

The Beta Decay of $^{79,80,81}\text{Zn}$ and Nuclear Structure around
the N=50 Shell Closure

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Dedication

To my father, who has always given me support in my academic pursuits and has never doubted my abilities.

Abstract

This dissertation reports on new information in the β^- [beta minus] decay of the neutron-rich nucleus ^{81}Zn , which populates states in its daughter nucleus ^{81}Ga . This includes new γ [gamma]-ray transitions in the daughter nucleus, ^{81}Ga , as well as a β [beta]-delayed neutron branching ratio. This isotope was produced at the Holifield Radioactive Ion Beam Facility of Oak Ridge National Laboratory through the Isotope Separation Online technique. They are fission fragments from proton-induced fission on a uranium carbide target. These fission fragments are ionized and both mass and isotopically separated before arriving at the Low Energy Radioactive Ion Beam Spectroscopy Station (LeRIBSS). The γ [gamma]-ray and β [beta] electron emissions from the decays are measured and analyzed in this work.

A new β [beta]-delayed neutron branching ratio is reported for this decay, which is in agreement with recent theoretical values. The core excited states in the daughter nucleus, ^{81}Ga , populated through allowed Gamow-Teller decays are analyzed. A trend in core excited states with other N=50 isotones indicates an increasing gap between a deeply bound neutron hole and the valence neutron above the N=50 gap upon moving towards doubly magic ^{78}Ni .

This dissertation also reports on additions to the decay schemes of ^{79}Zn and ^{80}Zn decays. Their decay level schemes have been expanded upon and an improved picture of the total allowed Gamow-Teller decay strength is known from ^{79}Zn to ^{81}Zn . This work presents an improved, albeit still incomplete, picture of the energy of states populated through Gamow-Teller decays from below to above the N=50 shell gap in zinc isotopes.

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1 Introduction

The atomic nucleus is a system of nucleons, protons with positive electrical charge and neutrons which are neutral, bound together by the strong nuclear force. Nuclear physics is the study of the structure and physical behaviors of atomic nuclei. Among more than 3000 known nuclei most nuclei are radioactive, among them 168 are stable and form the so-called valley of stability between atomic numbers $Z=1$ and $Z=83$.

For nuclei with relatively few protons and neutrons the strong force favors the formation of stable nuclei with an approximately equal number of protons and neutrons. However, for most nuclei, those with atomic number $Z \geq 20$, the stable nuclei are those with more neutrons than protons. This is because more neutrons are required to balance the Coulomb repulsion of the protons. Nuclei which are not stable will decay to nuclei closer to stability until eventually a stable nucleus is reached. The known decay modes of the nuclear ground states include atomic electron capture, β decay, α particle and cluster emission, β -delayed neutron(s) or proton(s) emission and two proton emission. Through observation of these decays it is possible to study nuclear structure in the resultant, or so-called daughter, nucleus.

One of the major goals in nuclear physics is to understand what determines the structure of atomic nuclei and how and why this structure changes with varying number of protons and neutrons. An alternative to nuclear structure studies through experimental decay are the methods using reactions between nuclei. This can also provide insight into how the nucleons are interacting with each other using, for example, transfer or fusion reactions. The advantage of reaction studies is that there is a possibility of determining what nucleus is to be studied and which of its properties will be studied through an appropriate choice of the nuclear reaction. However, this is at the cost of increased experimental complexities compared to decay spectroscopy experiments.

Another essential part of nuclear physics research is nuclear theory, which is concerned with generating models of the atomic nucleus with better predictive power of the experimental observables. These models are usually based on fundamental principles. One of the current trends is the development of nuclear models that are able to derive properties of the nuclei from the free nucleon-nucleon interactions, for example when nucleons scatter with each other outside of the nuclear medium.

In this work we present the results of the experimental studies of nuclei near doubly magic ^{78}Ni by using decay spectroscopy methods. We will relate these results to those obtained with

the use of nuclear reaction experiments. Finally, we provide interpretation of the data using theoretical calculations based on a contemporary nuclear model.

The nucleus ^{78}Ni is a neutron-rich nucleus with traditional “magic” numbers of nucleons. It consists of 50 neutrons and 28 protons. Nuclei with magic proton/neutron numbers have higher proton/neutron separation energies than their neighboring nuclei. The magic numbers are the result of the emergence of the shell structure of nuclei, as will be explained in the following section. Because one needs to understand how and if the shell structure changes away from stability, and hence if the magic numbers change, there is a strong motivation for studying nuclei which are both away from stability and contain traditional magic numbers of nucleons. Studies of nuclei with only one proton or one neutron outside of a doubly magic nucleus are particularly interesting because they may reveal the simplest excitation modes of a single “valence” nucleon in the mean potential, which is generated by the doubly magic nucleus. This is analogous to electron orbitals in alkali atoms. These data can provide important information on the persistence of nuclear magicity away from stability.

The nucleus ^{81}Zn and its decay product ^{81}Ga , which are the main subjects of this work, are both examples of such systems. The nucleus ^{81}Zn has 51 neutrons and 30 protons. In this study we provide the experimental data [1] which will determine the spin of its ground state. This was controversial and used previously [2] to make claims on the effects of shell evolution.

The nucleus ^{81}Ga has 50 neutrons and 31 protons. The excited states populated in this nucleus reveal the information on various modes of excitations involving the valence protons and excitations of the doubly magic core. These experimental data were combined with the experimental systematics of other $N=50$ and $N=51$ nuclei.

Finally, the decays of $N=49$ ^{79}Zn and $N=50$ ^{80}Zn are also discussed in this work. The decay level schemes for the daughter gallium isotopes are expanded to higher energies in comparison to previous literature. These expansions provide an improved, although still incomplete, picture of the energetics of decays of deeply bound neutrons from below to above the $N=50$ shell gap.

2 Theoretical Background

2.1 Shell model of the atomic nucleus

In the quest for scientific understanding, a model is necessary for describing the behaviors of a particular system. A model is a good description of the system if it consistently predicts behaviors and quantities which are verified to be true in experiments. In this case the quantum mechanical system of an atomic nucleus is the system to be probed. In nuclear physics the shell model is an elementary model in which many predictions are verified with experimental evidence. The following are derivations and observations from experiments which explain and support this model.

2.1.1 Central potential

Combined effects of nucleon-nucleon interactions generate a mean field which form a potential well [3]. The protons and neutrons reside in this potential. Approximate methods are used when solving for the energies and configurations of states in which an atomic nucleus exists. One example of a nuclear potential is the Woods-Saxon potential, shown in figure 1. The equation for this potential, as a function of radial distance r , is the following:

$$V(r) = \frac{-V_o}{1 + \exp(\frac{r-R}{a})} \quad (1)$$

Where V_o is the maximum depth of the potential well, r is the radial distance, R is the nuclear radius, and a is the value for the surface diffuseness. The nuclear radius $R = r_o A^{1/3}$, where A is the number of nucleons (mass number) in the nucleus. For short distances nucleons exist in a deep, attractive well.

The nuclear shell model provides a framework in which one analyzes nuclear structure and the population of various states in atomic nuclei. The presence of this central potential well gave rise to formulation of the nuclear shell model, analogous to the atomic shell model. The emergence of nuclear shells and therefore magic numbers can be illustrated by the simple example of the radial potential with $V=0$ at r less than R_o and $V = \infty$ at r greater than R_o , analogous to the one-dimensional infinite square well potential. This means that the nucleons cannot exist beyond this radius R . It also means that the wave function for the nucleus must be zero at this boundary where $r = R$. Using such a potential yields the following Schrödinger equation:

$$(\frac{d^2}{dr^2} + \frac{2m}{\hbar^2}(E_{n\ell} - \frac{\hbar^2 \ell(\ell+1)}{2mr^2}))u_{n\ell}(r) = 0 \quad (2)$$

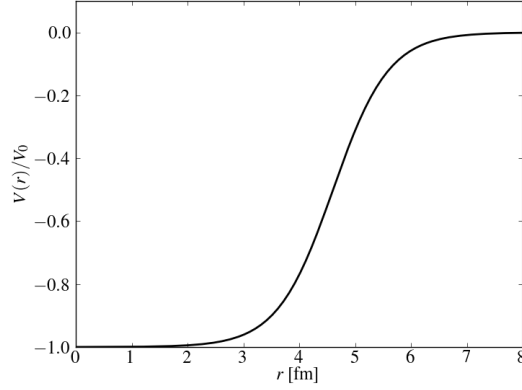


Figure 1: Woods-Saxon potential.

Where the solutions, $u_{n\ell}$, are spherical Bessel functions. Thus $u_{n\ell} = j(k_{n\ell}r)$, and

$$k_{n\ell} = \sqrt{\frac{2mE_{n\ell}}{\hbar^2}} \quad (3)$$

The n is the principal quantum number and the ℓ is the orbital angular momentum number. The n number can take on values of 1, 2, 3, etc. and ℓ can be 0,1,2,3,etc. The boundary condition that the wave function $u_{n\ell}$ must be zero at $r=R$ leads to quantization of energy levels. As each level is filled with protons or neutrons, this leads to the following magic numbers, or numbers of nucleons filling each closed shell:

2, 8, 18, 32, 50, etc.

The degeneracy of a level originates from the Pauli principle, where no two nucleons can have the same set of quantum numbers.

Here some of the empirical nuclear magic numbers such as 2, 8, and 50 are produced but not the numbers 20, 28, and 126. But another central potential may be attempted, that of the three dimensional harmonic oscillator potential. Here the potential which the nucleons experience is:

$$V(r) = \frac{1}{2}m\omega^2r^2 \quad (4)$$

Where m is the mass of the particle, ω is the frequency of oscillation, and r is radial distance. This yields the following Schrödinger equation:

$$\left(\frac{d^2}{dr^2} + \frac{2m}{\hbar^2}\left(E_{n\ell} - \frac{1}{2}m\omega^2 r^2 - \frac{\hbar^2 \ell(\ell+1)}{2mr^2}\right)\right)u_{n\ell}(r) = 0 \quad (5)$$

The solution of this differential Schrödinger equation depends upon Laguerre polynomials L:

$$u_{n\ell} \propto \exp\left(-\frac{m\omega r^2}{2\hbar}\right) r^{\ell+1} L_{\frac{\ell+1}{2}}^{\frac{n+\ell-1}{2}}\left(\sqrt{\frac{m\omega}{\hbar}}r\right) \quad (6)$$

The energy of each level with this harmonic oscillator potential is:

$$E_{n\ell} = \hbar\omega\left(2n + \ell - \frac{1}{2}\right) \quad (7)$$

The degeneracy in each level from using this potential is:

$$n_{\Lambda} = (\Lambda + 1)(\Lambda + 2) \quad (8)$$

Where the number Λ relates to the principle and orbital angular momentum quantum numbers as:

$$\Lambda = 2n + \ell - 2 \quad (9)$$

The “magic numbers” produced from using this potential are: 2, 8, 20, 40, 70, etc. Again, some of the real nuclear magic numbers are reproduced for the filled shells using this central potential, but not all of them.

2.1.2 Spin-orbit splitting

In the 1940’s Maria Goeppert Mayer, Hans Jensen, and others realized that a central potential alone was not reproducing the experimental magic numbers of nucleons in filled shells [4]. They realized they must use a spin-orbit interaction as well, just as is done in atomic physics with the potential experienced by electrons. The total potential felt by the nucleons would take the following general form:

$$V = V(r) - f(r)\vec{L} \cdot \vec{S} \quad (10)$$

Where V is the total potential, $V(r)$ is the central potential part, $f(r)$ is an arbitrary function based upon the radial distance, $\vec{L}$ is the orbital angular momentum vector, and $\vec{S}$ is the intrinsic

spin angular momentum vector. The total angular momentum operator $\vec{J}$ is given as the addition of the orbital and intrinsic spin operators:

$$\vec{J} = \vec{L} + \vec{S} \quad (11)$$

Therefore the derivation of the $\vec{L} \cdot \vec{S}$ term is as follows:

$$\vec{J}^2 = \vec{L}^2 + \vec{S}^2 + 2\vec{L} \cdot \vec{S} \quad (12)$$

Therefore:

$$\vec{L} \cdot \vec{S} = \frac{1}{2}(\vec{J}^2 - \vec{L}^2 - \vec{S}^2) \quad (13)$$

Because the orbital angular momentum and intrinsic spin angular momentum operators commute, the order of the multiplication above is of no importance. The expectation value of this spin-orbit multiplication is then:

$$\begin{aligned} \langle \vec{L} \cdot \vec{S} \rangle &= \langle \frac{1}{2}(\vec{J}^2 - \vec{L}^2 - \vec{S}^2) \rangle \\ &= \frac{\hbar^2}{2}[j(j+1) - \ell(\ell+1) - s(s+1)] \\ &= \frac{\hbar^2}{2}[j(j+1) - \ell(\ell+1) - \frac{3}{4}] \end{aligned}$$

Because $s = \frac{1}{2}$ for each nucleon. For $j = \ell + \frac{1}{2}$

$$= \frac{\hbar^2}{2}\ell$$

and for $j = \ell - \frac{1}{2}$

$$= -\frac{\hbar^2}{2}(\ell+1)$$

Thus the orbitals labeled with quantum numbers n and ℓ have split degeneracy with the introduction of the spin-orbit term into the potential. The new orbitals are labeled with the quantum numbers n , ℓ , and j , where n is the principle quantum number, ℓ is the orbital angular momentum, and j is the total angular momentum taking into account both ℓ and s , which is always $\frac{1}{2}$ for nucleons. This splitting of degeneracy and the orbitals in this nuclear shell model with the inclusion of the spin-orbit potential are shown in figure 2.

The energies of the new, split levels with respect to the degenerate levels are as follows:

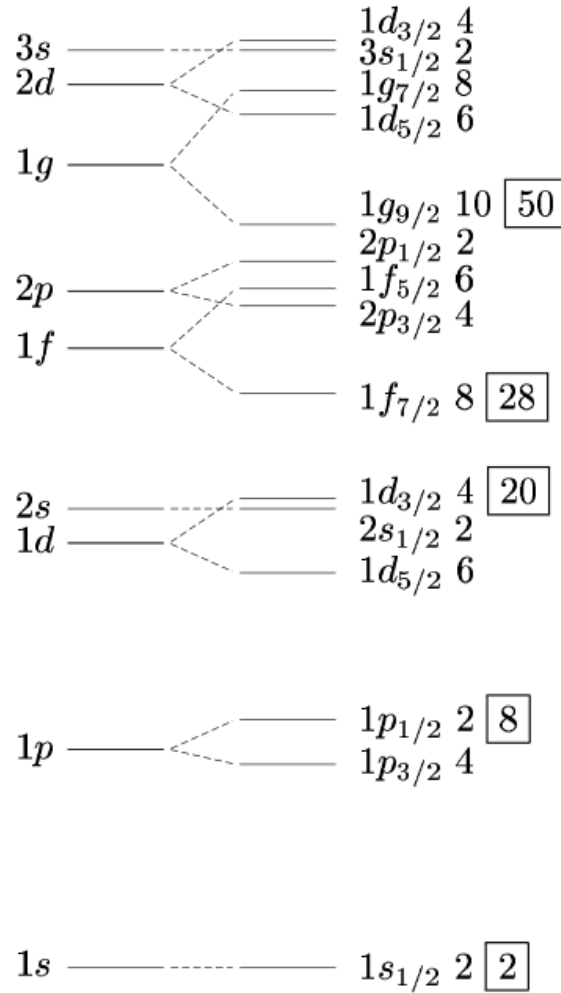


Figure 2: Shell model of the atomic nucleus. The splitting of degeneracy upon introduction of a spin-orbit term in the potential is shown. The final orbitals on the right are denoted in an $n\ell j$ notation, standing for principal quantum number n , orbital angular momentum ℓ in traditional spectroscopic letters s,p,d,f,g,etc. and the total angular momentum j of the orbital. The amount of nucleons, protons or neutrons, which can reside in each state is given by $2j + 1$.

$$E_{n\ell}(j = \ell + \frac{1}{2}) = -\frac{\hbar^2 \ell}{2} \int d^3r |\Psi_{n\ell}(\vec{r})|^2 f(r) \quad (14)$$

And

$$E_{n\ell}(j = \ell - \frac{1}{2}) = \frac{\hbar^2(\ell + 1)}{2} \int d^3r |\Psi_{n\ell}(\vec{r})|^2 f(r) \quad (15)$$

Thus the difference in energy between these two levels which split the degeneracy is:

$$\Delta = \hbar^2(\ell + \frac{1}{2}) \int d^3r |\Psi_{n\ell}(\vec{r})|^2 f(r) \quad (16)$$

The main point from this equation is that the larger the orbital angular momentum ℓ , the larger the spin-orbit splitting. Thus the energy difference between the two g levels, with $\ell=4$, is larger than the splitting between two p levels, with $\ell=1$. This can be seen in figure 2, where the energy difference between $1p_{1/2}$ and $1p_{3/2}$ is smaller than the splitting between $1g_{7/2}$ and $1g_{9/2}$ [5].

2.2 Experimental evidence of shell gaps and magic numbers

This spin-orbit splitting supplied a theoretical explanation for the previously observed magic numbers in nuclei. Those nuclei with specific numbers of protons and/or neutrons exhibited properties of enhanced stability. For example, the neutron separation energies in nuclei reach a local maximum at a magic number of neutrons and drop significantly with the addition of one more neutron.

Other examples which show shell structure in nuclei and contain observations of these magic numbers are charge distributions, deformations, and the energy of 2^+ states in even-even nuclei. This is shown in figure 3. A maximum in these energies indicates that a shell gap needs to be crossed in order to promote nucleons to the first excited states at these numbers of neutrons. Thus there are many examples experimental verifications of the existence of shell structure and magic numbers in nuclei.

2.3 Microscopic description of β decay

A nucleus may undergo one of three different known types of β decay. It may emit a positron and a neutrino after a proton turns into a neutron. This is known as β^+ decay. It may undergo electron capture, wherein an inner atomic shell electron is absorbed by a proton, thus the proton

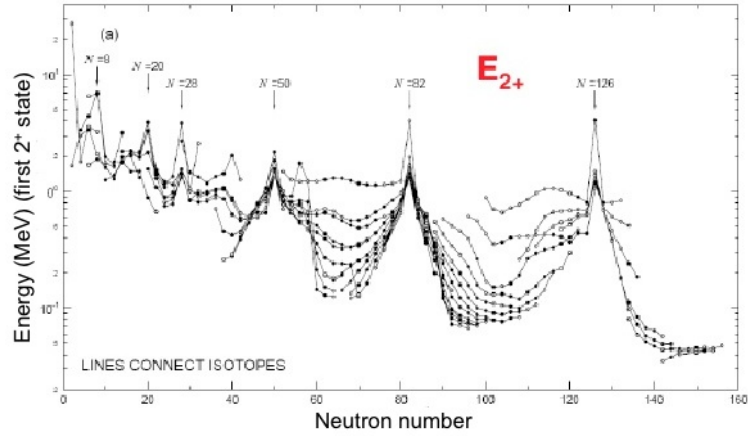


Figure 3: The first 2^+ state energies in even-even nuclei are at local maxima at “magic” numbers of neutrons. This means that at those nuclei more energy is required to excite nucleons to this first excited state than for the immediate neighboring even-even nuclei. This indicates an increased energy gap for nuclei with these particular neutron numbers.

changes to a neutron and a neutrino is also emitted. This is called electron capture (EC). A third possibility is that an electron and antineutrino are emitted and a neutron in the nucleus changes into a proton. This last example is known as β^- decay.

When a nucleus undergoes β decay this involves the weak interaction and a change of a quark flavor inside one of the nucleons. For example, in β^- decay a neutron will change into a proton and this involves a down quark decaying to an up quark. This is because a neutron consists of two down quarks and one up quark, but it transitions into a proton, which has two up quarks and one down quark. The change of flavor of the down quark to an up quark is mediated by a gauge boson for weak interactions. In β^- decay a W^- boson is emitted to carry negative charge for formation of the electron. This is depicted in the Feynman diagram in figure 4. This shows that a down quark in the neutron changes into an up quark and emits a W^- boson, which in turn changes into an electron and an anti-neutrino. The final result is a proton with the electron and anti-neutrino carrying away energy and momentum. Charge is conserved in the process.

The following equation simply restates the full changes that take place during β^- decay. A neutron changes into a proton and an electron and anti-neutrino are emitted.

$$n = p + e^- + \bar{\nu} \quad (17)$$

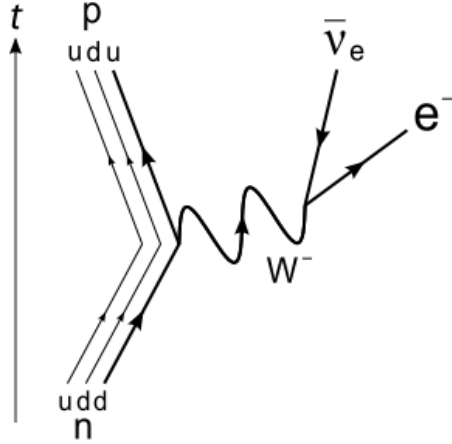


Figure 4: Feynman diagram of β^- decay, with time progressing upward on the vertical axis.

The quantity referred to as lepton number is also conserved in all β decays. The left hand side, consisting of just the neutron, has no lepton. The right hand side has an electron, with lepton number as positive one, and an anti-neutrino, with lepton number as negative one, thus conserving a net lepton number of zero on both sides.

The helicity of the leptons involved is another point important for nuclear β decay. Helicity is defined as the scalar product of spin angular momentum with linear momentum, divided by the scalar of linear momentum of the particle. Thus the helicity is as follows:

$$h = \frac{\vec{\sigma} \cdot \vec{p}}{p} \quad (18)$$

A lepton is labeled left-handed if its linear momentum and spin angular momentum are antiparallel. The helicity would be negative. A lepton is right-handed if the linear momentum and spin angular momentum are parallel and the helicity would be positive. In β decay all leptons are left-handed and all anti-leptons are right-handed. Experimentally the electron neutrino has a measured helicity of about -1, the electron anti-neutrino has a helicity of about +1. They would equal +1 or -1 exactly if the neutrino, or whichever particle is in question, were massless because the eigenvalue for helicity is actually $\frac{v}{c}$, the speed of the particle over the speed of light. The observation in β^- decay is that whether or not the total change in angular momentum is by one unit $\hbar$ or zero units $\hbar$, the linear momentum and the spin angular momentum of the neutrino must be parallel. This is due to conservation of linear and angular momentum for all particles involved [6]. Thus in β^- decay the neutrino must always be right-handed, with

positive helicity, and by lepton number conservation it must be an anti-neutrino. The opposite is observed in β^+ decay, where the neutrino emitted must be a regular neutrino by lepton number conservation but the linear and spin angular momentum vectors must always be anti-parallel by conservation of linear and angular momentum of all particles involved. Therefore, in β^+ decay the particle emitted is a neutrino which is also left-handed with negative helicity. One can rewrite the reactions for β^- and β^+ decay as the following equations, specifying right-handed anti-neutrinos and left-handed neutrinos.

$$n \rightarrow p + e^- + \bar{\nu}_R \quad (19)$$

$$p \rightarrow n + e^+ + \nu_L \quad (20)$$

In order to describe what the quarks experience as they undergo the necessary changes, one must note that in strong interactions the flavors of the quark are conserved. Thus the strong-interaction quark eigenstates are simply consisting of up, down, charm, strange, top, and bottom quarks. But in decays mediated by the weak interaction the flavor of the quark changes, as is discussed here with β^- decay changing a down quark into an up quark. Furthermore, in observed weak decays the quark flavor changes to a definite flavor, i.e. only one flavor is observed in the final product. But the real relationship between the initial and final quark states can be described in one way with the following example:

$$u \rightarrow d' = d * \cos\theta_c + s * \sin\theta_c \quad (21)$$

Where the final quark states is some mix of a down quark, d, and a strange quark, s, in this case. This mixing angle θ_c is called the Cabibbo angle. For β decay here we only concern ourselves with the transition between the up and down quarks. This alone is governed by:

$$G_v = G_F \cos\theta_c \quad (22)$$

The relationship between the G_v and G_F quantities from the quark flavor changing argument can be analyzed further with knowledge of helicity in nuclear β decay. Operators acted upon by a parity transformation, or reflection across the axes, for example from x to -x in cartesian coordinates, may or may not conserve the helicity. Specifically, the operators which are scalars, pseudoscalars, and tensors produce leptons with both helicities when this parity transformation is applied. Only vector and axial vector operators (axial vector operators are those which do not change sign upon reflection across the axes, such as angular momentum) obey what is observed

in nature. That is, they produce leptons of the correct helicity observed in β decay experiments. Because the vector and axial vector operators give different parities, they must be used together in a linear combination when treating β decay. This is part of the so-called V-A theory for β decay. That is, nuclear β decay is based upon the use of both a vector operator and an axial vector operator. The relation between the coupling constants for the vector and axial vector operators may be introduced as:

$$g_A = \frac{G_A}{G_V} = \frac{f_\pi g_{\pi n}}{M_N c^2} \quad (23)$$

This is referred to as the Goldberger-Treiman relation. The f_π is the pion decay constant and the $g_{\pi n}$ is the coupling constant between the pion and a nucleon. The pion is a meson consisting of some up and down quark linear combination, which mediates the residual strong force between neighboring nucleons, thus binding them together in the atomic nucleus. Also, M_N is the mass of the nucleon in question and c is the speed of light. This number g_A is experimentally measured to be [8]:

$$g_A = -1.259 \quad (24)$$

2.4 Fermi and Gamow-Teller operators

In nuclear β decay there are various quantum mechanical rules governing the probability of a transition between a state in a parent nucleus to one in a daughter nucleus. There are two general types of allowed β decays, Fermi and Gamow-Teller. The operators for both of these transitions act on nuclear wave functions to change neutron to proton or vice-versa. An isospin lowering operator for example will change a neutron to a proton and an isospin raising operator will change a proton to a neutron. An isospin lowering operator acting on a proton will yield a null wave function, essentially it annihilates the proton state. An isospin raising operator acting on a neutron does the same. Given these purely mathematical constructs of isospin raising and lowering operators, the Fermi and Gamow-Teller operators are as follows. First the Fermi operator:

$$\sum_i \tau^-(i) = T^- \quad (25)$$

And the Gamow-Teller operator:

$$\sum_i \hat{\sigma}(i) \tau^-(i) = Y^- \quad (26)$$

The Gamow-Teller operator also includes a Pauli spin matrix operator, $\hat{\sigma}(\iota)$, for change in spin. It is a vector operator which will change the spin of the old neutron to a new spin for the new proton. Because in Gamow-Teller decay the resultant electron and anti-neutrino couple to spin one, the possible spin projections for two particles coupled to spin one are one, zero, or negative one. It is a triplet state. This vector operator has three components because of this triplet nature:

$$\hat{\sigma} = (\sigma_1, \sigma_2, \sigma_3) \quad (27)$$

Where the σ_1 , σ_2 , and σ_3 are standard Pauli spin matrices in three dimensions. The minus sign in the Fermi and Gamow-Teller operators mean the isospin operators are lowering operators because this is for β^- decay, in which a neutron changes to proton. If it were β^+ decay then the - sign would be replaced by a + sign for isospin raising operators. The i index is present for adding over all of the decaying nucleons.

The transition probability for a nuclear β decay to occur comes from the calculation of a matrix element. That is, an expectation value of an appropriate quantum mechanical operator between the initial and final nuclear states in question. In Dirac notation:

$$M_{if} = \langle \Psi_f | \hat{O} | \Psi_i \rangle \quad (28)$$

Specifically here, the Fermi and Gamow-Teller reduced transition probabilities, which are scalars, are given by placing the Fermi and Gamow-Teller vector operators, respectively, as acting between the initial and final nuclear states in the decay. For Fermi transitions:

$$B_F^\pm = \frac{|\langle \Psi_f | \sum_i \tau^\pm(\iota) | \Psi_i \rangle|^2}{2J_i + 1} \quad (29)$$

And for Gamow-Teller transitions:

$$B_{GT}^\pm = \frac{c_A^2 |\langle \Psi_f | \sum_i \sigma(\iota) \tau^\pm(\iota) | \Psi_i \rangle|^2}{2J_i + 1} \quad (30)$$

Where J_i is the total angular momentum of the initial state and c_A is a constant pertaining to a renormalization in the Gamow-Teller decay [7].

In general, if a bound neutron in an atomic nucleus undergoes Fermi β^- from a spin projection (m_j) of $+\frac{1}{2}$, it will change into a proton with spin projection of $+\frac{1}{2}$ and the electron and anti-neutrino will couple to spin 0. However, in a Gamow-Teller decay that same neutron

would change into a proton of spin projection $-\frac{1}{2}$ and the electron and anti-neutrino will couple to spin 1. Thus a pure Gamow-Teller decay is a “spin-flip” decay because of the change in spin between the old and new nucleon.

To acquire the operators for forbidden transitions, one must simply derive them from a sequence of allowed transition operators. For example, an operator for a second forbidden transition between a state with $J^\pi = 0^+$ to $J^\pi = 2^+$ would first involve an allowed transition operator changing the 0^+ state to a virtual 1^+ state, and another allowed transition operator to transform the 1^+ state to a 2^+ state. First forbidden decays, and all others which require a change in parity, would also require the operator to change the orbital angular momentum by an odd amount of 1,3, or 5 etc. units to change the parity [8].

2.5 Decay lifetimes, energetics, and selection rules

A rule known as Fermi’s golden rule states that the decay rate from one state to another is given by the final density of states multiplied by this matrix element squared. Using Fermi’s golden rule of transition rates between initial and final states, one can calculate the decay rate. The decay rate per final electron energy interval is given by this expression:

$$\frac{dW}{dt dE_e} = G_\beta^2 \frac{d\rho}{dE_e} [B_F + B_{GT}] \quad (31)$$

The ρ factor here is the final density of states. The G_β is a coupling constant for the weak interaction, which is the force involved in mediating β decay at the quark level. This expression states that the rate at which a β decay occurs from initial state in the parent nucleus to a final state in the daughter nucleus depends upon the derivative of the density of states with the energy of each possible final state, multiplied by the addition of the Fermi and Gamow-Teller reduced transition probabilities.

To obtain the total decay rate for all decays from ground state of the parent nucleus to all possible states in the daughter nucleus, one must integrate the following new expression for decay rate per unit of final energy interval:

$$\frac{dW}{dt dE_e} = \frac{G_\beta^2}{2\pi^3} [F(E, Z) p_e E_e (\Delta_f - E_e)^2] [B_F + B_{GT}] \quad (32)$$

Here the p_e is final electron momentum at infinity, E_e is the final electron energy at infinity, Δ_f is the decay energy $E_{initial} - E_{final}$ and $F(Z, E)$ is the Fermi function. In general it is calcu-

lated as this ratio:

$$F(E, Z) = \left| \frac{\Psi(0)_{with}}{\Psi(0)_{without}} \right|^2 \quad (33)$$

That is, the ratio of the electron wave function at the nucleus with the coulomb interaction versus without the coulomb interaction all squared. For non-relativistic electrons and a point-like nucleus, the following is the Fermi function, given as:

$$F_{NR}(E, Z) = \frac{2\pi\eta}{1 - e^{-2\pi\eta}} \quad (34)$$

Where $\eta = \frac{\pm Ze^2}{v_e}$, Z is atomic number of the parent, e is the electron charge, and v_e is the velocity of the emitted electron (or positron) at infinity. This effects the shape of the final electron energy distribution. The β^- curve will have more counts at lower energy whereas the β^+ decay will have fewer counts at lower energy because of the coulomb interaction either attracting or repelling this lepton (electron or positron) to or from the nucleus. The endpoint of such a spectrum is the Q value, the energy difference between ground state of parent and ground state of daughter.

