Study of Nuclear Level Densities from Deuteron-Induced Reactions on $^{54,56,58}$Fe and $^{63,65}$Cu

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This dissertation titled
Study of Nuclear Level Densities from Deuteron-Induced Reactions on $^{54,56,58}$Fe and $^{63,65}$Cu

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Abstract

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The nuclear level density is known to be the major source of uncertainty in the calculation of compound cross sections. The problem is attributed to the fact that most experimental level density information is obtained from the neutron resonance data, which are confined to narrow ranges of excitation energy and angular momentum. In addition, the level density derived from the average neutron resonance spacing is strongly dependent on the spin cut-off parameter, which is not known experimentally at the neutron binding energy. To minimize these uncertainties, we study the nuclear level density using the analysis of evaporation spectra from compound nuclear reactions.

Double differential cross sections and angular distributions for the $(d,n)$, $(d,p)$, $(d,d)$, and $(d,\alpha)$ reactions have been measured with deuteron energies between 5 and 9 MeV. The angular distributions of outgoing particles have indicated that cross sections measured at backward angles are dominated by compound reaction mechanism. These cross sections were used to test known nuclear level density models commonly applied in nuclear reaction codes: (1) the Fermi gas model [1]; (2) the Gilbert-Cameron model [2]; and (3) the microscopic model using the Hartree-Fock-Bogoliubov method [3].

The nuclear level densities of $^{64}$Zn, $^{66}$Zn, $^{55}$Fe, $^{57}$Fe, $^{55}$Co, $^{57}$Co, $^{59}$Co, $^{52}$Mn, $^{54}$Mn, and $^{56}$Mn have been deduced from the experimental compound cross sections. Most of the best-fit parameters have been found to agree with the Gilbert-Cameron model with parameters according to Iljinov systematics [4]. The compound fractions of the total cross sections have been observed to decrease with increasing bombarding energy for all reactions. The contribution of deuteron breakup has negligible effect in the backward angles.
To my parents, Isaias and Fe
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1 INTRODUCTION

Nuclear reaction cross sections play a significant role in basic and applied nuclear physics. For example, the knowledge of cross sections for different types of nuclear reactions is essential for explaining astrophysical phenomena and for the development of future generation nuclear reactors. Due to the enormous amount of cross section data needed in applications, such as mentioned above, theoretical calculations are imperative to supplement the limitation that experimental cross section databases offer. For the past several decades, nuclear reaction codes have been developed to allow a comprehensive modeling of different reactions and predict reliable nuclear information. However, to accurately provide these information, accurate input parameters are required in the calculations.

The compound nuclear reaction is one of the major types of reaction mechanisms. It is often known to be the dominant mechanism at low bombarding energies \( \left( < 5 \frac{MeV}{A} \right) \) [6]. The concept of compound nuclear reaction arises from the compound nucleus model introduced by Bohr [7] in 1936. In this model, the interaction between a projectile and a target nucleus results in the formation of a compound nucleus, a relatively long-lived intermediate state, that subsequently decays by particle or \( \gamma \) emissions. The formation and decay stages are considered to be independent of each other. More details about the compound nucleus model will be discussed in Chapter 2. For reactions involving medium and heavy mass nuclei, it is necessary to treat the compound reaction statistically due to the large density of final states involved. In such cases, the compound cross section is calculated through the statistical Hauser-Feshbach theory, developed independently by Hauser and Feshbach [8], and Wolfenstein [9]. The key ingredients in this theory are the transmission coefficients and the nuclear level densities.
1.1 Nuclear level density

The nuclear level density is defined as the number of nuclear levels per unit of excitation energy. It is a fundamental property inherent to each nucleus. The earliest theoretical work on the nuclear level density was done by Bethe [10] in 1936 that led to the formulation of the Fermi gas model for level densities. Bethe’s work is based on the assumptions that the nucleons within the nucleus are non-interacting fermions (Fermi gas system) and the spacing between single particle states is equidistant and independent of the energy. Unfortunately, these assumptions fail to incorporate important nuclear features such as: (1) fluctuations of the single-particle density due to shell effects; (2) increase of the nuclear binding energy due to short-range residual interactions; and (3) collective effects (e.g. rotation and vibration), which resulted in numerous refinements of the Fermi gas model [2, 11–13]. Implementing these features in the form of semi-empirical models into the analytical level density formula has been found to improve the accuracy of the calculated cross sections, especially when employed with parameters from global systematics.

Complex theoretical approaches for calculating the level density at a microscopic level, such as shell model Monte Carlo (SMMC) calculations [14], finite temperature Hartree-Fock-Bogoliubov theory [3], and relativistic mean field [15] have made some progress in the past decades. Logically, the microscopic approach should provide a better description of the level density compared to phenomenological formulas. This is true especially at regions where no experimental information is available – such as at very high excitation energies or for nuclei far from the valley of stability.

There are several experimental methods to determine the nuclear level density. The application of a particular method depends on the excitation energy region being studied. Here, a short description of each method is discussed.
• **Counting of discrete levels.** The level density information is extracted through direct counting of the observed levels per unit of excitation energy. This method is applicable only at low excitation energies near the ground state, where nuclear levels are well separated and the average nuclear level width $\Gamma$ is much smaller than the average level spacing $D$ ($\Gamma \ll D$). The counting of resolved levels does not apply at high excitation energies due to the experimental resolution. As excitation energy increases, the level spacing rapidly decreases resulting in missed levels. Complete level schemes (including spin and parity) of most stable nuclei are well documented in the literature [16]. Levels for light nuclei are known to be complete up to 5-6 MeV. This value decreases down to 1-2 MeV for heavier mass nuclei where the level density is higher.

• **Counting of neutron resonances.** Neutron resonances are observed near the neutron binding energy. At this excitation energy region, levels can still be resolved through neutron capture reactions only when low angular momentum (such as s-wave) is transferred. The neutron energy resolution is required to be better than the average s-wave resonance spacing $D_0$ to resolve the individual levels. This is provided by the time-of-flight technique. Extensive neutron resonance compilations are available for s-wave resonances in Ref. [17]. The level density can be directly determined from the information of the average density of s-wave resonances, spin distribution, and parity ratio. Charged-particle capture resonances may also yield level density information in the same manner as neutron resonances, but these are limited to low mass nuclei because of the Coulomb barrier.

The uncertainties associated with this technique come from missing resonances, wrong spin assignments, and the ambiguity of the assumed spin and parity distributions. For the past decades, various statistical techniques [18–20] were
developed to estimate the fraction of missing resonances, and these were successfully carried out to correct old measurements [17, 21]. The spin and parity distributions, on the other hand, are far from being resolved. The latter might not be so important for heavy nuclei since the assumption that positive and negative parities are equally distributed might be true for excitation energies above 7-10 MeV, but it may likely be a factor for nuclei with $A < 70$ [22]. Currently, the largest uncertainty associated with the level density from neutron resonance data is related to the uncertainties of the spin cut-off parameter in the spin distribution function. This information lacks the experimental data to support the available models.

- **Analysis of evaporation spectra.** In the statistical regime, the spectra of outgoing particles from compound nuclear reactions are proportional to the level densities of the residual nuclei. By applying the statistical Hauser-Feshbach formalism [8] in analyzing the measured compound cross section, one can extract the experimental level density [23]. A more elaborate description of this method will be presented in Chapter 3.

- **Analysis of level widths through Ericson fluctuations.** This approach is performed at substantially higher excitation energies (such that $\Gamma \gg D$) where nuclear levels strongly overlap. Excitation functions measured with high energy resolution demonstrate fluctuations as a function of energy. Analysis of these fluctuations provides the average level width in the continuum $\langle \Gamma \rangle$, also known as the coherence width. This information can be further used to determine the level density [24, 25]. Since this method requires very high experimental resolution to observe the tiny fluctuations in the excitation function, it has only been applied to a few nuclei.

The majority of level density information is derived from the densities of discrete levels or neutron resonances. Recent parameter systematics for phenomenological level
density models are obtained by model interpolation between these two regions. However, the uncertainties associated with the neutron resonance data due to their strong dependence on the spin cut-off parameter and the parity ratio might present a concern. Exploring other methods that do not use the neutron resonance data should minimize these uncertainties and therefore, improve the nuclear level density parametrization. The evaporation method is the best candidate among the other techniques since it covers a wider range of excitation energy and it is less sensitive to spin and parity distributions. Despite these advantages, a relatively small number of nuclei have been explored by this technique.

1.2 Level density applications

A specific example of nuclear level density application is related to astrophysical nuclear reactions. The kinetic energy distribution of nuclei in an astrophysical plasma follows the Maxwell-Boltzmann function at a given temperature. As a result, a large fraction of nuclear reactions proceeds through the compound reaction mechanism [26]. Some of the significant reactions include the neutron capture for s-process and r-process and the proton capture for rp-process. Capture cross sections are crucial for understanding the stellar evolution and the nucleosynthesis of heavy elements. In particular, nuclei near the iron region play a special role in nucleosynthesis since they are the heaviest nuclei created by fusion inside the stars, and they act as seeds for the synthesis of much heavier elements. These reactions can be calculated through the Hauser-Feshbach model. Predictions are very important in these cases since cross sections of interest are either experimentally scarce or unavailable. Moreover, explosive nuclear burning in a stellar environment produces unstable nuclei, which can be targets for subsequent reactions. As of today, the study of nuclei far from the valley of stability is still a major experimental challenge, although some of the exotic light nuclei can now be studied in radioactive ion beam facilities.
Another relevant application of the level density is related to fusion and fission reactor designs. The information about the energy-angle correlations of the evaporated particles and the double differential cross sections are important to model the efficiency and safety of the structural designs by predicting neutron production, activation, heating, shielding requirements, and material damage. A more direct impact of this study results from the iron and copper being the primary constituents of the reactor walls [27, 28]. Moreover, fusion reactors involve d or d-T fuels. Accurate data for induced radioactivity assessment are of great interest, especially for experimental and evaluated data of deuteron-induced reactions, which are less comprehensive compared with nucleon-induced reactions.

1.3 Objectives of the study

The goal of this study is to determine the level densities of selected nuclei around the mass region $52 \leq A \leq 66$ with a technique based on compound nuclear reactions. The information obtained will be utilized to deduce the best level density parameters and test them with existing global parametrizations. The extraction of level densities will be made by analyzing the evaporation spectra from selected nuclear reactions.

1. The reactions of $d + ^{63}\text{Cu}$ and $d + ^{65}\text{Cu}$ are used to study the residual nuclei of $^{64}\text{Zn}$ and $^{66}\text{Zn}$ by measuring the outgoing neutron spectra.

2. The reactions of $d + ^{54}\text{Fe}$, $d + ^{56}\text{Fe}$, and $d + ^{58}\text{Fe}$ are used to study the residual nuclei of $^{55}\text{Fe}$, $^{57}\text{Fe}$, $^{55}\text{Co}$, $^{57}\text{Co}$, $^{59}\text{Co}$, $^{52}\text{Mn}$, $^{54}\text{Mn}$, and $^{56}\text{Mn}$ by measuring the outgoing neutron, proton, and $\alpha$-particle spectra.

In addition to studying the level density of residual nuclei, our results will also evaluate the practicality of using deuterons to induce compound nuclear reactions. As the lightest composite particle with a loosely bound state, deuterons easily break up due to an encounter of nuclear field or Coulomb field. This would result in several other processes apart from
those when only nucleon-induced reactions take place. Although a study [29] has shown that neutron spectra measured at backward angles with 7 MeV deuterons contain negligible non-compound contribution, this will be further verified by our experiments. Different outgoing channels will also be explored in the process.

1.4 Thesis outline

The dissertation is organized as follows. In Chapter 2, the theoretical background that reviews the fundamental concepts involved in this study is presented. It starts with the compound reaction mechanism and proceeds with the Hauser-Feshbach theory as its statistical interpretation. This is followed by the discussion of level density models generally employed in cross section calculations. Chapter 3 describes in detail the experimental setup and data analysis. Chapters 4 and 5 are devoted to the presentation and discussion of the results. Finally, the summary and conclusions are given in Chapter 6.
2 THEORETICAL BACKGROUND

2.1 Nuclear reaction mechanisms

Nuclear reaction mechanisms at low energies are classified mainly into three categories: (a) Direct, (b) Compound, and (c) Pre-equilibrium. The difference between them is based on the interaction time of the projectile and the target before emission of particles takes place (see depiction in Figure 2.2). For a given reaction, all reaction mechanisms may contribute to the final cross section. Therefore, to study the contribution of each specific process, one needs to consider the interaction energy, the shape of the emission spectrum, and the angular distribution. The following sections discuss these properties for each reaction process.

![Figure 2.1: Schematic of direct, compound, and pre-equilibrium nuclear processes. Figure is reproduced from Ref. [5].](image)
2.1.1 Direct reactions

Direct reactions are characterized by extremely fast reaction times with the time scale of $\sim 10^{-22}$ s. This is approximately equal to the time interval the projectile traverses the boundaries of the target nucleus. The contribution of direct reactions becomes more significant as the energy of the projectile increases. Since the associated De Broglie wavelength of the incident particle is comparable to the size of the target nucleus, localized interactions may take place where only a few surface nucleons are involved. Hence a strong correlation between the initial and final states of the reaction is expected. The direction of the outgoing particles favors the direction of the projectile resulting in a strongly forward-peaked angular distribution. Moreover, the angular distribution exhibits a diffraction pattern that enables the identification of spin and parity for the residual nuclear states.

For deuteron projectile, its low binding energy easily leads to breakup reaction where the deuteron fragments into proton and neutron due to interaction with the Coulomb or nuclear fields of the target. Since the process proceeds very fast, it is generally considered a direct reaction. The angular distributions of neutron and proton fragments are observed to be forward peaked, and their energy distribution is a continuous line with peak centered at approximately half of the incident deuteron energy.

Among the models used to describe direct reactions are the Distorted Wave Born Approximation (DWBA) and Coupled Reaction Channel (CRC) [30–32].

2.1.2 Compound reactions

In the compound nuclear reactions [7], the incident projectile is captured to form a long-lived intermediate state, called compound nucleus. The energy acquired by the compound nucleus is released in the form of particle evaporation or $\gamma$-radiation emission. Such a reaction can be represented as
where \( a, A, C^*, b, \) and \( B \) are the projectile, the target nucleus, the compound nucleus, the ejectile, and the residual nucleus, respectively. According to the compound nucleus model, the compound reaction starts with the fusion of the projectile and the target nucleus. The entrance channel energy is shared among the constituent nucleons of the compound system. The relatively long reaction time associated with compound nuclear reactions makes it an extremely slow process, that is several orders of magnitude longer than direct reactions. Figure 2.2 shows a simple diagram for the formation of a compound nucleus and its decay.

An important consequence of the compound nucleus model is the "loss of memory" on how the compound nucleus is formed. One can then assume that there exists no connection between the entrance and exit channels of the reaction except for the conserved quantities such as energy, angular momentum, and parity. This is the so-called Bohr independence hypothesis [7]. Several observable properties are used to experimentally validate the compound nuclear reactions:

- The decay probability of a compound nucleus (for a specific excitation energy) is independent of the entrance channel (see Refs. [33, 34]).

- The energy spectrum of emitted particles approximately follows a Maxwellian distribution [35].

- The angular distribution of evaporated particles from a compound nucleus decay is symmetric about 90° in the center-of-mass frame [36].

Weiskopf and Ewing [35] were the first to develop the statistical model for describing the spectra of evaporated particles. The derivation of this model is based on the Bohr independence hypothesis where the cross section of a compound reaction is expressed
Figure 2.2: Excitation of a compound nucleus and its decay, where $E_c$ is the excitation energy of the compound nucleus $C^*$ and $E_b$ is the energy of the evaporated particle for the exit channel $(b + B)$.

independently as a product of the compound nucleus formation cross section $\sigma_{CN}$ and its decay probability $P_b$

$$\sigma_{ab} = \sigma_{CN} P_b. \quad (2.2)$$

The Weiskopf-Ewing theory can be used to predict angle-integrated cross sections. This feature is rather limited especially when used in detailed analysis of the experimental cross
sections. The main disadvantage of this theory is that it does not take into consideration the conservation of angular momentum.

2.1.2.1 Hauser-Feshbach theory

An improvement of the Weiskkopf-Ewing model is the Hauser-Feshbach theory [8]. This theory explicitly considers the conservation of angular momentum and parity in the formation of the compound nucleus and its decay. Suppose a binary compound reaction given in expression (2.1) have angular momenta $i_a, I_a, J, i_b, I_b$ for the projectile, target, compound nucleus, ejectile, residual nucleus, respectively, the total cross section for the decay of an excited compound nucleus to a single final state is given as follows

$$
\sigma_{ab} = \sum_{J, \pi} \sigma_{ab}^{J, \pi},
$$

(2.3)

where $J$ and $\pi$ are the angular momenta and parities of the states populated in the compound nucleus.

From the Bohr independence hypothesis, equation (2.3) can be rewritten as a product of the formation cross section $\sigma_{CN}$ and decay probability for each $J, \pi$ state

$$
\sigma_{ab}(\epsilon_a, U_b, I_b, \pi_b) = \sum_{J, \pi} \sigma_{CN}^{J, \pi}(\epsilon_a) \frac{\Gamma_b(J, \pi, \epsilon_b, I_b, \pi_b)}{\Gamma(J, \pi)},
$$

(2.4)

where $\Gamma_b$ is the decay width for a particular channel $b$ with quantum numbers $J, \pi$ by emission of particle $b$ with energy $\epsilon_b$ and residual level with excitation energy $U_b$ and quantum number $I_b$ and $\pi_b$. $\Gamma$, on the other hand, is the total decay width, which is the sum over all decay widths ($\Gamma = \Gamma_a + \Gamma_b + \Gamma_c + ...$). The formation cross section of a compound nucleus with angular momentum $J$ and parity $\pi$ can be calculated in terms of the transmission coefficient [37] as

$$
\sigma_{CN}^{J, \pi}(\epsilon_a) = \frac{\pi}{k^2} \sum_{|S_a|=|l_a-i_a|} \sum_{|S_d|=|J-S_a|} \frac{(2J+1)T_{a}^{i_a}(\epsilon_a) \cdot \delta_{\pi a \pi}^{S_a}}{(2I_a+1)(2I_a+1)},
$$

(2.5)

where the $\delta$ function conserves the parity in the formation cross section.
Similarly, the decay width $\Gamma_b$ can be expressed in terms of the transmission coefficients summed over the contributions from allowed angular momenta by the conservation of $J$ and $\pi$ as

$$\Gamma_b(J, \pi, \epsilon_b, I_b, \pi_b) = \sum_{S_b=|I_b-i_b|}^{|I_b+S_b|} \sum_{l_b=-|J-S_b|}^{|J-S_b|} T_{l_b}^b(\epsilon_b) \cdot \delta_{\pi_b, \pi}^{(1)/b}$$

(2.6)

and $\Gamma$ is expressed as

$$\Gamma(J, \pi) = \sum_b \int_0^{U_{max}} \sum_{I_b, \pi_b} \rho_b(U_b, J_b, \pi_b) \cdot \Gamma_b(J, \pi, \epsilon_b, I_b, \pi_b) dU_b,$$

(2.7)

where the sum of $b$ represents the different exit channels, $U_b$ is the residual nuclear energy related to the particle emission energy as $U_b = E_{CN} - B_b - \epsilon_b$, $\rho_b(U_b, I_b, \pi_b)$ is the spin and parity dependent level density for the residual nucleus reached by emission of particle $b$. More details about the transmission coefficients can be found in the following section. The level density is typically assumed to be a product of the total level density $\rho_{tot}(U)$ and a normalized spin distribution function.

