40Ca core and the second of the pf5/2g9/2 orbits with a 60Ni core. We have found that the latter gives superior results and these are shown in figure 6.1. The effective interaction is a realistic G matrix whose monopole part has been phenomenologically adjusted and which has been used previously in 76Ge and 82Se [167]. Calculations are done with the shell-model code ANTOINE [168]. The difference between the T=0 and T=1 states is well reproduced, as is the band structure on top of the 1 state. The shell model gives levels with spin 9+, 11+ and 13+ at energies of 4728, 5339 and 6249 keV, respectively, which are close in energy to the levels observed above the 7 state. Moreover, as shown in table 6.1, the calculated value of B(E2;3+ → 1+) agrees well with that deduced from the measured half life of the 818 keV level.

Table 6.1: Values of B(E2;3+ → 1+) in e²fm⁴

<table>
<thead>
<tr>
<th>Shell Model</th>
<th>Shell Model</th>
<th>Experiment</th>
</tr>
</thead>
<tbody>
<tr>
<td>(pf5/2g9/2)</td>
<td>(pf5/2)</td>
<td></td>
</tr>
<tr>
<td>139</td>
<td>89</td>
<td>152(60)</td>
</tr>
</tbody>
</table>
To shed further light on its structure, and to pave the way for heavier $N \approx Z$ nuclei where a shell model treatment becomes intractable, the states of $^{62}$Ga have been studied in the framework of the Interacting Boson Model [169, 170]. For $N \approx Z$ nuclei, bosons with $T_z = 0$ (corresponding to correlated np pairs) should be included, and for odd-odd nuclei, both $T=0$ and $T=1$ bosons are of importance for the low-energy spectrum. A version of the IBM that includes these bosons has been proposed by Elliott and Evans [171] and is referred to as IBM-4 [172]. The bosons in IBM-4 are assigned orbital angular momenta $L = 0$ ($s$) and $L = 2$ ($d$), and they now carry spin and isospin labels $S$ and $T$ with the combinations $(ST) = (01)$ and $(10)$ being retained. This choice is motivated by the short-range character of the residual nucleon-nucleon interaction which energetically favours two-nucleon configurations which are spatially symmetric and hence antisymmetric in spin-isospin. The resulting total angular momenta of the bosons in IBM-4 are thus $J = 0, 2 (T=1, T_z = 0, \pm 1)$ and $J = 1^2, 2, 3 (T=0)$; the bosons are denoted as $^{2S+1}LJ$, and hence include $^1s_0$, $^1d_2$, $^3s_1$, $^3d_1$, $^3d_2$ and $^3d_3$.

A calculation for $^{62}$Ga has been performed by Van Isacker et al. [161] which represents the first full diagonalisation of an IBM-4 Hamiltonian. While full details of the calculation are beyond the scope of this thesis, they are presented in reference [160]. The method used follows the analysis of Halse et al. who studied even-even [173] and odd-odd [174] sd-shell nuclei in the context of IBM-4. In this approach, boson energies and interaction matrix elements are first derived from the shell model. To simplify the mapping procedure in this first application of the method, the model space is restricted to $p/(2p)$ and the appropriate single-particle energies and interaction parameters of the modified surface delta interaction (MSDI) are used [139]. The boson energies are determined from the In-1p nucleus $^{58}$Cu, taking $^{56}$Ni as the core.

The IBM-4 results for $^{62}$Ga are compared with the MSDI $(fp)$ shell-model calculation in figure 6.1. This represents a first confirmation of the validity of the IBM-4 insofar that a good agreement is obtained between the two sets of calculations for low angular momentum states. The IBM-4 approximation worsens as the angular
momentum of the states increases because of the absence of high angular momentum bosons. For example, the yrast $6^+$ at 2371 keV in the shell model presumably has only a very small overlap with the yrast $6^+$ at 4048 keV in the IBM-4.

Not surprisingly, both the MSDI, restricted-basis shell model calculation and the IBM-4 approximation derived from it do not match the quality of agreement with the data offered by the $fpg$ shell model calculation. Indeed, judging by the predicted $B(E2)$ values and the regularity and spacing of the $1^+$ band, it would seem that, even at this low mass, the $g_{9/2}$ intruder orbital is important in providing the required degree of collectivity. Nevertheless, with these caveats, it is now possible to use the IBM-4 calculation to analyse the pair structure of states in $^{62}$Ga. As an example, shown in figure 6.1 for each state is the expectation value of the number of $T=0$ bosons as a percentage of the total. While the $T=0$ states in $^{62}$Ga are clearly dominated by $T=0$ bosons, the calculation predicts that this dominance only extends to the level of roughly two-thirds, the remaining pairs have $T=1$. However, the value calculated for the ground state of 23% $T=0$ cannot be true since the $ft$ value of the beta decay is strongly dependent on this mixing, and the measured value would suggest at most, only a few percent $T=0$.