Finally, integrating over all possible final electron energies leads to the total decay rate from parent ground state to all daughter states.

$$\frac{dW}{dt} = \frac{G_\beta^2 m_e^5}{2\pi^3} f[B_F + B_{GT}] \quad (35)$$

Where f is the Fermi integral:

$$f = \frac{1}{m_e^5} \int_{m_e}^{\Delta_f} F(Z, E) p_e E_e (\Delta_f - E_e)^2 dE_e \quad (36)$$

This integral is over the energy range from the rest mass of the electron up to the total decay energy available. Table 1 summarizes allowed β transition properties.

The $\log ft$ value is used to factor out this Fermi integral factor. The f is from the Fermi integral discussed above. It includes the information on the overlap of the initial and final wave functions and the energy difference. The $t_{1/2}$ is for the partial half-life, or half-life of that particular transition from initial state to a specific final state.

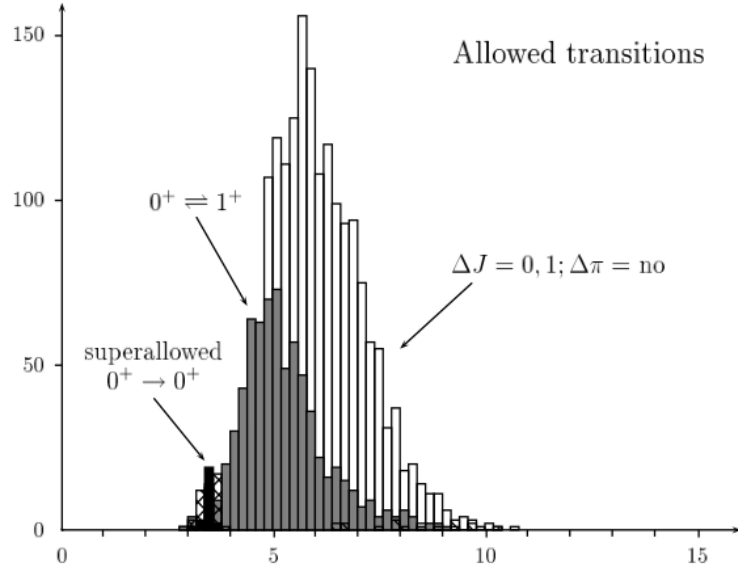


Figure 5: Empirical $\log ft$ values for allowed β decays [9].

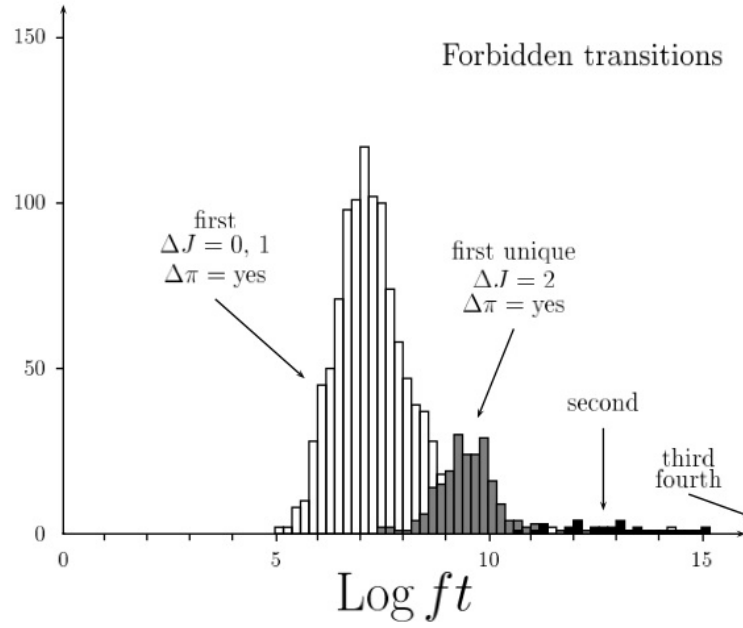


Figure 6: Empirical $\log ft$ values for forbidden decays [9].

Table 1: Table describing allowed β transition properties

Transition	Spin	orb. ang. mom.	parity change	isospin
Fermi	$\Delta J=0$	$\Delta L=0$	no	$\Delta T=0$
Gamow-Teller	$\Delta J=0,1$ not 0 to 0	$\Delta L=0$	no	$\ \Delta T\ =0,1$

$$t_{1/2} = [dW/dt]^{-1} \ln 2 \quad (37)$$

The full quantity $ft_{1/2}$ is thus given by:

$$ft_{1/2} = \frac{2\pi^3 \ln 2}{G_\beta^2 m_e^5} \frac{1}{B_F + B_{GT}} = \frac{4\pi D}{B_F + B_{GT}} \quad (38)$$

Here D is a constant in units of time. This $ft_{1/2}$ is the value for a transition to a particular state with an energy labeled E_f . But summing over all possible final states with E_f energy yields the full half life for the β decaying state to all of the possible final states.

$$T_{1/2}^{-1} = \sum f_f \frac{B_F(E_f) + B_{GT}(E_f)}{4\pi D} \quad (39)$$

As stated before, it is more common to take the log of the quantity ft and thus it becomes $\log ft$. The larger this number the weaker the transition because it means the reduced transition probabilities are smaller, and thus the smaller this number the stronger the transition because those same reduced transition probabilities are larger.

Table 1 displays the total angular momentum change, or change in total spin J, which occurs with each type of allowed β decay transition. However, a transition where the $\Delta J= 0, 1$, or 2 with a change in parity is called a first forbidden transition.

Other transition types include allowed transitions that are 0^+ to 0^+ and isospin does not change. These are superallowed decays and their $\log ft$ values range from 3 to 4. Also, there are so-called isospin forbidden decays, which are the same as allowed decays but isospin changes. Their $\log ft$ values range from 6.4 to 10.3.

The allowed transitions do not carry away orbital angular momentum, but forbidden transitions do. A transition with a $\Delta J= 2$ and parity change is called first forbidden unique. There are second, third, and fourth forbidden transitions as well. Each successive forbidden decay

Table 2: Table showing properties of forbidden β decays.

Transition	Spin	parity change	typical log ft
First forbidden	$\Delta J=0,1,2$	yes	5.5-9
Second forbidden	$\Delta J=2,3$	no	11-13
Third forbidden	$\Delta J=3,4$	yes	18

allows one more unit of angular momentum change and whether the parity changes depends on this relationship:

$$\Delta\pi = (-1)^n \quad (40)$$

where n is the order of the forbidden transition. For example, first forbidden is $n=1$, second forbidden is $n=2$, etc.

2.6 β decay branching ratios

When a parent nucleus undergoes β decay to its daughter, a certain percentage of all β decays will go to the ground state, each excited state available through selection rules, and to neutron unbound states. There are two general concepts which determine the β decay branching ratios distribution. This distribution is influenced first by energetics of the decay. Decay to the ground states and low energy states in the daughter are highly preferred. The second is that of nuclear structure. The wave function overlaps, between the parent state and each state in the daughter, determine the so-called reduced transition matrix elements for each transition.

The Fermi phase space function shown above determines how the population of states in the daughter nucleus is weighted just by the energy difference between the ground state of the parent and the state being populated in the daughter. As can be seen by the shape of the function in figure 7, there is a large preference for decay directly to the ground state, but this preference rapidly decreases with smaller decay energy higher up in the Q_β window.

The nuclear structure S_β function depends upon the overlap of initial and final wave functions populated in the decay. Thus it contains the nuclear structure influence on the decay as opposed to the energetics. It is both of these functions that are folded together to produce the final β decay intensity distribution, labeled I_β in figure 7. This S_β function also contains resonance features such as the Gamow-Teller resonance, usually centered around 15 MeV in or

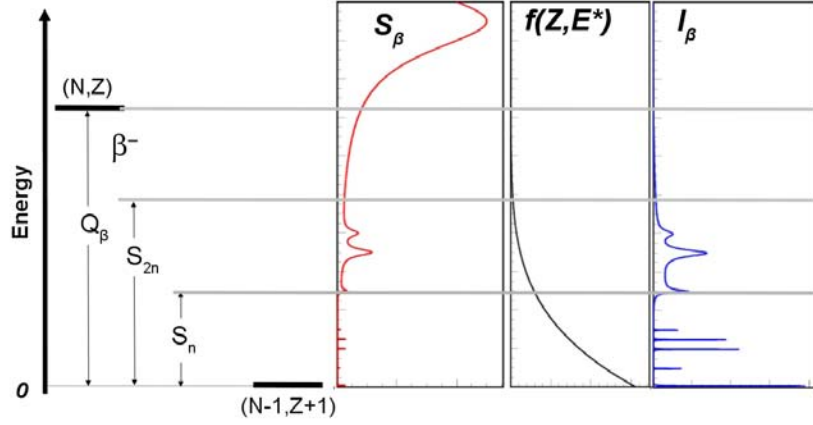


Figure 7: The S_β function and the Fermi statistical factor, both of which contribute to the β decay intensity distribution, I_β . The one neutron and two neutron separation energies are also shown, along with the full Q_β decay window value.

above the decay window, and a Pygmy resonance. The GT resonance in the daughter nucleus is a group of excited state such that all of the protons are in a group oscillating with respect to a group formed by all the neutrons. The Pygmy resonance is likened to the same oscillations but since it is at a lower energy in the decay window it is more similar to just surface protons oscillating with surface neutrons. However, they are still resonances in which the daughter nucleus may be populated after decay. Thus the S_β function has peaks at these resonances, which in turn enhances the final I_β values at those energy ranges.

The decay half-life can be written in terms of this S_β function and Fermi statistical factor folded together. It is:

$$T_{1/2} = \frac{1}{\sum_i S_{\beta i} f_i(Z, E^*)} \quad (41)$$

Here the Z is proton number of the parent nucleus and E^* is the energy of an excited state populated in the daughter. Therefore, this last equation relates the half-life of decay to both the nuclear structure factor on β decay, S_β , and the energetics influence in the Fermi statistical factor, $f(Z, E^*)$.

3 Experimental studies of the neutron-rich isotopes of Zn

3.1 Introduction

One major field of nuclear structure physics involves the observation of radioactive nuclear decay. Here we take advantage of the β^- decay to populate excited states in the daughter nucleus. When a nucleus undergoes β^- decay from the ground state, or possibly an isomeric excited state, it may decay directly to the ground state of the daughter, to an excited state, or to a neutron unbound state and thus emit a neutron as well. The decay radiation provides the experimenter with a probe into the structure of each populated nucleus. A detected γ -ray represents a transition between nuclear states because it was emitted when that nucleus transitioned from one state to another. Furthermore, different γ -ray transitions emitted in sequence represent a cascade of transitions in that nucleus. Time properties of a γ -ray emission also reveal the nucleus from which it decayed through half-life identification. It may also reveal whether an isomeric state is present because the half-life would be different outside of error bars and other evidence would have to show it as not originating from another nuclear decay. In general this is how decay spectroscopy is conducted, through analysis of decay radiation to discover new information about both nuclear structure and its influence on decay properties.

The success of decay spectroscopy experiments depends largely on the ability to produce nuclei of interest with large enough statistics above a background of radiation, which is always detected regardless of the experimental environment. The success of the experiment also requires little or no significant contaminants from unwanted isotopes arriving at the experimental setup. The main point is to separate measured decays of the desired ions of interest from background radiation. Also, since the implanted ions are far from stability and even the daughters are long-lived decaying nuclei, it is necessary to move the long-lived daughter activity away from the detectors to a location behind lead shielding.

Two methods for decay spectroscopy studies of neutron-rich nuclei at the HRIBF (Holifield Radioactive Ion Beam Facility) include studies at the Low Energy Radioactive Ion Beam Spectroscopy Station (LeRIBSS) and the Ranging-Out technique [10]. The Ranging-Out technique makes use of isobaric separation of ions but also post-accelerates those ions to 3 MeV/nucleon to then pass through an ionization chamber. The ionization chamber is used for further isobaric separation. Ions with a higher proton number Z are stopped in the gas whereas lower Z isobars are not. Energy loss through a gas is proportional to the atomic number of the ion squared,

Z^2 , according to the Bethe formula. Therefore, the more neutron-rich isobars pass through the ionization chamber. Isobaric purity is improved upon in this manner because only the low Z , more neutron-rich ions will pass to implantation site within view of the detectors.

Studies at LeRIBSS do not involve ions which are post-accelerated in the tandem. They are passed through the high resolution magnetic isobaric separator, but are implanted upon a moving tape collector or silicon detector at 200 keV energy. The general purpose of LeRIBSS is to study the decays of fission fragments after they are implanted on a tape. No reaction is required of the fission fragments once they arrive at LeRIBSS, nor are they required to pass through either an ionization chamber filled with a gas nor through any foils of significant thickness. The γ and β decay photons and particles from these fission fragments are measured and after some short time the long-lived activity can be transported away from view of the detectors.

In the particular case of zinc ions, it is advantageous to study their decays at LeRIBSS because zinc cannot make a negative ion. Hence it cannot be post-accelerated since the HRIBF tandem accelerator needs negative ions to be injected into it. The study of zinc, or any ion which cannot be made negative, at other experimental stations which require post-acceleration is simply not possible because of this reason.

3.2 Experimental details

3.2.1 ISOL technique at the HRIBF

The Isotope Separation Online Technique is utilized [11] [12] at the HRIBF to produce neutron-rich, unstable nuclei. The process starts with the acceleration of protons from the Oak Ridge Isochronous Cyclotron (ORIC) to an energy of 54 MeV and with a beam current of 10-15 μ A. The protons impinge upon a high-temperature target of uranium carbide, located on a high voltage platform. This creates fission fragments from the ^{238}U nuclei. The fission products cover a wide mass distribution and are transported, as a vapor, to a coupled ion source. The ion source is an electron beam plasma ion source. Fission products are extracted as positive ions off of the ion source, which is at 160 kV. The source is on a 40 kV potential platform with respect to the whole platform. The positive ions pass through the first stage dipole resolving magnet, located on the high voltage platform, to select a specific mass of ions. These positive ions are accelerated off of the platform to a total energy of 200 keV. They are then transported to the second-stage, high resolution magnet for isobaric mass separation.

The first challenge with producing neutron-rich nuclei through this ISOL technique is the small cross section for producing such nuclei through fission of Uranium (It should be noted that our cross sections of production are typically higher than at many fragmentation facilities). Often times the cross section for producing nuclei along an isotonic chain will drop by at least one order of magnitude with the addition of each neutron along the chain away from stability. Another problem exists because of the challenge of isobaric separation. If isobaric mass separation is not of sufficient strength and since the amount of neighboring isobars may be orders of magnitude greater in beam intensity than the desired ion, the experimenter may never be able to separate the large amount of unwanted isobars from the tiny amount of desired isobars in the beam. In the particular case of ^{81}Ga versus ^{81}Zn , the gallium arrives at 10^6 pps (with open slits, explained below) whereas the zinc arrives at LeRIBSS at 30 pps.

A second challenge with producing and obtaining neutron-rich nuclei from an ISOL source is that of high temperature chemistry. Some ions are retained on the walls of the source chamber because of chemical properties of this element with its surroundings. Even if this element eventually releases from the source, it may be after one or more half-lives have passed, thus too many of the ions will have decayed before reaching the experimental end station. Nickel is one such example, it is difficult to produce from ISOL sources because of this chemistry reason.

3.2.2 Charge exchange cell

There is a cell containing cesium vapor which may be used before the high resolution isobaric separator to change positive ions into negative ions. This is needed for ions which will be post-accelerated because the ions come out of the source as positive but must be negative before injection into the tandem accelerator. As the ions pass through the vapor they collide with the gas molecules and some will exchange charge. The cesium atoms will give up an electron to the ion passing through because they have an electron affinity favoring the release of an electron. Cesium has an electron affinity of +46 kJ/mole [13]. It is an alkali metal which gains energy by loosing its one valence electron. Two collisions are required, one to go from positive to neutral and another to go from neutral to negative. For some elements it is not possible to gain a negative charge state due to an opposite sign on their electron affinities, such as with zinc and cadmium [13]. These elements cannot form negative ions. Therefore, in some cases it may be advantageous because one can study an isotope whose isobaric neighbor does not form a

negative ion. Sending those ions through the tandem means the final beam will have the desired ion without the isobaric neighbor present. For example, copper isotopes have been studied with negative ions after post-acceleration through the tandem because the isobaric neighbor zinc is not present in the beam.

However, one major problem with using this charge exchange cell is that it reduces beam intensity because many of the ions will loose too much energy and not pass through the cell. It reduces the beam to about 5-10% of its original intensity, although this varies some depending upon the element and its electron affinity. Thus it is advantageous to study the decay of positive ions at the LeRIBSS when possible because such a beam will be of higher intensity than a negative ion beam.

3.2.3 High resolution isobar magnets

Magnets are often used to perform mass separation of ions. The Lorentz force experienced by a charged particle in motion inside of a magnetic field is utilized and is given by:

$$\vec{F} = q\vec{v} \times \vec{B} \quad (42)$$

Where q is the charge of the particle, $\vec{v}$ is the velocity vector of the particle, and $\vec{B}$ is the magnetic field vector. The force vector $\vec{F}$ is thus equal to this vector product between velocity and magnetic field. This magnetic Lorentz force provides a centripetal force on the charged particle, thus moving it along a circular path with some radius, ρ .

$$\vec{F} = q\vec{v} \times \vec{B} = \frac{mv^2\hat{\rho}}{\rho} \quad (43)$$

It is common to rewrite this equation such that a quantity called the magnetic rigidity, the field times the radius, has a magnitude of:

$$B\rho = \frac{mv}{q} \quad (44)$$

Thus a specific magnetic rigidity can be set in a magnet, a given field for a given radius is determined from the mass, charge, and velocity of the ion moving through the $\vec{B}$ field.

The High Resolution Isobar Magnet setup consists of a pair of large-radius, double focusing dipole magnets with a theoretical mass-resolving power of $\frac{m}{\Delta m} = 10000$ utilized to separate

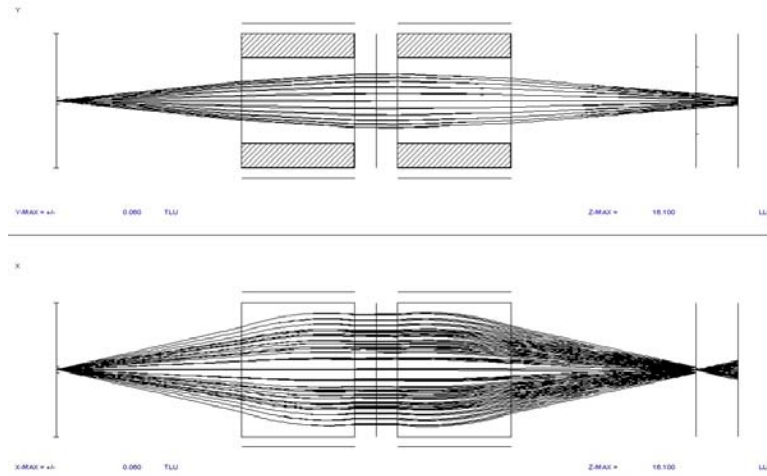


Figure 8: The vertical focusing of ions moving through the high resolution isobar magnet. Vertical focusing is an effect of the fringe magnetic fields at the edges of the magnet. Thus it is advantageous to have two magnets and four edges instead of one magnet with two edges. This helps to focus the beam vertically and improves resolution [14].

isobars from each other. Here m is the mass number of the isobars and Δm is the mass difference between the two neighboring isobars which need to be separated. In an experiment the mass-resolving power is limited by at least three factors: Beam emittance from the ion source, energy spread of the ions, and charge state of the ions. For example, a large beam emittance will create a large beam spot at the object position and thus the image beam spot will be too large to resolve the masses. A large spread in energy of the ions will create the same effect. Ions arriving at different energies into the magnet will bend at different radii, thus decreasing or destroying the mass resolution. Different charge states can also disrupt mass resolution, as explained in the next section.

Figure 8 shows how ions move through this high resolution isobar magnet in the plane perpendicular to the plane of the ground. One advantage of having two magnets with a space in between is that there is extra vertical space for transmission. It is also advantageous to have two magnets and four edges instead of one magnet with two edges because vertical focusing is only affected by the fringe fields at the edges. Thus better resolution is obtained with four edges and two magnets. Figure 9 shows how the ions move in the plane parallel to the plane of the ground. A longer bend radius also improves resolution and focusing through the magnets [14].

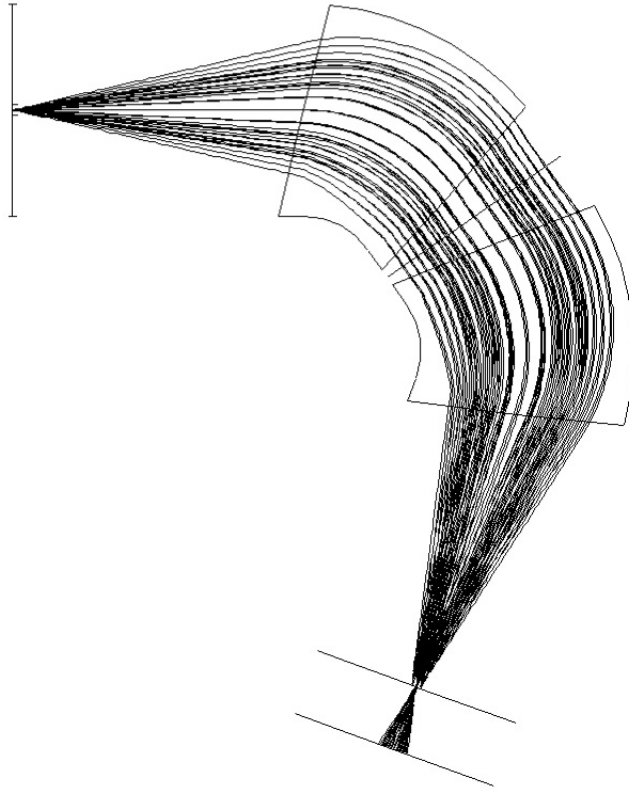


Figure 9: The horizontal focusing of ions moving through the high resolution isobar magnet. Another advantage of having two magnets instead of one is that a longer bend radius improves resolution [14].

3.2.4 Charge state contaminants

Most ions which are converted to a negative charge through the cesium charge exchange cell exit with a charge of -1. It is very rare for ions to acquire a charge of -2 or more. However, positive ions emerge from the ion source and have a +1 charge but may possess a +2 charge state instead. This is more likely than the creation of a -2 charge state in the charge exchange cell. The important point is that since the mass separators actually work on separating ions with the same mass over charge ratio, A/Q , it is possible for a heavier ion with charge +2 to come through if it has the same A/Q as the desired ion. For example, it is possible for ^{81}Zn with a charge state of +1 to come through the mass separator at the same time as say ^{162}Eu with a charge state of +2. Thus the spectrum would contain decays from the ^{162}Eu as well. This is not a problem if the intensity of europium ions is not obscuring the observation of the zinc decays (as was the case here, we did not observe such decays with an ^{81}Zn beam, but this description is included as a possible example of charge state contaminants). However, if there is too much intensity in the beam from the charge state contaminant then a solution must be found. A molecule containing ^{81}Zn would have a different mass and so this would be a solution, if it could be made and extracted from the ion source. It is guaranteed to have a different A/Q ratio if the charge state is still +1. In general, using molecular beams to obtain a specific isotope may be beneficial because of both this charge state contaminant issue and also because neighboring isobars may be too close in mass. Therefore creating a beam of molecules may eliminate this issue of resolving close masses, depending upon which element and isotope is desired.

3.3 LeRIBSS

After the ions leave the High Resolution Isobar Magnets they arrive at LeRIBSS, the Low Energy Radioactive Ion Beam Spectroscopy Station. They avoid the tandem, thus no post-acceleration occurs, and the ions arrive with an energy of 200 keV. The ions implant upon a tape, which is part of the Moving Tape Collector (MTC), designed to collect short lived activity in view of the detectors and move long lived activity away from the detectors behind a hanging wall of lead. The purpose of the lead wall is to shield the long lived activity from the γ -ray and β electron detectors. The ions may arrive as positive or negative since the tandem is not used, which requires negative ions for acceleration. For many isotopes it is beneficial to keep the ions positive because the cesium vapor charge exchange cell decreases the beam intensity and lowers emittance quality. Use of the tandem reduces the intensity of the beam by a fac-

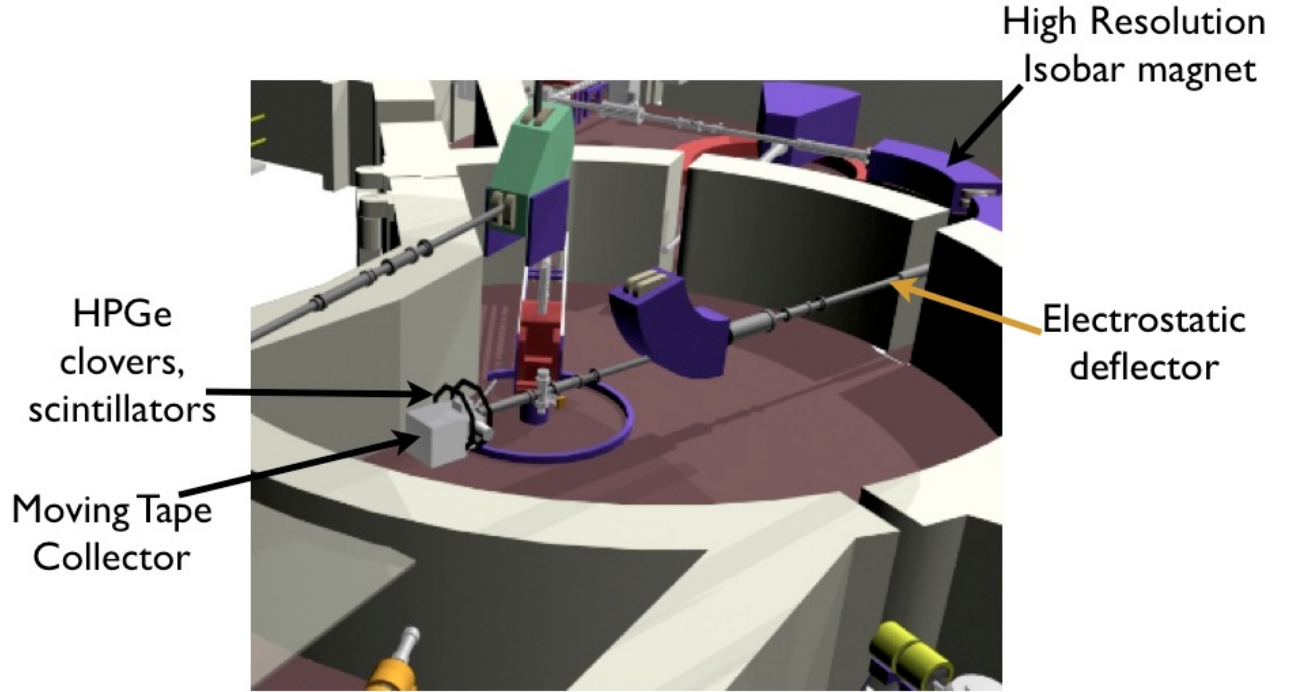


Figure 10: The LeRIBSS experimental setup. The beam comes in from right to left, first through the high resolution isobar separator, past an electrostatic deflector, and it bypasses the tandem because post-acceleration is not needed. The ions arrive at 200 keV energy and implant on a moving tape collector (MTC). The site of implantation is surrounded by four high purity germanium detectors to measure γ decays and two plastic scintillators for β electron detection.

tor of 10 and the charge exchange cell usage further reduces the beam intensity by a factor of 20.

The ions implant upon the tape at a site which is surrounded by two plastic scintillators for detection of β decay electrons and four high purity germanium clovers (HPGe) to measure the γ radiation, see figures 10 and 12. All of these detectors input their measurements into digital data acquisition boards for digitization and basic analysis of the signals into a raw data file.

3.3.1 Electrostatic deflector and tape system

As the ions approach the LeRIBSS end station they pass by an electrostatic deflector, which is designed to deflect the beam completely off of the implantation site such that only decays may be observed with no further implantation from a beam. There is a waveform generator

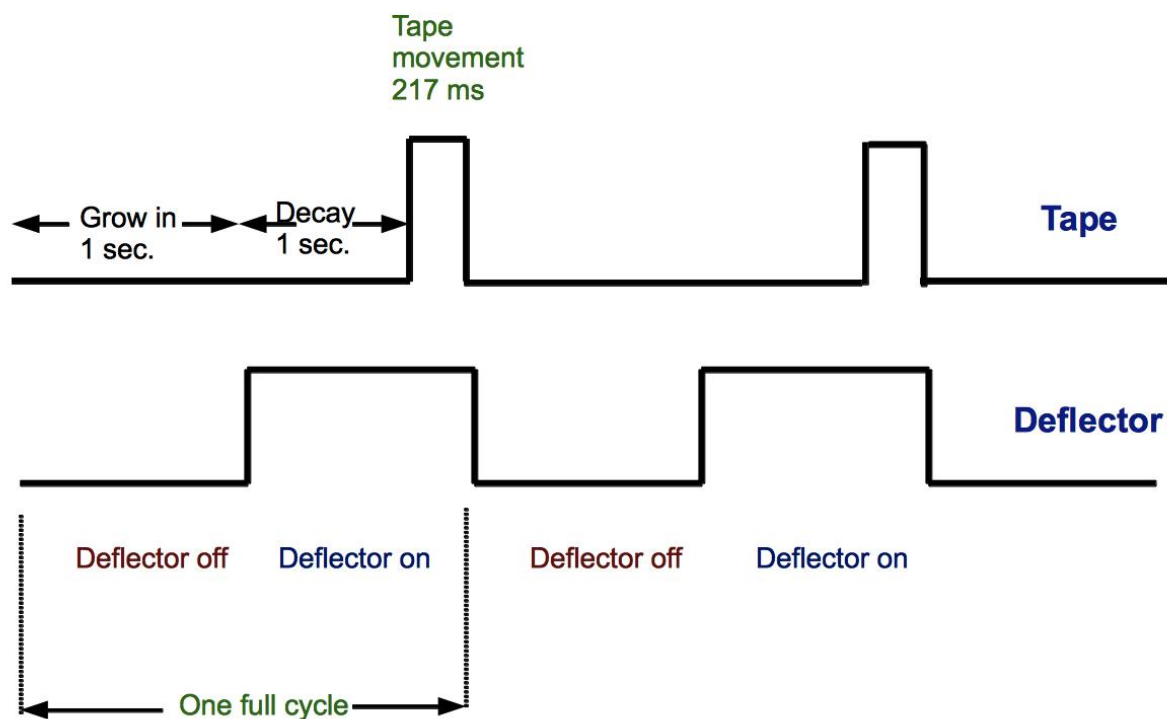


Figure 11: Schematic of cycles for both electrostatic deflector and the MTC. The deflector allows beam particles to collect for one full second (or whatever desired time), then it deflects away the beam for another full second, and finally the tape moves over a period of about 220 ms. At the end of the tape move the deflector goes off and the whole cycle starts over.

connected to the voltage supply for this electrostatic deflector, which controls the cycling of this deflector. It also controls the MTC movements. Thus a series of grow in and decay cycles may be observed from the activity of the implanted ions on the tape. Figure 11 contains an example schematic of the tape movement cycle overlaid with the electrostatic deflector cycle. The beam gathers for one second, for example, followed by deflection of the beam and so only decays are observed during the next second. Finally, the tape moves the activity away from the implantation site, which takes 217 ms or longer if desired, and the whole cycle begins again. It is possible to allow continuous beam implantation or a long free decay time over many half-lives. This helps to determine an upper and lower baseline in the counts versus time profile to improve upon the half-life determination. The time for continuous implantation, free decays, and the time taken for moving the tape are all variables which can be altered. It is possible to move the tape different distances as well. Thus the experimenter may put the tape under more stress by making it move further away in a shorter amount of time. Also, the deflector polarity is positive and the voltage pulse rises from 0 to +1900V. This is preferred to a negative voltage pulse because such a pulse would pull the positive ion beam the other direction during deflection, potentially contaminating the tape with isobaric contaminants each cycle.