Substituting equations (2.5), (2.6), and (2.7) back to equation (2.4) provides the Hauser-Feshbach expression for the angle-integrated cross section

$$\sigma_{ab}(\epsilon_a, U_b, I_b, \pi_b) = \frac{\pi}{k_a^2} \sum_{J, \pi} \frac{2J + 1}{(2I_a + 1)(2I_a + 1)} \frac{\sum_{S_{a,l_a}} T_{a}^{l_a}(\epsilon_a) \sum_{S_{b,l_b}} T_{b}^{l_b}(\epsilon_b) \rho(U, J, \pi)}{\sum_b \sum_{l_b, \pi_b} \Gamma}.$$ 

(2.8)

Extending equation (2.8) above for decay to final states corresponding to particle emission energies between $\epsilon_b$ and $\epsilon_b + d\epsilon_b$, equation (2.8) becomes

$$\frac{d\sigma_{ab}}{d\epsilon_b} = \frac{\pi}{k_a^2} \sum_{J, \pi} \frac{2J + 1}{(2I_a + 1)(2I_a + 1)} \frac{\sum_{S_{a,l_a}} T_{a}^{l_a}(\epsilon_a) \sum_{S_{b,l_b}} T_{b}^{l_b}(\epsilon_b) \rho(U, J, \pi)}{\sum_b \int_0^{U_{max}} \sum_{S_{b,l_b}} \Gamma_b(U_b) dU_b}.$$ 

(2.9)

2.1.3 Pre-equilibrium reactions

The time scale and number of intranuclear collisions for pre-equilibrium reactions are intermediate between direct and compound reactions. Particle emission in this process takes place after the first step of the projectile and target interaction occurs, but long before the composite system reaches statistical equilibrium. Hence it is also termed as
a pre-compound process. The pre-equilibrium contribution in the outgoing particle spectra appears as a structureless high energy enhancement, which is very distinct compared to both the Maxwellian-type distribution for a compound reaction and the strong population of selected low-lying states in the case of direct reactions. Moreover, the angular distribution of the pre-equilibrium particles is forward peaked. This is an indication that the memory of how the composite system was formed is still preserved.

Pre-equilibrium model calculations are based on time-dependent master equations that describe the probability of transitions to different particle-hole state (exciton) configurations. The exciton number is the primary parameter in these calculations. Integrating the master equations over time provides the energy-averaged emission spectrum. More details are found in Refs. [30, 32, 38]

### 2.2 Phenomenological nuclear level density

The nuclear level density represents the basic statistical information of the nuclear levels as a function of excitation energy. Since the experimental information on the nuclear level density is restricted, phenomenological level density models are generally employed to predict them over required excitation energy region. The models are parametrized based on the experimental information mainly by the low-lying discrete levels and the density of neutron resonances.

The nuclear level density for a particular \((J, \pi)\) is generally expressed as a product of the parity distribution \(P(E, \pi)\), the total level density \(\rho(E)\), and the spin distribution \(S(E, J)\)

\[
\rho(E, J, \pi) = P(E, \pi)S(E, J)\rho(E). \tag{2.10}
\]

The parity distribution provides the ratio between the numbers of negative \(n_-\) and positive \(n_+\) parity states as a function of excitation energy \(E\). This is expressed as

\[
P(E, \pi_\pm) = \frac{n_{\pi}(E)}{n_+(E) + n_-(E)}. \tag{2.11}
\]
It is often assumed that the parity distribution is independent of energy and is equally distributed with \( P(\pi) = 1/2 \). Ericson [39] showed that a few number of single particle levels with different parities would immediately approach equal distribution even for a small excitation energy. However, several theoretical studies [40–42] have shown the dominance of one type of parity at low energies and that the parity ratio only approaches to an equal distribution near the particle separation energy. Experimental measurements [43] have also confirmed such dependence.

Al-Quraishi et al. [22] proposed the following empirical formula for the parity distribution

\[
P(E) = \frac{1}{2} \left( 1 \pm \frac{1}{1 + \exp \left[ k (E - \delta) \right]} \right)
\]  

(2.12)

with \( k = 3 \text{ MeV}^{-1} \) and \( \delta = a + b/A^c \) where \( A \) is the nuclear mass number, and \( a, b, \) and \( c \) are constants derived from fitting discrete levels at low excitation energies for a set of nuclei. The + and - signs are used for \( P(0) \sim 1 \) and \( P(0) \sim 0 \), respectively.

The spin distribution function \( S(E, J) \) [39, 44, 45] in equation (2.10) is given by

\[
S(E, J) = \frac{2J + 1}{2\sigma^2} \exp \left[ -\left( J + \frac{1}{2} \right)^2 \right]
\]  

(2.13)

where \( \sigma \) is the spin cut-off parameter and the prefactor \( 2J + 1 \) is the statistical weight. The expression for the spin cut-off parameter \( \sigma \) is energy dependent, which makes the spin distribution function also energy dependent. More discussion on the spin cut-off parameter follows in the next section.

Finally, the level density \( \rho(E) \) takes several functional forms. Three commonly used level density expressions are: (1) the Fermi gas model, (2) the constant temperature model, and (3) the Gilbert-Cameron model, which is a combination of the first two.
2.2.1 Fermi gas model

The level density expression based on the Fermi gas approximation was derived by Bethe [10]. It assumes a non-interacting particle model where the nucleons move independently of each other under the influence of an overall central potential. The average potential replaces the complex nucleon-nucleon interaction in the system. Since nucleons are fermions, they are also subjected to the Pauli exclusion principle that prevents two identical fermions to occupy the same quantum state. Further assumption is that the single particle states are equally spaced. For a two-fermion system, the total Fermi gas level density is given by

\[
\rho(E) = \frac{1}{12 \sqrt{2} \sigma} \frac{\exp \left[ 2 \sqrt{aE} \right]}{a^{1/4} E^{5/4}}.
\]

This is obtained by summing the level density \( \rho(E, J, \pi) \) over spins and parities. Here, the level density parameter \( a \) and the spin cut-off parameter \( \sigma \) are model parameters.

Despite the crude assumptions, the Fermi gas model has been successful in describing the nuclear level density. However, it was not able to describe a large set of experimental data with a single set of parameters partly because it does not incorporate important nuclear properties. As an example, the assumption of non-interacting nucleons ignores the pairing effect and any collective enhancements due to nucleon-nucleon interactions. Moreover, the energy independent single particle states neglect the nuclear shell structure. These shortcomings were addressed by introducing one or several free parameters of the Fermi gas expression. The next subsections discuss the parameters of the Fermi gas level density model and how they were modified to account for its limitations.
2.2.1.1 Fermi gas level density parameters

- **Pairing correlations** \( \Delta \). Residual interactions between nucleons alter the nuclear level density at low excitation energies. An example is the pairing interaction that causes the staggering of the binding energies for even-even, even-odd, and odd-odd proton-neutron numbers. An explanation of the even-odd staggering is that an extra amount of energy is required to break a nucleon pair in the ground state. That extra energy is the pairing energy resulting in an energy shift \( \Delta \). The pairing energy shift \( \Delta \) may be obtained from the differences between separation energies or empirical masses of neighboring isotopes and isotones. It has been found that this parameter follows a general trend as \( \Delta \approx \frac{12}{\sqrt{A}} \) [46].

In the conventional shifted Fermi gas model, the odd-even effect is taken into account by subtracting the pairing energy shift from the excitation energy \( E \), such that \( U = E - \Delta \) [47, 48]. \( U \) is also known as the ”effective” excitation energy. The shifted Fermi gas model only has the level density parameter \( a \) as the free parameter even when \( U \) is used in equation (2.14) instead of \( E \). Adopting the same formula, some authors [44, 49, 50] proposed to treat both the level density parameter \( a \) and the pairing energy shift \( \Delta \) as free parameters. This formula provided reasonable fits with the experimental level density [11]. This model is called the back-shifted Fermi gas model since the experimental values of \( \Delta \) were found to shift towards more negative values compared to the ones used in the conventional shifted Fermi gas model.

- **Level density parameter** \( a \). The level density parameter \( a \) of equation (2.14) is related to the density of single-particle levels at the Fermi energy and reflects the properties of the single-particle potential. These relations are given by the equations of state expressed as [38]:

\[
U = aT^2, \quad S = 2aT, \quad \text{and} \quad \sigma^2 = IT. \tag{2.15}
\]
Here, $U$ is the effective excitation energy, $T$ is the nuclear temperature, $S$ is the entropy and $I$ is the nuclear moment of inertia. Assuming equidistant spacing model for a two-fermion gas of protons $\pi$ and neutrons $\nu$, the analytical expression for the level density parameter $a$ is related to the state densities $g_{\pi(\nu)}$ as

$$a = \frac{\pi^2}{6} (g_{\pi}(U_F) + g_{\nu}(U_F))$$

(2.16)

Here, $g_{\pi(\nu)}(U_F)$ is the proton (neutron) state density at the Fermi energy.

Extensive analysis of the experimental level density of neutron resonances shows that the general trend of $a$ varies linearly with the mass number $A$ \cite{51, 52} as

$$a = \text{const} \cdot A \ (\text{MeV}^{-1})$$

(2.17)

In some cases, an additional parameter is included to describe the experimental data better. Such expression is given by $a = cA + d$, where $c$ and $d$ are parameters adjusted to fit the experimental data at different mass ranges \cite{53}.

With the accumulation of neutron resonance data, it became apparent that $a$ significantly deviates from expression (2.17) for nuclei close to the magic shell regions. It was also observed that at higher excitation energies, the nuclear shell effects of the level density disappear \cite{54}. To account for both effects, Ignatyuk \cite{13} suggested to employ an energy-dependent level density parameter given by the expression

$$a(U) = \tilde{a} \left\{ 1 + \delta W \left[ 1 - \exp(-\gamma U) \right] \right\} U.$$

(2.18)

Here, $\tilde{a}$ is the asymptotic level density parameter, $\delta W$ is the shell correction energy, and $\gamma$ is the damping parameter. As seen in equation (2.18), the shell correction $\delta W$ dampens out at high excitation energies ($U \to \infty$) due to the term containing the
damping parameter $\gamma$. The shell correction $\delta W$ is determined as the difference between the experimental nuclear mass $M_{\text{exp}}$ and the mass from the liquid drop model calculation $M_{\text{LD}}$ as

$$\delta W = M_{\text{exp}} - M_{\text{LD}}. \quad (2.19)$$

$M_{\text{exp}}$ are taken from the Audi-Wapstra mass compilation [55] while the $M_{\text{LD}}$ is calculated from the Myers-Swiatecki [56] mass formula

$$M_{\text{LD}} = M_n N + M_H Z + E_{\text{vol}} + E_{\text{sur}} + E_{\text{coul}} + E_{\delta}. \quad (2.20)$$

Here, $M_n$ is the mass of the neutron, $M_H$ is the mass of the proton, $E_{\text{vol}}$ is the volume component, $E_{\text{coul}}$ is the coulomb component, $E_{\text{sur}}$ is the surface component, and $E_{\delta}$ is the pairing component.

The asymptotic level density parameter and damping parameter generally employs the systematics according to Ignatyuk [13] or Iljinov [4] as

$$\tilde{a} = 0.154A + 6.3 \times 10^{-5}A^2 \quad \text{and} \quad \gamma = -0.054, \quad (2.21)$$

$$\tilde{a} = 0.114A + 9.80 \times 10^{-2}A^{2/3} \quad \text{and} \quad \gamma = -0.051, \quad (2.22)$$

respectively.

Another formula to calculate $a$ was suggested by Al-Quraishi et al. [57]. They proposed two forms of $a$ which are thought to provide a better level density description for nuclei far from the valley of stability by considering the isospin dependence. These are

$$a = \alpha A/ \exp[\beta(N - Z)^2] \quad (2.23)$$

and

$$a = \alpha A/ \exp[\gamma(Z - Z_0)^2]. \quad (2.24)$$
The values of the parameters are given by $\alpha = 0.1062$ and $\beta = 0.00051$ for equation (2.23) while $\alpha = 0.1068$, $\gamma = 0.0389$, and

$$Z_0 = \frac{0.5042A}{1 + 0.0073A^{2/3}}$$

(2.25)

for equation (2.24). $Z_0$ is the approximate atomic number for a $\beta$-stable nucleus of mass number $A$. Equation (2.25) was derived from the semi-empirical formula for the determination of the masses of nuclei without the pairing term.

Equation (2.23) accounts for the effect of isospin conservation as one moves away from the valley of stability. This points out that the level density decreases as the value of $|N - Z|$ increases. From a microscopical point of view, equation (2.24) considers the role of single particle states, which become unstable when drip lines are reached. Both formula take their typical form of $a = \text{const} \cdot A$ for nuclei near the valley of stability.

- **Spin cut-off parameter.** The spin cut-off parameter $\sigma$ represents the width of the spin distribution of levels at a given excitation energy. This parameter can be obtained in two different ways. According to Ericson [39], the spin cut-off parameter is given by

$$\sigma^2 = gT\langle m^2 \rangle$$

(2.26)

where $g$ is the density of single particle states (see equation (2.16)), $T$ is the temperature such that $T = \sqrt{U/a}$, and $\langle m^2 \rangle$ is the average of the square of the spin projection for single particle states near the Fermi energy. According to Ref. [52], $\langle m^2 \rangle = 0.24A^{2/3}$. This gives the spin cut-off parameter as

$$\sigma^2 = 0.1461A^{2/3} \sqrt{aU}.$$  

(2.27)

The second method to obtain the spin cut-off parameter is by assuming that the nucleus is a rigid sphere. The excitation modes of the nucleus is then separated
into intrinsic excitation due to individual nucleons and rotational excitation. Such consideration leads to the spin cut-off parameter

\[ \sigma^2 = IT, \]

where \( I = \frac{1}{h^2} \frac{2MA^5r_0^2}{5} \), \( M \) is the nucleon mass, and \( r_0 \) is the reduced nuclear radius. Taking the reduced radius to be \( r_0 = 1.25 \text{ fm} \) and expression (2.15) give the expression for the spin cut-off parameter as

\[ \sigma^2 = 0.0145A^{5/3} \sqrt{\frac{U}{a}}. \quad (2.29) \]

### 2.2.2 Constant temperature model

It has been noticed [2] that the cumulative number of levels at low excitation energy is well described by a two-parameter equation given by

\[ N(E) = \exp \left( \frac{E - E_0}{T} \right) \quad (2.30) \]

where \( E_0 \) and \( T \) are free parameters. From equation (2.30), the constant temperature model (CTM) of the level density [2] is

\[ \rho(E) = \frac{dN}{dE} = \frac{1}{T} \exp \left( \frac{E - E_0}{T} \right). \quad (2.31) \]

\( T \) is known as the constant temperature parameter that does not vary with the excitation energy. This is in contrast to the Fermi gas model where \( T = \sqrt{(E - \delta)/a} \). Although the constant temperature level density was proposed for low excitation energies where it agrees well with experimental data, recent papers suggest that the expression (2.31) can be valid even up to 20 MeV [58].

### 2.2.3 Gilbert-Cameron model

The Gilbert-Cameron model (GCM) uses both the constant temperature model and the Fermi gas model to describe level densities in a wide range of excitation energies [2].
A matching energy \( E_M \) is defined such that the level density below it is described by the constant temperature while the level density above \( E_M \) is represented by the Fermi gas model

\[
\rho(E) = \rho_{CTM}(E) \text{ for } E \leq E_M, \quad (2.32)
\]

\[
\rho(E) = \rho_{FGM}(E) \text{ for } E \geq E_M. \quad (2.33)
\]

The \( E_M \) is determined such that both level densities and their derivatives are continuous at that point:

\[
\rho_T(E_M) = \rho_{FG}(E_M) \quad (2.34)
\]

and

\[
\frac{d\rho_T}{dE}(E_M) = \frac{d\rho_{FG}}{dE}(E_M). \quad (2.35)
\]

Hence, five parameters (matching point energy \( E_M \), nuclear temperature \( T \), energy shift \( E_0 \), level density parameter \( a \), and pairing energy \( \delta \)) are needed for the Gilbert-Cameron model.

### 2.3 Microscopic models

In addition to phenomenological level density models, microscopic models are also developed. Microscopic models are appropriate for application in experimentally unknown regions because these models are based on first principles and do not contain phenomenological parameters. Microscopic models of nuclear level density have been developed in the past decades but they suffer from a lack of accuracy to reproduce experimental data compared to that of highly parametrized analytical expressions. Microscopic calculations are based on the combinatorial approach, statistical model, spectral distribution, Monte-Carlo, and quantum Monte-Carlo approaches [3, 14, 59, 60].

In this work, we used the microscopic combinatorial approach of Ref. [3] in our calculations. This method consists of using the single-particle level scheme obtained from the axially symmetric Hartree-Fock-Bogoliubov (HFB) method based on the BSk14
effective interaction to construct incoherent particle-hole (ph) state densities as a function of the excitation energy, the spin projection, and the parity. BSk14 has been tuned with nuclear binding energies and the spectral gaps characterized by the pairing interaction to reproduce the experimental odd-even mass differences. The pairing interaction is of primary importance for a reliable prediction of the NLD. Once the incoherent particle-hole state densities are determined, the collective effects are accounted for to deduce the total level density. Collective rotational effects are included by building up rotational bands consistently. Phenomenological vibrational enhancement is incorporated by treating the phonon excitations explicitly through the boson partition function. The resulting NLD were found to reproduce the available experimental data both in the low energy region and in the region of neutron resonances.