In summary, the calculations predict that the $T=0$ states lie 386 keV above the $T=1$ ground state in $^{62}$Ga. This compares with an experimental gap between the first two observed states of 571 keV. This can be viewed in the context of the $T=0$ ground state in $^{58}$Cu and the higher lying (1006 keV) first observed $T=0$ state in $^{74}$Rb. Thus a picture is beginning to emerge of a gradual rise in the $T=0$ states as mass increases along the N=Z line in this shell, this rise being accompanied by a gradual reduction in the contribution of $T=0$ pairs. However, the beta decay of the N=Z=39 nucleus $^{78}$Y has recently been studied at the Argonne National Laboratory [43] using the reaction $^{40}$Ca($^{40}$Ca,pn)$^{78}$Y at 125 MeV. The beta decaying state was determined to be a $5^+$ $T=0$ state, with an excitation energy of not more than 500 keV above the $T=1 I^+=0^+$ ground state. It is hence concluded that the $T=1 np$ pairing is quenched compared to the next lowest N=Z odd-odd nucleus $^{74}$Rb.
Clearly, the degree of collectivity is limited in $^{62}$Ga, and more information on non-yrast states would be necessary to distinguish the presence of any degree of energy gap associated with coherent n-p pairing. Finally, while the full shell model calculation gives a very satisfactory result in this nucleus, such a method will become increasingly intractable as the number of valence nucleons, and hence collectivity, increases along the N=Z line. For this reason, the success of the first numerical application of the isospin invariant IBM formalism in reproducing its, albeit limited, shell model origins is particularly encouraging. It offers the hope that this method of truncating the problem can be applied in the near future to examine the role of isospin in the development of collective structure along the length of the N=Z line.

6.4 Future Prospects.

The beta decays of the next six odd-odd N=Z nuclei [40, 41, 42] up to $^{86}$Tc have identified $T=1$ ground states, paving the way for further structure investigations of this type. One transition is already known in the N=Z=43 nucleus $^{86}$Tc [175], although its position in a decay scheme is unknown. The isomer study by Grzywacz et al. [158] also observed the beta decay of $^{66}$As. A decay scheme is deduced with states of spin either $1^+$ or $2^+$ placed at energies 837 keV and 963 keV above the $0^+ T=1$ ground state [176]. The first excited state in the $T_x=1$ isobar $^{66}$Ge is at 957 keV, possibly implying that the 963 keV state has $T=1$, and the 837 keV state may be $T=0$. However, this is all supposition, especially when there is no firm experimental evidence that either of these transitions feed directly into the ground state. Therefore, an experiment has been performed at the Oak Ridge National Research Facility, using an array of five clover detectors and five single crystal 25% efficient detectors in conjunction with a recoil mass separator and ion-chamber to study the excited states of $^{66}$As via the reaction $^{40}$Ca($^{32}$S,$\alpha$p$n$) $^{66}$As. The analysis of this data is currently in progress, and it is hoped that this type of study will yield more information about the phenomenon of n-p pairing and isospin mixing in this revealing region of the table of isotopes.
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7. *An Yrast Study of the A=61 Isobars \(^{61}\text{Cu}\) and \(^{61}\text{Zn}\).*


8. *High Spin States in \(^{58}\text{Ni}\).*


9. *Discrete Line Spectroscopy of \(^{62}\text{Ga}\) with AYEBALL.*


10. *High Spin Studies of \(^{62}\text{Zn}, \(^{61}\text{Zn}, \(^{61}\text{Cu}\) and \(^{58}\text{Ni}\).*

ORAL PRESENTATIONS

1. The AYEBALL Gamma-ray Spectrometer on the Argonne Fragment Mass Analyser, North West Europe '96 Institute Of Physics Conference, Amsterdam, April 1996


3. Experimental Studies of N~Z Nuclei in the Mass 60 Region, EPSRC Postgraduate Summer School, University of Nottingham, September 1997.

4. First Observation of Excited States in the N=Z=31 Nucleus $^{62}_{31}$Ga using AYEBALL and PEX, Invited Seminar at Notre Dame University, IN. Aug 1997