Another point on the deflector is that it is an electrostatic deflector as opposed to a magnetic separator. This is because it avoids causing some form of hysteresis effect on the isobaric separator. Such constantly changing magnetic fields would alter the field strength from the isobaric high resolution separator, which is a danger that is unnecessary to introduce. All that is required is to move the ions away from implanting at LeRIBSS and an electrostatic deflector is sufficient.

3.3.2 Detector system

At the LeRIBSS setup there is the CARDS, the Clover Array for Radioactive Decay Studies. It contains the four clover detectors, each clover consisting of four high purity germanium crystals packed together in a tight geometry. Figure 12 shows a closeup of the clovers and two plastic scintillators around the site of implantation in the beam pipe. Each clover has four crystals, thus there are sixteen crystal detectors total in CARDS.

The site of implantation on the tape is surrounded, outside of the beam pipe, by the high purity germanium clover detectors. When ionizing radiation interacts with the electrons in the germanium atoms, the photoelectric effect occurs such that energy from the γ -ray photons is



Figure 12: The four high purity germanium clovers and two plastic scintillators (covered in black tape) surrounding the site of implantation, which is inside the beam pipe.

deposited and the charges are liberated as electrons and subsequent holes. The amount of electrons and holes liberated is directly proportional to the energy of the initial γ -ray photon which induced the photoelectric effect.

The collection of charges liberated from ionizing radiation requires that none of those charges recombine with the surrounding atoms. It is also important that the collection time be sufficiently fast enough. For an operating temperature of 77 Kelvin, the electric field must be a minimum of 10^5 volts/meter [15]. Setting an extra large voltage will ensure that the electrons and holes will reach the readout locations at the ends of the crystals and not recombine with nearby atoms. The signal goes through the preamplifier, which increases the amplitude of the signal to a voltage appropriate for reception into the digital data acquisition boards. These boards must accept a signal in the one to five volt range.

Another source of radiation which the germanium detectors will always measure is that of background γ radiation. In the environment there are many sources of background radiation, including concrete blocks, lead blocks, radon gas, etc. It is also possible to see γ -rays from neutron induced reactions in the germanium from any neutrons coming from the uranium carbide fissioning ion source, although ideally such neutrons should not move past appropriate shielding

around the source.

3.3.3 HPGe energy and efficiency calibrations

There are two types of calibrations which must be made with the germanium detectors. The first is an energy calibration. When a germanium detector measures a single γ -ray photon, it places this count in a channel. For the same energy photon this count will be in a different channel in different germanium detectors. Each germanium crystal detector has a different response to the same energy. Therefore calibration sources, which emit a decay pattern with known γ -ray peaks at known energies, are measured in the detectors so that one can match the channel numbers in each detector with the correct energies. The experimenter must plot the channel number at which each peak occurs vs. the known energy of said peak from the calibration source. This plot is then fitted with a polynomial, either linear or quadratic, which is then used when the raw data is scanned into a histogram. Thus the channel number in each germanium crystal detector is transformed into a correct energy number in keV. One must also create a residual plot showing the difference between real and calculated energy vs. the real energy. Such a plot should be a random scatter, otherwise it would suggest a bias in the fit between channel number and energy, which should not exist in a good fit.

The second type of calibration is for efficiency of detecting γ -ray photons at different energies. The detectors do not cover a total 4π solid angle, thus not all γ -ray photons can be intercepted from the activity on tape. But even if they were covering total solid angle, not all γ -ray photons will deposit their energy in the germanium. Some may pass through without interaction or Compton scatter all or part of their energy. There is a relationship between the energy of the photons and the efficiency with which the germanium detectors measure them. From zero keV up to approximately 80 keV, the efficiency of detection generally rises with increasing energy. This is because with higher energy photons they are less likely to be absorbed by materials which are housing the germanium crystals. Above this point the efficiency generally decreases, but the rate of decrease depends on the energy range. To determine this curve of efficiency vs. photon energy, γ -ray sources are utilized and a source whose date of creation and absolute activity at the time of creation must be measured. Thus the absolute rate of γ -rays per second at the time of measurement is calculated from knowing the half-life for each energy emitted from said source. Measuring how many decays per second occur in the detectors vs. the absolute decays per second yields an efficiency of detection at each of these energies. A plot of these efficiencies is

fitted with a polynomial to know the efficiency at any given energy in the available energy range.

An SRM source, or standard reference material source, is used because the absolute amount of γ decays per second are known for the time in which the source was created. In this particular experiment the experimenters were made aware of the existence of the SRM source after other sources were used at LeRIBSS. Therefore the SRM source was measured in front of a small germanium crystal detector which was previously energy calibrated. The other sources, which were utilized at LeRIBSS, were also remeasured in front of this small detector at the exact same position, with less than a millimeter error in position, along with the SRM source such that the counts in each of the peaks could be rescaled to match the SRM source intensity. An efficiency curve was fit to the points measured from this SRM source. The measured counts from all of the other sources are each divided by the efficiency at that energy derived from this fit to the SRM source. Therefore, from this the experimenter calculates the absolute γ -rays emitted per second in all of the sources. This is possible because all sources were used at the same position in front of the same germanium crystal for the same amount of time, exactly one hour. With knowledge of the absolute rate of γ -rays emitted per second at each energy from each source, the experimenter then returns to the measured source data at LeRIBSS to calculate the efficiency of detection at each energy emitted from the sources. Table 3 shows sample γ energies from each source used in this efficiency calibration.

With any γ -ray efficiency measurements there are at minimum two particular problems. The first is the measurement of summation peaks and thus photo-peak efficiency is lost. That is, some of the γ -ray photons will be measured at the same time in the germanium and this charge collection will occur at the same time, leading to a summation peak in the spectra. For example, in the ^{60}Co source there are two γ -ray energies emitted, one at 1173 keV and the other at 1333 keV. These two photons can sum together in the detectors and result in a measured peak at 2506 keV.

The other problem pertains to dead time in the data acquisition system. If a source of γ radiation is active enough, this can lead to times in which the data acquisition is overwhelmed and cannot record measured decays while it is busy processing previously measured decays. The end result is that fewer counts are recorded than what is really measured in the detectors. This is referred to as dead time. In the efficiency source measurements there is an observable effect from dead time on the measured efficiency of the ^{226}Ra source. Its activity was higher than the

Table 3: Table of information on the efficiency calibration sources used at LeRIBSS.

Source	E_γ (keV)	absolute rate (s^{-1})	counts (s^{-1})	efficiency at LeRIBSS (percent)
^{133}Ba	302.4	441.6	80.15(170)	18.15
^{137}Cs	661.1	14934.5	1396.37(40)	9.35
^{60}Co	1172.8	12211.3	768.09(60)	6.29
^{60}Co	1332.0	12094.1	702.67(60)	5.81
^{226}Ra	186.1	1432.6	315.89(245)	22.05

others used because it emits significantly more γ -rays per nucleus per unit time compared to, for example, ^{60}Co or ^{133}Ba . The ^{226}Ra source emits 12 γ -rays in conjunction with a higher activity than the other sources. Also, it is difficult to determine how much each of the γ transitions will form summation peaks. This can lead to higher counts in the high energy peaks but a loss in counts in the lower energy peaks. Dead time from high activity and summation of multiple γ peaks leads to a lower efficiency curve for ^{226}Ra . Therefore a simple readjustment factor is applied to the efficiencies of the ^{226}Ra points such that they match the efficiency curves for the other sources.

The total number of known γ energies vs. their efficiencies from the sources is 28. Therefore, an efficiency curve of energy in keV vs. efficiency in percent is fitted to these points in the energy range of 50 keV to 2204 keV. Any energies measured in the decay experiments at LeRIBSS will have efficiencies given as a function of energy by this fit.

This efficiency curve only extends to an energy of 2204 keV. In the decay spectroscopy data measured for the zinc isotopes, including ^{81}Zn , there are γ -ray energies at 4.2 and 4.8 MeV. Thus a calculation or fit to find the efficiency of detecting such high energies is needed. Data from the decay of ^{79}Zn was utilized. There is a peak at 3.9 MeV and the relative intensity is known in literature [16]. Knowing the relative intensity between this 3.9 MeV transition and transitions at energies with known absolute intensity, such as 865 keV, allows one to calculate the absolute intensity of the 3.9 MeV transition. Once the absolute intensity of such a transition is known, comparing to the measured intensity leads to a calculated efficiency from the simple ratio of measured intensity to absolute intensity per unit time.

An efficiency curve is measured for both γ singles, i.e. a spectrum of all γ -rays measured

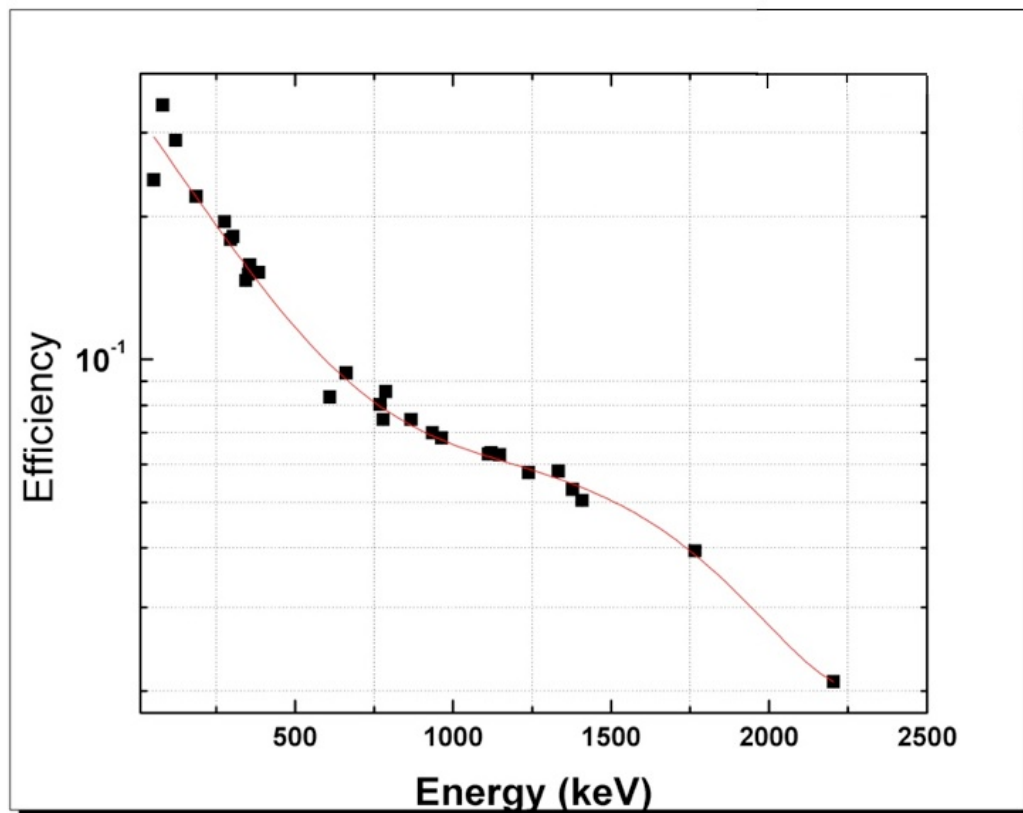


Figure 13: Efficiency curve for the clover array of 16 high purity germanium crystals for detecting γ radiation at LeRIBSS.

with no gates or conditions, figure 13, as well as an addback spectrum. The addback spectrum is constructed by adding together all γ -rays measured in the same event window ($2.5\mu\text{s}$) in the same clover. This is meant to “sum back” the Compton scattered, high energy photons. For higher energies there should be an enhancement in their efficiency of detection by performing an addback algorithm and this efficiency is measured with the same calibration technique described above.

The efficiency of detection in a single crystal is approximately $1/16$ the value at any energy for the whole array. However, two of the four clovers are closer to the site of implantation by about one centimeter, thus they have enhanced efficiencies, especially at lower energies, in comparison to the other two clovers.

3.3.4 Plastic scintillators for β decay detection

The two plastic scintillators at LeRIBSS surround the beam pipe around the site of implantation where they measure the electrons emitted from the implanted ions on the tape. The implanted nuclei on tape will β decay and emit electrons with energies between a few keV up to about 11 MeV, depending on decay energies of the ions. Most of those electrons have enough energy to pass through the 0.5 mm aluminum pipe and pass into the plastic scintillators. Some will have enough energy to pass through both of these undetected. Electrons at 1 MeV for example will travel on the order of a few centimeters in air, but in this plastic it will travel about 1-3 mm before inducing scintillation. The plastic has a model number of BC-404 (the basic chemical is called polyvinyltoluene), made by the company BICRON, and is designed to be transparent to the type of light it emits from the scintillation caused by incoming radiation. The range of wavelengths of scintillation light produced is from about 380 to 500 nm, with a peak at 410 nm. The detectors are about 15 cm in length and surround a beam pipe which is 0.5 mm thick and 5.12 centimeters in diameter. Thus they are about 2.5 centimeters away from the site of implantation at the perpendicular distance of closest approach. Each scintillator covers almost 2π solid angle around the site of implantation. Together the calculation of solid angle coverage is about 97% of the total 4π solid angle given these geometrical parameters.

The basic process of detection in plastic scintillators involves ionizing radiation entering the plastic and causing an excitation in the plastic molecules. This excitation will decay and release its energy via light emission. These photons of light travel through the plastic and enter photomultiplier tubes cemented onto the ends of the plastic. These photons then induce the

photoelectric effect in the phosphorescent glass entrance window in the attached photomultiplier tubes. The resultant photoelectrons travel along a series of dynodes, each at appropriate voltages such that electron multiplication is occurring with each subsequent dynode. This electron multiplication is the process by which a large enough signal is read from the photomultiplier tube.

An ideal scintillator will have a linear relationship between the energy of the particle detected and the amount of photons released after the deposition of the particle's energy. There should be a high ratio of amount of photons released to each unit of energy deposited in the material of the detector. Also, the emission of the photons should be prompt and have a quick decay time. Fluorescence is the ideal interaction in the material. Phosphorescence is another type of light emission but has a longer decay time and longer wavelength. Delayed fluorescence is just like fluorescence but with a longer decay time for the light emission. Therefore, these latter two processes actually contribute to background counts in the scintillator measurements due to their long decay times and only the prompt fluorescence is useful for signals [15]. An attempt to measure those longer decay time scintillation events may be ill advised because this would lead to increasing the event window time and thus record correlations between events which should be not truly in coincidence with each other.

The efficiency of detecting an electron in the plastic scintillators is very close to 100%. Sometimes an electron may scatter and only deposit part of its initial energy. However, the purpose of the scintillators at LeRIBSS is to detect the presence of an electron, not to measure its energy. The electrons can also backscatter after having left the plastic detector, possibly reentering the scintillator to deposit energy more than once and misleading the experimenter to believe they have measured more than one electron. However, backscattering increases quickly with atomic number, Z , of the scintillator material. Because organic scintillators, such as plastic, have a very low Z this is not a problem in these experiments. For this specifically sized detector, organic scintillators are suitable for detecting β decay electrons with energy of about 10 MeV or lower because of this reason.

3.3.5 Microchannel plate detector

The microchannel plate (MCP) detector setup is used for analyzing radioactive beams. It is inserted directly into the path of a radioactive ion beam without significantly disturbing the beam particles while they are in transit to the experimental station. The detector is shown in

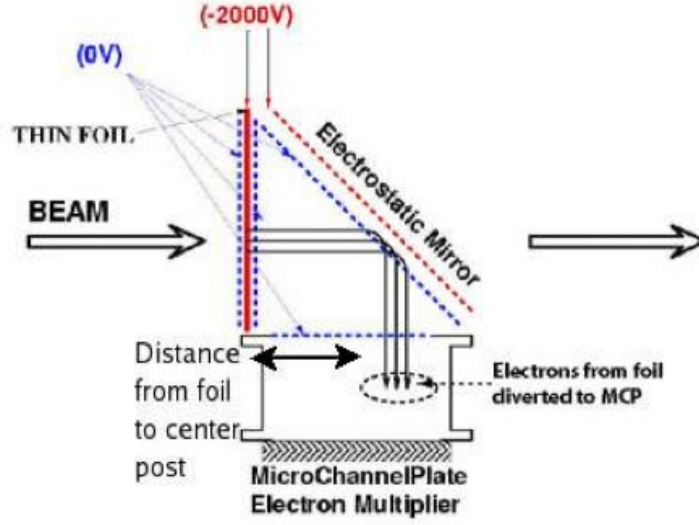


Figure 14: This is the illustration of the MCP setup. The beam comes in from left to right. The beam continues mostly undisturbed but electrons are liberated from the carbon in both directions. The electrons moving to the right are deflected down to the microchannel plate for multiplication and detection.

figure 14. The MCP detector, which will be used in future LeRIBSS experiments, consists of a triangular apparatus with very thin gold wire grids on two of the sides. The beam moves in from the left through a thin carbon foil of $5 \mu\text{g}/\text{cm}^2$. This small thickness is required because of the low energy of beam particles involved. Greater thicknesses would slow or stop the beam particles in this foil. The Bethe formula gives the linear stopping power, or differential energy loss per unit length through the material. This differential energy loss per unit length is defined as:

$$\frac{-dE}{dx} = \frac{4\pi e^4 z^2 N B}{m_o v^2} \quad (45)$$

Where e is the electron charge, z is atomic number of the beam particle, m_o is the rest mass of an electron, and v is the velocity of the beam particle. The B expression is given as follows:

$$B = Z \left[\ln \frac{2mv_o^2}{I} - \ln \left(1 - \frac{v^2}{c^2} \right) - \frac{v^2}{c^2} \right] \quad (46)$$

I is the mean excitation potential for the specific material, natural carbon in graphite form in this setup. The Bloch approximation says that this I is $10\text{eV} \cdot Z$, where eV is the energy unit electron-volts and Z is the atomic number of the target. In this case $Z=6$, so $I = 60 \text{ eV}$ by this

approximation. For smaller velocities, i.e. non-relativistic, it is reasonable to neglect the last two terms in the B equation, so only the first logarithmic term remains[15]. The important point is that the stopping power is inversely proportional to the particle's kinetic energy and particles with greater charge will have a greater stopping power because of the z in the numerator outside of the logarithmic term.

It is important to note that the above Bethe formula concerns electronic stopping power. This pertains to inelastic collisions between the ions and bound atomic electrons in the material. At lower energies and with heavy beam particles it becomes important to take into account nuclear stopping power. This comes from the elastic collisions between the ions and atomic nuclei in the material. This type of stopping power can be derived from the interatomic Coulomb repulsion between the atoms. Typically the nuclear stopping power peaks at an energy of 1 keV/nucleon. In this work the beams are around mass $A=81$ and the energy is 200 keV. Thus the energy is higher than 1 keV/nucleon (200 keV is higher than 81 keV) and so more of the total stopping power originates from the electronic type rather than the nuclear type.

When the beam particles pass through the foil electrons are liberated because kinetic energy from the beam is sufficient for ionization in the carbon atoms. The liberated electrons are repulsed from the carbon foil in both directions because it has -1000 volts of potential on it. This foil is surrounded on both sides by thin gold wire grids both at zero, or ground, potential to keep the thin foil from bending. The electrons move away from the foil towards wire grids on the hypotenuse of the triangular setup, which act as an electrostatic mirror and bend the electrons into the microchannel plate. This plate consists of three layers: an anode, a back, and a front. It has voltages across the anode and back while the front is kept at ground potential. It is made of a ceramic material but coated on one side with conducting metal. The surface contains periodically spaced holes of diameter about $4\mu\text{m}$ and depth $100\mu\text{m}$. Because an electrostatic potential exists across its depth and as the electrons come in to hit the inside walls of the holes, the amount of electrons is multiplied. All these electrons gather at the opposite side and a large enough signal is then sent to the data acquisition.

The MCP can be used to supply a timing signal for the arrival of a beam particle before it implants on tape. This timing signal can be used as a start signal for subsequent decays measured in the detectors. In this work the tape movement signals were used as a start signal for the grow and decay of activities.

3.3.6 Digital data acquisition

The information read out from HPGe detectors, plastic β decay scintillators, the electrostatic deflector, the tape movement, MCP, etc. are all sent into a number of Pixie16 digital data acquisition boards. Each board can take in up to sixteen different signals. Each signal has to pass through the analog conditioning part, which consists of an offset adjustment unit and analog Nyquist filtering [17]. This analog signal then goes through an Analog to Digital Converter, or ADC. This involves sampling the analog pulse once every 10 nanoseconds to measure its voltage. Thus the analog signal turns into a digital signal with a single voltage value every 10 nanoseconds. Thus the ADC has a sampling rate of 100 million samples per second.

After digitization, the signal passes through a digital filter for measuring time and energy with various parameters in this filter. Every four signals are grouped together and managed by an FPGA, or Field Programmable Gate Array. It is this array which manages those parameters for running filters over the digitized signals. A diagram showing the full layout of a pixie16 board is shown in figure 15.

A so-called trapezoidal filter is used to determine the energy of the digitized pulse, see figure 16. There are two regions over the time scale where summing of counts occurs. At each point in time the number of counts in each red summation region in the figure is recorded. Then these two summation regions are moved to the right in time to the next 10 ns bin and a record of counts in each red region is taken again. At each of these points in time the number of counts in the left region is subtracted from the number of counts from the right region. These differences are then plotted at each point in time, which results in the pulse seen in the middle row of the figure. This is how the “trapezoidal” pulse is formed from this trapezoidal filter. The lowest row in the figure shows another trapezoidal pulse from a trapezoidal filtering, but with a shorter time length in the two red summation regions, the integration times. Thus a digitized pulse is transformed into a pulse of trapezoidal shape, the amplitude of which is proportional to the energy of the signal. Thus the energy is extracted from the vertical position of the plateau in this resultant, trapezoidal pulse shape.

An example of experimentation performed in this research with Pixie16 trapezoidal filters is shown in figure 16. A parameter called the TRIGGER_THRESHOLD controls when a pixie channel will trigger on the incoming signals or ignore it as noise. Shown in figure 17 is the recorded efficiency for detecting a known pulser signal versus the trigger threshold parameter

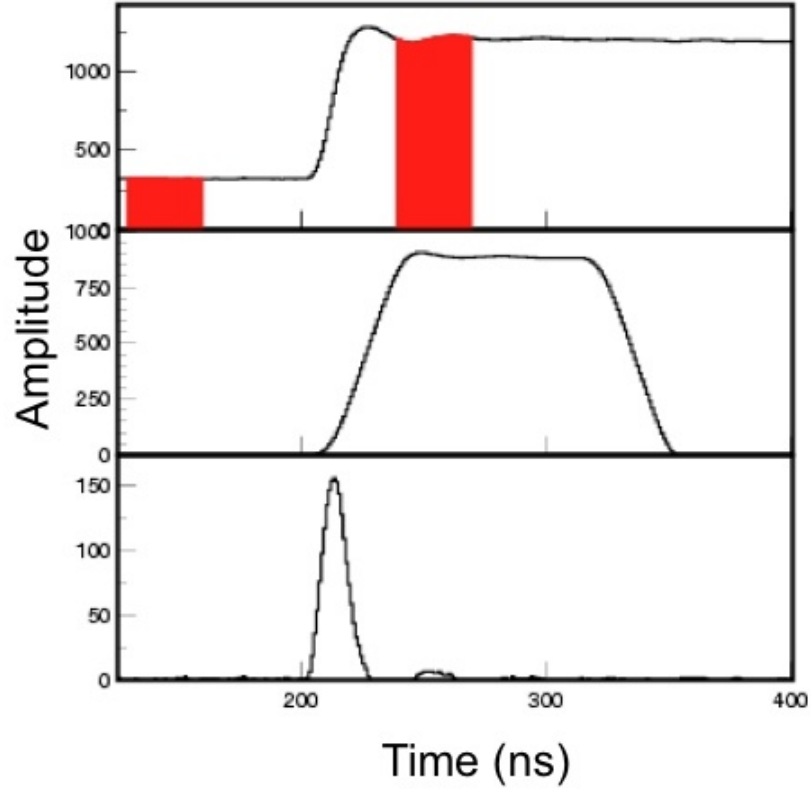


Figure 16: Trapezoidal filter for determining the energy from a digitized pulse in a Pixie16 channel. The top panel shows the digitized signal, which originally came as an analog signal from a preamplifier connected to a germanium crystal, with the two regions marked for summation. At each point in time the summations are subtracted one from another, left summation region subtracted from the right one, and this value from the subtraction is graphed at each 10ns bin in time. This results in the trapezoidal shaped signal shown in the center row. The bottom panel is also a resultant shape from the trapezoidal filtering but from shorter integration times, i.e. shorter times in those red summation regions.

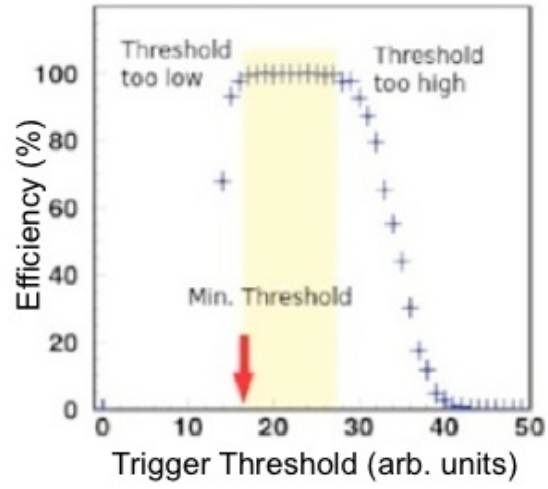


Figure 17: Trigger threshold vs. efficiency of detecting a pulser signal in a channel in Pixie16.

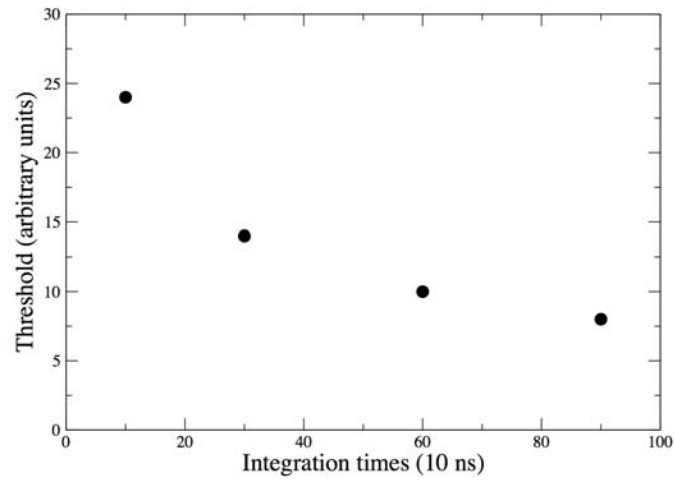


Figure 18: Integration time vs. minimum trigger threshold for obtaining 100% detection efficiency of a pulser signal.

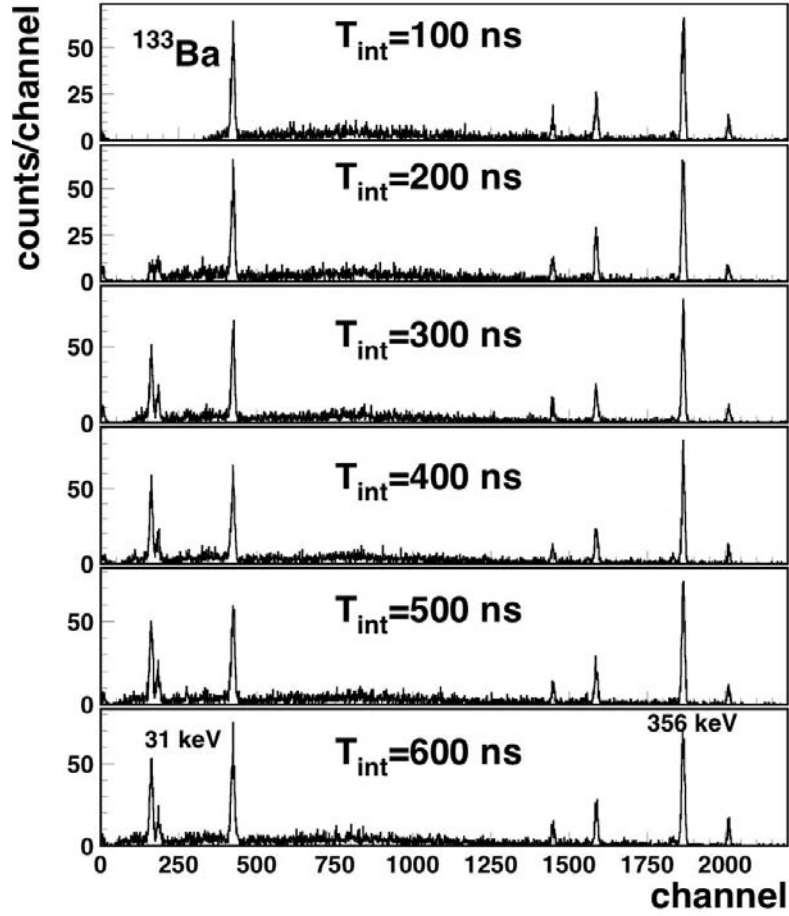


Figure 19: A ^{133}Ba source of γ -rays measured in the HPGe detectors. Increasing the integration time parameter in the trapezoidal filter increases sensitivity for low amplitude, low energy signals. Thus the 31 keV peak is more efficiently detected with longer integration times in the filter because the number of counts detected increases with increasing integration times in the filter.

on the horizontal axis, in arbitrary units. There is a range of trigger threshold values for which the pulser signal is recorded with 100% efficiency. A minimum threshold for detecting 100% of the signal is also indicated. Low trigger thresholds are important when measuring, for example, low intensity beams with low energy signals and the experimenter wishes to acquire all of these signals efficiently against an equal or nearly equal background of electrical noise. Given a signal from a pulser with a specific amplitude, one can measure the range of trigger thresholds needed to measure this signal with 100% efficiency, see figure 17. This is of course only valid when keeping all other parameters in the FPGA constant throughout the measurement [18].