2.4 Transmission coefficients

Another important ingredient for the Hauser-Feshbach calculations is the transmission coefficients \[37\]. As shown in equations (2.9) and (2.8), transmission coefficients are needed to calculate both the formation of the compound nucleus and its decay to various channels. Individual particle transmission coefficients are calculated by solving the Schrodinger equation with the optical potential that represents the particle-nucleus interaction. More details on optical potential are discussed in the following section. The particle transmission probability is typically expressed by the sum over all allowed partial waves as

\[
T_j^\mu(E, J, \pi, E_i^\mu, J_i^\mu, \pi_i^\mu) = \sum_l \sum_s T_{js}(E_i^\mu) \tag{2.36}
\]

where \(|J - s| \geq l \leq |J + s|\) and \(|J_i^\mu - J_j| \geq s \leq |J_i^\mu + J_i^\mu|\). The \(\gamma\)-transmission coefficients, on the other hand, are derived from the \(\gamma\)-strength functions given by \[61\]

\[
T_{XL}(E_\gamma) = 2\pi E_\gamma^{2L+1} f_{XL}(E_\gamma). \tag{2.37}
\]
The $f_{XL}$ is the strength function for the transition of multipolarity $XL$, which is directly related to the radiative width $\langle \Gamma_{\gamma i} \rangle$, the level spacing $D_i$ of initial states, and the transition energy $E_\gamma$ through

$$f_{XL} = \frac{\langle \Gamma_{\gamma i} \rangle}{E_\gamma^{2L+1}D_i}.$$  

(2.38)

In general, all the multipolarities contribute to the $\gamma$-ray emission probability but the largest portion of the strength is due to E1 or M1.

### 2.4.1 Optical model potential

The optical model [30] simplifies the description of the projectile and target interaction by replacing the complex nucleon-nucleon forces by the interaction of the projectile and a potential well. The expression of the optical model potential is hence based on the essential features of the nucleon-nucleon interaction (e.g. a short range interaction) and the general properties of nuclear matter. The optical potential affects the elastic scattering of the particles through its real component. It is known that as the energy increases, non-elastic reaction channels open up, reducing the flux of the elastic channel. To describe the absorption of the incoming flux, one can invoke a complex potential with a negative imaginary part in analogy to the absorption of light as it passes through an absorbing medium. Hence, a potential consisting of both real and imaginary parts is capable of calculating both the flux of the scattered incident particles and the flux absorbed, which leads to non-elastic reactions. The parameters of the potential are adjusted from global fits of the reaction and elastic scattering cross sections.

The optimized analytical expression of the optical potential is composed of three fundamental terms. The first one is the nuclear potential which is further subdivided into real and imaginary components to account for the elastic and non-elastic channels. The second term is the Coulomb potential which contributes if the incident particle is charged. Finally, the spin-orbit term can be added to the optical potential to enable the calculation
of the polarization of the scattered particles. A general radial-dependent OMP expression for the nucleon-nucleus interaction is given by

\[ V(r) = -V_V f(r) - iW_V f(r) - i4a_D W_D \frac{d}{dr} f(r) + \left( \frac{\hbar}{m_D c} \right)^2 \]

\[ \frac{V_{SO}}{r} \frac{d}{dr} f_{SO}(r) \cdot l \cdot \sigma + i \left( \frac{\hbar}{m_{SO}} \right)^2 \frac{W_{SO}}{r} \frac{d}{dr} f_{SO}(r) \cdot l \cdot \sigma + V_C(r) \]

(2.39)

where \( V_V \) and \( W_V \) are the depths of the real and imaginary components of the volume potential, \( W_D \) is the depth of the surface derivative term, \( V_{SO} \) and \( W_{SO} \) are the real and imaginary depths of the spin-orbit potential, and \( V_C \) is the Coulomb term. The quantity \( l \cdot \sigma \) is the scalar product of the intrinsic orbital and spin angular momentum operators. The radial-dependent form factor \( f(r) \) is usually taken to have the Woods-Saxon shape

\[ f(r) = \frac{1}{1 + \exp[(r-R)/a_d]} \]

(2.40)

It is parametrized by the effective nuclear radius \( R \) and the diffuseness parameter \( a_d \). The radius parameter is proportional to the mass number and it is defined as \( R = r_0 A^{1/3} \).

The optical model parameters for nucleon-nucleus reactions are well documented in the literature and these are successfully applied for a wide range of target nuclei and energy.
3 EXPERIMENTAL DETAILS

The experiments were all carried out at the Edwards Accelerator Laboratory of Ohio University. In this chapter, an overview of the accelerator laboratory and description of the experimental setup and data analysis are presented.

3.1 Overview of the Edwards Accelerator Laboratory

The Edwards Accelerator Laboratory has a tandem type electrostatic accelerator that accelerates light ions by using a Van de Graaff machine. It is capable of generating up to a 4.5 MV potential on the terminal. The schematic layout of the facility is shown in Figure 3.1. There are two types of ion sources in the laboratory: (1) the duoplasmatron and (2) the cesium sputtering ion sources. The duoplasmatron is used exclusively for producing a helium beam while the cesium sputtering source is used for all other light particles.

The particles extracted from the ion sources are negatively charged. They are deflected by the injection magnet and guided to the beam line by the Einzel lens and XY steerers. The acceleration of the charged particles is a two-stage process for a tandem type accelerator. The first stage is the acceleration of the injected negative ions towards the positively charged terminal. Inside the terminal, the accelerated ions undergo electron stripping by gas or carbon foil strippers converting them into positive charged states. This leads to the second stage of acceleration where the positively charged ions are repulsed from the terminal. The particles are then guided and focused through a series of slits, steerers, and quadrupole magnets located immediately after the high energy accelerator tube. The 90° analyzing magnet differentiates the particles from the beam according to their energy, mass, and charged state. The magnetic field is set through a nuclear magnetic resonance (NMR) probe. Once the ions pass through the analyzing magnet, the switching magnet redirects them to the preferred beam line. Five beam lines behind the switching magnet are available for use. Two are in the small target room and the other three are in the large target room.
Another beam line, called the swinger, is also available which is located in the vault area. Figure 3.1 shows the schematic layout of the Edwards Accelerator Laboratory.

The Edwards accelerator uses a pulsing and bunching system to convert a direct current (DC) beam to a pulsed beam. Unlike a DC beam, a pulsed beam allows the use of a time-of-flight method, which is necessary for neutron measurements. To pulse and bunch the beam, the DC beam is initially chopped by the deflection system driven by a sinusoidal voltage. The available frequencies for the chopper varies from $f = 5 \text{ MHz}$ to $f = 5/64 \text{ MHz}$. These frequencies correspond to beam periods from 200 to 12800 ns, respectively. The chopped beam is then bunched through the Klystron buncher that can make pulses with width as small as 1 ns for protons and deuterons.

Figure 3.1: Schematic layout of the Edwards Accelerator Laboratory. Three experimental rooms are: the vault area (high energy area), the small target room, and the large target room. Figure courtesy of Catalin Matei.
3.2 Swinger facility and neutron measurements

The swinger facility of the Edwards Accelerator Laboratory is designed for high-precision neutron physics experiments. Excellent energy resolution is achieved through long flight paths with the 30-m time-of-flight tunnel. A massive concrete wall shields the tunnel from background neutrons and γ rays produced in the vault area. Neutrons produced from the beam-target interaction enter the tunnel through a collimation system made up of a 30-cm diameter cylindrical hole in the wall. The neutron detector is positioned inside the tunnel such that it only sees the scattering target as the primary source of neutrons and γ rays.

Another important component of the swinger facility is the beam swinger line. This allows neutron measurements at different angles by rotating the beam line and the target to cover angles ranging from 0° to 162° while keeping the detector assembly fixed inside the tunnel area. The bending of the beam is achieved by the beam-swinger magnet that guides the accelerated ions inside the "C-shaped" beam line. The accelerated ions then exit the swinger beam line perpendicular to its original direction. It is important that the position of the scattering target is aligned with the axis of rotation of the swinger beam line so that the target remains on the same position regardless of the angle and direction of the beam. The target position is aligned using two telescopes oriented perpendicular with each other. Figure 3.2 shows a picture of the swinger beam line.

The targets are mounted on a circular target holder (as shown in Figure 3.4) housed in a cylindrical target chamber. The chamber is attached at the end of the swinger beam line. After alignment of the target position, the chamber is sealed and is subjected to high vacuum. A leak test is then performed to make sure that the system is airtight.

The scattering chamber is electrically isolated from the swinger beam line. Before the beam enters the chamber, it initially passes through two tantalum collimators with diameters 3.18 mm and 1.59 mm. Electrons liberated from any surface inside the chamber
when hit by the beam are restrained from exiting the scattering chamber by the charge suppression ring placed immediately after the collimators. The charges collected from the stopping plate, target, and scattering chamber are integrated and digitized by the beam current integrator (BCI) module. This value is used to determine the accumulated charge and to monitor the beam current.

### 3.2.1 Neutron measurements

Neutrons are uncharged particles, so they essentially interact with the atomic nucleus only. The range of interaction is very short resulting in neutrons to penetrate deeper in matter. A general mechanism involved in the detection of neutrons is to convert them to charged particles by interaction with the detector material. The charged particles are
Figure 3.3: Schematic layout of the swinger beam line, tunnel area, and the detector system.

Figure 3.4: Target holder inside the scattering chamber of the swinger beam line. The target holder is operated remotely.
responsible for exciting the active component of the detector (i.e. scintillator) producing the signal for data processing.

In the experiment, the neutron detector consisted of a NE213 liquid scintillator coupled to a photomultiplier tube (PMT). This yields an excellent neutron-γ pulse shape discrimination (PSD) which is a necessary property for neutron measurements. Due to the scintillator composition (CH\textsubscript{1,212}), several mechanisms are involved in registering the neutrons in the detector. The primary mechanism is the elastic scattering of neutron on a proton. Since the neutron mass is almost equal to the proton mass, the neutron energy is transferred effectively to the recoil proton in a collision. Another mechanism is the elastic scattering of neutron on a carbon atom. Since carbon is heavier than proton, the recoil energy is not enough to create a significant number of scintillation photons. However, neutron scattering with carbon may distort the response function of the scintillator due to energy loss. For neutron energy that exceeds 8.5 MeV, the non-elastic scattering of neutron on carbon should be taken into consideration because they start to dominate at this energy [16]. The two competing reactions are given by $^{12}$C(n,$\alpha$)$^9$Be and $^{12}$C(n,n’)$^3$α. It should be noted that $\alpha$ particle has a lower scintillation efficiency compared to recoil proton. The byproduct particles from the scattering process [e.g. recoil protons and $\alpha$-particles] are the ones that excite the scintillator molecules that subsequently produce the scintillation photons. The amount of photons is proportional to the energy deposited by the neutrons. The photons are then collected and converted to electrical signals by the PMT.

In addition to the detector array located inside the tunnel, other neutron-sensitive detectors are also installed in the vault area. These are stilbene detectors used to monitor the neutron flux primarily coming from the beam-target reactions. These detectors also detect neutrons from other sources, such as neutrons produced from the slits and collimators, and down-scattered neutrons. One of them is attached to the beam swinger line such that its
angle is always fixed relative to the incident beam, while the other one is positioned on the wall.

Three NE213 scintillation detectors, positioned in an array form, constitute the main detector assembly inside the tunnel area. The center of the assembly is aligned with the scattering target through an optical telescope. Figure 3.5 displays the neutron detector assembly and the data acquisition (DAQ) system. Lead bricks can also be set up to further reduce the $\gamma$-ray background.

Figure 3.5: The neutron detector array comprises three NE213 scintillation detectors arranged in triangular form and positioned such that the center of the array is aligned with the center of the scattering target.

### 3.2.1.1 Time-of-flight technique

The conventional method for measuring the neutron energy is the time-of-flight technique (TOF). This utilizes the fast-timing capability of the detector instead of the energy deposited by neutrons. For accelerator experiments, the time-of-flight technique requires the beam to be pulsed for timing information.
The relative flight time is determined as a time difference between start and stop signals. In the experiment, the start signal is produced by the beam pick-off (BPO) device, which generates a signal every time a beam passes through it. On the other hand, the stop signal is generated every time a neutron or γ ray is registered in the neutron detector. The BPO signal undergoes a long delay that it is registered in the DAQ system later than the detector signal. As a result, the channels in the raw TOF spectrum are arranged in a reverse order relative to the absolute time of flight (e.g. the neutrons with short flight times are registered at large channel numbers). Figure 3.6 shows the time-of-flight spectrum taken from the $^{27}$Al(d.n) reaction. To convert the TOF spectrum to the neutron energy spectrum, we used the relativistic time and energy relation given by the expression

$$E_n = E_0 \left[ \frac{1}{\sqrt{1 - \left( \frac{d}{c} \right)^2}} - 1 \right]$$

(3.1)

where $d$ is the flight path, $t$ is the flight time, $E_0$ is the neutron rest mass, and $c$ is the speed of light. The classical expression of the kinetic energy also works well for our energies ($E_n \ll E_0$).

The same time-of-flight concept applies for γ rays. Since the speed of light and the flight path are constant, the corresponding flight time is also constant. Hence, γ rays that originate from the beam-target reactions appear as a large narrow peak in the TOF spectrum (see Figure 3.6). The γ-peak position is utilized to determine the absolute time calibration.

The energy resolution associated with the TOF technique depends on several factors: (1) the uncertainty due to the target thickness $\delta E_t$; (2) the detector time resolution $\delta t_d$; (3) the uncertainty related to the beam time width $\delta t_b$; (4) the thickness of the detector $\delta l$; and (5) the uncertainty of the beam energy $\delta E_b$. The energy uncertainty can then be written in terms of the relative uncertainties as

$$\frac{\delta E}{E} = \sqrt{\left(\frac{2\delta l}{t}\right)^2 + \left(\frac{2\delta t}{t}\right)^2 + \left(\frac{\delta E_t}{E_t}\right)^2 + \left(\frac{\delta E_b}{E_b}\right)^2}.$$  (3.2)
Figure 3.6: Raw TOF spectra taken from one of the $^{27}$Al(d,n) calibration run. The blue line displays the spectrum including both neutrons and $\gamma$ rays. The red line shows the neutron-gated TOF spectrum.

Here, $\delta t^2$ is the sum of $\delta t_d^2$ and $\delta t_b^2$. In many cases, the energy resolution is dominated by the time resolution $\delta t$ or by the uncertainty due to the energy loss in the target. The former becomes significant for shorter flight paths while the latter for thicker targets. Hence, the flight path can be maximized by setting the relative uncertainty in time equal to the uncertainty in energy due to target thickness. It is also important to note that the uncertainty due to the thickness of the scintillator may also dominate for short flight paths and low neutron energies.

### 3.2.1.2 Neutron-$\gamma$ Identification

The NE213 liquid scintillator is sensitive to both $\gamma$ rays and neutrons. Since we are only interested with neutrons, it is essential to separate the events produced by each radiation before any further analysis is carried out. This is achieved by using the pulse shape discrimination (PSD) properties of the detector where events are distinguished based on the shape of the light pulse generated in the scintillator. The shape of the pulse depends
on the population of the two main components of the scintillator – fast and slow. The proportion depends on the type of radiation that deposits the energy. For instance, $\gamma$ rays absorbed by the detector induce larger fraction of the fast component due to the generation of fast electrons, whereas neutrons populate the slow component more due to the recoil protons. Hence, the light pulse produced by neutrons is expected to be longer in time compared to that from $\gamma$ rays. Using this idea, the discrimination between events from neutrons and $\gamma$ rays is implemented based on the ratio of the head charge (tail charge) and the total charge of the integrated electrical signal. The head charge (tail charge) is determined by integrating a time fraction of the pulse in the fast (slow) component [62]. An example of a PSD spectrum is shown in Figure 3.7.

![Neutron-$\gamma$ identification plot from $^{27}$Al(d,n) calibration run.](image-url)
3.2.1.3 Time and detector efficiency calibration

Converting each channel in the raw TOF spectrum to their equivalent absolute flight time is required to obtain the neutron energy spectrum. This is achieved by using the time calibration of the detector system, which is determined by the following steps. The first step is the accumulation of random events using either an artificial radioactive source or natural background radiation into a TOF spectrum. The start signal is prompted by a pulse registered in the neutron detector, while the stop signal comes from the pulse generator. This procedure yields an uncorrelated time spectrum that accounts for differential nonlinearity of the time-to-amplitude converter (TAC). The number of counts in each TAC channel is taken to be proportional to the width of that channel. For the second step, a pulse generator is used to generate the calibration peaks in the TOF spectrum spaced by specific time interval. Figure 3.8 shows both the random distribution and the time calibration peaks from one of the neutron detectors. Finally, in order to obtain the absolute time calibration, we employed the γ-peak position in the TOF spectrum to find the offset. The absolute time of the γ peak can be calculated accurately. Once the absolute TOF spectrum is obtained, we then apply equation (3.1) to further transform it to neutron energy spectrum.

NE213 scintillation detectors are not 100% efficient. In order to determine the absolute neutron yield from the nuclear reactions, the number of neutrons registered by the NE213 scintillation detector should be corrected by its detection efficiency. This correction is necessary for determining the absolute reaction cross sections. The detection efficiency of the NE213 scintillation detectors is obtained from the ratio of the neutron flux registered by the detector and the known "standard flux". In our experiments, the calibration neutron flux with energies from 300 keV to 15 MeV was produced by the $^{27}$Al(d,n) reaction on a stopping target at $E_d = 7.44$ MeV. Outgoing neutrons were detected at $\theta = 120^\circ$. For the "standard flux", the data were taken from Ref. [63] where neutrons from the same reaction
Figure 3.8: Time calibration peaks (top) and random distribution (bottom) from one of the NE213 scintillation detectors. The time calibration peaks are separated by a 50-ns time interval.
were detected using a fission chamber with uranium fission foils. The well-established neutron-induced fission cross sections of uranium yield an accurate neutron flux. From the “standard flux” \( N_{\text{standard}} \) and the one registered by NE213 detector \( N_{\text{NE213}} \), the detection efficiency as a function of neutron energy is calculated as

\[
\epsilon(E) = \frac{N_{\text{NE213}}(E) / (\Delta\Omega \cdot \Delta C \cdot \Delta E)}{N_{\text{standard}}(E)}
\]

(3.3)

where \( \Delta\Omega \) is the solid angle and \( \Delta C \) is the integrated incident beam current.

### 3.2.2 Data Acquisition (DAQ) system

In the TOF technique, the energy, the PSD, and the time information from the detector system are the important outputs to be processed. A block diagram of the data acquisition system for neutron measurements is shown in Figure 3.9. A bias voltage of -2700 V is supplied to the photomultiplier tube (PMT) by the high voltage power supply. Light pulses produced by the scintillator are guided to the PMT where they are converted to electrical signals. Two output signals are provided by the PMT: (1) the fast anode (timing) signal and (2) the slow dynode (linear) signal. The linear output is connected to a 113 Ortec preamplifier to produce a voltage pulse. The pulse is fed to an Ortec 572 amplifier for signal amplification and the output is passed to an analog-to-digital converter. On the other hand, the anode signal from the PMT is split into two. One signal goes to a mesytec MPD-4 module for pulse shape discrimination where it is carried out by a constant fraction discriminator. The output is then connected to an Ortec 427A delay amplifier and then to an ADC for the PSD processing. The second signal proceeds to a fast amplifier and then to an Ortec 934 discriminator. The output is fed to an Ortec 437A time-to-analog converter for the TOF information.
Figure 3.9: Block diagram of the electronic setup for neutron measurements.

3.3 Charged-particle measurements

Charged-particle measurements were performed with the charged-particle spectrometer installed on the 65° beam line located in the large target room (see Figure 3.1). Figure 3.10 displays a photo and a schematic layout of the spectrometer. The scattering chamber, which houses the targets and detectors, is about 32 cm in radius. Unlike neutrons, charged particles lose their energy even when they travel through air. Hence, it is essential to put the whole system including the target and detectors under high vacuum. The chamber is designed to hold 10 detectors through apertures positioned at 22.5° to 157.5° angles with an increment of 15°. The detectors can directly be mounted on the apertures for short flight
path. In case a longer flight path is needed, a TOF tube is attached to the aperture before mounting the detector (see Figure 3.10). For our experiments, the three most forward angles had approximately 2-m flight path and the remaining angles had approximately 1.5-m. The longer tubes were chosen to lessen the solid angles subtended by the detectors, thereby, reducing the high count rate in the forward angles due to elastic scattering. Also, longer flight path reduces the relative uncertainty in the TOF which leads to a better mass resolution for particle identification.

The targets were mounted on a target ladder capable of holding 7 target foils (see Figure 3.11). The ladder was positioned at the center of the scattering chamber, and it was operated by a rotating knob for raising and lowering it. The center of each target position on the ladder were properly calibrated through a laser installed inside the beam line.

The targets were oriented 40° relative to the plane perpendicular to the incident beam direction to optimize the detection of emitted particles from the reactions. Such orientation increases the effective thickness of the target by about 30%, which was accounted for in the calculation of the absolute cross sections. The advantage of an experiment where measurements at all angles are carried out simultaneously is that it eliminates the possible variation of the cross sections due to the target’s non-uniform thickness and the effect of beam profile fluctuations.

3.3.1 Charged-particle detection

There are several types of charged-particle detectors. Frequently employed ones for spectroscopic and nuclear reaction studies are solid-state detectors. In the experiment, we employed silicon surface barrier detectors, which are known to be efficient because of their high stopping power vital for energy absorption, low energy threshold for electron-hole pair creation, and a moderate band gap.
Figure 3.10: Photo of the charge particle spectrometer (top) and schematic layout of the scattering chamber (bottom) for charged-particle measurements. The target holder is located at the center of the chamber. The silicon detectors (Sid) are attached to the end of the extended tubes.
Solid state detectors typically consist of junctions made up of n-type and p-type semiconductors (germanium or silicon). Under reverse bias, a depletion region is formed along the layer where both semiconductors are in contact. This region acts as the active component of the detector. When a particle impinges the detector, it deposits its energy by generating electron-hole pairs. The voltage applied to the detector induces an electric field that sweeps the electrons and holes out from the depletion region and these are collected as electrical signals. The signal provides the arrival time of the particle and the pulse height.

3.3.1.1 Particle identification

Silicon surface barrier detectors are sensitive to any charged particles (e.g. protons, deuterons, alphas) incident on them. It is important to identify which specific particle type produced the event. In our experiments, the particle identification was performed by simultaneously measuring the energy deposited by the particle in the detector and the time it takes for the particle to reach the detector from the target. The flight time was obtained in a similar fashion with that in the neutron experiments as explained in section 3.2.1.1.
The distinction of events among particles is based on their mass $m$ and speed $v$. Incident particles with the same kinetic energy $E$ but of different masses travel with different speeds. Since the particles are produced almost simultaneously and the flight path $d$ is fixed, the amount of time $t$ it takes for a particle to reach the detector varies depending on its mass according to the classical relation

$$
t = \sqrt{\frac{1}{2} \frac{md^2}{E}}.
$$

(3.4)

By plotting both the time and the deposited energy for each event, one can see that particles of the same mass gather together to create their own band. Figure 3.12 shows a typical TOF vs pulse height plot. A two-dimensional gate is then applied to isolate the events generated by each particle type. Since the particle identification is determined by the mass, particles with similar masses such as $^3$He and $^3$H cannot be distinguished from each other with this method. These particles, however, are not detected in our experiments since their corresponding cross sections are negligible for the reactions presently studied.

### 3.3.2 Data Acquisition (DAQ) system

Figure 3.13 shows the block diagram of the data acquisition system for the charged-particle measurements. The silicon surface barrier detector is connected to an Ortec 142 preamplifier to provide bias to the detector through an external HV power supply and to amplify the weak electrical signals from the detector. The preamplifiers are fast-rise-time and charge-sensitive devices designed for optimum performance with charged-particle detectors. In order to improve the signal-to-noise ratio, the preamplifiers were positioned very close to the detectors. The output from the preamplifiers includes time and energy signals. The energy output is fed to a spectroscopic amplifier for signal amplification. The output is then connected to the Q1 input of the computer for digitization. The time signal, on the other hand, is outputted to a timing filter amplifier. The amplified signal is then connected to a constant fraction discriminator (CFD) and to a gate and delay generator.
Figure 3.12: Typical TOF vs energy spectrum used to identify charged particles registered in silicon detectors. The three bands correspond to proton, deuteron, and $\alpha$-particle events.

Finally, the output is then connected to the TOF input of the data distributed acquisition system (DDAQ). The signal from the time output is used as the start signal while the signal from the beam pick-off enters the DDAQ for the stop signal. The elapsed time between the two pulses provides the relative TOF information for each event registered in the detector.

### 3.4 Cross section calculation

The particle energy spectrum indicates only the number of particles registered in the detector. To compare the spectrum with theoretical calculations, the number of particles needs to be converted to reaction cross section. The differential cross section is determined by

$$\frac{d\sigma}{d\Omega} = \frac{N_S}{N_i \cdot \Delta t \cdot \Delta \Omega \cdot n_i \cdot \epsilon}$$  \hspace{1cm} (3.5)
where

\[ N_s \] is the number of scattered particles registered in the detector;

\[ N_i \] is the number of incident particles on the target;

\[ \Delta t \] is the thickness of the target;

\[ n_i \] is the number of atoms per unit volume of the target;

\[ \Delta \Omega \] is the solid angle subtended by the detector relative to the scattering target; and

\[ \epsilon \] is the detector efficiency.

The quantity \( N_i \) is calculated from the integrated beam current measured by the Faraday cup. The \( n_i \) is computed from the density of the target material and atomic weight of the target nucleus. The solid angle \( \Delta \Omega \) is determined from the effective area of the detector and the distance between the detector and the target position. Finally, the detection efficiency depends on the type of detector being used. For charged particles, silicon detectors are known to be nearly 100% efficient so \( \epsilon \) is approximately equal to 1. Neutron detectors, on the other hand, only detects a fraction of the incident neutron flux, hence \( \epsilon \) is less than 1.
In the experiments, the systematic uncertainty of the measured cross section is attributed to the BCI (1%), detection efficiency (5%), and target thickness (15%). Thus, the total systematic uncertainty of the measured cross section is about 16%. In order to get the total uncertainty, this is added in quadrature with the statistical uncertainty.

Another important factor that contributes to the calculation of cross section is the dead time. This is the time interval that the data acquisition (DAQ) system uses to process a single pulse. Signals generated during this time may be lost resulting in incorrect registered number of counts. To correct for the dead time, we used the expression

\[ DT_{\text{cor}} = \frac{\text{measured counts}}{1 - \tau_D} \]  

(3.6)

where \( \tau_D \) is the fractional dead time of the DAQ system and \( DT_{\text{cor}} \) is the dead time corrected counts.

Equation (3.6) accounts only for pulses that are lost when another pulse is being processed. In some cases, two pulses which are very close to each other in time may be summed together to represent a single pulse with an incorrect amplitude. Such events are called "pile-up", and they are not accounted for by the dead time correction. Hence, it is important to keep the count rate low to minimize the number of pile-up events. The detector located at 37.5° relative to the incident beam direction has the largest dead time due to the high count rate from elastic scattering of deuterons. The dead times for the other detectors were found to decrease as we moved to backward angles. During the experiment, the dead time was minimized by tuning the beam such that the count rate of the detector at 37.5° was less than 1.3 kHz. The average detector dead times measured for \( E_d = 5 \) MeV are tabulated in Table 3.1. These values were also found to decrease as the deuteron incident energy increased (e.g. average dead time for the detector at 37.5° with \( E_d = 9 \) MeV is 5%).

The targets used in the experiments are self-supporting foils. Details of the target foils are summarized in Table 3.2.
Table 3.1: Angular positions of the silicon detectors and their typical dead time values from nuclear reactions with $E_d = 5$ MeV.

<table>
<thead>
<tr>
<th>Detector</th>
<th>$\theta$ (°)</th>
<th>Dead time (%)</th>
<th>Detector</th>
<th>$\theta$ (°)</th>
<th>Dead time (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>37.5</td>
<td>21</td>
<td>5</td>
<td>127.5</td>
<td>1</td>
</tr>
<tr>
<td>2</td>
<td>52.5</td>
<td>7</td>
<td>6</td>
<td>142.5</td>
<td>1</td>
</tr>
<tr>
<td>3</td>
<td>67.5</td>
<td>3</td>
<td>7</td>
<td>157.5</td>
<td>1</td>
</tr>
<tr>
<td>4</td>
<td>97.5</td>
<td>2</td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

Table 3.2: Target information.

<table>
<thead>
<tr>
<th>Element</th>
<th>Isotopic abundance (%)</th>
<th>Thickness (mg/cm²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{63}$Cu</td>
<td>99.8</td>
<td>2</td>
</tr>
<tr>
<td>$^{65}$Cu</td>
<td>99.7</td>
<td>2</td>
</tr>
<tr>
<td>$^{54}$Fe</td>
<td>96</td>
<td>0.61</td>
</tr>
<tr>
<td>$^{56}$Fe</td>
<td>99.9</td>
<td>0.54</td>
</tr>
<tr>
<td>$^{58}$Fe</td>
<td>81.9</td>
<td>0.64</td>
</tr>
</tbody>
</table>

3.5 Experimental level density

The Hauser-Feshbach theory is described in section 2.1.2.1. There it shows the necessary steps in calculating the average cross section for compound nuclear reactions. The formation of a compound nucleus is expressed in terms of the transmission coefficients of the incoming particle, and its decay is expressed in terms of both the transmission coefficients of the outgoing particles and the level density of the residual nuclei. In
this study, we applied the Hauser-Feshbach theory to obtain the level density from the experimentally measured compound reaction cross sections. This will be discussed in the sections that follow.

3.5.1 Extraction of the level density

The level density of a residual nucleus is deduced using the method suggested by Vonach [23]. This involves direct bin by bin renormalization of the input level density in Hauser-Feshbach calculations until the calculated spectrum reproduces the experiment. The method starts with an assumed initial level density function for the residual nucleus of interest. The transmission coefficients have a minor influence on the calculations so uncertainties associated with them can be neglected at this point. The cross section for the first-step particle emission is computed and compared with the experimental cross section. In order to extract the level density, we use the Hauser-Feshbach given by equation (2.9) to calculate for the cross section \( \frac{d\sigma}{dE} \). By comparing the calculated (model) value of \( \frac{d\sigma}{dE}_{\text{mod}} \) with the measured value \( \frac{d\sigma}{dE}_{\text{meas}} \), the level density \( \rho_{\text{meas}}(E) \) can be derived as [23]

\[
\rho_{\text{meas}}(E) = \rho_{\text{mod}}(E) \left( \frac{\frac{d\sigma}{dE}_{\text{meas}}}{\frac{d\sigma}{dE}_{\text{mod}}} \right).
\] (3.7)

The data points \( \rho_{\text{meas}} \) obtained are fitted by a level density function to extract the best-fit parameters. These parameters are then applied to generate the new input level density of that residual nucleus for calculating the new cross section. The method is repeated until the extracted level density converges. Finally, the obtained level density is required to match the density in the discrete level region to get the final experimental level density.

The level density extracted by this technique is restricted only to excitation energies up to the particle separation energy of the nucleus being studied. Beyond this energy, second-step particle emission may likely contribute to the measured cross section. The second-step emission occurs when the excitation energy of the first compound nucleus is large enough to permit a sequential emission of more than one particle. Since the second
stage particle emission originates from another compound nucleus, using the corresponding cross sections may provide incorrect level density information.

### 3.5.2 Non-compound contamination in the evaporation spectra

The Hauser-Feshbach formulas given by equations (2.8) and (2.9) assume only a compound reaction mechanism. However, direct, including breakup reaction, and pre-equilibrium mechanisms contribute to the experimental total cross section. To effectively use the Hauser-Feshbach theory for obtaining the level density, it is necessary to minimize the contribution of non-compound processes.

The cross section of non-compound reactions strongly depends on the energies of the incident and outgoing particles. As discussed in the beginning of Chapter 2, direct and pre-equilibrium emissions dominate at high projectile energies. Hence, it is desirable to use low bombarding energies to produce compound reactions. It should be noted that the contamination from non-compound emission affects more the high energy region of the particle spectrum [23]. In addition to using low bombarding energy, one can also test the contribution of each mechanism through the angular distribution of the outgoing particles. An angular distribution with symmetry at 90° should indicate the dominance of compound emission [36] while a forward-peaked distribution implies direct and pre-equilibrium emissions. This suggests that spectra measured at backward angles and at low bombarding energies are mainly due to compound reaction mechanisms.
4 Experimental results from d + $^{63,65}$Cu measurements

In this chapter, we present the results obtained from the analysis of measured evaporated neutrons from d + $^{63}$Cu and d + $^{65}$Cu reactions. The experiments were carried out with the swinger facility of the Edwards Accelerator Laboratory. The detector array consisting of three NE213 detectors was positioned 6.92 m away from the target, inside the well-shielded TOF tunnel. Since the diameter of an NE213 detector is 20.5 cm, the solid angle subtended by a single detector was calculated to be $6.87 \times 10^{-4}$ sr. Outgoing neutrons were detected at 20°, 40°, 60°, 80°, 100°, 120°, 130°, 140°, and 150° angles.

In the following sections, we discuss the detector efficiency measurements, data analysis, and results.

4.1 Neutron detector efficiency

Detection efficiency accounts for the unregistered neutrons that passed through the detector. The correction of neutron spectra with the detection efficiency is necessary for calculating the absolute cross sections. Experimental details for obtaining the detection efficiency of the NE213 scintillation detectors from $^{27}$Al(d,n) reaction have already been described in Chapter 3. Figure 4.1 shows the neutron emission spectrum of the d + $^{27}$Al reaction as measured by one of the NE213 scintillation detectors. The detection efficiency of the NE213 detector is calculated as the ratio between measured neutron flux and standard neutron flux taken from Ref. [63] (see equation (3.3)). Figure 4.2 displays the detection efficiencies obtained for the NE213 scintillation detectors. The noticeable difference in low energy cut-off among the detection efficiencies is due to different low energy threshold applied for each detector. Setting an energy threshold for the detector was done to reduce wrap-around neutrons from being registered as an event. Wrap-around neutrons are low energy neutrons registered on a time stamp derived from a later reaction causing them to appear as signals in the high energy region. The elimination of low energy neutrons by
increasing the energy threshold of the detector, however, decreases the detection efficiency. This is illustrated in Figure 4.2 where the discriminator thresholds of each detector were set at slightly different values. The peaks in the detection efficiency were found to originate from neutron interaction with nuclei present in air, aluminum in the scattering chamber, and carbon in the NE213 liquid scintillator [64]. Hence, these effects are automatically taken into account once the detection efficiency is incorporated in determining the experimental cross sections.

4.2 Neutron cross sections

Neutron emission spectra were measured from \( d + {}^{63}\text{Cu} \) and \( d + {}^{65}\text{Cu} \) reactions to study the residual nuclei \( {}^{64}\text{Zn} \) and \( {}^{66}\text{Zn} \). The \( {}^{63}\text{Cu} \) and \( {}^{65}\text{Cu} \) foils were bombarded by deuterons with energies of 6 and 7.5 MeV. The choice of having two bombarding energies was motivated to check for contribution from pre-equilibrium processes, which are known to increase with bombarding energy. Figure 4.3 shows the sample TOF spectrum and its equivalent neutron energy spectrum measured at 150°.
Before the TOF spectrum was converted to neutron energy spectrum, background events were subtracted. First, γ-ray events were removed through pulse shape discrimi-
nation as discussed in section 3.3. Second, events due to background neutrons which are uncorrelated in time (e.g. down-scattered neutrons), were subtracted as uniform distribution in the TOF spectrum. Finally, events from neutrons due to interaction of beam with the surrounding materials other than the target were removed by subtracting target-in and target-out spectra normalized by their respective beam current integration (BCI).

4.2.1 Double differential cross sections

The neutron energy spectrum was converted to cross sections in the laboratory frame of reference with equation (3.5). Experimental cross sections were further transformed to the center-of-mass (CM) frame of reference, where the target particle and projectile are moving opposite in direction with equal momenta. The conversion of our cross sections from laboratory system to CM system was done in order to compare them with theoretical calculations, where results are generally expressed in the CM system. Since the transformation was assumed only for a binary (d,n) reaction, energies of neutrons from secondary emission are slightly different from their correct center-of-mass energies. For neutron peaks due to reactions on target contaminants, such as carbon and oxygen, a constant shift to lower energy with respect to angle in the CM frame is expected. Such variation can be used to spot contamination peaks.