Another important parameter includes the integration time in the trapezoidal filter. This refers to the amount of time each summation region covers in the filter, seen in figure 16. The relationship we measured between this integration time and ability to detect a pulser signal is shown in figure 18. The conclusion from this measurement is that with longer integration time the efficiency of detecting the pulser signal improves. It improves from 100 to 600 ns but from 600 to 900 there is little improvement. Thus it appears that summing over a larger time range improves detection efficiency, but this improvement is limited to some maximum amount of time, about 600 ns.

Another example of this can be seen when measuring a source of γ radiation in front of the HPGe detectors, see figure 19. As integration time in the trapezoidal filter is increased, the γ -ray transition of 31 keV in a ^{133}Ba source is measured more efficiently. An increase in the measured counts in this peak is seen in figure 19 as integration time is increased from the top panel to the bottom.

The conclusion here is that a longer integration time was helpful in improving the acquisition of a small signal out of a nearly equal sized electrical noise background. This results in improved ability to detect low γ -ray energies amongst electrical background noise. One final note on integration times in the trapezoidal filter: Integration times for the different detectors (HPGe, scintillators, etc.) are determined mainly by their different rise times. The shorter the rise time, for example with scintillators on the order of 1.8 ns, a smaller integration time will be needed compared to the other detectors.

4 Experimental Analysis and Results: ^{81}Zn

4.1 Data reduction methods

A raw data file, which contains all the energies, time stamps, etc. for each measurement in the detectors must be scanned using computer code to produce useful spectra for analysis. The first example is that of a singles γ -ray spectrum. This is a one-dimensional spectrum which contains all counts of γ -rays measured in the HPGe crystals. Before moving on to other spectra which need a “gate”, or a requirement of more than one measurement in coincidence, one must define what is referred to as the temporal width of an event window. This is the amount of time in which all measurements taken are considered to be part of the same “event.” The time length of an event is determined from a histogram which shows the time differences between all measurements. This histogram has a two large peaks, one at time zero to 10 ns and another around $1.33\ \mu\text{s}$. This means that most events (physics processes such as a β decay followed promptly by γ -ray emission) are within either 10 ns of each other, the smallest time interval which can be measured in pixie16 digital acquisition boards used here. The other peak indicates a spike of events with time difference $1.33\ \mu\text{s}$ between them. However, there are many other events with time differences of 20ns all the way up to $2.5\ \mu\text{s}$. The goal here is to choose a time length to define an event which is not too small such that many valuable measurements are excluded, but not too long such that certain measurements are correlated unduly. This would lead to false coincidences in the analysis. Thus in this experiment the event length was chosen to be $2.5\ \mu\text{s}$ because it includes all of the appropriate peaks in this spectrum.

The next important γ -ray spectrum is the β -gated γ spectrum. A measurement of a γ -ray in the HPGe cannot be placed into this spectrum unless a β decay electron is also measured in the same event window time ($2.5\ \mu\text{s}$) along with it. The point is that all of the γ -rays in this spectrum must be in coincidence with a β decay that they followed, thus only β -delayed γ -ray transitions are placed in this spectrum. The γ -rays seen here should only originate from a β decaying nucleus from the implanted beam ion and not from background sources.

Another important spectrum is a two dimensional one of time vs. energy of γ -rays, see figure 20. The vertical axis contains the time at which counts of a γ -ray are collected for all the γ -ray energies. This time is in reference to a specific movement of the moving tape collector (MTC). The cycle begins with the signal which stops the tape movement. It ends after the tape movement ceases and that same signal is given again. The horizontal axis is contains the energies of

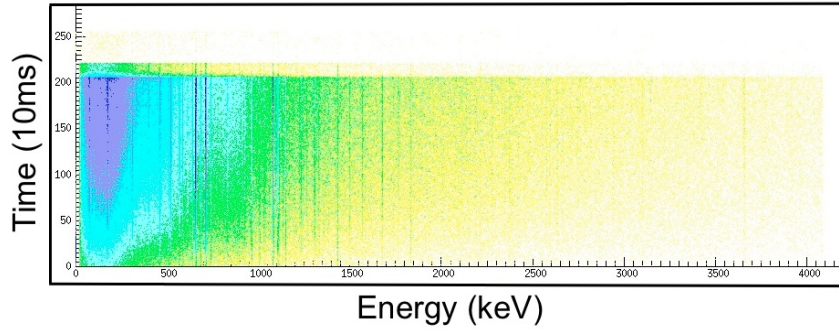


Figure 20: The two-dimensional spectrum where a “cut” can be made on a specific γ -ray energy on the horizontal axis to create a one dimensional spectrum showing a grow-in and/or decay out of activity for that one specific γ -ray energy.

those γ -rays. A “cut” in this two dimensional spectrum on a specific energy peak will produce a one-dimensional spectrum showing counts versus time for that particular γ -ray energy. Thus it is likened to making a slice in the vertical direction but only at a specific energy on the horizontal axis. The resultant one dimensional plot will contain a grow-in and/or decay curve. The experimenter can then fit a function to this curve and obtain a half-life for the specific γ -ray transition in question. The half-life obtained will only be altered by the amount of statistics that can be gathered in said time window. The more statistics the smaller an error bar which can be obtained on the half-life measurement. However, if a long decay out time is recorded then the full γ -ray spectrum will contain decays from the other isobars’ decays (long lived activity from daughter, granddaughter, etc.), quickly reducing the ability to observe weaker transitions from the parent decay in the spectrum.

Furthermore, another series of important two dimensional spectra pertain to coincidences between different γ -ray emissions occurring in cascade with each other. A γ - γ matrix must be created to obtain this information. The basic idea is that a count is placed in the spectrum if a certain γ -ray is measured within the same event window as another γ -ray. The count will be placed at a position such that its x-coordinate is the energy of the first γ -ray detected and the y-coordinate is the energy of the other γ -ray measured in the same event. This is further explained in the section on β - γ - γ coincidences below. Thus a count is placed where two different γ -rays were emitted in cascade of each other from within the same nucleus after β decay.

Finally, one more useful spectrum is that which shows the energy of each γ -ray vs. the time delay between detection of the β decay electron and its subsequent γ -ray emission. Every measured event has a time stamp on it, recorded in the digital electronics. Thus such correlations are possible through the digital data acquisition system. This spectrum is able to reveal isomers in the daughter nucleus because of this time delay. The typical range of isomers detectable with this setup ranges from about 100 ns to a few microseconds. In this particular data set the event window length was set at $2.5\mu\text{s}$. Therefore, $2.5\mu\text{s}$ is the maximum amount of time for an observable isomer in this manner. Isomers with longer decay time should be detectable by measuring a specific γ -ray emission from the isomeric state's decay and observing a significantly different half-life than from the parent or daughter ground state decays.

4.2 Construction of the decay level scheme

The β decay of ^{81}Zn populates either the ground state, an excited state, or a neutron unbound state in ^{81}Ga . Populating a neutron unbound state leads to neutron emission and population of states in ^{80}Ga . A decay scheme of ^{81}Zn is constructed based upon measurements of β and γ radiation taken at LeRIBSS. When a nucleus undergoes β decay to a daughter nucleus, a number of processes can occur. A one step process takes place where the parent nucleus decays directly to the ground state of the daughter and no γ transitions are emitted until this daughter nucleus itself eventually decays. Also possible is a two step process wherein the parent nucleus decays to an excited state of the daughter and the daughter subsequently decays to either another, lower excited state or its own ground state via γ -ray emission. The energy of this γ -ray photon is equal to the energy difference between the excited state from which it decayed and the lower energy state it decays to in the daughter. The nucleus may also decay to a higher excited state and γ decay multiple times to reach the ground state of the daughter nucleus. Another process that is possible is where the parent nucleus undergoes β^- decay to a state above the neutron separation energy. That is, it decays to a neutron unbound state and thus the daughter nucleus strongly prefers emitting a neutron as opposed to a γ -ray. The emission of a neutron leads to population of an excited state (or ground state) of a nucleus with one more proton but two less neutrons.

We use three concepts as guidelines to construct a decay level scheme showing these different processes: Intensities of the γ transitions, $\beta-\gamma-\gamma$ coincidence data, and half-lives of each γ -ray transition.

4.2.1 Intensity balance

The intensity is defined as the measured counts in a peak divided by the efficiency of detecting the γ -ray at that particular energy. The counts come from a gaussian fit to the peak, here done with the DAMM spectroscopy computer program [19]. This gaussian peak fit includes a fitting of the background in the γ -ray spectrum. The definition of intensity may need to include the efficiency of any other detector used to detect those counts. For example, the efficiency of detecting β decay electrons in the scintillators would be required while using counts in a β -gated γ spectrum.

There is a summing rule of intensities needed while building a decay scheme. The intensity of all γ transitions feeding a nuclear state must be less than or equal to the total amount of intensity from γ -ray transitions decaying from that nuclear state. Extra intensity which is decaying from the nuclear state and which cannot be accounted for from the γ -ray transitions feeding it from higher energy states must originate from β decay intensity directly to said state. A schematic of this rule for intensity balance is shown in figure 22. The intensity of the γ transitions labeled B and C must add together to be either greater than or equal to the intensities of γ transitions D and E. The net difference is labeled F, the intensity that must directly feed level A through β decay of the parent nucleus to the daughter. Thus the intensity of the β decay directly feeding level A can be either positive or zero, within error bars. This is a rule which cannot be violated simply because an extra intensity feeding a state from above cannot simply disappear, it must appear as some other form of decay away from that state to lower energy states. This rule is then used as guidance in constructing the decay level scheme with γ -ray transitions and their known intensities.

4.2.2 The $\beta - \gamma - \gamma$ coincidences

Two or more γ -ray transitions that are emitted from the decay in coincidence with each other, i.e. in the same event window of $2.5\mu\text{s}$, must be transitions in cascade within the decay level scheme. A two-dimensional histogram is created for acquiring these coincidences after scanning the raw data file. A count is placed at the appropriate coordinates to show which two γ -ray transitions were emitted in coincidence with each other. For example, in figure 22 if γ -rays C and D are measured in coincidence then a count is placed at the coordinates for the energy of γ -ray C on the horizontal axis and at the energy of γ -ray D on the vertical axis. Also, a count is placed at the energy for γ -ray D on the horizontal axis and one at the energy for γ -ray C on the vertical axis. Thus a “cut” can be made in the two-dimensional spectrum around the desired

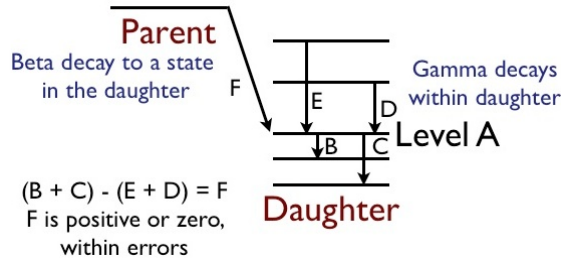


Figure 22: The intensity decaying into level A is less than or equal to that of the intensity decaying out. The intensity of direct decay, or side feeding, to that state is given by the amount decaying out minus the amount decaying into it from higher energy states in the decay window.

energy peak such that a one-dimensional spectrum of counts vs. energy is created. Therefore it will show all the γ -rays in coincidence with said γ energy peak. Here a “cut” on the energy for γ -ray C should produce a one-dimensional spectrum with a peak at the energy for γ -ray D, thus showing their coincidence relationship. A cut in the two-dimensional spectrum on γ -ray D will produce a similar one-dimensional spectrum, but with a peak at the energy of γ -ray C. These two γ -ray transitions must be placed in a direct cascade with each other. The order is then determined by the intensity balance and any other γ -ray transitions seen in coincidence with C and D.

4.2.3 Half-life measurements

It is important to identify the origin of a γ -ray transition by measuring its decay half-life. This will match it with the correct decaying nucleus, for example the desired parent decay versus the daughter or granddaughter decay.

The two-dimensional spectrum for energy of counts on the horizontal axis and time of measurement for each count on the vertical axis has the purpose of providing grow-in and/or decay curves for γ -ray measurements. Thus a “cut” in this spectrum at a desired energy for a specific γ -ray will yield a time profile for the transition. Therefore, an appropriate grow-in of activity and/or decay out activity curve may be fit to it. This allows for the determination of a half-life for that specific transition. The grow-in and decay equations are shown with the half-life kept as the free parameter along with an amplitude A to the fit. Here t is the independent variable time, T is the time at which the grow-in part ends and the decay out part begins, the decay constant λ is defined with the half-life in it, and A is the amplitude of the curve. All transitions de-exciting

states in the daughter nucleus should have a half-life matching the decay half-life of the parent. These half-life measurements help to identify γ -rays as transitions belonging to either the parent decay, daughter decay, β -delayed neutron emission, or as a γ -ray from background. Thus these three measurements of $\beta - \gamma - \gamma$ coincidences, intensity balance, and half-life all provide means by which to identify and properly place γ -rays in the decay scheme or elsewhere.

The following equations are simple Bateman equations for the free decay of activity, where the different variables are explained above and shown below. The first equation applies to the grow-in of activity. The second one applies to the decay out of activity, where the amplitude is given by the same expression for the grow-in half but at the time T in which the free decay part begins. Thus it is the amplitude of the free decay part. The third equation below is the decay constant λ , depending upon both half-life and using the constant $\ln(2)$.

$$f(t) = A * (1 - \exp^{-\lambda t}) \quad (47)$$

$$f(t) = A * (1 - \exp^{-\lambda T}) * \exp^{-\lambda(t-T)} \quad (48)$$

where

$$\lambda = \frac{\ln(2)}{T_{1/2}} \quad (49)$$

4.2.4 β efficiency calculation

Many of the γ -ray transitions seen in this ^{81}Zn decay experiment are only observable with a gate on a β decay as well. Therefore, not only is the efficiency of γ -ray detection needed, as explained before, but efficiency for detection of the β decay electron is also needed. When a neutron-rich parent nucleus undergoes β^- decay, it releases an electron and an anti-neutrino. The energy of both of those particles must equal the energy difference between the decaying state in the parent and the populated state in the daughter nucleus. Thus the possible energies of the outgoing electron resulting from the population of a certain state in the daughter form a non-symmetric gaussian like distribution of energies. The amount of electrons at any given energy available to it after decay to the daughter state and release of the anti-neutrino forms this distribution. Furthermore, any specific level in the daughter which decays via a γ -ray can be populated either through direct β decay from the parent or through a γ decay from a higher energy state in the daughter. Thus the experimenter must take into account both direct population through

β decay as well as β decay to higher energy states in the daughter with subsequent γ decays populating the state of interest. It is the energy of these β decay electrons which determine the efficiency by which they are detected in the scintillators. A calculation of a so-called Q effective value, Q_{eff} , is needed for each γ -ray transition observed in the β -gated γ spectrum. It is mostly this spectrum which is used for building the level scheme since most of the transitions are only observable in the β -gated γ spectrum.

This value should take into account both the actual Q value down to said state in the daughter as well as the Q value for direct β decays to states above the state of interest, which feed it via γ decays. For example, in figure 22 a Q_{eff} value is needed for the γ transition A. If the level A is only populated through direct β decay and no γ decays from higher energy levels, then the Q_{eff} is equal to the Q value, or energy difference between the ground state of the parent and that particular state in the daughter. However, if there are β decays to states above level A and thus subsequent γ decays down to level A, then the Q_{eff} value will be a lower value than just the Q value, taking into account lower energy electrons detected in the scintillators in coincidence with the γ -ray transition A. The two equations below show how the Q_{eff} value is calculated as the difference between the Q value to that state minus an E' value, which takes into account feeding from γ -ray transitions above and total decay out of the level in question. [20]

$$Q_{eff} = Q - E' \quad (50)$$

$$E' = \frac{[\Sigma I_{\gamma out} - \Sigma I_{\gamma in}]E_{level} + \Sigma I_{\gamma} E'_{level}}{\Sigma I_{out}} \quad (51)$$

Here the E prime is calculated with two terms in the numerator and one in the denominator. The first term in the numerator is the total γ intensity de-exciting the nuclear level minus the total γ intensity feeding the level from energy levels above. Then this difference is multiplied by the energy of the level. The second term in the numerator sums all of the γ intensities feeding the level from above, each multiplied by the energy of the level from which it decays. The term in the denominator refers to the total γ intensity de-exciting the level. Thus this E prime value is the amount by which the normal Q value down to a state should be modified by the β decay to states above this level in question. All γ transitions de-exciting this level will have the same efficiency of β electron detection in the scintillators because they have the same weighting of the Q value, i.e. the same Q_{eff} . Once a Q_{eff} value is obtained for a given γ -ray transition de-exciting the level, there is a relationship between the Q_{eff} value and the efficiency

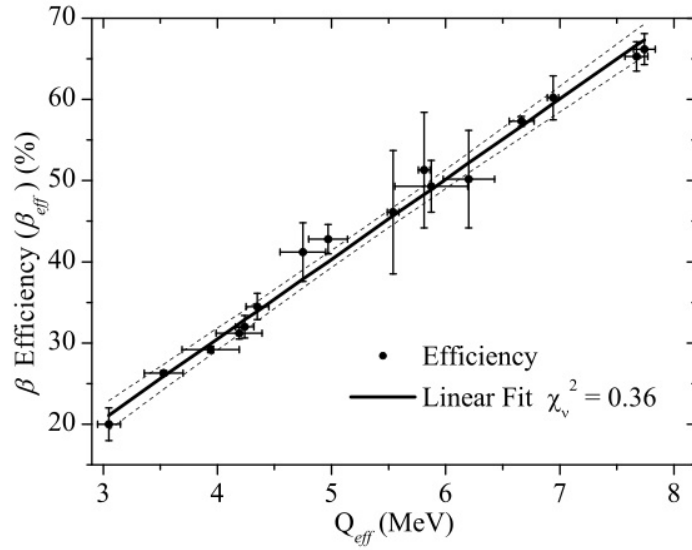


Figure 23: The experimental values for β detection efficiencies vs. the Q_{eff} values on the horizontal axis. A linear fit is also shown here for a shorter range of Q_{eff} values in the work by J.A. Winger and S. Ilyushkin [20]. However, to accommodate a wider range of Q_{eff} values needed in the ^{81}Zn decay, the non-linear equation discussed in their work and this text is utilized instead.

of detecting the β decay electron emitted from the decays to that level. This relationship was determined experimentally for the setup at LeRIBSS [20] from γ -rays that could be measured in both a singles γ spectrum as well as the β -gated γ spectrum. The efficiency for those transitions which can be observed in both types of spectra is simply the ratio of counts in the β -gated spectrum over the counts observed in the singles spectrum. Then a fit is made to these data points such that the β detection efficiency of transitions which can only be measured in the β -gated γ spectrum can be determined. The experimental points for the Q_{eff} values and the β detection efficiencies for various γ -ray transitions are shown in figure 23.

The equation used to fit these points is the following:

$$\beta_{eff} = 100 * (1 - 0.8315^{Q_{eff}-1.993}) \quad (52)$$

Where the β efficiency is in percent and the Q_{eff} is in MeV. The constants originate from the fit performed to known β detection efficiencies by S. Ilyushkin *et al.* [20]. Table 4 shows the γ -ray transitions in ^{81}Zn decay with their calculated efficiency of detection of the β decay

Table 4: Table showing the actual Q_β value down to each level in ^{81}Ga as well as the calculated Q_{eff} values. The efficiencies of detecting β decay electrons emitted after population of each state are shown. The intensities for the 4294 and 4880 keV transitions were extracted from the singles addback spectrum and thus their β electron detection efficiencies are not needed.

Level (keV)	Q value (MeV)	Q_{eff} (MeV)	β eff. (%)
351	11.077	10.784	80.3(9)
802	10.626	10.626	79.7(10)
1267	10.161	10.161	77.8(40)
1458	9.970	9.970	77.1(40)
1936	9.492	7.486	63.7(190)
4294	7.134	7.134	–
4880	6.548	6.548	–

electrons. Also, the total Q_β and Q_{eff} values down to the levels which depopulated via those γ -rays are shown.

4.2.5 Determination of isobaric contamination

It is important to determine the amount of isobaric impurity present in the beam of particles implanting on tape. This is because the amount of intensity in the daughter decays is needed to calculate a ground state to ground state branching ratio. If there are two sources of the daughter to granddaughter decay radiation, then this ground state to ground state branching ratio will be inflated from the real value. This is because decays of isobaric contaminants will be indistinguishable from the parent to daughter to granddaughter decay radiation. This is the main motivation for determining the amount of isobaric contamination. It must be calculated in order to readjust for the ground state to ground state branching ratio.

To determine the percent of beam particles which are the neighboring isobar, in this case the undesired ^{81}Ga versus the desired ^{81}Zn , one must fit the grow-in curves of γ -ray transitions from both the parent and the daughter decays. The following differential Bateman equations are used for the parent and daughter decay activities, which include the implantation rates for the parent zinc and the daughter gallium.

$$\frac{dA_1}{dt} = I_1 - \lambda_1 A_1(t) \quad (53)$$

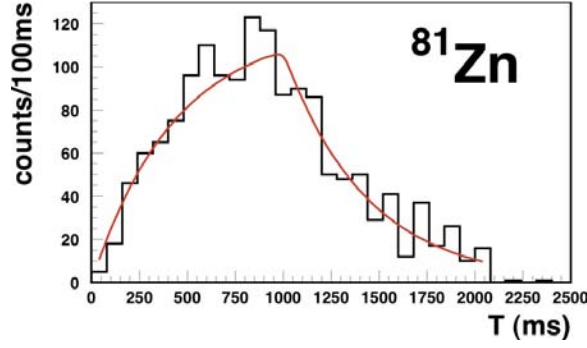


Figure 24: The fit of the Bateman equation solution (A_1) to the parent decay activity for the grow-in portion of its time profile. A simple exponential decay is also fitted to the second half of the time profile. The second half is simply a free decay out because the beam was deflected away from the site of implantation.

$$\frac{dA_2}{dt} = I_2 - \lambda_2 A_2(t) + \lambda_1 A_1(t) \quad (54)$$

The solutions to the Bateman equations are the following:

$$A_1(t) = C_1 \left(\frac{I_1}{\lambda_1} - \frac{I_1 \exp^{-\lambda_1 t}}{\lambda_1} \right) \quad (55)$$

$$A_2(t) = C_2 \left(\frac{1}{\lambda_2(\lambda_1 - \lambda_2)} \right) (-I_2 \lambda_2 + I_2 \lambda_1 - I_1 \lambda_2 + I_1 \lambda_1 + I_1 \lambda_2 \exp^{-\lambda_1 t} + I_2 \lambda_2 \exp^{-\lambda_2 t} - I_2 \lambda_1 \exp^{-\lambda_2 t} - I_1 \lambda_1 \exp^{-\lambda_2 t})$$

Where A_1 is the number of parent decays per unit time as a function of time, A_2 is the number of daughter decays per unit time as a function of time, and λ_1 and λ_2 are the decay constants for parent and daughter, respectively. The t is time and I_1 and I_2 are the implantation rates for parent and daughter, respectively. The solutions also come from initial conditions such that at time $t=0$ the functions N_1 and N_2 are also zero. The constants C_1 and C_2 are the multiplication of the three following quantities: Efficiency of γ detection, the efficiency of β electron detection for the particular γ -ray utilized, and the percent of the total β decay intensity that this particular γ -ray carries away in the decay.

These two solutions to the above differential Bateman equations are fitted to the grow-in profiles of the strongest γ -rays observed from the decay of the parent (the equation for N_1) and the daughter (the equation for N_2). The decay constants (λ_1 and λ_2) contain half-lives, which

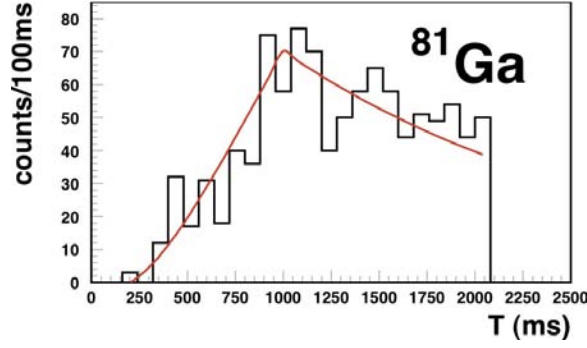


Figure 25: The fit of the Bateman equation solution (A_2) to the daughter decay activity for the grow-in portion of its time profile. A simple exponential decay is fitted to the second half of this time profile. The second half here is also a free decay out because of deflection of the beam away from the site of implantation.

are placed in as constant numbers and not as parameters varying with the fit. The only varying parameters are the implantation rates, I_1 and I_2 .

The rates obtained are as follows: In the fit shown in figure 24 the rate of zinc implantation was 22.91(75) ions per second. The rate of zinc implantation obtained from the daughter curve is 23.12(19) ions per second. The fit to the daughter activity grow-in is shown in figure 25. From this fit to the daughter activity the I_2 , or rate of implantation for the daughter gallium, is 0.00219(2) ions per second, essentially zero within error bars. Thus this initial result shows that the rates of zinc implantation obtained from both the parent and daughter activities agree with each other. Furthermore, the rate at which gallium is directly implanting on tape is zero, within errors.

More specifically, a closer inspection of this fit reveals an upper limit on the isobaric contamination rate. At three standard deviations in the fit, the I_2 is 5% that of the I_1 . This means that at most the isobaric impurity is 5%. That is, for every 100 zinc ions implanting on the tape there are at most 5 ions of gallium implanting on tape. Thus this is the conclusion reached, that the isobaric impurity is 5%. As a result, in this experiment it can be stated that a nearly pure beam of ^{81}Zn was obtained with little ^{81}Ga contaminants. This information will be utilized to adjust the ground state to ground state branching ratio obtained when the decay level scheme and β decay branching ratios are calculated.

Table 5: The γ -ray transition details for ^{81}Zn . Each energy and relative intensity is shown. Lifetime column states whether fits match the lifetime of ^{81}Zn decay. A column for coincidences is shown along with previous literature energies and intensities.

E_γ keV	I_γ	lifetime	coincidence (keV)	literature E_γ	lit. I_γ
350.9(3)	100(4)	yes	451, 916, 1107, 1585, 2358	351.1	100 [2]
451.6(4)	19(3)	yes	351	451.7	20 [2]
802.2(7)	6(3)	yes	—	no	—
916.2(8)	5(3)	no	351	no	—
1107.6(9)	4(3)	yes	351	no	—
1458.3(12)	4(2)	yes	—	no	—
1585.4(13)	9(5)	yes	351	no	—
1936.3(17)	10(6)	yes	—	no	—
2358.4(20)	17(7)	yes	351, 1585	no	—
4294.0(40)	13(4)	no	—	no	—
4880.0(42)	19(5)	no	—	no	—

4.3 Level scheme for ^{81}Zn decay

The information on γ -ray transitions placed in ^{81}Zn decay is shown in table 5. The transitions at 350.9 keV and 451.6 keV were also previously measured [2]. The nine other transitions are all new ones measured in this experiment. They were first identified in the β -gated γ -ray spectra followed by $\beta - \gamma - \gamma$ coincidence data to verify their placements in the level scheme, along with the intensity balance. Finally, half-lives were measured for these transitions to further verify whether or not they belong to the decay of ^{81}Zn . Because of low statistics many of these transitions were difficult to fit for half-life determination. However, the important point is whether or not they are significantly closer to the half-life of ^{81}Zn , 304(13) ms, instead of the ^{81}Ga daughter decay, 1.217(5) seconds [21]. These are the only two possible nuclei each transition can decay from, once the possibility of a transition belonging to background has been eliminated.

The 351 keV transition can only be analyzed with its peak in the β -gated γ spectrum because there is an omni-present background peak at 351.9 keV from radium decay to ^{214}Pb . Such background lines will not appear in the β -gated γ spectrum because the majority of β decay electrons from γ -rays in the background will not be detected in the plastic scintillators. The 451 keV is also observed along with its coincidence with the 351 keV transition, just as in the previous

publication [2]. However, here a crossover transition of 802 keV is also observed. First one has to verify that this peak at 802 keV is more than just the sum of a 351 keV γ -ray and a 451 keV γ -ray. This involves using the efficiency of detection for each γ -ray and multiplying them together to calculate the efficiency of detecting both γ -rays summed together. Then from this efficiency and the measured intensity of each individual transition added together, one can calculate the expected number of counts in this crossover transition. Here the expected number of counts is 11.5 above background. However, we measure 57.5 counts above background. Therefore, we conclude that the 802 keV peak represents a real transition and more than just the summation of the 351 and 451 keV γ -ray transitions.

4.3.1 The 802 keV transition and M1 hindrance

The transition from the second excited state to the ground state (this 802 keV transition) is postulated to be an M1 transition. This assignment is based upon shell model calculations. According to these calculations the second excited state in ^{81}Ga should have $J^\pi = 3/2^-$. The first excited state also has a $J^\pi = 3/2^-$, although the wave function there may be more mixed, as discussed below. The ground state is predicted to have $J^\pi=5/2^-$ and the spin was measured to be 5/2 in an experiment at ISOLDE-CERN [23]. Thus a transition from the second excited state to the ground state and another from the first excited state to the ground state should involve the same change in spin J. However, one involves more energy than the other. The 802 keV M1 transition is greater in energy than the 351 keV M1 transition. But a mathematical relation for the decay constant for transitions between nuclear states is needed in order to discuss any of these transitions. It is given by equations pertaining to Weisskopf estimates [24].