The cross sections of $^{63}\text{Cu}(d,Xn)$ and $^{65}\text{Cu}(d,Xn)$ reactions for given emission angles and bombarding energies are shown in Appendix A. The cross sections are all inclusive where all emitted neutrons from all open Xn channels are included. The open Xn channels with their reaction Q-values are tabulated in Table 4.1.

4.2.2 Angular distributions

Level density study through the evaporation technique requires the measured particles to emanate from the decay of compound nuclei. As already discussed in Chapter 2, direct and pre-equilibrium reactions favor the emission of particles in the forward direction
Table 4.1: Q-values for open neutron channels. $E_{b(min)}$ corresponds to the minimum excitation energy of the residual nucleus (RN) where neutrons from the reaction product starts to contribute in the cross section.

<table>
<thead>
<tr>
<th>Reaction products</th>
<th>Q-value (MeV)</th>
<th>$E_{b(min)}$ (MeV)</th>
<th>RN</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{64}$Zn+n</td>
<td>5.49</td>
<td>0</td>
<td>$^{64}$Zn</td>
</tr>
<tr>
<td>$^{60}$Ni+n+$\alpha$</td>
<td>1.53</td>
<td>3.95</td>
<td>$^{62}$Ni</td>
</tr>
<tr>
<td>$^{63}$Cu+n+p</td>
<td>-2.22</td>
<td>7.71</td>
<td>$^{63}$Zn</td>
</tr>
<tr>
<td>$^{56}$Fe+n+2$\alpha$</td>
<td>-4.76</td>
<td>10.24</td>
<td>$^{58}$Fe</td>
</tr>
<tr>
<td>$^{63}$Zn+2n</td>
<td>-6.37</td>
<td>11.86</td>
<td>$^{64}$Zn</td>
</tr>
</tbody>
</table>

while those from compound emissions exhibit symmetry about 90° in the CM system. Hence, the analysis of the angular distribution of emitted particles can be used as a tool to determine angles where compound reaction mechanisms dominate. Angular distributions were evaluated by integrating the double differential cross sections over selected energy intervals.

Figure 4.4 shows the angular distributions for both $^{63}$Cu(d,n) and $^{65}$Cu(d,n) reactions at 6- and 7.5-MeV bombarding energies. The forward-peaked feature is apparent for all angular distributions. This is an indication that a fraction of emitted particles was produced by non-compound reaction mechanisms, particularly from direct and pre-equilibrium emissions. It is also observed that the relative fraction of the forward-peaked component increases with the outgoing neutron energy. Conversely, the flat angular distributions in the backward angles imply the dominance of compound emissions. The isotropic angular distribution is in agreement with the Hauser-Feshbach calculations performed for the angular distribution of (d,n) reaction. It is predicted that nucleons emitted from
deuteron-induced compound reactions exhibit a symmetric angular distribution at 90° with anisotropy of less than 10%. Since the cross sections in the backward angles are almost equal, the statistics was further improved by averaging the data from 120° to 150°. From this point forward, compound cross sections are represented by the average cross sections over backward angles (see Figure 4.5).

The experimental angular distributions have been fitted by the phenomenological Kalbach formula [65]. This is expressed as

\[
\sigma(\theta) = C(\exp(a_d \cos \theta) + R[\exp(a_c \cos \theta) + \exp(-a_c \cos \theta)]).
\] (4.1)

The first term represents the forward-peaked component of the angular distribution known to be from direct and multi-step direct (MSD) while the second term describes the symmetric component from equilibrium and multi-step compound (MSC) contributions.
It should be mentioned that both MSD and MSC are classified under pre-equilibrium emissions according to the quantum mechanical theory of pre-equilibrium reactions by Feshbach, Kerman, and Koonin (FKK). MSD emission occurs during the early stage of the formed composite nucleus where only a few degrees of freedom are excited. In the MSC emission, the composite nucleus has already undergone a large cascade of two-body...
interaction but the system has not attained equilibrium state yet. In equation (4.1), \( C \) is the overall normalization factor; \( R \) is the coefficient for the term containing the MSC and equilibrium part; \( a_c \) and \( a_d \) describe the steepness of the slope for both MSC and MSD, respectively; and \( \theta \) is the emission angle in the CM frame. Here, we assume that the symmetric component is due to compound reactions only.

We also employed equation (4.1) to quantitatively estimate the compound fraction of the total reaction cross section. Here, the double differential cross sections are integrated over the energy interval where only the first stage neutron emission occurs. The compound fractions obtained are then used as reduction factors for the theoretically calculated cross sections, which are based from the compound model only. Hence, a compound fraction acts as a factor to account for the incoming flux that proceeds through the non-compound reaction mechanisms.

The estimated fractions are tabulated in Table 4.2. The decrease of the compound nuclear fraction as the beam energy increases is in agreement with the observation from Ref. [66]. This dependence is a direct consequence of the redistribution of the entrance channel energy. With higher projectile energies, nucleons involved in the intranuclear collisions are more probable of receiving sufficient energy for them to be emitted in the continuum even before the composite system achieves equilibrium. The larger non-compound fraction for the deuteron energy of 7.5 MeV may suggest a larger pre-equilibrium MSC contribution compared with the data from 6 MeV. It should be mentioned that the uncertainties of the compound fractions given in Table 4.2 are obtained from the fit.

4.3 Compound neutron cross sections

Figure 4.6 shows the measured cross sections averaged over backward (120°, 130°, 140°, and 150°) angles. Theoretical calculations were carried out in the framework of
Table 4.2: Estimate of compound fractions obtained from the experimental angular distributions based on equation (4.1). Uncertainties are given in parentheses.

<table>
<thead>
<tr>
<th>Energy (MeV)</th>
<th>Compound (%)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>d + $^{63}$Cu</td>
</tr>
<tr>
<td>6</td>
<td>78(4)</td>
</tr>
<tr>
<td>7.5</td>
<td>73(3)</td>
</tr>
</tbody>
</table>

the statistical Hauser-Feshbach (HF) theory using the nuclear reaction code EMPIRE [31]. In order to account for the non-compound component, the calculated cross sections were multiplied by the compound fractions found in Table 4.2. The optical model potentials (OMP) required for the calculation of the transmission coefficients were taken from Ref. [67] for neutrons and protons, Ref. [68] for $\alpha$-particles, and Ref. [69] for deuterons. For the nuclear level density inputs, we employed two phenomenological models: (1) the Fermi gas model (FGM) using the parameters obtained from the von Egidy systematics (FGM (Egidy sys)) [1]; (2) the Gilbert-Cameron model (GCM) [2] using the level density parameter systematics of Iljinov (GCM (Iljinov sys)) [4] and Ignatyuk (GCM (Ignatyuk sys)) [13]; along with (3) the microscopic model using the combinatorial approach based on the microscopic Hartree-Fock-Bogoliubov method (HFBM) [3]. These are some of the commonly employed level density models in HF theoretical calculations.

The spin cut-off parameter used in the GCM for the Fermi gas component above the matching energy is given by equation (2.29) while the constant temperature component uses the spin cut-off parameter as a linear interpolation between the discrete levels and equation (2.29) at the matching energy. The Fermi-gas level density parameter $a$ is energy dependent as suggested in Ref. [13] given by equation (2.18). The temperature $T$ for the constant temperature part is adjusted such that the level density at low energies is consistent
with the density of low-lying levels and that it should be continuous at the matching energy where the transition from constant temperature model to Fermi gas model occurs.

For the FGM parametrized by the von Egidy systematics of Ref. [1], the spin cut-off parameter utilized the form

\[
\sigma^2 = 0.0146 A^{5/3} \frac{[1 + \sqrt{1 + 4a(U - E_0)}]}{2a}. \quad (4.2)
\]

This equation is adequate for low energy reactions according to Ref. [70]. Both the level density parameter \(a\) and the backshift energy parameter \(E_0\) are empirically determined in terms of its correlation with the deuteron pairing energy \(P_d\) (such that \(P_d = 1/4(-1)^{Z+1} [S_d(A + 2, Z + 1) - 2S_d(A, Z) + S_d(A - 2, Z - 1)]\)) and the shell correction \(S(Z, N)\) (see equation (2.19)). Ref. [1] has given the equations

\[
a/A = p_1 + p_2 S'(Z, N) + p_3A \quad (4.3)
\]

and

\[
E_0 = \begin{cases} 
  p_1 - 0.5P_d + p_4 \frac{dS(Z, N)}{dA}, & \text{(even-even)} \\
  p_2 - 0.5P_d + p_4 \frac{dS(Z, N)}{dA}, & \text{(odd-A)} \\
  p_3 + 0.5P_d + p_4 \frac{dS(Z, N)}{dA}, & \text{(odd-odd)} 
\end{cases} \quad (4.4)
\]

where \(p_1, p_2, p_3, \text{ and } p_4\) are constants summarized in Table 4.3. It is important to mention that the level density parameter \(a\) in all parameter systematics is dependent on the shell correction.

We first consider the cross sections obtained from the 6-MeV bombarding energy. Theoretically calculated cross sections using GCM and FGM level density models are in reasonable agreement with each other and with the experimental data points for neutron energies between 2 and 7 MeV as shown in Figure 4.6. Slight deviations appear only at the lowest and highest neutron energies. The latter feature would pose a problem when applied
Table 4.3: Constants for von Egidy level density parameter systematics. The constants $p_1$, $p_2$, $p_3$, and $p_4$ are distinct for equations (4.3) and (4.4).

<table>
<thead>
<tr>
<th></th>
<th>$p_1$</th>
<th>$p_2$</th>
<th>$p_3$</th>
<th>$p_4$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$a/A$</td>
<td>0.127</td>
<td>4.98\times10^{-3}</td>
<td>-8.95\times10^{-5}</td>
<td></td>
</tr>
<tr>
<td>$E_0$</td>
<td>-0.468</td>
<td>-0.565</td>
<td>-0.231</td>
<td>0.438</td>
</tr>
</tbody>
</table>

Table 4.4: Parameters for the Gilbert-Cameron (GCM) and Fermi gas (FGM) level density models. Best GCM parameters are underlined.

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>$\tilde{\alpha}$</th>
<th>$U_x$</th>
<th>$\delta$</th>
<th>$E_0$</th>
<th>$T$</th>
<th>ref</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{64}$Zn</td>
<td>10.11</td>
<td>8.08</td>
<td>3.0</td>
<td>-0.30</td>
<td>1.07</td>
<td>[13]</td>
</tr>
<tr>
<td></td>
<td>8.86</td>
<td>9.49</td>
<td>3.0</td>
<td>-0.93</td>
<td>1.23</td>
<td>[4]</td>
</tr>
<tr>
<td>$^{66}$Zn</td>
<td>10.43</td>
<td>6.62</td>
<td>2.95</td>
<td>0.42</td>
<td>0.93</td>
<td>[13]</td>
</tr>
<tr>
<td></td>
<td>9.12</td>
<td>7.92</td>
<td>2.95</td>
<td>-0.14</td>
<td>1.09</td>
<td>[4]</td>
</tr>
</tbody>
</table>

FGM Parameters (Egidy systematics [1])

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>$a$</th>
<th>$\delta$</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{64}$Zn</td>
<td>7.21</td>
<td>0.80</td>
</tr>
<tr>
<td>$^{66}$Zn</td>
<td>8.14</td>
<td>1.02</td>
</tr>
</tbody>
</table>

to nuclear reactions that involve a compound nucleus at very high excitation energies where these models may depart substantially from each other. Hence, additional level density information, e.g. from Ericson fluctuations, is needed to test the models at that energy domain. It should also be noted that the microscopic model also reproduces the
The cross sections calculated from GCM are in close agreement with the experimental data for both $^{63}\text{Cu}(d,n)^{64}\text{Zn}$ and $^{65}\text{Cu}(d,n)^{66}\text{Zn}$ reactions. However, to consistently reproduce the cross sections, different level density parameter systematics have been employed. The GCM parameters used for $^{63}\text{Cu}(d,n)^{64}\text{Zn}$ were taken from the Iljinov level density parameter systematics [4] while those of $^{65}\text{Cu}(d,n)^{66}\text{Zn}$ were from the Ignatyuk systematics [13] (see underlined values in Table 4.4). Using either Iljinov or Ignatyuk systematics for both nuclei would give disagreement with the data points for one of them. The same systematics would predict a $\sim 16\%$ difference in $T$ and $\sim 13\%$ in $a$ for $^{64}\text{Zn}$ and $^{66}\text{Zn}$ while the data indicate the same $T$ within $2\%$ difference. The use of different parameter
systematics in GCM only demonstrates that global level density parametrizations used in Ignatyuk and Iljinov systematics do not always work for all nuclei.

When the cross sections for $^{63}$Cu(d,n)$^{64}$Zn and $^{65}$Cu(d,n)$^{66}$Zn reactions with bombarding energy of 7.5 MeV were calculated using the best GCM level density parameters obtained from the 6-MeV data, they were observed to slightly underestimate the experimental cross sections in the high neutron energy region (as shown in Figure 4.6). This may be attributed to the presence of pre-equilibrium process, which is known to enhance the higher energy portion of the neutron energy spectrum more than the lower energy part.

### 4.4 Level densities of $^{64}$Zn and $^{66}$Zn

The level densities of $^{64}$Zn and $^{66}$Zn have been extracted from the double differential cross sections measured at backward angles by the method described in section 3.5. The deduced level densities are shown in Figure 4.7. Here, we show only the statistical uncertainties of the data points. Good agreement is observed between level densities obtained from the cross sections measured with 6- and 7.5-MeV deuterons at excitation energies above 5.5 to 6.5 MeV. For lower excitation energies populated by high energy outgoing neutrons, we only present the data points from 6-MeV deuterons. The level densities produced by 7.5-MeV deuterons have not been deduced for low excitation energies due to the contribution of non-compound reaction mechanisms. The extracted level densities also display step structures at excitation energies $\sim$3 and $\sim$2.5 MeV for $^{64}$Zn and $^{66}$Zn, respectively. Studies have shown that such structure is characterized by a sudden increase of the level density and it can be interpreted as the breaking of nucleon Cooper pairs [71]. The increase in level density is related to the creation of new degrees of freedom resulting from the broken pairs.

The experimental level densities of $^{64}$Zn and $^{66}$Zn were fitted by the constant temperature model function to obtain the best-fit parameters. It was observed that the
slope for a specific isotope slightly changed with bombarding energies. The temperatures for $^{64}$Zn were found to be 1.10(2) and 1.19(2) for 6- and 7.5-MeV deuteron energies, respectively. Similarly, the temperatures for $^{66}$Zn were found to be 1.08(2) and 1.18(2) MeV. The slight increase in temperature may be attributed to the contribution of non-compound reaction mechanisms obtained with the higher deuteron bombarding energy. Direct and pre-equilibrium mechanisms are known to enhance the neutron cross sections at high energies resulting in a less steep slope.

Figure 4.7 also presents the comparison of the experimental level densities to known NLD models. No available global systematics has provided parameters for specific NLD model to consistently reproduce the data points for both $^{64}$Zn and $^{66}$Zn. The GCM model which best describes the data points used different parameter systematics for each nucleus. In addition, the reasonable agreement between the experiment and calculated cross sections using FGM (Egidy sys) (see Figure 4.6) is not apparent in Figure 4.7. FGM (Egidy sys) level densities exhibit an offset relative to the experimental data points. The shape of the level density calculated from the von Egidy systematics is not consistent with the experimental data particularly at high excitation energies. The difference in shape and slope of $^{64}$Zn and $^{66}$Zn as predicted by von Egidy, Iljinov, and Ignatyuk parameter systematics are due to the shell correction $\delta W$ for these nuclei according to formula (2.19).
Figure 4.7: Level densities of $^{64}$Zn (left) and $^{66}$Zn (right). Points are experimental data. Lines are calculations. Histogram is the density of discrete levels.
5 EXPERIMENTAL RESULTS FROM D+\(^{54,56,58}\)Fe MEASUREMENTS

This chapter presents the results obtained for d + \(^{54}\)Fe, d + \(^{56}\)Fe, and d + \(^{58}\)Fe reactions. Four of the major exit channels, namely neutron, proton, deuteron, and \(\alpha\)-particle channels, have been studied to provide better understanding of the reaction mechanisms and level densities. Charged particles were measured by the charged-particle spectrometer at selected angles 37.5°, 52.5°, 67.5°, 97.5°, 127.5°, 142.5°, and 157.5°. Registered particles were identified based on their TOF information and deposited energy in the silicon detectors (see section 3.3 for further details). The experimental technique and the data analysis for the neutron channel are similar to the ones presented in Chapter 4, thus, these will not be repeated here.

In the following sections, we discuss the energy calibration of the charged-particle detectors and determination of their subtended solid angles, data analysis, and results.

5.1 Detector calibration and solid angle determination

The energy calibration of the silicon detectors is required to convert the ADC spectrum to particle energy spectrum. In the experiment, the calibration peaks were provided by \(^{12}\)C(d,p)\(^{13}\)C, \(^{12}\)C(d,d)\(^{12}\)C, and \(^{12}\)C(d,\(\alpha\))\(^{10}\)B reactions with incident deuteron beam energies of 5 and 7 MeV. These reactions produced calibration peaks between 2 and 9 MeV. For energies above 9 MeV, we employed the peaks that populate the ground state and first few excited states of \(^{55}\)Fe through the \(^{54}\)Fe(d,p)\(^{55}\)Fe reaction. Table 5.1 summarizes the states populated by these reactions. The excitation energies corresponding to these peaks are available in the Evaluated Nuclear Structure Data File (ENSDF) database [16]. Figure 5.1 shows the spectrum from d + \(^{12}\)C reaction with \(E_d = 7\) MeV. The centroid of each peak was identified by Gaussian curve fitting. The equivalent energies in the laboratory frame were determined from the reaction kinematics based on the excitation energy, the detection angle, and the Q-value of the reaction. The energy calibration constants were
extracted from the coefficients of the linear regression fit between the centroid channels and the equivalent energies of the peaks (see Figure 5.2).

Table 5.1: Levels of residual nuclei used to calibrate the silicon detectors.

<table>
<thead>
<tr>
<th>Reaction</th>
<th>Q-value (MeV)</th>
<th>Level (J^π)</th>
<th>Energy (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>^{12}\text{C}(d,p)^{13}\text{C}</td>
<td>2.72</td>
<td>1/2^-</td>
<td>0.0</td>
</tr>
<tr>
<td></td>
<td></td>
<td>1/2^+</td>
<td>3.089</td>
</tr>
<tr>
<td></td>
<td></td>
<td>3/2^-</td>
<td>3.684</td>
</tr>
<tr>
<td></td>
<td></td>
<td>5/2^+</td>
<td>3.853</td>
</tr>
<tr>
<td></td>
<td></td>
<td>5/2^+</td>
<td>6.864</td>
</tr>
<tr>
<td>^{12}\text{C}(d,d')^{12}\text{C}</td>
<td>0</td>
<td>0^+</td>
<td>0.0</td>
</tr>
<tr>
<td></td>
<td></td>
<td>2^+</td>
<td>4.438</td>
</tr>
<tr>
<td>^{12}\text{C}(d,\alpha)^{10}\text{B}</td>
<td>-1.34</td>
<td>3^+</td>
<td>0.0</td>
</tr>
<tr>
<td></td>
<td></td>
<td>1^+</td>
<td>0.718</td>
</tr>
<tr>
<td></td>
<td></td>
<td>1^+</td>
<td>2.154</td>
</tr>
<tr>
<td>^{54}\text{Fe}(d,p)^{55}\text{Fe}</td>
<td>7.073</td>
<td>3/2^-</td>
<td>0.0</td>
</tr>
<tr>
<td></td>
<td></td>
<td>1/2^-</td>
<td>0.411</td>
</tr>
<tr>
<td></td>
<td></td>
<td>5/2^-</td>
<td>0.931</td>
</tr>
<tr>
<td></td>
<td></td>
<td>7/2^-</td>
<td>1.316</td>
</tr>
</tbody>
</table>

The solid angle subtended by each detector was established by measuring the elastic scattering of deuterons on a gold foil. At bombarding energy of 5 MeV, the elastic scattering cross section can be calculated accurately using the Rutherford scattering formula. The number of elastically scattered deuterons was determined from the integral of the elastic scattering peak in the spectrum. Peaks measured at angles 37.5° and 52.5° were corrected
Figure 5.1: Energy spectrum from $^{12}\text{C}(d,p)^{13}\text{C}$ (blue), $^{12}\text{C}(d,d)^{12}\text{C}$ (red), and $^{12}\text{C}(d,\alpha)^{10}\text{B}$ (green) with deuteron beam energy of 7 MeV measured at 37.5°.

Figure 5.2: Energy calibration fit of the detector located at 37.5° relative to the beam line.
for contaminations due to silver and copper impurities present in the gold foil. The solid angle subtended by each detector was calculated according to equation (3.5) using the information of the elastic cross section, the elastic scattering yield, the intensity of the beam, and the thickness of the target.

To validate the calculated solid angles, we initially utilized them to obtain the elastic scattering cross sections for deuterons on $^{56}\text{Fe}$ with incident energies of 5 and 7 MeV. The obtained ratios of elastic and Rutherford cross sections were then compared to previously published data reported in Ref. [66] (see Figure 5.5). Both are in excellent agreement.

Table 5.2: Solid angles determined from the elastic scattering cross sections of deuterons on gold.

<table>
<thead>
<tr>
<th>Detector Angle (degree)</th>
<th>Solid angle (sr)</th>
<th>Statistical uncertainty (sr)</th>
</tr>
</thead>
<tbody>
<tr>
<td>37.5°</td>
<td>7.27×10^{-5}</td>
<td>9.38×10^{-8}</td>
</tr>
<tr>
<td>52.5°</td>
<td>6.42×10^{-5}</td>
<td>1.55×10^{-7}</td>
</tr>
<tr>
<td>67.5°</td>
<td>6.40×10^{-5}</td>
<td>2.43×10^{-7}</td>
</tr>
<tr>
<td>97.5°</td>
<td>1.94×10^{-4}</td>
<td>7.66×10^{-7}</td>
</tr>
<tr>
<td>127.5°</td>
<td>2.15×10^{-4}</td>
<td>1.12×10^{-6}</td>
</tr>
<tr>
<td>142.5°</td>
<td>1.88×10^{-4}</td>
<td>1.17×10^{-6}</td>
</tr>
<tr>
<td>157.5°</td>
<td>1.83×10^{-4}</td>
<td>1.23×10^{-6}</td>
</tr>
</tbody>
</table>

5.2 Elastic cross sections and optical model potential

Elastic scattering data for deuterons are not as extensive compared to protons and neutrons. Moreover, deuteron optical model potentials (OMP) available in the literature have been optimized using data derived from energies above 20 MeV and angles below 90°. Recent studies have shown that cross sections calculated from known deuteron optical
model potentials differ from experimental data especially for energies below 20 MeV, and that adjustments of the parameters are necessary in order to make them agree [72, 73]. Since the deuteron optical potential is necessary for calculating the fusion cross section, we employed our elastic scattering data to test known optical potential parametrizations at incident deuteron energies considered in this study. Parameters that closely reproduce the experimental elastic angular distributions are the ones utilized for the compound model calculations.

![Elastic scattering spectrum](image)

Figure 5.3: Elastic scattering spectrum in the laboratory frame from $^{54}$Fe(d,d)$^{54}$Fe reactions measured at 157.5°.

The deuteron elastic scattering cross sections on $^{54,56,58}$Fe nuclei have been obtained by integrating the elastic scattering peaks. Figure 5.3 shows the elastic scattering spectrum for deuteron on $^{54}$Fe measured by one of the detectors. The additional peaks in the spectrum suggest the presence of contaminants in the $^{54}$Fe foil. The doubling of the contamination peaks indicates that the contaminants accumulated on both the front and back surfaces of
the target foil and the energy gap between the peaks resulted from the energy loss of the
deuterons in the target foil.

The Woods-Saxon form factor (see formula (2.40)) was implemented all throughout
our calculations. Deuteron optical potential parameters of Haixia et al. [69], Yinlu et
al. [74], Bojowald et al. [75], and Lohr et al. [76] were tested. The OMP parameters of
Haixia [69] and Yinlu [74] are applicable for a wide variety of nuclei with deuteron energies
between 1 and 200 MeV, whereas the OMP parameters of Bojowald [75] are appropriate for
bombarding energies between 20 and 200 MeV. The OMP parameters of Lohr, on the other
hand, are locally parametrized suitable only for deuteron energies between 8 and 13 MeV.
Therefore, in our analysis, the OMP parameters of Bojowald and Lohr were extrapolated
down to deuteron energies of 5, 7, and 9 MeV. The parameters were retrieved from the
Reference Input Parameter Library (RIPL-3) [30], and these are listed in Appendix B.
The smooth variation of the parameters with respect to the mass number is expected since
the optical potential varies slowly between neighboring nuclei. Any rapid variation in the
parameters is a consequence of nuclear structure effects such as the presence of collective
states.

Figures 5.4 - 5.6 show the comparison between the present data and the optical model
calculations for bombarding energies of 5, 7, and 9 MeV. All optical potentials describe the
data reasonably well at forward angles. The optical potentials of Haixia [69] and Yinlu [74]
exhibit distinct minima at backward angles, which are not observed with the potentials from
Bojowald [75] and Lohr [76]. For this reason, the best fit procedure was not performed on
our data points because they have large angular intervals that prevent us from making an
unambiguous fit.

As tabulated in Table 5.3, the total absorption cross sections from different deuteron
optical potential parametrizations vary by about 20% for a given bombarding energy. Since
the deuteron OMP of Lohr [76] consistently agrees with our elastic angular distribution
data, it was used in the HF analysis of our compound cross section data. It should be noted
that the total absorption cross sections calculated with the OMP parameters of Haixia and
Lohr differ within 2% only (see Table 5.3). Therefore, using either one of them in the
data analysis should give results with minimal difference. Data points for $^{56}$Fe(d,d)$^{56}$Fe
reactions with deuteron bombarding energies of 5 and 7 MeV from Ref. [66] are also
plotted for reference (see Figure 5.5).

Table 5.3: Deuteron absorption cross sections as predicted from different optical model
potentials.

<table>
<thead>
<tr>
<th>Nucleus</th>
<th>Energy (MeV)</th>
<th>Haixia [69] (mb)</th>
<th>Yinlu [74] (mb)</th>
<th>Lohr [76] (mb)</th>
<th>Bojowald [75] (mb)</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{54}$Fe</td>
<td>5</td>
<td>536</td>
<td>483</td>
<td>542</td>
<td>617</td>
</tr>
<tr>
<td></td>
<td>7</td>
<td>985</td>
<td>912</td>
<td>994</td>
<td>1093</td>
</tr>
<tr>
<td></td>
<td>9</td>
<td>1244</td>
<td>1156</td>
<td>1260</td>
<td>1368</td>
</tr>
<tr>
<td>$^{56}$Fe</td>
<td>5</td>
<td>553</td>
<td>501</td>
<td>558</td>
<td>636</td>
</tr>
<tr>
<td></td>
<td>7</td>
<td>1008</td>
<td>936</td>
<td>1015</td>
<td>1117</td>
</tr>
<tr>
<td></td>
<td>9</td>
<td>1270</td>
<td>1182</td>
<td>1282</td>
<td>1395</td>
</tr>
<tr>
<td>$^{58}$Fe</td>
<td>5</td>
<td>571</td>
<td>514</td>
<td>573</td>
<td>655</td>
</tr>
<tr>
<td></td>
<td>7</td>
<td>1031</td>
<td>950</td>
<td>1034</td>
<td>1142</td>
</tr>
<tr>
<td></td>
<td>9</td>
<td>1296</td>
<td>1198</td>
<td>1303</td>
<td>1421</td>
</tr>
</tbody>
</table>

The optical potential parameters for neutrons and protons were taken from the OMP of
Koning et al. [67]. These are the same parameters employed in Chapter 4 for the calculation
of $^{63,65}$Cu(d,n) cross sections. For $\alpha$ particles, instead of using the OMP of Avrigeanu et al.
[68], we employed the OMP of McFadden et al. from Ref. [77]. It was done because the
calculated cross sections based on Avrigeanu’s OMP overestimated our experimental $(d,\alpha)$ cross sections at low emission energies. In addition, McFadden’s OMP has been suggested to be reliable for level density studies as indicated in Ref. [23]. Although Avrigeanu’s OMP was used to calculate the transmission coefficients of outgoing $\alpha$ particles in Chapter 4, it has a minor effect on the calculated neutron cross section since the $\alpha$ channel’s contribution to the total absorption cross section is less than 7% for both $d + ^{63,65}$Cu reactions.

5.3 Experimental data of non-elastic cross sections

5.3.1 Double differential cross sections

The inclusive neutron, proton, and $\alpha$-emission cross sections in the CM system for $d + ^{54}$Fe, $d + ^{56}$Fe, and $d + ^{58}$Fe reactions are displayed in Appendixes A.3 - A.9. Deuteron beam energies of 5, 7, and 9 MeV were used for the charged-particle measurements while only 7 MeV was used for the neutron measurements. All the exit channels that may open up for a given bombarding energy can be determined according to their reaction Q-values shown in Table 5.4.

The distinct population of low-lying levels of the residual nuclei is evident in the emission spectra due to the good energy resolution of silicon detectors. As shown in Appendixes A.3 - A.5, proton emission spectra above 12 MeV measured at 37.5° and 52.5° are not available since the thickness of the detectors used in these angles is not enough to stop the highly energetic protons. The energy loss of the incident deuterons in each target foil was estimated using the program SRIM (Stopping and Range of Ion in Matter) [78]. For $E_d = 5$ MeV, the energy loss was found to be less than 20 keV for all targets. For higher incident deuteron energies, the energy loss was lower than 20 keV. Since the cross sections were subsequently binned to 100 keV for the charged particles and 200 keV for the neutrons, the effect of energy loss in the target was neglected.
Figure 5.4: Ratio of deuteron elastic scattering and Rutherford cross sections in the center-of-mass system for $E_d = 5$ MeV (top), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (bottom) for the $^{54}\text{Fe}(d,d)^{54}\text{Fe}$ reaction. Points are experimental data, lines are optical model calculations.
Figure 5.5: Ratio of deuteron elastic scattering and Rutherford cross sections in the center-of-mass system for $E_d = 5$ MeV (top), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (bottom) for the $^{56}$Fe(d,d)$^{56}$Fe reaction. Points are experimental data, lines are optical model calculations.
Figure 5.6: Ratio of deuteron elastic scattering and Rutherford cross sections in the center-of-mass system for $E_d = 5$ MeV (top), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (bottom) for the $^{58}\text{Fe}(d,d)^{58}\text{Fe}$ reaction. Points are experimental data, lines are optical model calculations.
Table 5.4: Table of open reaction channels for 9 MeV deuteron-induced reactions on $^{54,56,58}\text{Fe}$. $E_{b(\text{min})}$ is the minimum excitation energy of the residual nucleus (RN) where the reaction product starts to contribute in the cross section.

<table>
<thead>
<tr>
<th>Reaction</th>
<th>Reaction Products</th>
<th>Q-value (MeV)</th>
<th>$E_{b(\text{min})}$ (MeV)</th>
<th>RN</th>
<th>Reaction Products</th>
<th>Q-value (MeV)</th>
<th>$E_{b(\text{min})}$ (MeV)</th>
<th>RN</th>
<th>Reaction Products</th>
<th>Q-value (MeV)</th>
<th>$E_{b(\text{min})}$ (MeV)</th>
<th>RN</th>
</tr>
</thead>
<tbody>
<tr>
<td><strong>Neutron</strong></td>
<td></td>
<td></td>
<td></td>
<td><strong>Proton</strong></td>
<td></td>
<td></td>
<td></td>
<td><strong>$\alpha$ particle</strong></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>d$+^{54}\text{Fe}$</td>
<td>$^{55}\text{Co}+n$</td>
<td>2.84</td>
<td>0</td>
<td>$^{55}\text{Co}$</td>
<td>$^{55}\text{Fe}+p$</td>
<td>7.07</td>
<td>0</td>
<td>$^{55}\text{Fe}$</td>
<td>$^{52}\text{Mn}+\alpha$</td>
<td>5.16</td>
<td>0</td>
<td>$^{52}\text{Mn}$</td>
</tr>
<tr>
<td></td>
<td>$^{54}\text{Fe}+n+p$</td>
<td>-2.22</td>
<td>5.06</td>
<td>$^{54}\text{Mn}+2p$</td>
<td>$^{54}\text{Fe}+n+p$</td>
<td>-2.14</td>
<td>9.21</td>
<td>$^{51}\text{Cr}+p+\alpha$</td>
<td>$^{51}\text{Mn}+n+\alpha$</td>
<td>-1.38</td>
<td>6.54</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$^{55}\text{Mn}+n+\alpha$</td>
<td>-5.37</td>
<td>8.21</td>
<td></td>
<td></td>
<td>-2.22</td>
<td>9.29</td>
<td></td>
<td></td>
<td>-3.7</td>
<td>10.53</td>
<td></td>
</tr>
<tr>
<td>d$+^{56}\text{Fe}$</td>
<td>$^{57}\text{Co}+n$</td>
<td>3.80</td>
<td>0</td>
<td>$^{57}\text{Co}$</td>
<td>$^{57}\text{Fe}+p$</td>
<td>5.42</td>
<td>0</td>
<td>$^{57}\text{Fe}$</td>
<td>$^{54}\text{Mn}+\alpha$</td>
<td>5.66</td>
<td>0</td>
<td>$^{54}\text{Mn}$</td>
</tr>
<tr>
<td></td>
<td>$^{56}\text{Fe}+n+p$</td>
<td>-2.22</td>
<td>6.02</td>
<td>$^{56}\text{Fe}+p+n$</td>
<td>$^{56}\text{Mn}+2p$</td>
<td>-2.22</td>
<td>7.64</td>
<td>$^{53}\text{Cr}+p+\alpha$</td>
<td>$^{53}\text{Mn}+n+\alpha$</td>
<td>-1.90</td>
<td>7.56</td>
<td></td>
</tr>
<tr>
<td></td>
<td>$^{56}\text{Co}+2n$</td>
<td>-7.61</td>
<td>11.37</td>
<td></td>
<td></td>
<td>-5.14</td>
<td>10.56</td>
<td></td>
<td></td>
<td>-3.28</td>
<td>8.94</td>
<td></td>
</tr>
<tr>
<td>d$+^{58}\text{Fe}$</td>
<td>$^{59}\text{Co}+n$</td>
<td>5.14</td>
<td>0</td>
<td>$^{59}\text{Co}$</td>
<td>$^{59}\text{Fe}+p$</td>
<td>4.35</td>
<td>0</td>
<td>$^{59}\text{Fe}$</td>
<td>$^{56}\text{Mn}+\alpha$</td>
<td>5.66</td>
<td>0</td>
<td>$^{56}\text{Mn}$</td>
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<td>7.36</td>
<td>$^{58}\text{Fe}+p+n$</td>
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<td>6.58</td>
<td>$^{55}\text{Mn}+n+\alpha$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td>$^{58}\text{Co}+2n$</td>
<td>-5.31</td>
<td>10.45</td>
<td></td>
<td></td>
<td>-3.62</td>
<td>7.98</td>
<td>$^{55}\text{Cr}+p+\alpha$</td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>


5.3.2 Angular distributions

The angular distributions of outgoing neutrons, protons, and $\alpha$-particles for selected energy intervals are displayed in Figure 5.7. No neutron measurement was done for the $^{56}\text{Fe}(d,n)$ reaction. So we employed the neutron spectrum from Ref. [66] for the same reaction with deuteron energy of 7 MeV in the analysis. Although no angular distribution is presented in that paper, it is stated that a flat distribution was observed at backward angles starting from 80°. Similarly for our $d + ^{54,58}\text{Fe}$ reactions, the outgoing neutrons exhibited a flat distribution at backward angles. Hence, neutrons emitted in the backward hemisphere were primarily from the decay of compound nuclei.