According to Weisskopf estimates higher energy difference between the states leads to a shorter lifetime and this 802 keV transition should be more intense than the 351 keV transition based upon these arguments. However, the direct β decay intensity to the first excited state is higher than the second excited state. Thus a better comparison would be between the 451 and the 802 keV transitions. They both de-excite the same state and thus see the same side feeding. The 802 keV transition should be more intense than the 451 keV transition simply because of the greater energy difference, but this is not what is observed. This is explained by so-called M1 hindrance. The idea is that this 802 keV transition must be forbidden by its $\Delta\ell$, or change in orbital angular momentum, it experiences. Transitions between nuclear states are forbidden if the $\Delta\ell$ is greater than or equal to two. This suggests that the $J^\pi = 3/2^-$ second excited state may be a pure $\pi p_{3/2}$ state whereas the first excited state would not be

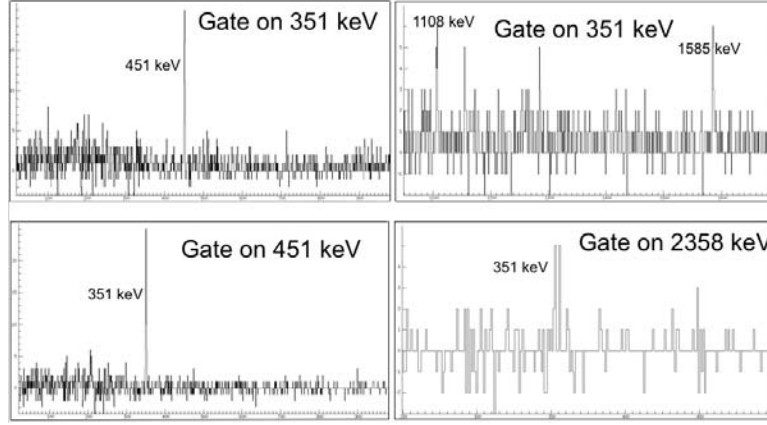


Figure 26: The $\beta - \gamma - \gamma$ coincidences are shown from gating on the 351, 451, and 2358 keV transitions.

a pure $\pi p_{3/2}$ state. Therefore, the first and second excited states may have wave functions with larger $\pi p_{3/2}$ components, with $\ell=1$, thus the 451 keV transition is not hindered by this $\Delta\ell$ -forbidden concept because its a transition between states which both have $\ell=1$. However, the transition from the second excited state to the ground state may involve a transition from a more pure $\pi p_{3/2}$ state, with $\ell=1$, to a more pure $\pi f_{5/2}$ state, with $\ell=3$. Thus this is a possible explanation as to why the 802 keV transition is more hindered than the 451 keV transition.

4.3.2 Coincidences and crossover transitions

To continue with the discussion of adding transitions to the decay level scheme, the 916 keV transition is placed in coincidence with the 351 keV transition, directly feeding the 351 keV level from a level at 1267 keV. The 1108 keV transition was also seen in coincidence with 351 keV and so it is placed as feeding that level from de-excitation of a level at 1458 keV. A crossover transition was observed at 1458 keV and was verified to be a transition which does not result from 351 and 1108 keV summation in the germanium detectors. This determination was accomplished with the same method as with the 802 keV transition.

The 1585 keV transition is seen in coincidence with both the 351 and the 2358 keV transitions. They are placed in a cascade with 1585 keV feeding the 351 keV level and the 2358 keV transition feeding the 1936 keV level. It then de-excites through the 1585 keV transition. The order is determined through the intensity balance. In addition, there is a transition observed at

1936 keV, thus further support is given to the establishment of a level at 1936 keV. The 2358 keV transition decays from a level at 4294 keV, where this crossover transition is also observed and decays from this level directly to the ground state. Finally, a transition of 4880 keV is observed and is placed as decaying from a level at 4880 keV to the ground state.

4.3.3 β decay branching ratios

As stated before, the percent of total β decay from the ground state of the parent nucleus to each observed state in the daughter nucleus is determined from the γ -ray intensity decaying from the state, then subtracted from the total γ -ray intensity feeding said state. The β decay branching ratio is defined as the amount of intensity feeding said state directly from the parent decay divided by the total amount of intensity observed feeding all daughter states from the parent. Here the β decay branching ratios are shown in figure 27 along with the completed decay level scheme.

The resulting β decay branching ratios also provide the ability to calculate a $\log ft$ value to each level in the daughter nucleus. This calculation is based upon the Q value down to each state, the proton number of the daughter, and the half-life of the parent decay. There are empirically established ranges for $\log ft$ values belonging to each type of β decay, whether it is a Fermi allowed decay, Gamow-Teller allowed, first forbidden, first forbidden unique, second forbidden, etc. Figure 6 shows the ranges of empirical $\log ft$ values assigned to different types of β decays [9], including first forbidden and first forbidden unique transitions.

4.3.4 Neutron emission branching ratio

The parent nucleus decays to a number of excited states in its daughter, including direct population of the ground state. But in the daughter nucleus there are energies high enough such that a single neutron will become unbound. All states above this so-called neutron separation energy are referred to as neutron unbound states. Thus direct β decays to these states will result in β -delayed neutron emission more often than γ emission. A neutron is emitted and thus populates a state in a nucleus with one more proton but two less neutrons. The reaction is represented in this equation:

$$Parent(A, Z, N) \rightarrow e^- + \bar{\nu} + Daughter(A - 1, Z + 1, N - 2) + n \quad (56)$$

The parent emits an electron and an antineutrino, as typical in β^- decay, but it also emits a neutron, n . The daughter in this case has a mass number one less than the parent and one

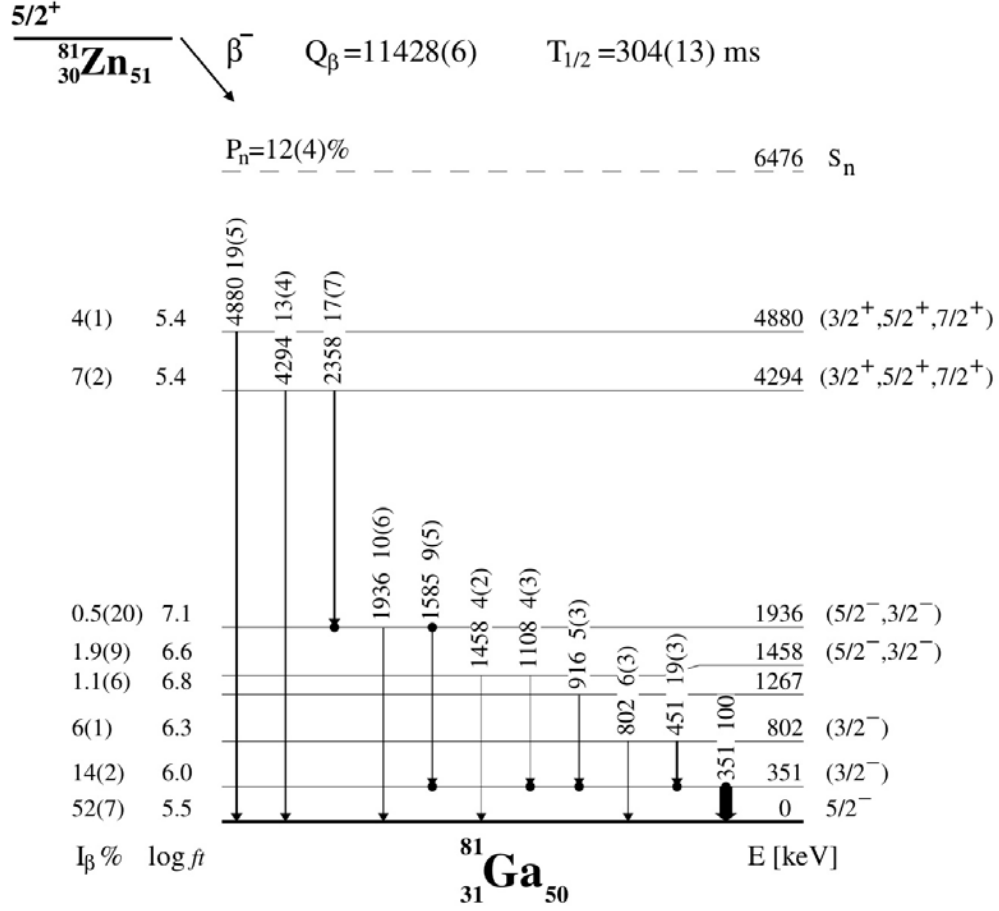


Figure 27: The complete decay level scheme for ^{81}Zn decay to ^{81}Ga . The 351 and 451 keV transitions, along with the first two excited states, were observed in a previous publication on ^{81}Zn decay [2]. The nine new transitions and five new levels are also shown. The population of the lowest states up to 1.936 MeV are through first forbidden decays of the parent and the population of the two observed states above 4 MeV are through allowed Gamow-Teller decays. Those two states above 4 MeV are identified as core excited states.

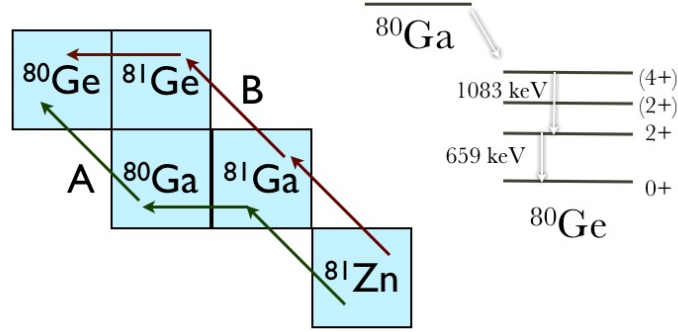


Figure 28: Derivation of the β -delayed neutron branching ratio from ^{81}Zn decay. The γ -ray transitions inside ^{80}Ge are from [25]. The neutron branching ratio is extracted from the 1083 keV transition because this state is less likely to be populated through β -delayed neutron emission from ^{81}Ga .

neutron is changed to a proton. However, the second neutron is emitted from the nucleus because it is in its neutron unbound states. The identification of β -delayed neutron emission must involve identifying a γ -ray emitted from either an excited state of this nucleus with (A-1, Z+1, N-2) or from this nucleus' subsequent β decay to its own daughter with (A-1, Z+2, N-3). Since it is β -delayed neutron emission from ^{81}Zn this involves observing the population of an excited state in either ^{80}Ga or a γ -ray transition from the decay of ^{80}Ga to ^{80}Ge . In this experiment a transition was only observed from the latter, the β decay of ^{80}Ga to ^{80}Ge .

The population of a state in ^{80}Ge decaying via γ -ray emission can occur via two different paths. These are shown in the figure 28. The 659 and 1083 keV transitions in ^{80}Ga to ^{80}Ge decay are possible transitions to utilize for calculating this branching ratio. However, one must calculate which of these two strongest transitions in ^{80}Ga decay is only populated through the β -delayed neutron emission of ^{81}Zn and subsequently ^{80}Ga β decay. The undesired problem is that one or more of these transitions are populated through the other path, the β -delayed neutron emission of ^{81}Ga . When the neutron is emitted from the neutron unbound states of ^{81}Ga it must obey the following angular momentum change rule.

$$|J_f - J_i| \leq \ell + s \quad (57)$$

The total angular momentum the neutron carries away between the initial and final states

must go into both the orbital angular momentum ℓ and the intrinsic angular momentum of the neutron s , which is $1/2 \hbar$.

The orbital angular momentum, ℓ , and the intrinsic angular momentum, s , of the neutron can couple to a number of values. These values are integer or half integer values between $|J_f - J_i|$ and $J_f + J_i$. For example, if the initial $J_i = 3/2$ and the final $J_f = 1$, then the orbital plus intrinsic angular momenta can both couple to $J=1/2$ or $J=3/2$ or $J=5/2$. Since the intrinsic angular momentum of a neutron is always $1/2$, this means that the ℓ value is $\ell=0$ for $J=1/2$, $\ell = 1$ for $J=3/2$, or finally $\ell=2$ for $J=5/2$. However, parity conservation also plays a role in determining the ℓ value which the neutron carries away. In cases where the parity does not change, the ℓ is an even number. When parity does change, the ℓ value is odd.

The amount of orbital angular momentum ℓ the neutron carries away will affect the potential barrier it must pass through. The potential, as a function of radial distance from the parent nucleus, is given by:

$$\frac{\hbar^2}{2m} \frac{\ell(\ell+1)}{r^2} \quad (58)$$

Here m is the mass of the neutron and r is the radial distance. Since a neutron has no charge it does not experience a Coulomb barrier, such as what a proton or alpha particle would experience. One must decide which orbital angular momentum values, ℓ , the neutron will carry away to determine the strength of this barrier. Thus there are a number of possible angular momentum barriers a neutron will experience as it is emitted from the neutron unbound states of ^{81}Ga to populate one of two states in ^{80}Ge . It can populate the lowest 2^+ state and subsequently γ decay through the 659 keV transition. To transition from the $5/2^-$ state in ^{81}Ga to a 2^+ in ^{80}Ge , the $\ell + s$ quantity (total J) must equal $1/2$, $3/2$, $5/2$, $7/2$, or $9/2$. This requires $\ell = 0, 1, 2, 3$, or 4 . However, because parity changes between these states the ℓ value can only be odd. Thus it is either 1 or 3 . In the case of a decay to a 4^+ state in ^{80}Ge the possible $\ell + s = \text{total } J$ values are $3/2$, $5/2$, $7/2$, $9/2$, $11/2$, or $13/2$. These correspond to $\ell=1, 2, 3, 4, 5$, or 6 . However, a change in parity means the possible ℓ values can only be 1 , 3 , or 5 .

Relying on the proposed spin values of the states in ^{80}Ge allows the experimenter to decide that population of the 1083 keV transition will be less “polluted” by this undesired path shown in figure 28 than the 659 keV transition. The possible angular momentum barriers are the same for both transitions except that the population of the 4^+ state (decaying through the 1083 keV)

has an extra value of ℓ which is higher than the ones available for populating the 2^+ state (decaying through the 659 keV). Thus on average there will be a smaller angular momentum barrier to populate the 2^+ state than the 4^+ state.

Thus the 1083 keV transition will be less contaminated with this β -delayed neutron branching from ^{81}Ga and so it is more trustworthy for obtaining the neutron branching ratio for ^{81}Zn decay.

The calculation of the neutron branching ratio utilizes the intensity of the 1083 keV peak, counts divided by its efficiency of detection in the γ spectra. Then the appropriate integrated Bateman factor is used to adjust for the half-life difference between ^{81}Zn decay and ^{80}Ga decay. Finally, the percent of total β decay strength which the 1083 keV transition carries away from ^{80}Ga is from literature. This yields an intensity, which upon dividing by the total intensity seen decaying to all available channels (to the ^{81}Ga ground state, excited states, and neutron emission), the percent of all ^{81}Zn β decay through β -delayed neutron emission is 12(4)%. Thus this analysis on the neutron branching ratio relies on data from previous literature on ^{80}Ga to ^{80}Ge decay [25].

4.3.5 Neutron branching ratio: Compare to theoretical predictions and prior values from literature

One set of calculations for the neutron emission branching ratios for neutron-rich zinc isotopes is included in figure 29. Our experimental value is within errors of the value calculated by I.N. Borzov in 2005 [26]. There he calculates 10(4)% for the neutron branching ratio of ^{81}Zn . Our value is slightly higher than the value reported by Verney *et al.* [2] of 7(3)%. Another recent value reported for the probability of neutron emission from ^{81}Zn is given by Hosmer *et al.* as 30(13)% [27]. This value was obtained from a β -ion-neutron correlation experiment at the National Superconducting Cyclotron Laboratory. Such differences are not uncommon when two different experimental techniques are utilized for measuring the same quantity. It should be noted that the value obtained in the work of this thesis is from a γ - β coincidence with the condition that the time between those events is no more than 2.5 μsecs . However, in the Hosmer *et al.* paper the correlation time between a β decay in the silicon detectors and a subsequent neutron detection in the NERO array is 200 μsecs . Thus the time window in the work in this thesis is much shorter and includes a γ -ray which can be uniquely identified with the nucleus from which it decayed. However, in a 200 $\mu\text{sec.}$ window in the Hosmer paper it is possible to

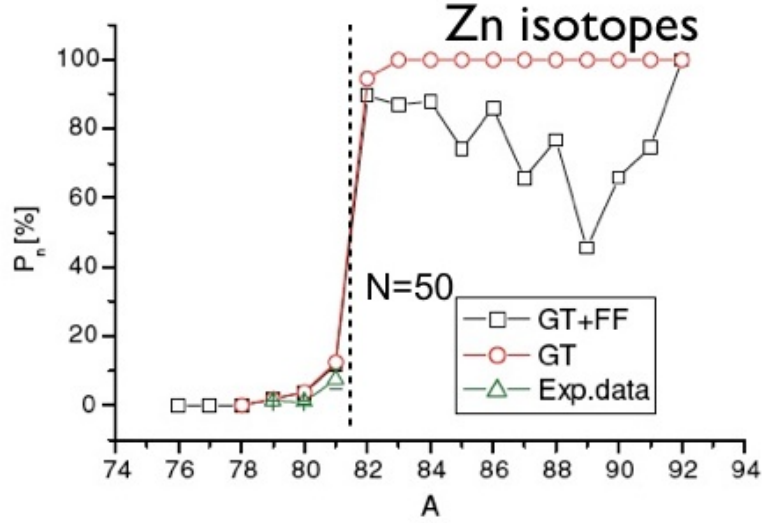


Figure 29: Predictions for the neutron branching ratios in neutron-rich zinc isotopes [26].

measure neutrons from other sources such as background and daughter decays in addition to the parent decay. This issue is addressed in their paper, however the difference in time window length is an important point providing more support for the number in this work. It is also noted that their neutron detection efficiencies may not be understood well enough. This is another contentious point which we avoid by detecting γ -rays as opposed to detecting neutrons directly.

4.4 Pandemonium effect and ground state branching ratio

For nuclei with a large Q_β or any large Q value for its decay, there is another difficulty present for calculating the ground state to ground state decay branching ratio. The calculation of the branching ratio to any state involves, as stated before, the total γ -ray intensity decaying out of the level of interest minus the total γ -ray intensity feeding into said level from above. The branching ratio to the ground state is calculated using the same principle, but the intensity decaying out of this level is taken from a γ -ray transition emitted from the decay of the daughter to the granddaughter nucleus. The only extra information needed is the percent of total β decay intensity carried away by whichever γ -ray transition is utilized, usually the strongest one in that decay. In this work the 216 keV transition from ^{81}Ga to ^{81}Ge is used and it carries away 37% [21] of the total β decay intensity of ^{81}Ga .

The problem with this transition exiting from the decay of the ^{81}Ga ground state is accounted

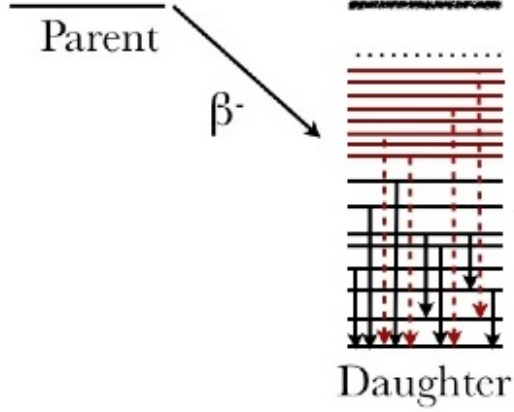


Figure 30: Generic description of the problem of pandemonium in β decaying nuclei. The solid, black arrows are the observed γ -ray transitions and the dashed, red arrows are the unobserved ones. The horizontal, dotted line near the top indicates a neutron separation energy and a solid black line above it indicates the top of the Q_β decay window.

for in the calculation of isobaric impurities in the beam. Now the problem pertains to all of the transitions feeding into this ground state from above in the ^{81}Ga decay scheme. Specifically, the transitions which are not observed are the origin of the problem. Because of low efficiency for detecting high energy γ -rays and because many of those transitions decaying from higher in the decay window may be of weaker intensity, it is not possible to observe all of the γ -ray intensity feeding the ground state. Therefore, the calculated ground state to ground state branching ratio may be artificially inflated because of that unobserved intensity feeding into it from higher states in the decay window.

4.4.1 Pandemonium: Basic information

This is the concept of so-called Pandemonium, which was addressed with a simulation by J.C. Hardy and others [28]. Pandemonium is addressed in later papers as well [29] [30]. In the first paper the authors put forward a fictional nucleus with the same proton and neutron numbers as ^{145}Gd , and so it would decay to ^{145}Eu . This is a fictional nucleus because a number of its characteristics are simulated with numerical methods. For example, the distribution of spin and parity values (J^π values) are simulated with a Monte-Carlo simulation. Next, the energy spacings between states are distribution with a Wigner semicircle distribution. This ensures a reasonably random distribution of energy gaps between all of the simulated excited states in this

nucleus. They note that this tends to overlook specific energy spacings which will be inherent to the unique nuclear structure of any particular nucleus, but this may be overlooked for the sake of a generic simulation, meant to illustrate the general problem of pandemonium. Finally, the distribution of allowed Gamow-Teller β decay strength to each of the available states is given by another random number distribution in this paper. Using this nucleus with its randomly distributed characteristics, the authors then calculate the full γ -ray intensities that would be emitted from such a nucleus. Energies below 1.7 MeV were ignored because the problem of the pandemonium effect deals with missing intensity from weak, high energy γ -ray transitions, not low energy ones. They then calculated a likely observable β -delayed γ -ray spectrum, if such a decay were to be measured with Ge(Li) detectors. An energy dependent efficiency is taken into account as well as 3 keV Gaussian FWHM resolution and double escape peaks. Also, a number of different decay intensities are used, such as 0.1 times the original amount and 10 times the original amount in order to measure the photo-peak intensities in the simulated γ -ray spectra. But even with these various intensities used, only 10% of all 1000 γ -ray transitions are observed and approximately 14% of the total γ -ray intensity above 1.7 MeV was completely undetected in each scenario. This illustrates how a nucleus with a high enough level density and a large enough Q_β decay window may have a significant amount of high energy but weak γ -ray transitions that are not observed, at least with high-resolution HPGe detectors. Upon comparison to the level scheme of ^{145}Gd decay at that time, the amount of counts observed makes the comparison to the 1 or 0.1 times simulated intensity data sets. Here they conclude that approximately 20% of the real decay of ^{145}Gd is completely unobserved in the real experiments. Thus they cannot be placed in the level scheme. The conclusion of this paper is simply to call into question the standard analysis of complex β decay level schemes because of this large amount of unobserved intensity from high energy states.

4.4.2 Estimation of missing β decay intensity

Because there is most likely unobserved β decay intensity in the experiment here, the experimenter must make some form of an approximation to take into account the γ -ray transitions and their intensities which are not observed decaying from high energy states. One possibility is to compare this decay to the closest nucleus with a decay which has very similar physics. The decay of the $N=51$ ^{83}Ge to $N=50$ ^{83}As involves nuclei which only contain two extra protons than ^{81}Zn and ^{81}Ga , respectively. The physics is similar because that decay must involve first forbidden decays to low energy states in the daughter and allowed Gamow-Teller decays to high energy states above about 3 MeV. This is because in the shell model picture for ^{83}Ge there is

only one valence neutron above the N=50 gap, in an s or d orbital with positive parity, but ^{83}As contains 5 protons (instead of 3 as in ^{81}Ga) in negative parity f and p orbitals. An allowed Gamow-Teller decay may only occur through the decay of a deeply bound neutron in the νf and νp orbitals, but this will only populate high energy, core excited states in the daughter.

The decay of ^{83}Ge to ^{83}As has been measured by J.A. Winger *et al.* [31]. There are 13 states above 3 MeV inside the ^{83}As level scheme, seven of which are said to be populated in allowed Gamow-Teller β decays from the parent because their $\log ft$ values are in the range of 5.5 to 6.0. The population of these states is compared to the allowed Gamow-Teller decays in ^{136}I decaying to N=82 ^{136}Xe , which is also singly-magic just like ^{83}As and ^{81}Ga . J.A. Winger also notes another experiment which showed there to be states with neutron particle-neutron hole configurations around 4.5 MeV in ^{136}Xe , populated through the reaction $^{136}\text{Xe}(p,p')$ [32]. The similar physics is the main reason an approximation is made through this comparison. In the decay of ^{81}Zn to ^{81}Ga there are only two states observed above 3 MeV, namely those seen at 4.294 and 4.880 MeV. In this comparison the assumption is made that the two states most strongly populated through Gamow-Teller decay are comparable to the two high energy states seen in ^{81}Zn decay. Therefore, there should be a minimum of 11 other states in this energy range of 3 to 5 MeV populated with similar decay strengths. The decay scheme for ^{81}Zn is simply amended by adding β decay strength to 11 fictitious states in this energy range, where the decay strength to these states is taken from those 11 states in ^{83}Ge decay. Table 6 shows all 13 states populated through decays of ^{83}Ge , each with Q_β values, $\log ft$ values, and β decay branching ratios shown.

Thus 11 fictitious states are added to the ^{81}Zn decay scheme with branching ratios equivalent to all of the above states, except the two at 4130.1 and 4191.4 keV, since those two have the strongest branching ratios and are compared to the two ^{81}Ga states observed above 4 MeV. It is assumed that these 11 “new” states will γ decay to the ground state and account for a minimum amount of the unobserved γ decay intensity to the ground state. When this new intensity is added, the ground state to ground state branching ratio in ^{81}Zn β decay changes from 52(7)% to 40(6)%. Tables 7 and 8 show how this missing intensity was accounted for in the tabulation of total β decay strength. First the intensities are shown to each state, including the ground state and β -neutron emission branch. Then there are 11 new states added with each having intensity equivalent to 2% of the total β decay strength. This 2% is approximately the average branching ratio to each of these 11 states in ^{83}As . This approximation slightly overestimates the

Table 6: The core excited states populated through allowed Gamow-Teller decay in ^{83}Ge decay to ^{82}As . The γ -ray intensities from these states are added to the ^{81}Ga level scheme as fictitious states which should exist there but are not observed in the current work. This added intensity is a partial approximation for missing intensity not observed because of the pandemonium effect.

Level energy (keV)	Q_β (keV)	$\log ft$	Branching ratio (%)
3522.5	5437.5	6.3	1.45
3733.4	5226.6	6.0	2.45
3956.4	5003.6	6.1	1.58
3999.6	4960.4	6.3	0.97
4030.5	4929.5	6.7	0.37
4130.1	4829.9	5.5	5.4
4191.4	4768.6	5.7	3.2
4221.2	4738.8	6.0	1.55
4228.3	4731.7	5.9	1.95
4364.3	4595.7	5.7	2.7
4405.6	4554.4	5.9	1.65
4434.1	4525.9	5.5	4.0
4841.6	4118.4	5.6	2.1

Table 7: The branching ratios in ^{81}Zn decay to ^{81}Ga before the adjustments for the small amount of isobaric contamination and the approximation for the pandemonium effect.

Level (keV)	Branching ratio (%)	$\log ft$
0	52(7)	5.5
351	14(2)	6.0
802	6.0(10)	6.3
1267	1.1(6)	6.8
1458	1.9(9)	6.6
1936	0.5(20)	7.1
4294	7.0(20)	5.4
4880	4.0(10)	5.4
β -neutron	12(4)	—

total decay strength to these high energy states, thus it is a worst case scenario approximation and is safer. Table 8 shows the new branching ratios after taking into account the intensities for these 11 fictitious states.

Table 8: The branching ratios in ^{81}Zn decay to ^{81}Ga after the adjustments for isobaric contamination and the pandemonium effect.

Level (keV)	Branching ratio (%)	$\log ft$
0	40(6)	5.6
351	10(1)	6.1
802	4.3(7)	6.4
1267	0.8(5)	7.0
1458	1.5(7)	6.8
1936	0.4(20)	7.2
4294	5.2(10)	5.5
4880	3.3(1)	5.5
11 “new” states	2.0 each	—
β -neutron	11(2)	—

5 Interpretations and Calculations for ^{81}Zn Decay

5.1 Theoretical shell model calculations

There are various theoretical tools which can be utilized to predict energy levels, J^π values, and $B(\text{GT})$ values in the β decay of atomic nuclei. These calculations involve using the shell model for atomic nuclei. Calculating the energy spectrum for an atomic nucleus is essentially a problem of solving the Schrödinger equation for the nuclear many-body problem.

$$\hat{H}\Psi = E\Psi \quad (59)$$

The wave function Ψ of the nuclear state is characterized by the quantum numbers in the shell model picture, described in the theoretical background chapter. The $\hat{H}$ is the Hamiltonian energy operator and E is the energy eigenvalue. Solving such an eigenvalue problem entails the diagonalizing of the matrix $\hat{H}$. However, the creation of that matrix with appropriate matrix elements is not trivial. It requires a detailed understanding of the specific nucleon-nucleon interactions involved within the particular nucleus in question.

The ideal theoretical calculation would involve calculating nuclear properties from each nucleon-nucleon interaction in the nucleus until all nucleon-nucleon pairings are taken into account. This would require an immense amount of computing power and an amount of time

which may be completely unavailable impractical, depending upon both the calculations and the computing system available. Thus there are a number of approaches to approximate the calculations. The shell model calculations use an inert core and the remaining nucleons outside of this core are treated as valence nucleons, which interact with only other valence nucleons through a residual interaction. This interaction is meant to approximate whatever real, but computationally insurmountable, interactions exist in said nucleus. For example, for calculations on ^{81}Zn and other $N=51$ nuclei, a ^{78}Ni inert core is utilized. The ^{78}Ni is a doubly magic nucleus according to traditional magic numbers and so it should be a good, inert core. That is, there are no interactions between nucleons inside the core with valence nucleons outside of the core. This only applies to energy levels low enough in the nuclei of interest such that core excited states are not populated. If such states are populated in the experiment, they will not be shown in the theory calculations because such states reside outside of the assumed model space.

The residual interaction between the valence nucleons is based upon some form of derived nucleon-nucleon interactions. A more modern handling of nucleon-nucleon interactions are based upon experimental data from scattering between two nucleons. However, renormalization for these interactions inside of the nuclear medium in the nucleus are needed. Experiments measuring phase shifts from scattered nucleons ultimately leads to mathematical formation of interactions such as CD-Bonn, Nijmegen, Argonne V_{18} , etc. [33].

5.2 $N=51$ systematics

Information on the changes in energy of states throughout an isotopic or isotonic chain often reveals information on how nucleon-nucleon interactions are driving the physics of those nuclei. For example, moving along an isotonic chain from lower to higher Z simply involves an addition of a proton with each new nucleus while the number of neutrons remains constant. As protons fill one orbital the behavior of the ground states or excited states may follow one trend, but upon filling the next available proton orbital this behavior may change. This would suggest a different proton-neutron interaction is controlling that change in behavior from one orbital to the next. Often times this description may be an oversimplification of what is really occurring because of the large amount of valence nucleons involved. However, near a doubly magic nucleus this description may be more valid as there are fewer valence nucleons interacting with each other.

The systematics of the ground states and first excited states in $N=51$ isotones, which includes

^{81}Zn , is the relevant example in this work. The other odd mass $N=51$ isotones which have been measured through decay and/or particle spectroscopy are ^{83}Ge , ^{85}Se , ^{87}Kr , ^{89}Sr , ^{91}Zr , etc.

5.2.1 Previous experiments on $N=51$ systematics

The nucleus ^{83}Ge is the latest addition to this isotonic chain. It was studied through the reaction $^2\text{H}(^{82}\text{Ge}, p)^{83}\text{Ge}$ [34]. It is an inverse kinematics reaction, meaning that the beam is the heavier particle, ^{82}Ge , and the target is the lighter particle, a deuteron from the deuterated plastic target. The ejectile is a proton and the resultant nucleus is ^{83}Ge in some excited state or its ground state. This is a so-called (d,p) reaction because it is essentially the introduction of a deuteron into the system but a proton is ejected. The outgoing protons carry information about both the energy and J^π of the populated state in the ^{83}Ge resultant nucleus. The energies and angular distributions of this outgoing proton are measured in the ORRUBA (Oak Ridge Rutgers University Barrel Array) silicon detectors.