Similar trend was observed for the proton angular distributions as shown in Figure 5.7. For the $\alpha$-particle exit channel, cross sections demonstrated nearly compound emissions for all angles (see bottom graphs of Figure 5.7). The most energetic $\alpha$ particles exhibited slight forward-peaked distribution while the low energy $\alpha$ particles displayed symmetric distribution about 90°. The degree of anisotropy can be used to evaluate the spin cut-off parameter $\sigma$ at a given excitation energy. Ericson [36] gave the semi-classical expression for this relation

$$\sigma_{CS}(\theta) = A \left( 1 + \beta \cos^2 \theta \right) \quad (5.1)$$

where

$$\beta = \frac{I_b^2 I_n^2}{2(2\sigma^2)^2 \hbar^4} = \frac{\mu_b E_b \mu_n E_n}{2(2\sigma^2)^2 \hbar^4} \quad (5.2)$$

$I_b$ and $I_n$ are the orbital angular momenta in the entrance and exit channels, respectively; $\sigma$ is the spin cut-off parameter; $\mu_b$ and $\mu_n$ are the masses of the projectile and ejectile, respectively; $E_b$ and $E_n$ are their corresponding energies. Hence, the evaporation technique using (d,$\alpha$) reaction would be a good candidate for spin cut-off parameter studies. A more general treatment of the angular distribution using quantum mechanical description is given in Ref. [79].
Compound and non-compound fractions of the total cross section have been estimated using the Kalbach formula given by equation (4.1). Due to the steep slope of the proton angular distributions, the Kalbach formula produced strongly correlated values for the compound and non-compound fractions. Instead of using the symmetric component of the Kalbach formula to identify the compound fractions, we assumed that the compound
cross section dominates at backward angles. We then estimated the compound component of the total cross section by multiplying the cross section measured at backward angle by $4\pi$. The total cross section was computed from the integral of the Kalbach formula. The compound fraction was then determined based on their ratio.

Table 5.5 tabulates the values of the compound fraction. Since (d,n), (d,p), and (d,α) channels almost exhaust the total reaction cross section, taking into account all these channels in estimating the compound fraction should provide more reliable values compared to when only one exit channel is considered. The three outgoing channels were available in our data for both d + $^{54,58}$Fe reactions with $E_d = 7$ MeV. For deuteron bombarding energies of 5 and 9 MeV, only the (d,p) and (d,α) channels were considered. The compound fractions have been used as reduction factors for the calculated cross sections based on the compound model only to account for the non-compound cross sections.

Table 5.5: Estimate of compound fractions obtained from the analysis of angular distributions.

<table>
<thead>
<tr>
<th>Deuteron Energy (MeV)</th>
<th>$d + ^{54}$Fe</th>
<th>$d + ^{56}$Fe</th>
<th>$d + ^{58}$Fe</th>
</tr>
</thead>
<tbody>
<tr>
<td>5</td>
<td>84</td>
<td>76</td>
<td>73</td>
</tr>
<tr>
<td>7</td>
<td>72</td>
<td>65</td>
<td>64</td>
</tr>
<tr>
<td>9</td>
<td>68</td>
<td>57</td>
<td>39</td>
</tr>
</tbody>
</table>

Previous studies of deuteron-induced reactions on $^{56}$Fe have been done by Mishra et al. [80] and Al-Quraishi et al. [66]. These authors were able to provide estimates of the compound fractions for this reaction at different bombarding energies. From Mishra
result, the compound fractions were obtained for deuteron energies of 4, 5, 6, 7, and 8 MeV given by 51%, 58%, 70%, 78%, and 77%, respectively. Al-Quraishi, on the other hand, found the compound fractions for deuteron energies of 5 and 7 MeV to be 85% and 70%, respectively. The compound fractions from Mishra were determined only from the (d,n) channel while the compound fractions of Al-Quraishi were acquired from the angular distributions of the (d,n), (d,p), and (d,α) channels. The results of Mishra indicate that the compound component for the d + 56Fe reaction increases with increasing bombarding energy. Our data and that of Al-Quraishi agree that the compound fraction decreases with increasing energy.

5.4 Comparison with theoretical calculations

In the following sections, the particle emission cross sections measured at backward angles and the deduced level densities of the residual nuclei were compared with theoretical predictions using the EMPIRE code. The following input level density models: the Gilbert-Cameron model (GCM) [2]; the Fermi gas model (FGM) [1]; and the microscopic approach based on the Hartree-Fock-Bogoliubov method (HFBM) [3] were tested. Both the Ignatyuk (Ignatyuk sys) [13] and Iljinov (Iljinov sys) [4] level density parameter systematics given by equations (2.21) and (2.22), respectively, were used to predict the parameters for the GCM, while the von Egidy systematics (Egidy sys) [1] was applied to provide the parameters for the FGM. The best-fit parameters for the constant temperature model (CTM) and Fermi gas model (FGM) were then obtained by fitting the experimental level densities.

Before fitting the deduced level densities, data points were initially matched with the density of low-lying discrete levels. In some cases, the experimental data points were not available in the discrete level region. To obtain the absolute normalization of the experimental level density, we used model extrapolation based on the constant temperature model (CTM) function. The model extrapolation procedure is as follows. A CTM fit
was applied to the data to extract the preliminary parameters. The model function was then integrated from the ground state up to the energy $E_{max}$, where levels are known to be complete, to get the value $\rho_M$. $E_{max}$ values are available from the RIPL database [30]. Similarly, the density of known low-lying levels was also summed from the ground state to $E_{max}$ to obtain the value $\rho_L$. The experimental points were then multiplied by the scaling factor $\rho_M/\rho_L$. Using the new absolute values of the experimental level density, another fit procedure was performed to extract the final best-fit parameters. In some instances, several iterations were needed to make sure that the parameters converge.

5.4.1 $^{55}$Co residual nucleus

![Graphs showing cross sections and level density](image)

Figure 5.8: $^{54}$Fe(d,n) cross sections for $E_d = 7$ MeV (left) and extracted level density of $^{55}$Co (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.
Figure 5.8 (left) shows the cross sections for the $^{54}$Fe(d,n) reaction with $E_d = 7$ MeV. Irregularities in the experimental cross sections between 2 and 6.5 MeV are observed. According to the spectra from target-in and target-out runs, the irregularities were attributed to unclean subtraction of peaks produced from d + $^{12}$C reaction. Based on the location of the $\gamma$ peaks, the carbon atoms were deposited on the surfaces of the collimators and beam stop. The fluctuations due to the contaminants cannot be removed cleanly from the spectrum because the fraction of beam hitting each component varied with every run.

Cross sections calculated with the GCM (Iljinov sys) level density are found to agree with the experimental cross sections. Calculations according to HFBM also describe the overall magnitude of the experimental cross sections but exhibit sinusoidal fluctuation due to collective effects incorporated in the model. The cross sections predicted by GCM (Ignatyuk sys), on the other hand, display a steeper slope compared to the data points.

Experimental level density is plotted in Figure 5.8 (right). The experimental data points are reproduced by GCM (Iljinov sys) well. The level density based on von Egidy parameter systematics is also in reasonable agreement with the experimental points. The magnitude of the level density according to HFBM overestimates the data. The best-fit parameters for CTM and FGM are: $T = 1.44$ MeV, $E_0 = -0.95$ MeV; and $a = 5.74$ MeV$^{-1}$, $\delta = 0.72$ MeV, respectively.

### 5.4.2 $^{55}$Fe residual nucleus

Figure 5.9 (left) shows the cross sections for the $^{54}$Fe(d,p) reaction with $E_d = 5$, 7, and 9 MeV. Distinct population of low-lying levels is evident in the spectra for all bombarding energies. The calculated cross sections also display contribution of second-step protons found at low proton energies.

The experimental level densities derived from the cross sections for $E_d = 7$ and 9 MeV are plotted in Figure 5.9 (right). The data points below 6 MeV are not included due to
Figure 5.9: $^{54}\text{Fe}(d,p)$ cross sections for $E_d = 5$, 7, and 9 MeV (left) and extracted level density of $^{55}\text{Fe}$ (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.

contamination of direct reaction mechanism in the cross sections. Hence, the experimental level density is normalized based on the extrapolation of CTM fit to the resolved region. The predicted level density according to GCM (Iljinov sys) is in good agreement with the data while the level densities of the GCM (Ignatyuk sys) and the HFBM are found to have steeper slope compared with the experiment. The best-fit parameters for CTM and FGM are: $T = 1.47$ MeV, $E_0 = -1.70$ MeV; and $a = 5.95$ MeV$^{-1}$, $\delta = 0.54$ MeV, respectively.

5.4.3 $^{52}\text{Mn}$ residual nucleus

Figure 5.10 (left) shows the cross sections for $^{54}\text{Fe}(d,\alpha)$ reaction with $E_d = 5$, 7, and 9 MeV. The $(d,\alpha)$ cross sections terminate at a higher $\alpha$ energies than the $(d,p)$ and $(d,n)$ cross sections due to large Coulomb barrier of the $\alpha$ channel. Theoretical calculations with
Figure 5.10: $^{54}\text{Fe}(d,\alpha)$ cross sections for $E_d = 5$, 7, and 9 MeV (left) and extracted level density of $^{52}\text{Mn}$ (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.

The level density from the GCM (Iljinov sys) describe the data points from all bombarding energies. Calculations with level densities according to GCM (Ignatyuk sys) and HFBM are found to underestimate the experimental cross sections.

The experimental level density for $^{52}\text{Mn}$ derived from cross sections measured with $E_d = 9$ MeV is plotted in Figure 5.10 (right). The comparison between the density of low-lying levels and experiment shows very good agreement for excitation energies below 3.5 MeV. The level density models based on GCM (Iljinov sys) and HFBM describe the data well. The FGM level density parametrized by von Egidy systematics seems to follow similar trend as GCM (Iljinov sys) level density until 3 MeV and then they diverge beyond that energy. The best-fit parameters for CTM and FGM are: $T = 1.30$ MeV, $E_0 = -1.81$ MeV; and $a = 6.30$ MeV$^{-1}$, $\delta = -0.50$ MeV, respectively.
5.4.4 $^{57}$Co residual nucleus

Figure 5.11: $^{56}$Fe(d,n)$^{57}$Co cross sections for $E_d = 7$ MeV (left) and extracted level density of $^{57}$Co (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.

Figure 5.11 (left) shows the measured cross sections (from Ref. [66]) for the $^{56}$Fe(d,n) reaction with $E_d = 7$ MeV. The calculated cross sections with level density parameters from GCM (Iljinov sys) show good agreement with the data only in terms of shape. The absolute magnitude overestimates the experiment by a factor of about 1.3. The HFBM and GCM (Ignatyuk sys) predictions show overestimation in the low energy region and exhibit steeper slope compared to the data.

The experimental level density of $^{57}$Co is plotted in Figure 5.11 (right). Previous level density measurement of Ref. [80] extracted using the combination of evaporation and Ericson fluctuation techniques is also plotted in the figure. A very good agreement
between the present and the previous experimental data is observed over a wide range of excitation energy. The level density according to the von Egidy systematics is in reasonable agreement with our experimental level density, but relative to the previous measurement, disagreement at excitation energies above 10 MeV is observed. The level density predicted by the HFBM is found to follow the trend of the experiment but with a larger magnitude. The best-fit parameters for CTM and FGM are: $T = 1.41$ MeV, $E_0 = -1.91$ MeV; and $a = 5.93$ MeV$^{-1}$, $\delta = -0.23$ MeV, respectively.

5.4.5 $^{57}$Fe residual nucleus

Figure 5.12: $^{56}$Fe(d,p)$^{57}$Fe cross sections for $E_d = 5, 7,$ and $9$ MeV (left) and extracted level density of $^{57}$Fe (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.

Figure 5.12 (left) shows the measured cross sections for the $^{56}$Fe(d,p) reaction with $E_d = 5, 7,$ and $9$ MeV. Theoretical calculations indicate more prominent contribution of the
second-step protons found in the low energy regions for all bombarding energies. However, this is not observed in the experimental cross sections. Calculated cross sections with the level density from the GCM (Iljinov sys) describe closest the data compared to the other level density models.

The experimental level density of $^{57}$Fe is derived from the cross sections measured at $E_d = 9$ MeV. This is plotted in Figure 5.12 (right). We also included the level densities extracted from the analysis of (d,α) evaporation method reported by Voinov et al. [29], and the one measured using the γ-ray folding technique via $^3$He inelastic scattering by Schiller et al. [81]. Comparison between the experimental level densities of the present result and those from the previous measurements does not agree. This might be explained by the uncertainty of the absolute normalization procedure performed in the present study based on model extrapolation towards the low energy region. According to Ref. [81], the experimental level density of $^{57}$Fe exhibits some structure at $\sim 3$ MeV of excitation energy which cannot be described by the extrapolation technique. The best-fit parameters for CTM and FGM are: $T = 1.44$ MeV, $E_0 = -2.62$ MeV; and $a = 6.28$ MeV$^{-1}$, $\delta = -0.23$ MeV, respectively.

5.4.6 $^{54}$Mn residual nucleus

Figure 5.13 (left) shows the measured cross sections for the $^{56}$Fe(d,α) reaction with $E_d = 5$, 7, and 9 MeV. Calculated cross sections with the level density from GCM (Iljinov sys) describe both the overall shape and the absolute values of the measured cross sections well.

The experimental level density extracted from cross sections measured with $E_d = 9$ MeV is plotted in Figure 5.13 (right). The level density is normalized with the density of discrete levels. We also plotted the previous result obtained from the (p,n) reaction of Ref. [82]. Agreement between both sets of data points is observed at higher excitation energies.
Figure 5.13: $^{56}$Fe(d,α)$^{54}$Mn cross sections for $E_d = 5, 7$, and 9 MeV (left) and extracted level density of $^{54}$Mn (right): Points are experimental data; lines are calculations; histogram is density of discrete levels.

only. The structures seen in Ref. [82] are not observed in our data except for the one at 4 MeV. HFBM does not agree with the experimental data in terms of both absolute values and shape. The best-fit parameters for CTM and FGM are: $T = 1.29$ MeV, $E_0 = -2.50$ MeV; and $a = 6.42 \text{ MeV}^{-1}$, $\delta = -0.96$ MeV, respectively.

5.4.7 $^{59}$Co residual nucleus

Figure 5.14 (left) shows the measured cross section for the $^{58}$Fe(d,n) reaction with $E_d = 7$ MeV. Prediction based on the level density parameters from GCM (Iljinov sys) describes the measured cross sections well.

The experimental level density is plotted in Figure 5.14 (right). The level density based on GCM (Iljinov sys) seems to describe the data best compared to the other level density
models. The FGM level density according to the von Egidy systematics is in reasonable agreement with the overall shape of our data but slightly larger in magnitude while level densities from GCM (Ignatyuk sys) and HFBM exhibit steeper slope compared with the data. The best-fit parameters for CTM and FGM are: $T = 1.29$ MeV, $E_0 = -1.85$ MeV; and $a = 6.79$ MeV$^{-1}$, $\delta = 0.12$ MeV, respectively.

### 5.4.8 $^{56}$Mn residual nucleus

Figure 5.15 (left) shows the measured cross sections for the $^{58}$Fe(d,\(\alpha\)) reaction with $E_d = 5$, 7 and 9 MeV. Calculated cross sections based on the level density parameters from the GCM (Iljinov sys) are found to best describe the experimental cross sections compared with the other level density models.
Figure 5.15: \(^{58}\text{Fe}(d,\alpha)^{56}\text{Mn}\) cross sections for \(E_d = 5\), 7, and 9 MeV (left) and the extracted level density of \(^{56}\text{Mn}\) (right): Points are experimental data; lines are calculations; histogram is the density of discrete levels.

The experimental level density derived from the data with \(E_d = 9\) MeV is plotted in Figure 5.15 (right). The FGM, HFBM, and GCM (Iljinov) are found to agree with the present data within experimental uncertainties. The best-fit parameters for CTM and FGM are: \(T = 1.32\) MeV, \(E_0 = -3.55\) MeV; and \(a = 6.31\) MeV\(^{-1}\), \(\delta = -2.0\) MeV, respectively.

### 5.5 Discussion

The theoretical calculations based on the Gilbert-Cameron model using the Iljinov parameter systematics (GCM (Iljinov sys)) have been shown to reproduce the experimental cross sections very well for almost all reactions studied. This is supported by the close agreement of the best-fit temperatures \(T\) of the level density with the ones predicted by the systematics listed in Table 5.6. From the table, the average value of the transition energy
between the constant temperature model and the Fermi gas model is about 8.4 MeV for the Iljinov systematics. This suggests that the FGM component has almost no influence in the predicted GCM (Iljinov sys) level densities shown in Figures 5.8 - 5.15. Also, the best-fit temperatures of our experimental level densities have been observed to vary slowly with mass number. This is even more evident when temperatures from different isotopes are compared.

The GCM parameters from Ignatyuk systematics (GCM (Ignatyuk sys)) are also presented in Table 5.6. The predicted temperatures for all considered nuclei are consistently lower than those from Iljinov systematics by about 20%. It should be noted that both Ignatyuk and Iljinov systematics provide only the asymptotic level density parameter \( \tilde{a} \) as expressed in equations (2.21) and (2.22), respectively. The difference in their parametrizations arises from the experimental data that were used to obtain the coefficients of the asymptotic level density parameter. Ignatyuk [13] reported that the parametrization for the asymptotic level density parameter was optimized by fitting the experimental neutron resonance data only. Iljinov [4], on the other hand, also indicated the use of neutron resonance data, but the data required the criteria that it should have at least 5-10 observed resonances per nucleus and that the measurements for a given nucleus should be performed by at least two different groups. In addition, they also included experimental level density data measured at much higher (through Ericson fluctuations data) and low \( \rho(U \geq 2 \text{ MeV}) \) excitation energies. The latter includes discrete levels known from level scheme. Hence, it is safe to say that the parameters predicted by the Iljinov systematics provide more reliable values due to the constraints imposed by the experimental data.

Best-fit parameters for FGM have also been determined from our experimental level densities. Before the FGM best-fit parameters were obtained, our data points were normalized using the extrapolation function based on the constant temperature function. The extracted best-fit parameters have been compared with the parameters predicted by
the von Egidy systematics (FGM (Egidy sys)) as presented in Table 5.7. The table also includes the deduced parameters from previous measurements for some of the nuclei. The values of the level density parameter $a$ extracted from our experimental level densities are in close agreement with the predictions from the von Egidy systematics. As shown in Figures 5.8 - 5.15, the predictions almost resemble the correct shape of the data points. It can also be observed that some of the FGM (Egidy sys) level densities are slightly larger in magnitude compared to the experimental data points. The problem on the magnitude have been indicated in the paper of von Egidy [1]. They have shown that for some nuclei, the predicted neutron resonance spacing deviate from the experimental ones by 25-100% in the mass region presently studied ($A \sim 55$).