Because the protons are emitted after populating states with different J^π values, they will come away from the reaction with different orbital angular momenta and thus arrive on the silicon detectors with different angular distributions. From these different angular distributions that experiment showed the first excited state in ^{83}Ge to have a J^π of $1/2^+$ and the ground state has a $J^\pi = 5/2^+$.

An experiment on decay spectroscopy at the HRIBF, using the so-called ranging-out technique and β -ion-neutron tagging, verified the energy of this first excited state inside of ^{83}Ge . The decay of ^{84}Ga , studied by J. A. Winger *et al.* [35], populates the first excited state of ^{83}Ge through β -delayed neutron emission. The energy in the γ -ray spectra is shown to be 247.3(5) keV from the ^{84}Ga β -neutron emission and 247.05(17) keV from the ^{83}Ga β decay. Thus this first excited state is at 247 keV. These two experiments serve as examples of how the experimental information on the systematics is amended. In this experiment we plan on using the observed β decay strength distribution along with β decay selection rules to add the J^π of the ground state of ^{81}Zn to the systematics.

5.2.2 ^{81}Zn ground state J^π

These $N=51$ isotones all possess one neutron above the $N=50$ shell gap. The nucleus ^{81}Zn has two protons above the ^{78}Ni core, ^{83}Ge has four, ^{85}Se has six, etc. With the addition of each pair

of protons there is a change in energy of the first excited state as well as changes in spin-parities for those states. With few valence nucleons outside of this ^{78}Ni core it is possible to investigate the data with the physics being driven by proton occupation number. All of these $N=51$ isotones mentioned have a $1/2^+$ first excited state and a $5/2^+$ ground state. The next isotone in the sequence, from higher to lower Z , is ^{81}Zn . The $1/2^+$ first excited state is continuously moving lower in energy from ^{89}Sr down to ^{83}Ge . If this trend continues to ^{81}Zn then this $1/2^+$ first excited state may stay a first excited with lower energy than in ^{83}Ge , or it may invert with the $5/2^+$ state to become the ground state itself instead. An opposite possibility is that it may reverse trend and move up in energy upon approaching ^{78}Ni . In the experiment discussed in this work, there is no direct measurement of the spin-parity of the ground state of ^{81}Zn . However, the measured β decay strength distributions in combination with β decay selection rules may allow for a determination of this ground state spin-parity. According to the systematics, the value can only be either $5/2^+$ or $1/2^+$. In order to use β decay selection rules, one has to look at a decay to a state in the daughter nucleus where the spin-parity is known with certainty. This way the choice is only between two options: A $5/2^+$ state decaying to a state in the daughter of known spin-parity, or a $1/2^+$ state decaying to that same final state. The only state in the daughter ^{81}Ga with known spin-parity is that of the ground state. The spin was measured to be $5/2$ through a magnetic moment measurement at ISOLDE-CERN [23]. The parity must be negative because gallium only possesses three protons above the $Z=28$ shell gap in f and p orbitals. Therefore, the ground state should have negative parity. The ground state J^π of ^{81}Ga is $5/2^-$.

Therefore, the two options are either a $5/2^+$ to $5/2^-$ decay or a $1/2^+$ to $5/2^-$ decay. The Gamow-Teller allowed decay is not possible because both options involve a change in parity, which is simply not allowed by Gamow-Teller selection rules. The first viable option is the $1/2^+$ decay to the $5/2^-$. By the selection rules explained in the theoretical background chapter, this is classified as a first forbidden unique transition because the change in angular momentum is by two units of $\hbar$ and the parity changes. Each type of β decay has a corresponding range of empirical $\log ft$ values, which are directly related by a simple equation to the decay branching ratio. The lowest $\log ft$ for a first forbidden unique is 7.5, see figure 6. That is, the ground state to ground state branching ratio would have to be less than 0.1%. However, here a branching ratio for the ground state to ground state decay is large enough such that the empirical $\log ft$ is about 5.6. This is after taking into account the small amount of isobaric contamination present and an approximation for unobserved β decay intensity to high lying states.

It is unlikely that even after these two approximations that the measured feeding intensity to the ground state to ground state decay is two orders of magnitude off from the real value. Therefore, a $1/2^+$ ground state decaying to a $5/2^-$ ground state in the daughter is not possible. The only option available is that it is a $5/2^+$ ground state decaying to a $5/2^-$ ground state. This is classified as a first forbidden transition because the $\Delta\ell=0$ and the parity changes. The empirical $\log ft$ values for first forbidden decays include 5.6 in their range, see figure 6. Therefore, when presented with a choice between first forbidden and first forbidden unique along with this large size of the β branching ratio, the $5/2^+$ to $5/2^-$ scenario is the only viable option. Therefore, it is concluded that the ground state J^π of ^{81}Zn is $5/2^+$. The first excited state is inferred to be $1/2^+$ based upon the systematics of the $N=51$ isotones.

5.2.3 Calculations on $N=51$ isotones

The energy of the first excited state in each of these isotones is shown in figure 31. The experimental values for these energies are known for zirconium to germanium isotones, but not yet for zinc. The points at the bottom are the $5/2^+$ ground states, whereas all of the first excited states have a J^π of $1/2^+$. The values in black are the theoretical energies of the first excited states from a shell model calculation. This calculation is based upon a ^{78}Ni core and a model space containing proton orbitals $f_{5/2}$, $p_{3/2}$, $p_{1/2}$, and $g_{9/2}$ whereas the neutron orbitals are $d_{5/2}$, $s_{1/2}$, $d_{3/2}$, and $g_{7/2}$ [36], where the nucleon-nucleon interactions are based upon the N3LO approach [37]. Thus the valence protons and one valence neutron in these $N=51$ nuclei can reside in these orbitals in any type of excited or ground state configuration. The main point of these calculations is to analyze the systematics of the ground states and first excited states in $N=51$ nuclei. Both energies and J^π values of the first excited states and ground states are of interest. The first calculation was performed with the canonical single particle states, meaning that the order and energy of the single particle states are not modified. It is possible to rerun the calculations but with modified energies. It is also possible to change each matrix element, which quantifies the strength of each nucleon-nucleon interaction. For example, the interaction between protons in the $f_{5/2}$ and neutrons in the $d_{5/2}$ single particle states can be strengthened or weakened or made to be zero, at least naively within this shell model calculation. Performing these changes leads to knowledge about how particular nucleon-nucleon interactions are affecting the systematics of these states.

The first investigation involves an interaction between the protons in the $\pi f_{5/2}$ and $\pi p_{3/2}$

Table 9: Shell model calculation information. These effective single particle state energies are in relation to the Z=28 shell gap for protons and N=50 gap for neutrons. Thus the $\pi f_{5/2}$ state is the lowest proton state and $\nu d_{5/2}$ state is the lowest neutron state.

Effective single particle state	Energy
$0\pi f_{5/2}$	0.000
$1\pi p_{3/2}$	1.100
$1\pi p_{1/2}$	2.500
$0\pi g_{9/2}$	4.500
$1\nu d_{5/2}$	0.00
$2\nu s_{1/2}$	1.30
$0\nu g_{7/2}$	1.80
$0\nu d_{3/2}$	2.40
$0\nu h_{11/2}$	3.00

orbitals with neutrons in the $\nu s_{1/2}$ orbital. In the theoretical predictions for the energies of the $1/2^+$ first excited states there is an obvious change in trend for nuclei below the Z=34 selenium to atomic numbers above it. From ^{79}Ni to ^{85}Se the energy is moving down from 1.27 MeV to 0.4 MeV. As atomic number Z increases, this energy reverses trend above Z=34 selenium and moves up in energy towards stability at ^{91}Zr . In the shell model picture the $\pi f_{5/2}$ orbital is filled from ^{79}Ni until ^{85}Se is reached. From ^{85}Se to ^{89}Sr the new protons are filling the $\pi p_{3/2}$ orbital. This change in which proton orbitals are being filled exactly overlaps where the change in trend occurs for the energies of these $1/2^+$ first excited states. It is believed that a proton-neutron interaction may be responsible for this change. The proton-neutron interaction between protons in the $\pi f_{5/2}$ orbital with neutrons in the $\nu s_{1/2}$ orbital must be of opposite sign than the interaction between protons in the $\pi p_{3/2}$ orbital and neutrons in the $\nu s_{1/2}$ orbital. Thus the first scenario, for nuclei with Z below 34, is an attractive interaction since the first excited state is pulled towards zero. However, the second scenario has a proton-neutron interaction which is repulsive for nuclei with Z above 34.

To further investigate this explanation for the change in trend, a modified shell model calculation was performed. Here the matrix elements controlling the interaction between $\pi f_{5/2}$ and $\nu s_{1/2}$ are set to zero and the calculation is rerun. The calculation is repeated but the matrix elements controlling interactions between $\pi f_{5/2}$ and $\nu d_{5/2}$ are set to zero. The later

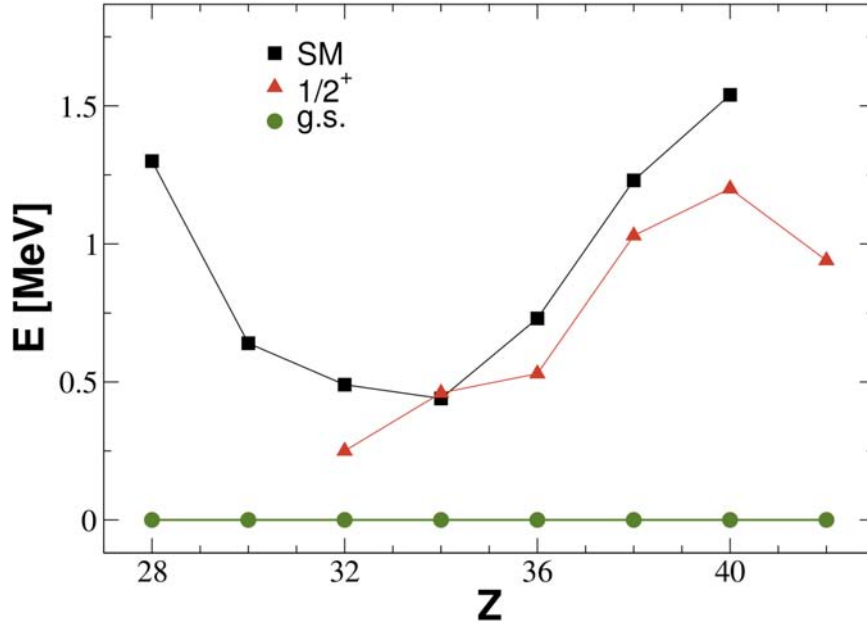


Figure 31: The systematics of states in N=51 isotones with the shell model calculation included.

interaction produces little to no change in the theoretical systematics of this first excited state's energy. However, the first interaction turned off ($\pi f_{5/2}$ with $\nu s_{1/2}$) does produce a noticeable result. It nearly “flattens” the trend in energies from ^{79}Ni to ^{85}Se . This strongly suggests that the proton-neutron interactions involving the $\nu s_{1/2}$ are more important than ones involving the $\nu d_{5/2}$ orbital. This is a vital question since all shell model calculations show that the ground state and first excited states in N=51 isotones have wave functions almost evenly mixed between $\nu d_{5/2}$ and $\nu s_{1/2}$ orbitals. These states are not single particle states, according to these calculations. The important proton-neutron interactions only involve the $\nu s_{1/2}$ orbital and not the $\nu d_{5/2}$ orbital, at least for those nuclei from ^{79}Ni to ^{85}Se which have fewer valence nucleons than the other N=51 isotones.

Finally, another calculation is performed where the matrix elements controlling the interaction between the $\pi p_{3/2}$ and $\nu s_{1/2}$ orbitals are set to zero. This results in little change in the systematics. This is expected because there are more valence nucleons for nuclei above Z=34, thus the nucleon-nucleon interactions in that region are more complicated. Therefore, this is why it is understandable that a change occurs for first excited states in nuclei below Z=34 but not much change occurs above Z=34. These extra calculations are shown in figure 32.

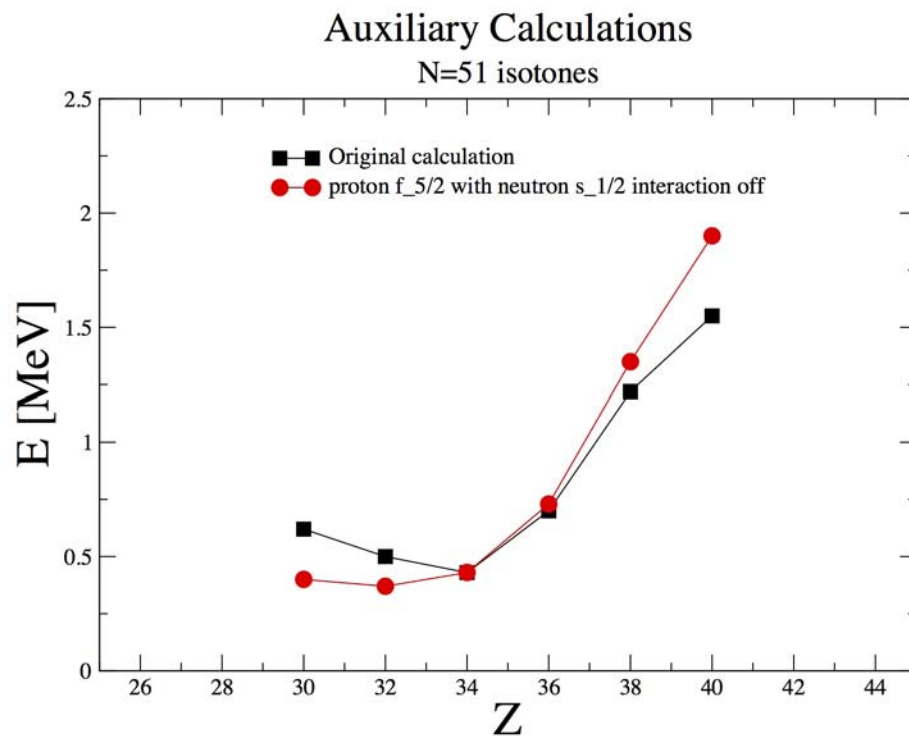


Figure 32: An example of changes to the shell model calculations. The original calculations are shown in black. The red points show the predicted energies of the first excited states when the $\pi f_{5/2}$ with the $\nu s_{1/2}$ interaction is set to zero.

5.3 Core excitations of ^{78}Ni

There are many configurations of nucleons in an atomic nucleus which will lead to the various excited states it can possess. States which have a low energy will usually only consist of the valence nucleons residing and coupling together in single particle states above the shell gaps. This ignores the cases where an intruder single particle state may exist in a major shell where it usually is not present. However, states which are high enough in energy may consist of nucleon holes residing in single particle states below these shell gaps. These are generally referred to as core excited states. In fact they are in the model space of the core nucleus. For example, obvious nuclei in which to look for such states is with the first excited state of a doubly magic nucleus. This is because the only way to achieve an excited state with such nuclei is to excite the nucleons beyond the shell gap. Since shells are filled for both protons and neutrons, the only states available for populating an excited state are for nucleons to reside beyond the shell gaps. The isotope ^{16}O is doubly magic and its first excited state is at 6.05 MeV. This is a relatively high first excite state energy, even compared to other doubly magic nuclei such as ^{40}Ca (first excited state at 3.74 MeV) and ^{208}Pb (first excited state at 2.61 MeV) [38].

5.3.1 Core excited states in $N=50$ isotones

In ^{81}Zn there is only one neutron outside of the core. In the shell model picture this neutron resides in a d or s orbital. The only energetically available proton single particle states for it to decay to are of opposite parity. This means that only first forbidden decays are possible via the decay of this valence neutron in ^{81}Zn because first-forbidden decays require a change in parity. Decay to states which are higher in energy in the daughter nucleus, that is decays forming a deeply bound hole in neutron orbitals below the $N=50$ gap, are damped heavily by the lower Fermi statistical factor (see theoretical background chapter). Decays to such states are not energetically favored. However, if allowed Gamow-Teller decays can occur through the decay of one of those neutrons, then the reduced transition matrix element would be about two orders of magnitude stronger and would provide a counterbalance to the lack of decay energy available. Therefore, population of core excited states in $N=50$ ^{81}Ga must involve the decay of a deeply bound neutron in the f and p orbitals, which in turn will decay to a proton in its appropriate spin-orbit partner in the proton f and p orbitals, see figure 33. The core excited states in ^{81}Ga must have J^π values of $3/2^+$, $5/2^+$, or $7/2^+$. These are the only states which can be populated through the Gamow-Teller decay of the $5/2^+$ ground state of ^{81}Zn because of selection rules.

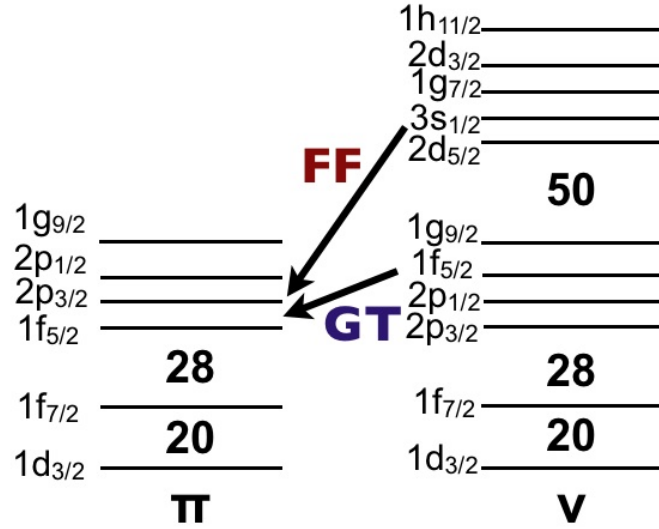


Figure 33: In the decay of ^{81}Zn and nearby $N=51$ isotones with one or a few neutron above the $N=50$ gap, only FF decays are possible from this valence neutron to decay to a proton in the low lying proton f and p orbitals in $N=50$ nuclei. The only energetically available allowed GT decays must occur via the decay of a deeply bound neutron in the neutron f and p orbitals to their spin-orbit partners in the proton f and p orbitals.

One way to investigate the behavior of core excited states across the $N=50$ isotonic chain is to calculate the average energy of the core excited states in these isotones. This is done through the use of the reduced transition matrix element for allowed Gamow-Teller decays, the $B(\text{GT})$ values. The $B(\text{GT})$ weighted energy calculated for each state populated in each of these isotones is needed. The following formula calculates this value:

$$\bar{E} = \frac{1}{N} \sum_{i=1}^N E_i B_i(\text{GT}) \quad (60)$$

The average comes from this weighted summation. The energy of each state populated in the decays to each $N=50$ daughter nucleus is multiplied by its $B(\text{GT})$ value. Thus it is a $B(\text{GT})$ weighted average of energies of the states. The $\bar{E}$ is the average energy of all states populated through only Gamow-Teller decay. Because the core excited states in ^{81}Ga and the nearby $N=50$ isotones are only populated through Gamow-Teller decay, the $\bar{E}$ is also an average energy for all of the core excited states. A graph of this $\bar{E}$ value in each of the $N=50$ isotones of interest shows an upward trend in the core excitations upon moving towards ^{78}Ni . In the data in this

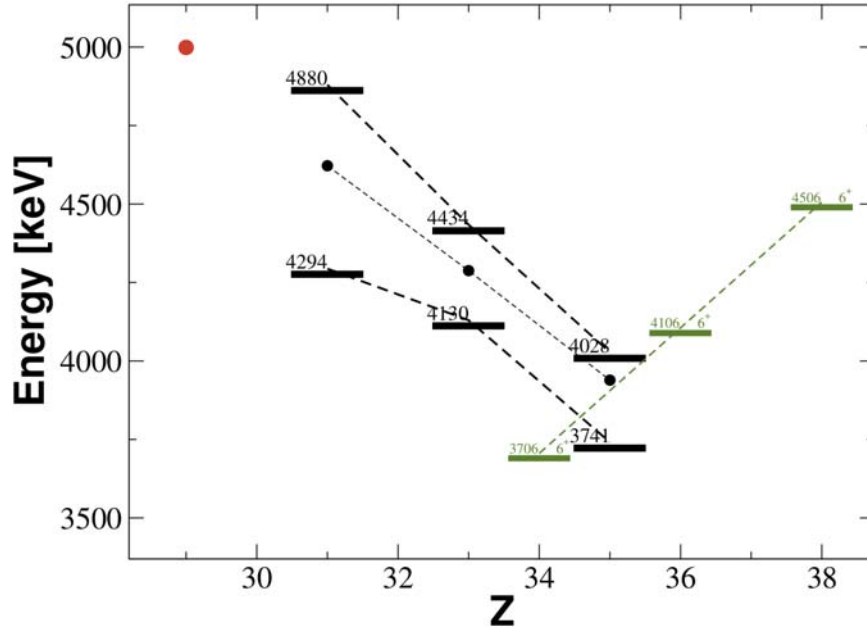


Figure 34: Core excited states in N=50 isotones. The black circles show the average energy of states populated through only Gamow-Teller decays in these isotones. They are the average energy for core excited states. The two states with strongest population in GT decay are shown in each of the three odd Z isotones. The green points show the energy of 6^+ yrast states in even Z, N=50 isotones [39]. The red point is the extrapolated energy for core excitations in ^{79}Cu , only one proton outside of the ^{78}Ni core. Those core excited states should reside around 5 MeV.

work the increase in the B(GT) weighted, core excited state energy for odd Z, N=50 isotones suggests an increase in the gap between the νf and νp orbitals and the $\nu d_{5/2}$ orbital above the N=50 gap. That is, the energy between the deeply bound neutron hole and the valence neutron above N=50 is increasing upon moving towards ^{78}Ni . This trend of the average energy for core excited states in N=50 isotones is shown in figure 34.

5.3.2 Possible N=50 and N=40 shell gap changes

Also shown in figure 34 is the systematics of yrast states in even Z, N=50 isotones [39]. The yrast states are defined as states which have the lowest energy for a given total angular momentum. The specific states shown here are from an experiment in which prompt γ -ray transitions are observed from the spontaneous fission of ^{248}Cm to populate states in ^{82}Ge [39]. Here there are 5^+ and 6^+ yrast states which are interpreted as having the following configuration: A hole in the $\nu g_{9/2}$ single particle state and a valence particle above the N=50 gap in the $\nu d_{5/2}$ single particle

state. This particle-hole pair then couples to the valence protons to form the yrast state. The main point made here is that the energy of these 5^+ and 6^+ yrast states are moving down with lower Z upon moving towards ^{78}Ni . If the interpretation of the configuration of these states is correct, it would suggest that the gap between the $\nu g_{9/2}$ and the $\nu d_{5/2}$ state is decreasing. That is, the $N=50$ shell gap may be decreasing towards ^{78}Ni . It was recently suggested by Porquet [40] as well, that ^{78}Ni is not a good doubly magic nucleus. Porquet states that the $N=50$ shell gap is decreasing by about 700 keV along the $N=50$ isotones from $Z=38$ (^{88}Sr) down to $Z=34$ (^{84}Se).

However, the interpretation of the drop in energy for these yrast states may not be correct. They may not originate from the coupling of a hole in $\nu g_{9/2}$ with a neutron in $\nu d_{5/2}$ and paired protons. The systematics of these yrast states may simply be an effect of configuration mixings or from particular proton-neutron interactions which are not yet fully understood. Also, even if these yrast states allow for direct probing of the $N=50$ shell gap, it is not yet known where these states reside for nuclei with Z less than 32. They may go back up in energy, thus indicating a strengthening of the $N=50$ shell gap. Therefore, although Porquet and Rzaca-Urban and others may suggest that the $N=50$ shell gap is decreasing because of the yrast behaviors in even Z , $N=50$ isotones, this is yet unconfirmed experimentally.

However, if this interpretation is correct and the trend in core excited states observed in this work is taken into account, it may indicate a decrease in the $N=50$ shell gap strength but an increase of a gap at $N=40$. If the $N=50$ gap is decreasing and the neutron particle-neutron hole gap just described is increasing, then this automatically signifies that the $N=40$ subshell gap between $\nu g_{9/2}$ and νf , νp orbitals must be decreasing. Again, this is unconfirmed because the magnitude of the $N=50$ gap change is itself unknown. Such confirmations can only come with the probing of the single particle nature of states in nuclei which have a single nucleon outside of the doubly magic core.

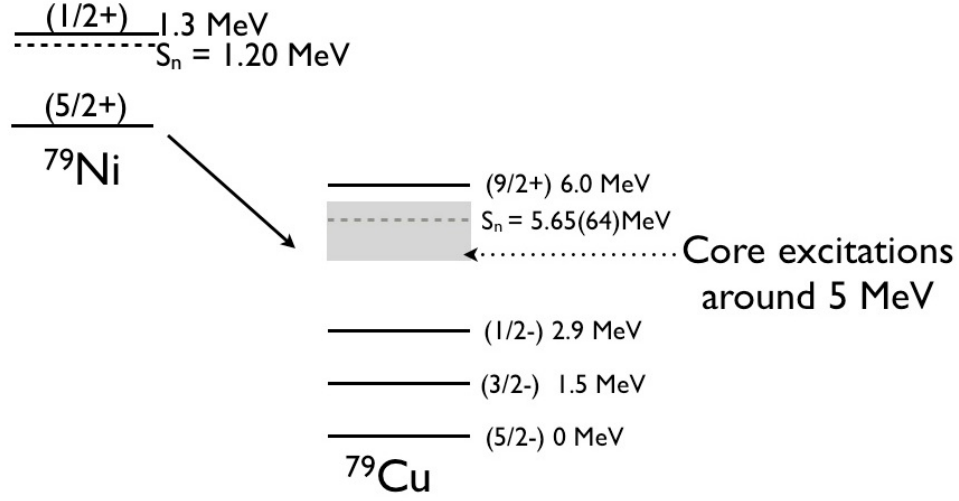
5.4 Predictions for ^{79}Ni and ^{79}Cu

The atomic nucleus ^{79}Ni contains only one neutron outside of the ^{78}Ni core and it is an $N=51$ isotone. The nucleus ^{79}Cu contains only one proton outside of ^{78}Ni and it is an $N=50$ isotone. Based upon previous discussions for both $N=50$ and $N=51$ isotones, various predictions may be made concerning these nuclei. The nucleus ^{79}Ni should have a $5/2^+$ ground state and a $1/2^+$ first excited state, according to the same shell model calculations shown in figure 31. The

ground state should be a pure $\nu d_{5/2}$ state and the first excited state should be a pure $\nu s_{1/2}$ state, located at 1.3 MeV according to the energy of this single particle state in the calculations. They should be single particle states because there is only this 51st neutron present with no coupling to the other nucleons, which reside beyond shell gaps.

The daughter in the β^- decay of ^{79}Ni , which is ^{79}Cu , should possess low lying states which would be single particle states for the single proton it has sitting outside of the ^{78}Ni core. Thus the ground state should be a $5/2^-$ because it is a $\pi f_{5/2}$ state. The first excited state would have $J^\pi = 3/2^-$ because of its $\pi p_{3/2}$ single particle nature. This state should be at 1.5 MeV according to the calculations. Finally, the second excited state would contain $J^\pi = 1/2^-$ because it would be a $\pi p_{1/2}$ single particle state, located at 2.9 MeV. There would also be a $9/2^+$ state for the $\pi g_{9/2}$ single particle state, located at 6 MeV. It may or may not be neutron unbound depending on where the neutron separation energy is located. According to mass extrapolations it is at 5.65(64) MeV [41].

As stated before, the extrapolation for the average energy of core excited states to ^{79}Cu puts them at around 5 MeV for this nucleus which is only ^{78}Ni plus one proton. Also, the first excited state in ^{79}Ni may or may not be neutron bound as well. This may affect neutron capture rates of the nucleosynthesis r-process. If a neutron captures onto this nucleus and populates that state, this would lead to a specific direction for the next step in the r-process. However, if that state is unbound then this will lead to a different direction for the r-process path. Whether or not this state is bound could affect the population of proposed r-process waiting point nuclei ^{79}Cu and ^{80}Zn [42] [43]. See figure 35 for a description of these predictions for ^{79}Ni and ^{79}Cu .



Excitations of single proton in fp orbitals

Figure 35: Predictions for the decay of ^{79}Ni to ^{79}Cu , derived from the calculations discussed with figure 31. The $N=51$ ^{79}Ni is predicted to also have a $5/2^+$ ground state, just as the case with ^{81}Zn , ^{83}Ge , etc., and a $1/2^+$ first excited state. The $N=50$ ^{79}Cu will have single particle states with its single proton outside of ^{78}Ni and the core excited states should be centered at 5 MeV. The first excited state in ^{79}Ni may or may not be bound, depending upon where the neutron separation energy truly resides. The S_n value shown here for ^{79}Ni is a predicted value from [44].

6 The Decay Scheme of ^{79}Zn

The atomic nucleus ^{79}Zn contains 30 protons and 49 neutrons. From the shell model point of view it has two protons and one neutron hole with respect to the ^{78}Ni core. It undergoes β^- decay to populate states in the $Z=31$, $N=48$ nucleus ^{79}Ga . This decay has been previously studied [16].

In this work the ^{79}Zn ions were produced through the same method as the ^{81}Zn with the ISOL technique, where protons induce fission in the uranium carbide target. The fission fragments released from the ion source were first mass separated and then isobarically separated from other mass 79 isobars. They are then implanted on tape at the LeRIBSS setup. The γ decays were measured with the four HPGe clovers and β decay electrons are detected in the plastic scintillator detectors. Most of the measurements of the ^{79}Zn decay were performed with a two second grow-in of activity. Another data set was collected with one second grow-in and one second decay. Total beam time for measuring ^{79}Zn was about 4 hours, although more beam time was taken in an attempt to see more neutron-rich $A=79$ isobars such as copper. The beam intensity for ^{79}Zn is about 900 ions per second. This is approximately the same order of magnitude as the previous publication, but it is improved upon because of the detection of multiple new transitions in this decay.

This previously established level scheme has been verified with the data analysis in this study. The γ - γ coincidence data along with intensity balance and half-life measurements of most of the transitions support the same placement of γ transitions as is shown in previous literature [16]. The only exception is the ordering of two of the transitions. The order of the 1100 keV and 1212 keV transitions are switched because their intensities are observed to be different outside of error bars, whereas the previous literature did not observe this difference. It was only observed that these two transitions are in coincidence with each other. In this work the decay scheme has been expanded upon by the addition of nine new transitions and six new levels.

6.1 Verification of original ^{79}Zn decay scheme

The placement of each transition is supported at minimum by β - γ - γ coincidence data and intensity balance, but often times a half-life measurement also verifies its origin as being from the decay of the parent ^{79}Zn .