For some of the nuclei studied, the influence of shell and pairing effects should be apparent in the experimental level densities. Let us consider the isotopes of cobalt. The valence protons and neutrons of cobalt occupy the single-particle $f^{7/2}$ orbital and the $pf$ shell. A full occupation of the $f^{7/2}$ orbital corresponds to a closed shell with proton or neutron number equal to the magic number 28. This is the case for $^{55}$Co where 8 neutrons are in the $f^{7/2}$ orbital. On the other hand, $^{57}$Co and $^{59}$Co have 2 and 4 valence neutrons in the $pf$ shell, respectively. The latter nuclei require extra energy to break the neutron pairs. Since the level density depends on the number of quasi-particle and available single-particle orbitals, one would expect a lower level density for $^{55}$Co compared to $^{57}$Co and $^{59}$Co. Breaking the neutron pairs of $^{59}$Co would provide more degrees of freedom for further excitation. Figure 5.16 shows the superimposed plots of the nuclear level densities for the cobalt isotopes.

Figure 5.17 shows the experimental compound cross sections compared with theoretical calculations based on the CTM best-fit parameters. The calculated cross sections have already been multiplied by the reduction factors found in Table 5.5 to account for the non-compound fractions in the total cross sections. The overall shape of each
Figure 5.16: Superimposed level densities of $^{55}$Co, $^{57}$Co, and $^{59}$Co (left); solid lines are best-fit constant temperature function; broken lines are density of discrete levels.

Experimental spectrum is reproduced by the calculation well, particularly the region where first stage emission occurs. For $d + ^{54}$Fe reaction, the magnitude of the neutron and $\alpha$-particle emission cross sections are found to be slightly underestimated by the calculations while the proton cross section is reproduced well. As we move to heavier mass iron isotopes, the calculated neutron emission cross sections reproduced our data for $^{56}$Fe(d,n) reaction but slightly overestimated the cross sections for $^{58}$Fe(d,n) reaction. For the proton channel, the trend is reversed. The calculated cross sections underestimated the data for $^{56}$Fe(d,p) reaction and the underestimation becomes severe for $^{58}$Fe(d,p) reaction. It appears that the ratio of calculated emission cross sections for neutron relative to proton is well reproduced for $d + ^{54}$Fe reaction, but it gradually deviates as we move to reactions with heavier mass iron isotopes. This feature is not yet understood and thus requires further attention in future experiments.
Figure 5.17: (Points) Emission spectra of neutrons, protons, and $\alpha$-particles. Theoretical calculations using CTM are shown in solid line.
Table 5.6: GCM parameters from Iljinov and Ignatyuk level density parameter systematics. Best-fit parameters for CTM is also shown.

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Table 5.7: FGM level density parameters from other measurements.

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5.6 Pre-equilibrium and breakup contributions

The probability of non-compound emission is known to depend on the type of incident and outgoing particles involved in the reaction. For the deuteron, its low binding energy easily leads to breakup into proton and neutron. Deuteron breakup may result in considerable contamination of the compound spectra depending on the energy, type of particle detected, and the angle of measurement. In addition, the neutron and proton fragments may initiate further reactions. This is the reason why in most evaporation experiments, if complex projectiles are used, particles with large binding energy (e.g. $\alpha$ particles) are preferred.

For deuteron, fragments from breakup have been shown to favor the forward angle direction. They have been observed to produce a narrow peak on top of the evaporation spectrum as shown in Refs. [86, 87]. Kalbach [88] proposed a phenomenological model to describe this peak using a Gaussian function with width

$$\Gamma = 1.15 + 0.12E_a - \frac{A_A}{140}$$  \hspace{1cm} (5.3)

and centroid energy (in MeV)

$$\epsilon_0 = \frac{A_b}{A_a} \left( \epsilon_a - B_{a,b} - \frac{Z_u Z_A}{9.5} \right) + \frac{Z_b Z_B}{9.5}$$  \hspace{1cm} (5.4)

where $\epsilon_a$ represents the channel energy, $B_{a,b}$ is the binding energy, and $Z_{b,B(a_A)}$ and $A_{b,B(a, A)}$ are the atomic numbers and mass numbers of the ejectile and residual nucleus (projectile and target nucleus), respectively. The phenomenological differential breakup cross section is given by

$$\frac{d\sigma_{BU}}{dE} = \sigma_{BU} \frac{1}{\Gamma \sqrt{2\pi}} \exp \left[ -\frac{(\epsilon_0 - E)^2}{\Gamma^2} \right]$$  \hspace{1cm} (5.5)

where

$$\sigma_{BU} = K \left( A_A^{1/3} + 0.8 \right)^2 \frac{1}{1 + \exp \left( \frac{13 - E_a}{6} \right)}$$  \hspace{1cm} (5.6)
$K$ is 18 for outgoing neutrons and 21 for outgoing protons [32], $E_a$ and $E_b$ are the laboratory energies of the projectile and ejectile, respectively.

The TALYS reaction code [32] has been utilized to estimate the contribution from breakup process where the Kalbach’s breakup model is integrated with the pre-equilibrium and other non-compound reaction mechanisms. The TALYS non-compound cross section is given by the sum of cross sections due to the following reaction mechanisms: the pre-equilibrium mechanism (PE) calculated with the two-component exciton model (EM), the nucleon transfer (NT) which includes stripping and pick-up, knock-out (KO) and breakup (BU) mechanisms expressed as follows

$$
\frac{d\sigma^{NC}}{dE} = \frac{d\sigma^{EM}}{dE} + \frac{d\sigma^{NT}}{dE} + \frac{d\sigma^{KO}}{dE} + \frac{d\sigma^{BU}}{dE}.
$$

(5.7)

The angular distribution of the non-compound double-differential cross section has the form given by equation (4.1). The pre-equilibrium component utilized in the calculation is the two-component exciton model.

Figure 5.18 displays the double differential cross sections of $^{65}$Cu(d,n) and $^{54}$Fe(d,p) reactions in the CM system compared with TALYS calculations with level densities based on the best-fit parameters given in Table 5.6. Stripping reaction is presumably the evident as a selective population of some states from the $^{54}$Fe(d,p) reaction. The decomposition of the predicted double differential cross section according to different reaction mechanisms is also displayed in the figure. The bumps (shown in red line) centered at about 6 MeV and 2 MeV for the proton and neutron spectra, respectively, are primarily due to breakup process.

In the calculations, the angular distribution for the compound cross section is assumed to be isotropic. The experimental proton spectrum measured at backward angles (e.g. 157.5°) has been found to have the least contribution of non-compound reactions except for the highest proton energies where direct reactions may have already contributed (see Figures 5.8-5.15). It seems that there is no structure present in the experimental data that would
represent the breakup contribution above 97.5° even though the calculations suggest so. This structure starts to become visible in the proton spectra only for angles below 67.5°. The magnitude of the bump continues to grow as the angle decreases. For neutron spectra, the peak for the breakup contribution is difficult to observe since it lies below the low neutron energy cut-off. For the available data, no enhancement has been observed even for the spectrum measured at 40°. However, enhancement of the cross sections for the most energetic outgoing neutrons is prominent.

Angle-integrated cross sections are best described by pre-equilibrium calculations using the exciton models. This is tested with our data using two-exciton model of TALYS and the PCROSS module of EMPIRE. The PCROSS module considers the one-component exciton model and, in addition, it incorporates the Iwamoto-Harada model for the pre-equilibrium emission of clusters of particles. TALYS offers both the two-component and one-component exciton models. Both were used to calculate the cross sections for the 64Cu(d,n) reaction and no significant difference between their results was observed.

The angle-integrated cross sections for 64Cu(d,n) with bombarding energies of 6 and 7.5 MeV are shown in Figure 5.19. Noticeable increase in the cross sections for high neutron energies relative to the compound component is apparent for both incident deuteron energies, although experimental data points from 7.5 MeV are clearly more enhanced relative to the data points from 6 MeV. This is expected as pre-equilibrium reaction mechanism plays a more important role at higher bombarding energies. The compound component shown in Figure 5.19 is calculated according to the best-fit parameters of the experimental level density for 64Zn. When the experimental points were compared with the pre-equilibrium models, neither EMPIRE nor TALYS was able to reproduce the data points. As shown in Figure 5.19, EMPIRE overestimates the angle-integrated cross sections while TALYS underestimates it. The EMPIRE results can be improved by adjusting the mean-free path parameter of the PCROSS module from its default value of 1.3 to 0.5.
Figure 5.18: Double differential cross sections for $^{54}$Fe(d,p) with $E_d = 9$ MeV (left) and $^{65}$Cu(d,n) with $E_d = 6$ MeV (right). Experimental data (points) are compared with calculations shown in solid lines with compound (green), non-compound (red) and total (blue) contributions.
Figure 5.19: Angle integrated cross sections of the $^{65}$Cu(d,n) reaction measured at deuteron energies of 6 MeV (left) and 7.5 MeV (right). Points are experimental data. Lines are calculations with pre-equilibrium exciton (PE) and compound (comp) models. The red lines are EMPIRE calculations using the default parameters, while the black lines are calculations with adjusted parameters.
6 SUMMARY AND CONCLUSIONS

The double differential cross sections for the $^{63}$Cu(d,Xn), $^{65}$Cu(d,Xn), $^{54}$Fe(d,Xn), $^{58}$Fe(d,Xn), $^{54}$Fe(d,Xp), $^{56}$Fe(d,Xp), $^{54}$Fe(d,Xα), $^{56}$Fe(d,Xα), $^{58}$Fe(d,Xα), $^{54}$Fe(d,d)$^{54}$Fe, $^{56}$Fe(d,d)$^{56}$Fe, and $^{58}$Fe(d,d)$^{58}$Fe reactions have been measured in the angular range between 20° and 160° with deuteron energies of 5, 6, 7, 7.5, and 9 MeV. The outgoing neutron spectra were obtained using the time-of-flight technique with the swinger facility of the Edwards Accelerator Laboratory while the outgoing charged-particle spectra were directly determined from both the time-of-flight and energy deposited in the silicon detector of the charged-particle spectrometer setup.

Measured differential elastic cross sections and their angular distributions have been compared to optical model calculations using the parametrizations of Refs. [69, 74–76]. The calculated elastic angular distributions are shown to underestimate experimental data points at backward angles by about 4-17%. Our experimental data are best described by the elastic angular distributions according to the deuteron OMP parameters of Lohr et al. [76]. This deuteron OMP parametrization has been used to calculate the transmission coefficients for the entrance channels of the reactions studied.

The cross sections measured at backward hemisphere have been found to have the least contamination from non-compound emissions. This is demonstrated by the flattening of the angular distribution for angles beyond 90° of the observed outgoing particles. The cross sections obtained in this angular range have been compared with theoretical calculations performed in the framework of the compound reaction model through the EMPIRE code. Different input level density models were tested including the Fermi gas model, the Gilbert-Cameron model, and the microscopic level density model based on the Hartree-Fock-Bogoliubov method. Among the phenomenological models, the Gilbert-Cameron model with the Ilijinov parameter systematics of Ref. [4] is consistent with the experimental results except for the $^{64}$Zn and $^{66}$Zn isotopes. The slopes of the experimental level densities
for these isotopes are found to be the same within 2% whereas the model predicted a
difference of about 16%. The microscopic level density model based on the Hartree-Fock-
Bogoliubov method [3] has also predicted the cross sections with a precision comparable to
the phenomenological models, but the predictions frequently overestimate the experimental
data points. Also, the predicted oscillation of the level density that resulted from the
inclusion of collective effects in the model is not supported by the experimental data for
some nuclei.

Nuclear level densities for $^{64}$Zn, $^{66}$Zn, $^{55}$Fe, $^{57}$Fe, $^{55}$Co, $^{57}$Co, $^{59}$Co, $^{52}$Mn, $^{54}$Mn, and
$^{56}$Mn have been deduced from the analysis of the experimental compound cross sections
measured at backward angles. Phenomenological GCM (see equations (2.32) and (2.33))
based on Iljinov [4] (GCM (Iljinov sys)) and Ignatyuk [13] (GCM (Ignatyuk sys)) systemat-
ics, FGM (see equation (2.14)) according to von Egidy [1] (FGM (Egidy sys)) systematics,
and the microscopic HFBM were compared to the experimental results. The GCM (Iljinov
sys) of equation (2.22) was found to describe all experimental level densities consistently
well. Experimental values of the temperature that describe the slope of the level density
function shows a good consistency for neighboring nuclei. The difference in level densities
among different isotopes is related to pairing and shell effects which are reflected by the
general energy shift $E_0$ rather than by the temperature $T$ itself. The best-fit level density
parameters $a$ of the Fermi gas model are in reasonable agreement with those predicted by
the von Egidy systematics. Our data have also agreed with level densities obtained from
previous experiments which either employed different reaction channels or entirely differ-
ent experimental technique. This is another indication that the presently measured level
densities have been deduced from compound reactions. This is a good indication that the
particle spectra from deuteron-induced reactions measured at backward angles are due pri-
marily to compound mechanisms at these bombarding energies.
The presence of direct and pre-equilibrium non-compound contributions are manifested in the particle spectra measured at forward angles. This is evident as forward-peaked asymmetry in the angular distributions observed for all reactions considered. The amplitude of the asymmetry becomes larger for higher outgoing particle energies indicating the increasing contribution of non-compound reaction components. The compound fraction for each reaction studied has been estimated from the angular distribution of the outgoing particles. It has been found that the compound fraction decreases with increasing bombarding energy. The deuteron breakup process has also been shown to have a negligible contribution in the cross section measurements at backward angles. The main mechanisms of the non-compound components are found to be from direct (presumably from pick-up or stripping reactions) and pre-equilibrium mechanisms, which enhance the cross sections in the high energy region of the outgoing particle spectra. The pre-equilibrium component can be well reproduced by the exciton model of reactions, however, model parameters need to be adjusted.

6.1 Outlook

In this study, we have shown that deuterons can be used as a projectile to obtain the level densities of medium to heavy mass nuclei. The angular distributions of emitted particles exhibited the dominance of compound reactions at backward angles. However, some of our experimental data do not agree with previous measurements of Refs. [29, 81]. A noticeable discrepancy in the magnitude of the calculated cross sections relative to our data for $^{56,58}\text{Fe(d,p)}$ reactions is also observed. To quantitatively estimate the uncertainty due to non-compound processes, it is important to perform experiments using different entrance channels that would lead to the same compound nucleus and excitation energy.

The quest for studying the nuclear level density far from the valley of stability has recently become a priority. The progress in the theoretical aspects of the nuclear level
density needs experimental verification which is lacking as of today. The emerging radioactive ion beam facilities would be an important component for such studies.
REFERENCES


APPENDIX A: EXPERIMENTAL DOUBLE DIFFERENTIAL CROSS SECTIONS

A.1 Cross sections of $^{63}$Cu(d,Xn) reactions

Figure A.1: Double differential cross sections of $^{63}$Cu(d,Xn) reactions for $E_d = 6$ MeV (left) and $E_d = 7.5$ MeV (right). The arrows indicate the energies where neutrons from second-step emissions start to contribute to the cross sections.
A.2 Cross sections of $^{65}$Cu(d,Xn) reactions

Figure A.2: Double differential cross sections of $^{65}$Cu(d,Xn) reactions for $E_d = 6$ MeV (left) and $E_d = 7.5$ MeV (right). The arrows indicate the energies where neutrons from second-step emissions start to contribute to the cross sections.
A.3 Cross sections of $^{54}$Fe(d,Xp) reactions

Figure A.3: Double differential cross sections of $^{54}$Fe(d,Xp) reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right). The arrows indicate the energies where protons from second-step emissions start to contribute to the cross sections.
A.4 Cross sections of $^{56}$Fe(d,Xp) reactions

Figure A.4: Double differential cross sections of $^{56}$Fe(d,Xp) reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right). The arrows indicate the energies where protons from second-step emissions start to contribute to the cross sections.
A.5 Cross sections of $^{58}$Fe(d,Xp) reactions

Figure A.5: Double differential cross sections of $^{58}$Fe(d,Xp) reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right). The arrows indicate the energies where protons from second-step emissions start to contribute to the cross sections.
A.6 Cross sections of $^{54}$Fe(d,$\alpha$) reactions

Figure A.6: Double differential cross sections of $^{54}$Fe(d,$\alpha$) reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right).
A.7 Cross sections of $^{56}\text{Fe}(d, X\alpha)$ reactions

Figure A.7: Double differential cross sections of $^{56}\text{Fe}(d, X\alpha)$ reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right).
A.8 Cross sections of $^{58}\text{Fe}(d,X\alpha)$ reactions

Figure A.8: Double differential cross sections of $^{58}\text{Fe}(d,X\alpha)$ reactions for $E_d = 5$ MeV (left), $E_d = 7$ MeV (middle), and $E_d = 9$ MeV (right).
A.9 Cross sections of $^{54}$Fe(d,Xn) and $^{58}$Fe(d,Xn) reactions

Figure A.9: Double differential cross sections of $^{54}$Fe(d,Xn) (left) and $^{58}$Fe(d,Xn) reactions (right) for $E_d = 7$ MeV.
### APPENDIX B: DEUTERON OPTICAL MODEL PARAMETERS

**B.1 Optical model parameters for d scattering on $^{54,56,58}$Fe nuclei with $E_d = 5$ MeV**

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<td>$^{58}$Fe</td>
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Table B.1: Deuteron optical model parameters for $^{54}$Fe, $^{56}$Fe, and $^{58}$Fe with 5 MeV incident energy.

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<td>0.0</td>
<td>Haixia [69]</td>
</tr>
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<td>0.81</td>
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B.2 Optical model parameters for d scattering on $^{54,56,58}$Fe nuclei with $E_d = 7$ MeV

Table B.2: Deuteron optical model parameters for $^{54,56,58}$Fe nuclei with $E_d = 7$ MeV incident energy.

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<th>$V_{SO}$</th>
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<td>Yinlu [74]</td>
</tr>
<tr>
<td>$^{58}$Fe</td>
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<th>$r_{WSO}$</th>
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<td>0.0</td>
<td>Haixia [69]</td>
</tr>
<tr>
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<td>1.17</td>
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<td>1.56</td>
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<td>0.64</td>
<td>1.23</td>
<td>0.81</td>
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<td>0.0</td>
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<td>0.69</td>
<td>0.75</td>
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B.3 Optical model parameters for $d$ scattering $^{54,56,58}\text{Fe}$ nuclei with $E_d = 9$ MeV

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<tbody>
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Table B.3: Deuteron optical model parameters for $^{54}\text{Fe}$, $^{56}\text{Fe}$, and $^{58}\text{Fe}$ with 9 MeV incident energy.

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<tr>
<th>$r_C$</th>
<th>$r_V$</th>
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<th>$a_{WV}$</th>
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<th>$a_{WS}$</th>
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<th>$r_{WSO}$</th>
<th>$a_{WSO}$</th>
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<tbody>
<tr>
<td>1.30</td>
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<td>Haixia [69]</td>
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