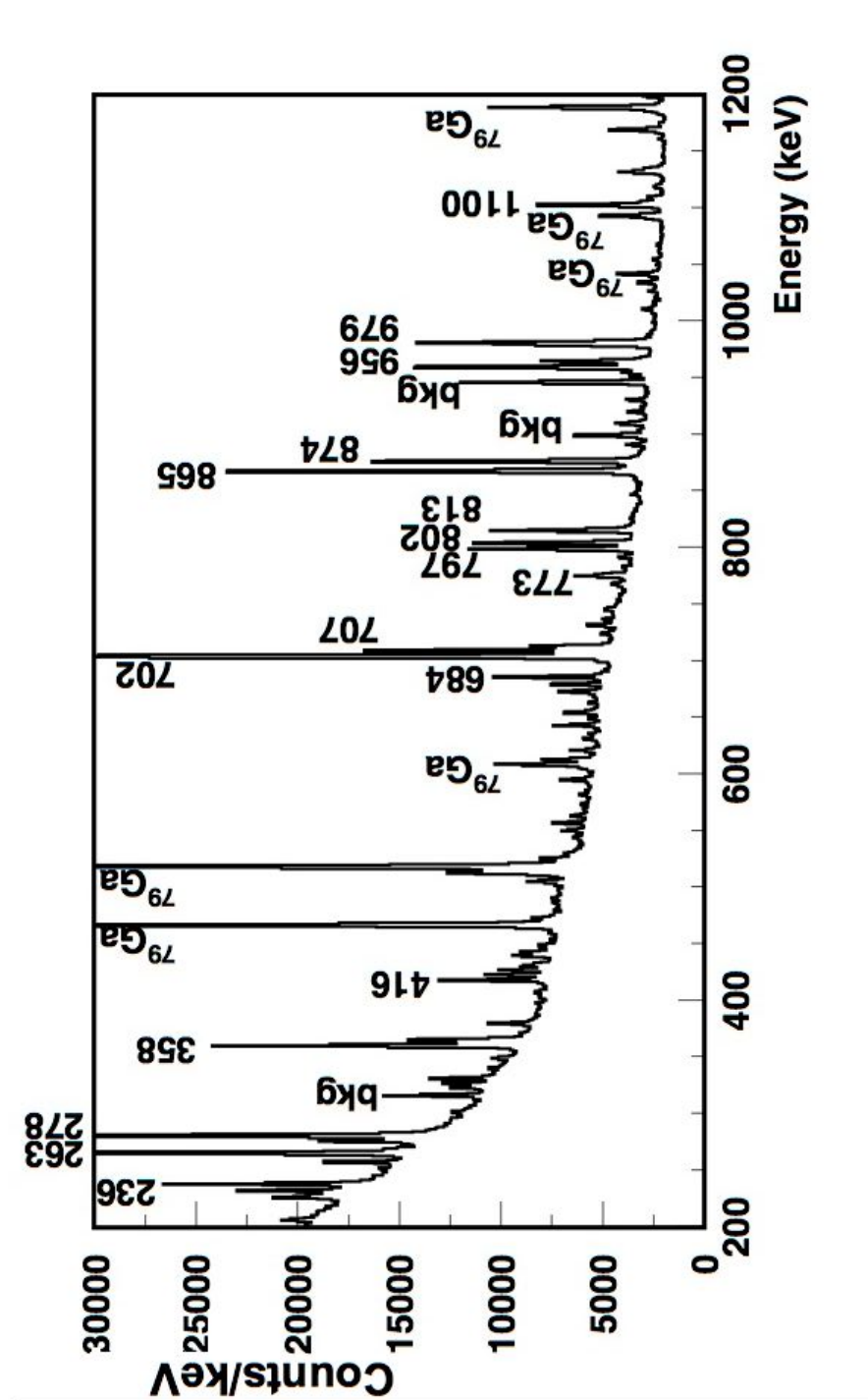


Figure 36: A sample of the β -gated γ -ray spectrum for ^{79}Zn to ^{79}Ga decay.

The first excited state was previously reported at 5.1 keV. This level is verified by the same transitions both feeding into it and those which are bypassing it to populate the ground state. The transitions of 273.8, 702, 797.1, 865.9, and 956.9 keV decay into this level at 5.1 keV. However, the transitions 278, 707.1, 802, and 962 keV all decay directly to the ground state. Thus the overlaps of these transitions reveal that a level must exist at 5.1 keV. Most of these transitions have other β - γ - γ coincidences and intensity summing which support their placements lower in the decay scheme as opposed to somewhere higher in energy. For example, the 278 and 273.8 keV transitions are observed in coincidence with 684 and 1778 keV transitions.

The second excited state is kept at 278.9 keV because of the transition of 278.9 keV decaying from it to the ground state, as previously discussed. The third excited state is at 707.1 keV because of the transition of 707.1 keV decaying from this level to the ground state. This transition is observed to be in coincidence with the transitions 874, 979, 1130, 1100, and 1212 keV. It is at the bottom of this cascade because of intensity balance. The fourth excited state at 802 keV is depopulated by a transition of 802 keV directly to the ground state. This transition is observed in coincidence with 263 and 813 keV, both of which are weaker in intensity and so they decay into the 802 keV level instead of out of it.

Table 10 lists all of the levels in the updated level scheme for ^{79}Ga along with the observed γ -ray transitions both decaying into and out of each level. The relevant $\gamma - \gamma$ coincidences are included.

6.2 Placement of new transitions in the ^{79}Zn decay scheme

There are new γ -ray transitions verified to be part of the ^{79}Ga decay scheme, populated in the decay of ^{79}Zn . These transitions are 730, 1802, 2156, 2419, 1575, 1723, 3205, 3933, and 3114 keV.

A transition of 1802 keV is placed as decaying from the new level established at 1807 keV down to the 5.1 keV level. This provides further support to the placement of this level in the level scheme. Furthermore, this transition is observed in coincidence with the 1212 keV transition and vice-versa. The 1212 keV transition is in coincidence with this 1802 keV γ -ray.

The 730 keV transition was found to be in coincidence with 358, 874, 979, and 702 keV transitions. Therefore, it is placed on top of this cascade of γ -ray transitions and it decays from a new level at 3649 keV.

Table 10: The decay of ^{79}Zn to ^{79}Ga information.

Level (keV)	Transitions into level (keV)	Transitions out of level (keV)	Relevant coincidences (keV)
5.1	273.8, 702, 797.1, 956.9, 1802, 3933	—	278.9, 273.8 with 684, 1778
278.9	684, 3205	278.9, 273.8	684 with 278.9, 273.8
707.1	874, 1130, 1100, 1941, 2034	707.1, 702	707.1 with 874, 979, 1130, 1100, 1212
802	813, 236	802, 797.1	802 with 263, 813
871	710, 1343	865.9	865.9 with 710, 1343
962	653, 1778, 2156, 2419	956.9, 962, 684	684, 956.9 with 653, 1778
1066	3114	263	263 with 802, 3114
1581	979	874, 710	710 with 856.9. 979 with 874
1615	—	653, 813	813 with 802, 797. 653 with 956, 962
1807	1212, 1575	1100, 1802	1212 with 1100. 1575 with 1100
1837.1	—	1130	1130 with 702, 707.1
2214	1723	1343	1723 with 1343
2561	773, 416, 358	979	979 with 773, 416, 358
2648.1	—	1941	1941 with 702
2741.1	236	2034, 1778	236 with 2034, 1778
2919	730	358	730 with 358
2977	—	416, 236	416 with 979 and 236 with 1778
3020	—	1212	1212 with 1100
3118	—	2156	2156 with 956.9
3335	—	773	773 with 979
3381	—	1575, 2419	1575 with 1100. 2419 with 956.9
3483	—	3205	3205 with 278
3649	—	730	730 with 358, 979, 874
3937	—	3933, 1723	1723 with 1343
4180	—	3114	3114 with 263

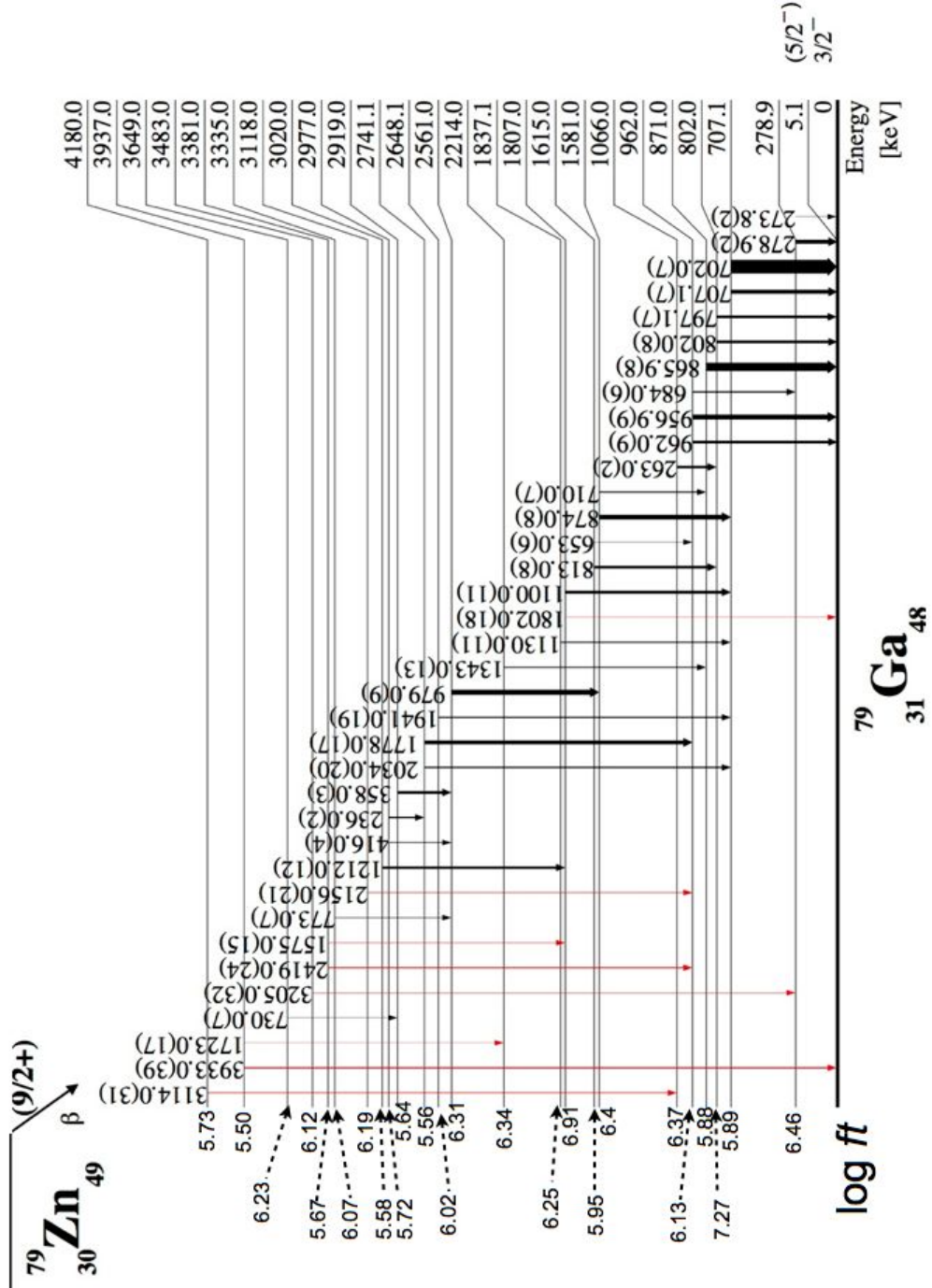


Figure 37: The level scheme for ^{79}Zn to ^{79}Ga decay. The $\log ft$ values for the β decays populating each state are shown. The newest transitions are marked in red arrows.

The transition of 2156 keV is seen in coincidence with 956.9, 962, and 684 keV. As a result, it is placed above these transitions and thus the nineteenth level at 3118 keV is a new state. The transitions 2419 and 1575 keV both decay from a new level at 3381 keV. The 2419 keV transition decays to the level at 962 keV and it is seen in coincidence with 956.9 and 962 keV. The 1575 keV transition decays to the level at 1807 keV and it is in coincidence with 1100 and 702 keV.

The 3483 keV level decays through a γ -ray transition of 3205 keV to the level at 278.9 keV. It is seen in coincidence with this 279 keV transition and the 278 keV transition is in coincidence with 3205 keV.

The 3937 keV level decays through the 3932.9 keV transition down to the level at 5.1 keV. The placement of this transition is given further support from the new transition 1723 keV and its placement. It decays from this level at 3937 keV down to the level at 2214 keV because it is seen in coincidence with the cascade of 1343 keV followed by 865.9 keV.

The 4180 keV level decays via the transition of 3114 keV, which goes to the level at 1066 keV. This 3114 keV transition is in coincidence with the 263 keV transition and the opposite coincidence of 263 keV with 3114 keV is also observed.

Thus this updated level scheme contains six new levels and nine new γ -ray transitions. The old level at 1919 keV has been replaced with the new 1807 keV level because the order of the 1100 and 1212 keV transitions was switched with respect to previous literature [16].

6.3 ^{79}Zn decay scheme discussion

The first excited state in the daughter, ^{79}Ga , is at 5.1 keV. A 5.1 keV γ -ray cannot be detected with HPGe γ -ray detectors as it is too low in energy and the efficiency is too low. Also, there is a significant amount of electrical noise in this energy range. As stated before, the existence of this level is known because of all the other γ -ray transitions either decaying to it directly or bypassing it to the ground state. Because the 5.1 keV transition is not observed directly, it is not possible to calculate a direct feeding to that level through β decay from the parent nucleus. As a result, the same problem exists with measuring direct feeding to the ground state. The experimenter here cannot distinguish between decay strength feeding the ground state from the

5.1 keV transition and direct feeding from β decay. Therefore the β decay branching ratios are shown to each individual excited state and to the neutron unbound states. However, feeding directly to the ground state will have to include feeding to the 5.1 keV first excited state as well.

Because this first excited state is so low in energy, it is possibly a β decaying isomer which decays to the granddaughter nucleus ^{79}Ge . The determination of its isomeric decay lifetime depends on both the energy of the state and the difference in intrinsic angular momentum between itself and the ground state of the ^{79}Ga . The ground state of ^{79}Ga has been measured to be $3/2^-$ [23]. The parity is likely to be negative because this nucleus contains three protons in the πf and πp orbitals with two neutron holes in the $\nu g_{9/2}$ orbital. In the shell model picture this means that all the particles and holes pair off until the net parity is negative. Thus the ground state of ^{79}Ga has $J^\pi = 3/2^-$. The first excited state may have $J^\pi = 5/2^-$ because other theoretical calculations show the ground state and first excited states to be either $J^\pi = 3/2^-$ or $5/2^-$. Furthermore, the systematics of less neutron-rich odd mass gallium isotopes all have either $J^\pi = 3/2^-$, $5/2^-$, or $1/2^-$ for the ground states. Here the ground state J^π is shown to be $3/2^-$ but the first excited state is tentatively labeled $5/2^-$ because it has not been directly measured. It is only inferred from systematics and theory predictions. The systematics for low-lying states in odd mass gallium isotopes are shown in figure 38. This is cited from the previously mentioned work where the ground state J^π ^{79}Ga and other nuclei were measured [23]. Given the facts that the ground state has $J^\pi = 3/2^-$ and the first excited state is most likely $J^\pi = 5/2^-$ with an energy of 5.1 keV, the possible multipolarity and lifetime of its γ decay can be calculated. The 5.1 keV transition is likely to decay via an M1 transition with a lifetime on the order of one to ten nanoseconds.

Furthermore, the ground state to ground state branching ratio must be zero. The first reason for this is from the shell model picture. Because the ground state of ^{79}Zn is $9/2^+$ and it decays to a $3/2^-$ ground state in the daughter, this is considered a third forbidden transition and as such there would be an unobservable and very small amount of decay strength directly feeding the ground state of the daughter. Also, there appears to be a significant amount of isobaric contamination in this beam according to both the intensity seen in the daughter γ decays and from attempts to fit the isobaric purity equations (see analysis chapter) to the grow-in of activity of ^{79}Zn and ^{79}Ga γ -ray transitions. These are the three reasons why it is concluded that there is essentially zero direct feeding of β decay strength to the ground state of ^{79}Ga .

Finally, the neutron branching ratio is calculated from a γ -ray transition emitted in the decay of ^{78}Ga to ^{78}Ge . The counts in its γ -ray peak are taken from the β -gated γ -ray spectrum because a small peak at this same energy appears in the background measurement. The calculated β -delayed neutron branching ratio for ^{79}Zn decay is 2.86(6) %. This is close to two previous values from literature: Hosmer *et al.* measured 2.2(1.4) % [27] and Pfeiffer *et al.* measured 1.3(4) % [45].

6.4 Isomers in N=49 nuclei

The nucleus ^{79}Zn is an N=49 isotone. There are certain isomers that appear in various parts of the chart of nuclei because of single particle states in the shell model picture, including for N=49 nuclei. The isomers of concern here are called spin-gap isomers. They normally form when two low energy states are significantly close to each other in energy and they possess a large difference in spin. The N=49 nuclei near ^{78}Ni contain neutrons in the $\nu g_{9/2}$ orbital. The possibility of isomers forming here exists because of the close proximity in energy of the $\nu g_{9/2}$ to the neutron orbitals $p_{1/2}$, $p_{3/2}$, and $f_{5/2}$. Various nucleon pairing scenarios may lead to states such as a $9/2^+$ close to either a $1/2^-$, $3/2^-$, or a $5/2^-$ state. Given the same energy difference between states, a $9/2^+$ state below or above a $1/2^+$ state with no others in between would lead to the longest lived, β decaying isomers. The equations noted in the theory section relate energy and spin difference for a γ -ray transition to its lifetime for decay.

The isomers in N=49 nuclei exist in the odd masses where an unpaired neutron creates a $9/2^+$ state next to a $1/2^-$ state. Thus the transition from the isomer to the ground state is simplistically viewed as a transition from a single hole in $\nu p_{1/2}$ to a single hole in $\nu g_{9/2}$. Table 11 lists these isomers.

In this work the only way to observe an isomer in ^{79}Zn is to measure the half-lives for the γ -ray transitions in its decay into ^{79}Ga and to note a significant difference in the half-lives. Some γ -ray transitions should be emitted from states populated in the decay of the ground state of the parent, with one half-life, and other transitions would decay from the isomer with a different half-life. However, in this work no such second half-life has been observed. This does not invalidate the existence of an isomer in ^{79}Zn , but it simply does not confirm the existence either because of the lack of evidence in this particular measurement. It is also possible that the half-life for this isomer is close enough to the half-life of the ground state decay such that the difference cannot be discerned in this experiment. They could be within error bars of each other.

Table 11: Isomers in N=49 isotones. Most of the odd A, N=49 isotones have a $9/2^+$ ground state and a $1/2^-$ first excited state, which is an isomer. The energies and lifetimes of these states are shown. The origin of most of these isomers, in a simplistic view, comes from a state with a hole (or particle) in $\nu p_{1/2}$ residing close in energy to a state with a particle (or hole, respectively) in $\nu g_{9/2}$. The nucleus ^{81}Ge is the exception with a $1/2^+$ state lower than the $1/2^-$ state, which has a half-life of about 0.5 ns.

Nucleus	Ground state J^π	Isomer J^π	Isomer energy (keV)	$T_{1/2}$ g.s.	$T_{1/2}$ isomer
^{81}Ge	$(9/2^+)$	$(1/2^+)$	679.1	7.6s	—
^{83}Se	$9/2^+$	$1/2^-$	228.5	22.3m	70.1s
^{85}Kr	$9/2^+$	$1/2^-$	305	10.7y	4.48h
^{87}Sr	$9/2^+$	$1/2^-$	388.5	stable	2.8h
^{89}Zr	$9/2^+$	$1/2^-$	587.8	78.4h	4.16m
^{91}Mo	$9/2^+$	$1/2^-$	653	15.49m	64.6s
^{93}Ru	$(9/2^+)$	$(1/2^-)$	734.4	59.7s	10.8s

A measurement of ^{79}Cu decay to ^{79}Zn may be the best option, or perhaps some reaction which populates the states in ^{79}Zn . In that latter case particle spectroscopy could be used to determine the J^π values of such states, which would in turn help to determine if a spin-gap isomer is present.

Table 12 shows some of the half-lives and their errors measured for various γ -ray transitions in ^{79}Zn decay into ^{79}Ga .

6.5 Measured S_β and calculated B(GT) values in ^{79}Zn decay

The investigation of allowed β decay strength distribution in ^{79}Zn decay is important because it may reveal where groups of states with J^π in ^{79}Ga which are close to the ground state J^π in the parent, ^{79}Zn , may reside. Some of those states may result from excitations of the ^{78}Ni core. Where in the Q_β window these states reside can contribute to the ability to discern where single particle states in the shell model picture reside for such neutron-rich zinc and gallium isotopes near ^{78}Ni .

A theoretical calculation was performed with a ^{56}Ni core and the model state has all traditional single particle states for the neutrons and protons between ^{56}Ni and ^{78}Ni . First this calculation was performed to determine the energies and J^π values of states in the daughter,

Table 12: Measured half-lives for transitions from the decay of ^{79}Zn to ^{79}Ga . The first three are previously measured transitions. The last two shown are examples of new transitions observed in this decay.

Transition (keV)	half-life (ms)	error(ms)
707.1	863	23
702.0	891	12
865.0	884	16
3933.0	810	74
1575.0	907	81

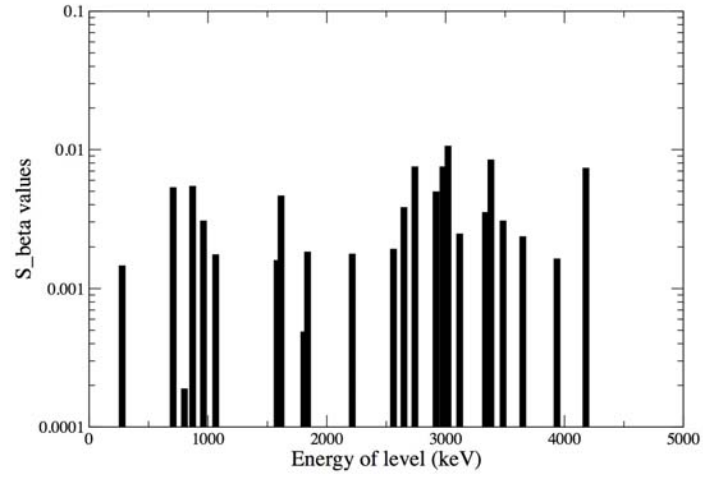


Figure 39: The experimental S_β values in ^{79}Ga , populated in the decay of the proposed $9/2^+$ ground state of ^{79}Zn .

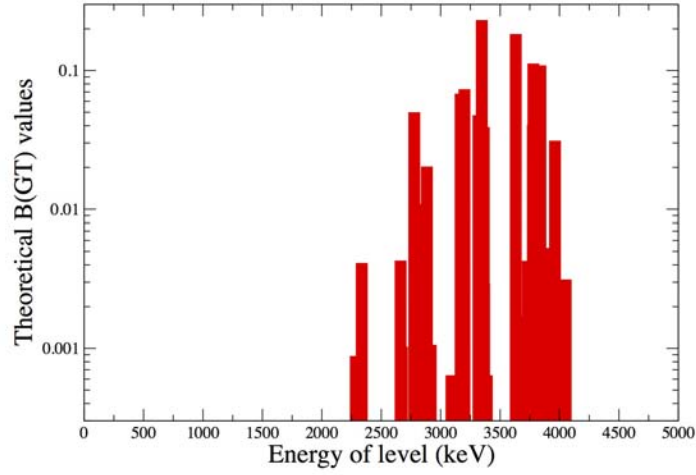


Figure 40: The theoretical B(GT) values in ^{79}Ga to predicted states populated in the decay of the proposed $9/2^+$ ground state of ^{79}Zn .

^{79}Ga . Then using the calculated wave functions for both these states in the daughter as well as that of the ground state in the parent nucleus, the reduced transition matrix elements for allowed Gamow-Teller transitions between these states are calculated. This calculation involves the overlap of the initial and final wave functions, as described in the equation for the B(GT) value in the theoretical background section. In the case of ^{79}Zn decay the ground state J^π is proposed to be $9/2^+$ because it has only one hole in the $\nu g_{9/2}$ orbital with two paired protons. Therefore the allowed Gamow-Teller decays from this state will only populate states in ^{79}Ga with J^π of $7/2^+$, $9/2^+$, and $11/2^+$.

The experimental B(GT) values here can only be called experimental S_β values because it is not known which observed states are truly populated in allowed Gamow-Teller decay vs. first forbidden decays. Some of the observed states will be populated through first forbidden decays as well. Thus these values for the population of states in ^{79}Ga are referred to as the experimental β decay strength distribution values, or experimental S_β values, and are shown in figure 39. The theoretical B(GT) values are shown adjacent to them in figure 40. A rudimentary comparison between the experimental and theoretical values shows that there may be some peaks in the theory values which align with the peaks in the experimental values.

Thus the evolution of the Gamow-Teller strength distribution from below to above the N=50 shell gap will be discussed in the next chapter once similar B(GT) values are displayed for ^{80}Zn decay.

6.6 ^{79}Zn decay conclusions

In conclusion, the ^{79}Zn to ^{79}Ga decay scheme has been expanded from 3.3 MeV to 4.18 MeV, in comparison to previous literature [16]. Nine new transitions and six new levels have been added to the decay scheme. One of those levels replaced a previous one because the order of the 1212 and 1100 keV transitions were also switched with respect to the previous literature. The observed β decay branching ratios allow for the calculation of experimental $\log ft$ values and S_β values to each observed level. Also, a theoretical shell model calculation was utilized to predict the B(GT) values to theoretical levels in ^{79}Ga as well. Finally, the evolution of ground state spin-parities in odd mass gallium isotopes is discussed along with possible isomerism in the parent, ^{79}Zn , as well as other N=49 isotones.

7 The Decay Scheme of ^{80}Zn

The decay of ^{80}Zn was also studied at the LeRIBSS setup. The same production method and experimental setup used for the ^{79}Zn and ^{81}Zn decays was also utilized here for ^{80}Zn decay measurements at LeRIBSS. Three files with two second grow-in of activity were taken and one file with four seconds of grow-in followed by four seconds of decay out was also measured. Approximately eight hours of beam time was taken. This decay has been studied by Ekström and Fogelberg [16] and J.A. Winger [48] and Nuclear data sheets [25] show the latest information on the decay scheme. The latest measurements on the β -delayed neutron emission branching ratio are included in the work by Hosmer *et al.* [27]. Compared to the latest measurement by J.A. Winger [48] there are about 3 to 4 times the statistics of ^{80}Zn decay in this work compared to previous work. This is based upon a comparison of the 712 and 716 keV peaks in their γ singles spectrum versus the one in this work. Finally, the half-life of ^{80}Zn decay is measured to be 540(20) ms based on the strongest transition at 712 keV, which is within error bars the same as the value from J.A. Winger [48] of 550(20) ms.

The nucleus ^{80}Zn is an N=50 nucleus and has 30 protons. Thus it has one magic number of nucleons. It undergoes β^- decay to populate states in the nucleus ^{80}Ga as well as some states in ^{79}Ga through β -delayed neutron emission. The latest decay level scheme [25] has been improved upon in this work. The following information verifies the original decay scheme. A sample spectrum taken during the measurement of ^{80}Zn decay is shown in figure 41.

7.1 Verification of original ^{80}Zn decay scheme

The first excited state was previously reported at 311.9 keV, with a γ -ray transition decaying from that level to the ground state. It is seen in coincidence with the 74 keV transition, and this transition is also seen in coincidence with the 311.9 keV. Therefore, the 74 keV transition is placed as feeding the 311.9 keV level and decaying from a level at 386.8 keV. The transition 1117 keV decays from a level at 1503.9 keV down to this level at 386.8 keV. This transition is seen in coincidence with 74 keV and 74 keV is itself in coincidence with 1117 keV.

The 174 keV transition was tentatively placed as decaying from a level at 485.9 keV because the same coincidence of 174 keV with 311.9 keV and 404 keV transitions are observed here, as is the case in the previous literature [48]. A transition of 486 keV is tentatively added as decaying from the level at 485.9 keV. This level and the two transition decaying from it are all

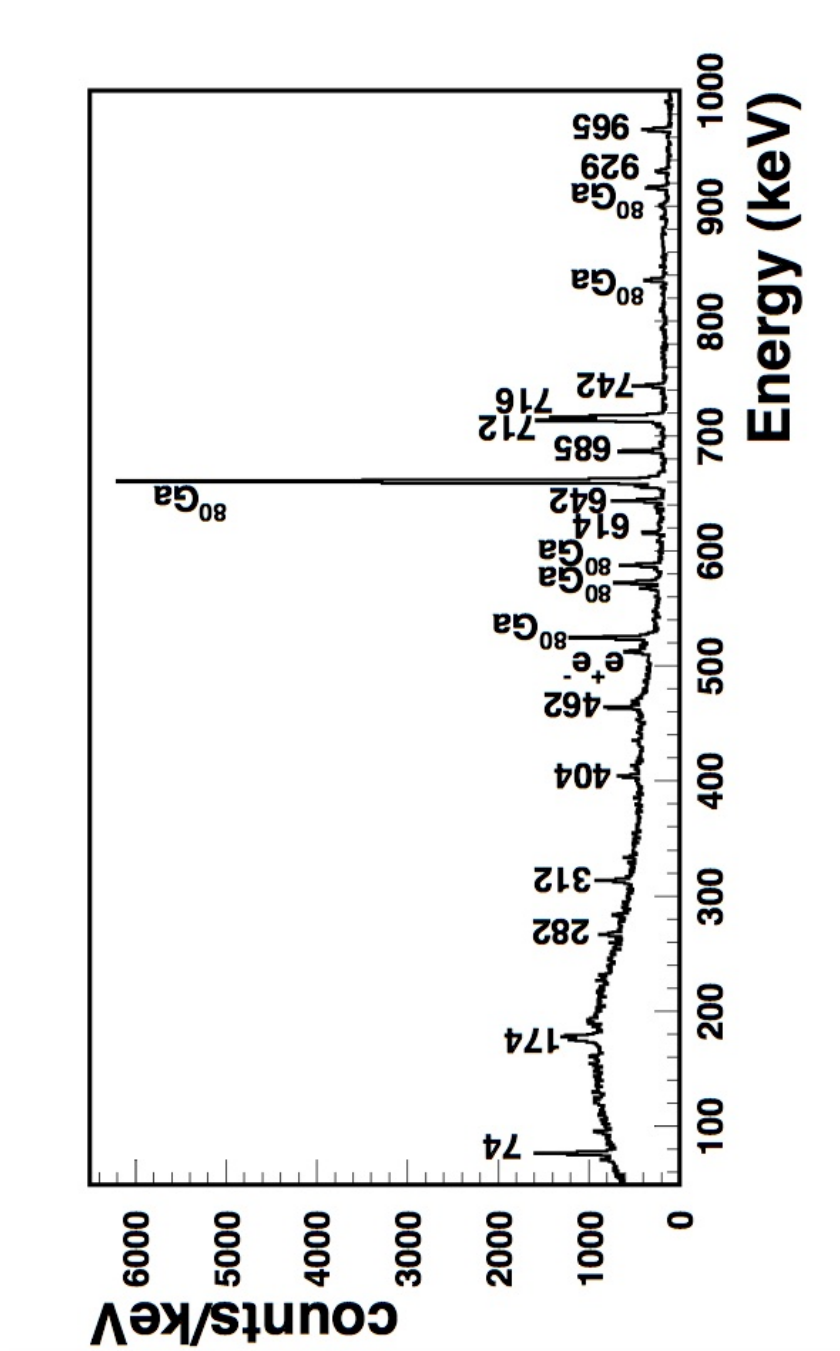


Figure 41: A sample from the β -gated γ -ray spectrum for the ^{80}Zn decay measurement.

Table 13: The decay information for ^{80}Zn to ^{80}Ga decay.

Level (keV)	Transitions into level (keV)	Transitions out of level (keV)	Relevant coincidences (keV)
311.9	74, 174	311.9	74 with 311.9
386.8	1117	74	74 with 1117
485.9 (tent.)	404	174, 485.9	174 with 311.9 and 404 with 174
685.6	742	685.6	685.6 with 742
712.6	716, 177, 791	712.6	712.6 with 716, 791
889.3	282, 614	889.3, 404, 177	889.3 with 614 and 282 with 177
929	2378	929	929 with 2378, 190
965	462	965	965 with 462
1171	332	1171, 282	1171 with 332
1334.1	94, 736, 169	1334.1	1334.1 with 94, 169, 736
1428	642, 1226	94, 462, 716, 742	642 with 94, 462, 716, 742
1503.9	567, 1152	169, 332, 614, 791, 1117, 1503.9	567 with 74, 311.9, 614, 791, 1117, 1503.9
2070.2	—	567, 642, 736	736 with 1334.1 and 567 with 614
2655	—	1226, 1152	1226 with 462, 716 and 1152 with 614, 791
3307	190	2378	190 with 2378, 929
3497	—	190	190 with 2378, 929

tentative here because the amount of intensity observed as decaying directly to the 485.9 keV level is slightly negative and is not quite zero within error bars. It is thus possible that an extra amount of very weak γ -ray intensity decaying from this level down to say the 386.8 keV level is present but not observed here. This would adjust the decay branching ratio to the 485.9 keV level to zero within error bars.

7.2 Placement of new transitions in the ^{80}Zn decay scheme

The first two transitions which are added to the decay scheme are 929 keV and 2378 keV. They are in coincidence with each other. The 929 and 2378 keV transition are also in coincidence with a 190 keV transition. Therefore, because of the coincidences and because of the intensities

of these three transitions, they are placed in cascade with 190 keV feeding into the 2378 keV, in turn feeding the 929 keV transition. The 929 keV decays into the ground state. This also establishes new levels at 929 keV, 3307 keV, and 3497 keV.

Another transition of 1226 keV is discovered and is placed into the scheme in figure 42. It decays from the level at 2655 keV into the level at 1428 keV. This transition is seen in coincidence with the 462, 712.6, and 716 keV transitions. In conjunction with this the 712.6 and 716 keV transitions are both observed in coincidence with the 1226 keV transition.

7.3 Isomeric state in ^{80}Ga

According to Cheal *et al.* a long-lived, low energy isomer exists in the nucleus ^{80}Ga [49]. The main conclusion reached in that work is that there is an isomeric state present from the hyperfine structure (hfs) observed in the atomic transitions because there is more than one set of hfs peaks, corresponding to more than one state with more than one nuclear spin. However, the energy and lifetime of this isomeric state is undetermined. It is also not possible for them to determine whether it is the 3^- state that is the ground state and the 6^- state is the isomer, or reversed. Previous publications put the ground state J^π of ^{80}Ga as 3^- [21] [48].

In the data presented in this work, no evidence for this isomeric state is present. It is possible that this state decays mostly by β decay to ^{80}Ge instead of γ decay to the ground state. This would explain why no transition in this energy range is observed in the γ -ray spectra. A β decaying isomer from ^{80}Ga could have a half-life different enough from the half-life of the ^{80}Ga ground state such that the experimenter could measure it outside of error bars. However, there are no such γ -ray transitions in the decay of ^{80}Ga to ^{80}Ge which have half-lives that differ outside of error bars.

Another reason why this isomer is not observed in this work is because it probably was not populated. In this experiment ^{80}Zn is produced and it decays from a 0^+ ground state. According to Cheal and others, the isomeric state is either 3^- or 6^- , both of which are extremely difficult to populate from a 0^+ state because of β decay selection rules and a huge change in spin and parity. Thus it is highly likely that this isomeric state was simply not populated in the β decay observed here, whereas in the experiment by Cheal it is the gallium itself which is directly extracted from the ion source and studied with a laser to probe the atomic spectroscopy. Thus

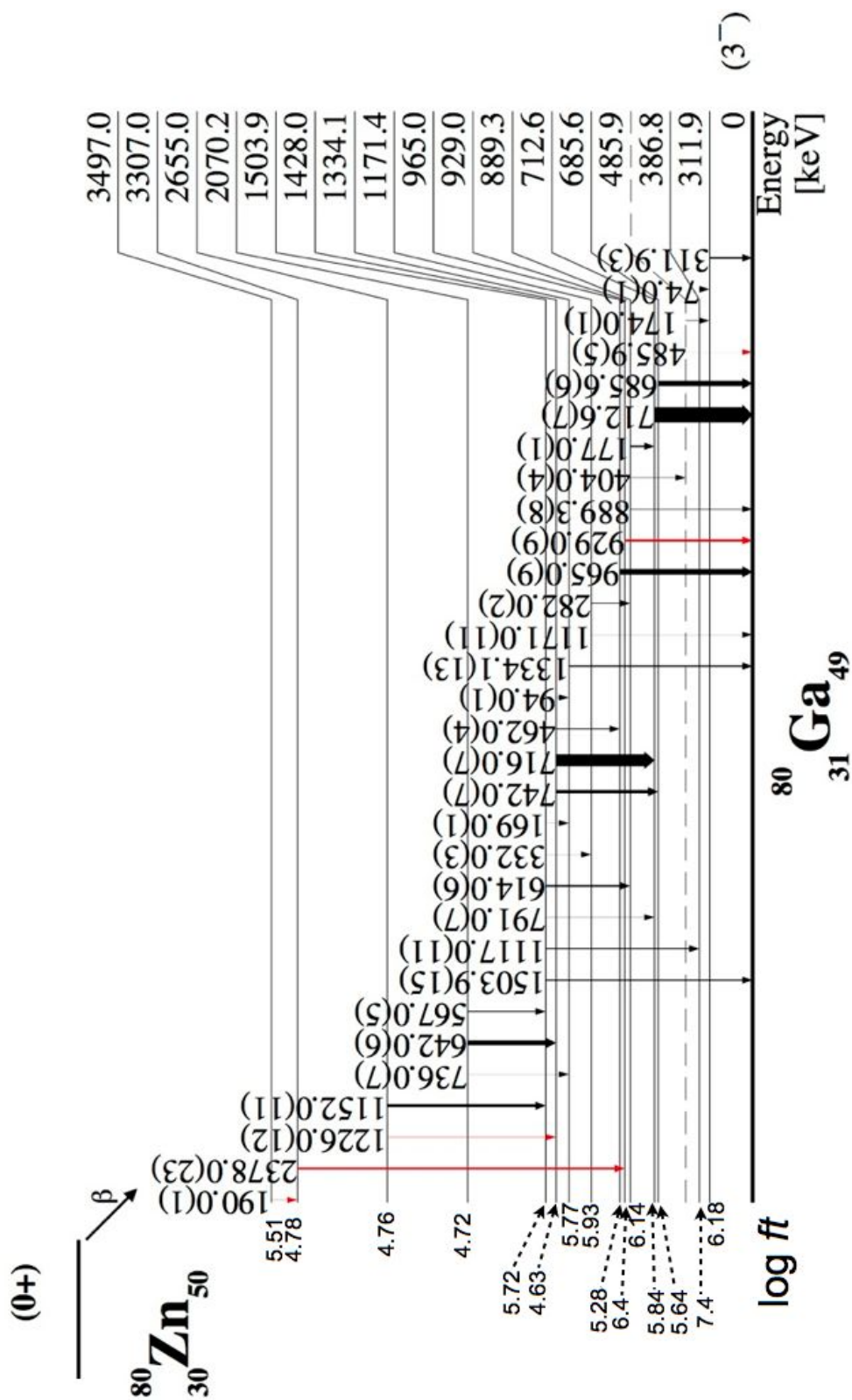


Figure 42: The decay scheme for ^{80}Zn to ^{80}Ga .

the isomeric state was populated in that experiment but probably not in the work presented here.

7.4 Discussion on ground state branching ratio

The ground state of ^{80}Ga is most likely either a 3^- or a 6^- state according to the isomeric state publication discussed above [49], as well as [21] and [48]. In the calculation of all β decay branching ratios to each excited state and the ground state in ^{80}Ga , this fact plays an important role in determining at least the ground state to ground state branching ratio. The parent nucleus ^{80}Zn is an even-even nucleus and therefore the spin and parity of the ground state must be 0^+ . The ground state of ^{80}Ga is either a 3^- state or a 6^- state, depending on which one is the isomeric state. A transition from a 0^+ to 3^- state is a third forbidden transition by selection rule definitions. The lowest $\log ft$ value on record for such a transition is 17.5 [9]. This is approximately 11 or more orders of magnitude weaker in transition strength than allowed β decay transitions.

The beam intensity for ^{80}Zn was about 100 ions/second. In order to see any of this direct feeding appear in the γ spectra the beam intensity would have to be increase to about 10^{13} ions/second, the amount of change needed to move from observing only allowed decays to observing third forbidden decays. An increase that large is currently not possible at any radioactive ion beam facility.

A 0^+ to 6^- transition is almost impossible. There is not a forbidden category labeled for it. If it existed it would be on the order of a fifth forbidden transition and a transition strength about 20 or more orders of magnitude weaker than allowed β decay transitions. It is safe to conclude that such a transition will not occur. Even if it ever did, it would not occur with high enough frequency that any experiment performed today or in the near future would be able to produce enough beam intensity to see this direct feeding.

7.5 Measured S_β and calculated B(GT) values in ^{80}Zn decay

The experimental S_β values for the population of states in ^{80}Zn are shown in figure 43. The theoretical B(GT) values for predicted levels in ^{80}Ga are shown adjacent, in the next figure down 44. A rough comparison between the experimental and theoretical values shows an agreement that there is an enhancement of Gamow-Teller decay strength to a level at about 1.4 MeV. At

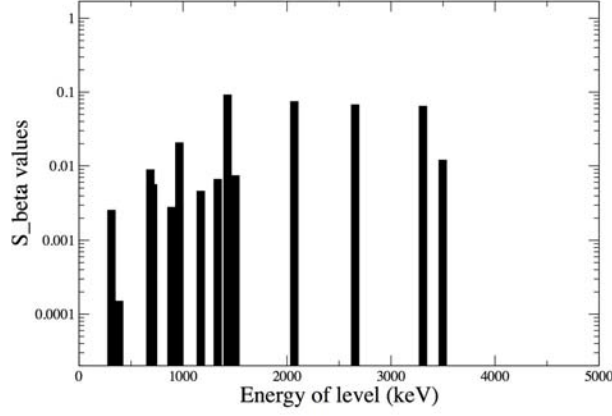


Figure 43: The experimental S_β values for levels populated in ^{80}Ga .

higher energies the experimental S_β values seem to have a plateau effect from 1.4 MeV to about 3.3 MeV. There are four predicted states in this energy range which show a similar behavior. Thus although there is not complete agreement between experimental and theoretical values, there are comparisons to be made between relative sizes and locations of enhanced $B(\text{GT})$ values in this decay.

As stated before, the allowed Gamow-Teller decay in ^{81}Zn to ^{81}Ga decay populates core excited states in the daughter because of the creation of a neutron hole in the f and p orbitals below the $N=50$ gap. The same Gamow-Teller decay can occur in ^{79}Zn to ^{79}Ga and ^{80}Zn to ^{80}Ga via the decay of a neutron in f and p orbitals, but in these nuclei the hole is not paired with any neutron above the $N=50$ gap. Thus they are not core excited states as in the case of ^{81}Zn decay. However, this comparison shows that the average energy of states populated through the decay of a neutron in f and p orbitals rises in energy from ^{79}Ga to ^{81}Ga . There is an increase from about 3 MeV in ^{79}Zn decay to about 4.4 MeV in ^{81}Zn decay. However, it should be noted that many predicted states above 3.3 MeV in ^{80}Zn decay to ^{80}Ga are not seen. Therefore this is an incomplete picture, especially for ^{80}Ga . This must be because of a drop in our sensitivity limits above this value. Given our efficiencies for detecting γ -ray transitions higher than 3.3 MeV means that the branching ratios to states above that energy are too small to detect given the current clover array setup and beam intensity.

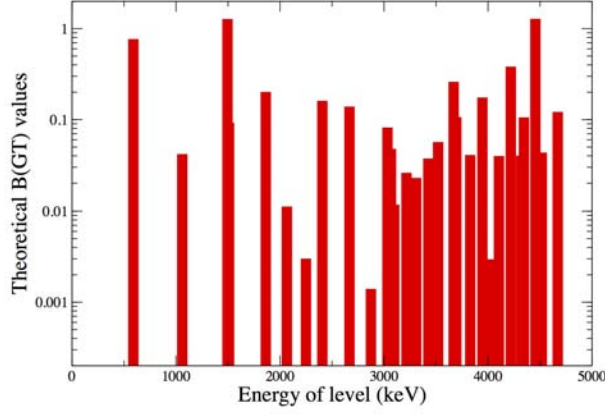


Figure 44: The theoretical B(GT) values for predicted levels populated in ^{80}Ga .

7.6 ^{80}Zn decay conclusions

The decay level scheme for ^{80}Zn to ^{80}Ga has been expanded from about 2.6 MeV to above 3.3 MeV. It was produced and studied at LeRIBSS through the ISOL technique, the same as with ^{79}Zn and ^{81}Zn decays. There is essentially no neutron branching ratio observed because no γ -ray transitions are observed which would be populated from such decay, whereas the previous literatures report 1% or less [48] [25]. An upper limit of 0.8% is stated in this work because this is a typical error on measuring such branching ratios in this analysis. The ground state to ground state branching ratio is approximately zero because of the large change in spin which would be needed for such a decay. Finally, the experimental S_β and theoretical B(GT) values for the populations of levels in ^{80}Ga are shown and discussed in comparison with ^{79}Zn and ^{81}Zn . This comparison shows a rise in the energy of states populated through the decay of neutrons in f and p orbitals from ^{79}Zn to ^{81}Zn .

8 Conclusions and Summary

The study of nuclear structure away from stability is accomplished through a number of different types of experimental techniques. Decay spectroscopy is the general method used in the experiments discussed in this work. Here the decay spectroscopy of $^{79,80,81}\text{Zn}$ nuclei is carried out and additions to the decay schemes are shown and discussed.

All three of these zinc ions are produced at the Holifield Radioactive Ion Beam Facility at Oak Ridge National Laboratory. They are produced through the Isotope Separation On-line Technique, starting with the acceleration of protons to induce fission in a uranium carbide target. The fission fragments are first mass analyzed to produce a beam of particles with the same mass number. Next they are sent through a high resolution isobaric separator to separate neighboring isobars from the desired ion. An isotopically pure or nearly pure beam is sent to the LeRIBSS, the low energy radioactive ion beam spectroscopy station. The ions of interest implant upon a tape, part of a moving tape collector. The site of implantation is surrounded by four high purity germanium clovers to measure γ decays and two plastic scintillators for detecting the β decay electrons.

There are nine new γ -ray transitions and five new levels added to the ^{81}Ga level scheme compared to previous literature [2]. The β decay branching ratios are also reported to each level. It is concluded that the lowest energy states observed up to 1936 keV are populated through first forbidden decays from the ground state of the parent, ^{81}Zn . However, the two states above 4 MeV placed in the level scheme are populated through allowed Gamow-Teller decays. This is because the branching ratios to those two states above 4 MeV are higher than most of the other branching ratios in the decay scheme and the $\log ft$ values are lower than the values for all of the other states. A neutron branching ratio of 12(4) % is measured here and it is in agreement with the value predicted by Borzov [26] of 10(3) %. It is however in disagreement with the value reported by Hosmer *et al* [27]. This is a disagreement which will have to be settled in future experimentation. However, our data comes from the condition of “gating” on a specific γ -ray transition in a shorter time window, as opposed to “gating” on a neutron event in a longer time window where said neutron may or may not come from daughter decays or is a background neutron. This lends more support to our measurement, as explained before.

The ground state to ground state branching ratio in ^{81}Zn decay involves calculating any isobaric impurity that may be present. In an attempt to determine the isobaric purity of the

beam, appropriate equations are fitted to the grow-in of activity from the strongest γ -rays of the parent and daughter decays. These fits reveal a 5 % impurity in the beam. Thus the outgoing intensity from the ^{81}Ga ground state decay is adjusted for this amount of extra intensity observed in the decay spectra.

The other contribution to the ground state to ground state branching ratio is from the pandemonium effect. The decay of an atomic nucleus with a large Q_β decay window to its daughter can lead to the emission of high energy, weak γ -ray transitions. The germanium detectors have a low efficiency of detecting high energy γ -rays. The problem increases if the decay strength is fragmented amongst a high density of states, leading to weaker transitions. Thus it is possible that the array of germanium detectors is not measuring many transitions feeding into the ground state of the daughter, thus leading to a ground state branching ratio which is larger than the real value. To approximate for this missing intensity from this pandemonium effect, a comparison is made to the decay of $N=51$ ^{83}Ge to $N=50$ ^{83}As . The physics of this decay is similar because according to J.A. Winger [31] the population of low energy states is through first forbidden decays and the population of high energy states centered around 4 MeV are through allowed Gamow-Teller decays. The assumption is made that at minimum the ^{81}Ga level scheme will contain as many states in this 3 - 5 MeV range which are seen in ^{83}As . Only two levels are currently seen in this range in ^{81}Ga , thus the remaining 11 states are added as “fictitious” states to this level scheme, with branching ratios to match those seen in ^{83}Ge to ^{83}As decay. Therefore, after adjusting for this new intensity feeding into the gallium ground state, a lower limit of 34% is set on the ground state to ground state branching ratio from ^{81}Zn decay. Thus the ground state to ground state branching ratio is artificially inflated by at most 18%, according to this estimation.

This modified branching ratio along with $N=51$ systematics, β decay selection rules, and the known ground state J^π of ^{81}Ga allow for the determination of the ^{81}Zn ground state J^π . There are only two options for this J^π value, according to $N=51$ systematics. It is either a $1/2^+$ state or a $5/2^+$ state. The ground state of ^{81}Ga is known to have $J^\pi = 5/2^-$ because the spin was directly measured by Cheal *et al.* [23] and the parity must be negative because gallium contains three protons in the f and p orbitals above $Z=28$ shell gap. If ^{81}Zn decays from a $1/2^+$ state to a $5/2^-$ state in the gallium, this would be a first forbidden unique decay. However, for this decay to be first forbidden unique the ground state to ground state branching ratio would have to be less than 0.1% and its $\log ft$ value would be 7.5 or higher. However this branching ratio

is at least 34% or higher and the $\log ft$ is about 5.6 instead. This is a two orders of magnitude difference. Therefore, the only option left is that the decay of ^{81}Zn involves a $5/2^+$ state decaying to a $5/2^-$ state, which is a first forbidden decay. Such decays are within the range of the $\log ft$ value of 5.6 [9], and so this is the conclusion reached: The ground state J^π of ^{81}Zn is $5/2^+$.

This information is added to the $N=51$ systematics knowledge. Theoretical shell model calculations are used with N3LO effective interactions and a ^{78}Ni core. A calculation for the ground states and first excited states in $N=51$ isotones shows that the ^{81}Zn is expected to have a $1/2^+$ first excited state and a $5/2^+$ ground state, in agreement with the experimental conclusion in this work.

These shell model calculations show a change in trend in the energies of the first excited states from below to above $Z=34$. It is believed that this change in trend is driven by proton occupation number. As protons fill the $\pi f_{5/2}$ orbital from ^{79}Ni to ^{85}Se , the first excited states are moving down in energy. However, from ^{85}Se to ^{89}Sr this first excited state is moving up in energy while the protons are filling the $\pi p_{3/2}$ orbital. Thus different proton-neutron interactions are taking place below and above $Z=34$, exactly where the change in trend of first excited state energies occurs. It is believed that the main proton-neutron interaction driving the behavior of these states below $Z=34$ is the $\pi f_{5/2}$ with the $\nu s_{1/2}$. Above $Z=34$ the main proton-neutron interaction driving the first excited states' energies is believed to be the $\pi p_{3/2}$ with the $\nu s_{1/2}$, which has a change in sign from the previous proton-neutron interaction, thus explaining the change in trend of those states.

The behavior of core excited states in $N=50$ isotones is also shown and discussed. For the decays which populate the $N=50$ ^{85}Br , ^{83}As , and ^{81}Ga isotones, a B(GT) weighted energy is calculated from experimental information. This is the average energy of core excited states, states which are only populated in allowed Gamow-Teller β decays. The trend in these core excited state energies indicates an increasing gap between the deeply bound neutron hole and the valence neutron above $N=50$ gap upon moving towards ^{78}Ni . An extrapolation was shown for ^{79}Cu , which only contains one proton outside of the doubly magic ^{78}Ni , and is expected to have core excited states centered around 5 MeV.

Using information on 5^+ and 6^+ yrast states in even Z , $N=50$ isotones allows for a discussion of possible shell gap changes around $N=40$ and $N=50$. According to Rzaca-Urban *et al.* [39], a

lowering in energy of these states upon moving towards ^{78}Ni indicates a weakening of the N=50 shell gap. In conjunction with the conclusion reached in this work about the neutron f and p orbitals moving away from the neutron $d_{5/2}$ orbital above N=50 gap, this must mean that the N=40 subshell gap is increasing upon moving towards ^{78}Ni . However, these are not definite conclusions because the interpretation from the yrast states may not be correct. Even if they are correct, these states may move up again in energy at isotones with Z lower than 34. In that case an increase in the N=50 gap would be indicated.

The decays of ^{79}Zn and ^{80}Zn have been previously studied but measured again at LeRIBSS. In this work the decay scheme of ^{79}Zn was expanded from about 3.3 MeV to above 4.1 MeV. It now has nine new transitions and six new levels. The possibility of isomerism in ^{79}Zn is discussed, along with a comparison to isomers in other N=49 isotones. Also, the evolution of J^π values for the ground and low lying states in odd mass gallium isotopes is shown from previous work, which includes both the empirical values and two sets of theoretical values. The important point is that the ground state was previously measured to be of $J^\pi = 3/2^-$, whereas the low lying first excited state at 5.1 keV must be $5/2^-$, inferred from systematics. Thus very little ground state to ground state branching is expected here because it would be a third forbidden transition, $9/2^+$ to $3/2^-$ decay. Also, this beam was not isotopically pure or nearly pure, thus it is not possible to extract a reliable ground state branching ratio.

The decay scheme for ^{80}Zn was expanded from about 2.6 MeV to above 3.4 MeV. There are four new transitions and three new levels in this decay scheme. The experimental S_β and theoretical B(GT) values for population of states in ^{80}Ga are also shown.

9 My Contributions

My main, personal contribution to this work is the analysis of the data sets for the decays of these three nuclei. I managed the data analysis code needed to produce appropriate spectra for extracting the necessary spectroscopic information. Based upon analysis of these spectra I was able to build the level schemes for these decays. Ultimately this is what led to obtaining neutron branching ratios, the J^π value of the ^{81}Zn ground state, core excited state systematics, etc. This led to the publication on ^{81}Zn decay: S. Padgett *et al.*, Phys. Rev. C **82**, 064314 (2010).

I learned how to use shell model calculation programs from both M. Hjorth-Jensen and A. Brown. I used them to investigate the behavior of $N=51$ systematics in light of our new experimental information on ^{81}Zn . Using the Nushellx package from A. Brown and others allowed me to calculate both energy spectra of nuclei around ^{78}Ni and $B(\text{GT})$ values in the decays of those nuclei. I also made use of Nushellx to calculate the energy spectra of neutron unbound states in nickel decays for VANDLE (Versatile Array for Neutron Detection at Low Energy) proposals.

I contributed heavily to the construction of the miniaturized MCP detector setup, now used at the LeRIBSS setup. I of course also participated in multiple experiments at the LeRIBSS setup and also at the RMS (Recoil Mass Separator) as well as experiments in the National Superconducting Cyclotron Laboratory.

References

References

- [1] S. Padgett *et al.*, Phys. Rev. C **82**, 064314 (2010).
- [2] D. Verney *et al.*, Phys. Rev. C **76**, 054312 (2007).
- [3] T.T.S. Kuo and G.E. Brown. Nuclear Physics **85** (1966).
- [4] M.G. Mayer and J.H.D. Jensen. Elementary theory of nuclear shell structure. Wiley New York (1955).
- [5] A. Das and T. Ferbel, Introduction to Nuclear and Particle Physics, Second Edition. World Scientific Publishing Co. Pte. Ltd., Hackensack, NJ (2003).
- [6] C.S. Wu *et al.*, Phys. Rev. **105**, no. 4, 1413-1415 (1957).
- [7] K. Grotz and H.V. Klapdor, The Weak Interaction in Nuclear, Particle, and Astrophysics. IOP Publishing Ltd (1990).
- [8] Nuclear physics class notes, Professor Witek Nazarewicz.
- [9] B. Singh *et al.*, Nuclear data sheets 84, 487 (1998).
- [10] C.J. Gross *et al.*, Eur. Phys. Journal A25, 115 (2005).
- [11] B. A. Tatum, Nucl. Instrum. Methods Phys. Res., Sect. B 241,926 (2005).
- [12] J.R. Beene *et al.*, J. Phys. G: Nucl. Part. Phys. **38**, 024002 (2011).
- [13] Nivaldo J. Tro, Chemistry: A Molecular Approach (2nd Edition). New Jersey: Pearson Prentice Hall. ISBN 0-13-100065-9. pp. 348 - 349. (2008) (via Wikipedia.org).
- [14] A.J. Mendez II, private communication.
- [15] Glenn F. Knoll, Radiation Detection and Measurement, third edition. John Wiley and Sons, Inc. New York (2000).
- [16] B. Ekström, B. Fogelberg, P. Hoff, E. Lund, and A. Sangariyavanish, Physica Scripta Vol. 34, 614-623 (1986).
- [17] Proakis, J.G. and Manolakis, D.G., Digital signal processing: principles, algorithms and applications, Macmillan (1996).
- [18] R. Grzywacz. Applications of digital pulse processing in nuclear spectroscopy. Nuclear Instruments and Methods in Physics Research, B, 204:649-659 (2003).

- [19] Display Analysis and Manipulation Module reference manual, August 28, 2005.
- [20] S. Ilyushkin *et al.*, Phys. Rev. C **80**, 054304 (2009).
- [21] P. Hoff and B. Fogelberg, Nucl. Phys. A **368**, 210 (1981).
- [22] A. Lisetskiy *et al.*, Phys. Rev. C **70**, 044314 (2004).
- [23] B. Cheal *et al.*, Phys. Rev. Lett. **104**, 252502 (2010).
- [24] R. Grzywacz. RIA summer school lecture notes (2005).
- [25] Nuclear Data Sheets 105, 223 (2005).
- [26] I.N. Borzov, Phys. Rev. C **71**, 065801 (2005).
- [27] P. Hosmer *et al.*, Phys.Rev. C **82**, 025806 (2010).
- [28] J.C. Hardy, L.C. Carraz, B. Jonson, and P.G. Hansen, Physics Letters, Vol. 71B, No. 2 (1977).
- [29] R.B. Firestone, Physics Letters, Vol. 113B, No. 2 (1982).
- [30] J.C. Hardy, B. Jonson, and P.G. Hansen, Physics Letters, Vol. 136B, No. 5,6 (1984).
- [31] J. A. Winger, J. C. Hill, F. K. Wohn, R. L. Gill, X. Ji, and B. H. Wildenthal, Phys. Rev. C **38**, 285 (1988).
- [32] P.A. Moore, P.J. Riley, C.M. Jones, M.D. Mancusi, and J. L. Foster, Jr., Phys. Rev. C. **1**, 1100 (1970).
- [33] R. Machleidt and I. Slaus, J. Phys. G: Nucl. Part. Phys. **27** R69-R108 (2001).
- [34] J.S. Thomas *et al.*, Phys. Rev. C **76**, 044302 (2007).
- [35] J.A. Winger *et al.*, Phys. Rev. C **81**, 044303 (2010).
- [36] M. Hjorth-Jensen, T. T. S. Kuo, and E. Osnes, Phys. Rep. 261, 125 (1995).
- [37] R. Machleidt, arXiv:0704.0807.
- [38] R. F. Casten, Nuclear Structure from a Simple Perspective, second edition. Oxford University Press (2005).
- [39] T. Rzaca-Urban, W. Urban, J.L. Durell, A.G. Smith, and I. Ahmad, Phys. Rev. C **76**, 027302 (2007).

- [40] M.G. Porquet *et al.*, Eur. Phys. J. A **39**, 295 (2009).
- [41] G. Audi, A. H. Wapstra, and C. Thibault, Nucl. Phys. A **729**, 337 (2003).
- [42] K.L. Kratz *et al.*, Astrophys. J. 403, 216 (1993).
- [43] K.L. Kratz *et al.*, Z. Phys. A, Hadrons and Nuclei, **340**, 419-420 (1991).
- [44] V.I. Isakov *et al.*, Physics of Atomic Nuclei, Vol. **63**, No. 10 pp. 1740-1746 (2000).
- [45] B. Pfeiffer *et al.*, Prog. Nucl. Energy 41, 39 (2002).
- [46] M. Honma *et al.*, Phys. Rev. C **80**, 064323 (2009).
- [47] B.A. Brown (unpublished).
- [48] J.A. Winger *et al.*, Phys. Rev. C **36**, 2 (1987).
- [49] B. Cheal *et al.*, Phys. Rev. C **82**, 051302 (2010).
- [50] J. Hakala *et al.*, Phys. Rev. Lett. **101**, 052502 (2008).
- [51] Photomultiplier Tubes and Assemblies for Scintillation Counting and High Energy Physics, Hamamatsu.
- [52] M. A. Preston, Physics of the Nucleus. Addison-Wesley Publishing, Inc. Reading, Massachusetts (1962).

10 VITA

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He went on to major in physics at the Florida State University. These years included two summers performing research as an undergraduate at both Florida State University and Michigan State University. The summer after undergraduate school was spent performing research at Oak Ridge National Laboratory before beginning graduate school at the University of Tennessee - Knoxville. Here he participated in various nuclear physics experiments at both Oak Ridge National Laboratory as well as the National Superconducting Cyclotron Laboratory of Michigan State University. During graduate school he also decided to learn the Spanish language. After three years of progress he is at an “advanced” level, according to modern foreign language professors. In his spare time he enjoys gardening, reading philosophy, and building model railroads.