Evidence for a Resonance at $E_p = 127$ keV in the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ Reaction

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Dissertation submitted in partial fulfillment of
the requirements for the degree of
Doctor of Philosophy in the Department of
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Abstract

(Physics – TUNL)

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Evidence for a Resonance at E_p = 127 keV in the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ Reaction.

A new resonance at $E_p = 127$ keV has been discovered in the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction which has which has a large (in some cases, more than an order of magnitude) effect on the rate of hydrogen burning in hot stars. The contribution of this previously undiscovered resonance to the rate of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction, which is the limiting reaction in CNO cycle fusion for most stars, is of critical importance in models of stellar evolution and calculations of the relative abundances of the CNO isotopes. The measurement was made possible by the development of a new target chamber which could be biased (to achieve the necessary beam energies), allowed γ-ray detectors to be placed close to the target (to achieve the necessary detector solid angle), and held a target of frozen, deuterated ammonia (ND_3). Yields of γ-rays from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ and $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ reactions indicate that the resonance strength, $\omega\gamma, is 4.5 \times 10^{-6}$ eV. Other parameters of the new resonance, such as the spectroscopic factor, have been derived for various possible values of the angular momentum (J) and parity (π) of this new state. This resonance corresponds to a previously unknown excited state in $^{15}\mathrm{O}$ at 7.414 MeV. Previous experimental stud-

ies which do not observe the resonance are re-examined in light of the present work, and two shell-model calculations which predict a "missing" energy level in ^{15}O are discussed. The techniques developed for the relative S-factor measurement are also applied to measure the ratio of the astrophysical S-factors of the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions at low energies in a single measurement, thus reducing systematic error in the experimental S-factor ratio. The carbon proton capture S-factor ratio is shown to be in agreement with the ratio calculated from the separate S-factor measurements of previous studies.

Acknowledgements

The data displayed herein could not have occurred without the assistance of a large set of people, particularly within the capture group. John Dunham, Richard O'Quinn, Paul Carter, and Sidney Edwards went the extra mile to provide technical assistance for this project on several occasions. In addition, a "cast of thousands" were involved in the many attempts (both successes and failures) to obtain the data which have led us to where we are now. Graduate students like Eric Wulf, Eric Schreiber, and Shane Canon, as well as postdocs like Sally Gaff and John "Ned" Kelley, faced up admirably to the daunting task of teaching a new generation of students enough experimental techniques to see this project through. Konstantin Sabourov has always been there, as we struggled together on many of the same experimental challenges. Konstantin has also, with the help of Rob Macri, provided valuable moral support during the years of pursuit of this final result. The younger generation of students, Brent Perdue, Brendan Crowley, Kaveri Joshi, and Amanda Sabourov helped tirelessly with the constant building and dismantling of lead shields and experimental setups. New postdocs Anton Tonchev and Mohammed Ahmed kept me rational during the final analysis of the data and through a few of the wilder experiments, and Caroline Nesaraja took on the responsibility for using our new techniques for a new experiment. Our visiting professors Mark Spraker and Kara Keeter added their own good advice, experience, and data-taking assistance to the project.

Presentation of the results contained within this dissertation has proven particularly difficult, because they represent a controversial addition to astrophysics. The

entire research group has stood behind this research, particularly the pillars of guidance and experience provided by Dick Prior and Henry Weller. Their assistance in both designing experiments and promoting the results has been invaluable.

I must finally acknowledge the large background of support provided to me by my family, friends, and Julie Roy. Their presence and support have been especially helpful and motivating during the writing of this dissertation.

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Chapter 1

Introduction

1.1 Overview

The cross section for the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction between E_p = 10-160 keV is of particular importance because it controls the rate of the primary mode of energy production in stellar plasmas approximately 30% hotter than those found in the Sun [Cla83]. From an astrophysical standpoint, we are motivated to study the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction at low energies because it is crucial to our understanding of stellar evolution models, and our understanding of such models is directly linked to our understanding of the origin of most elements heavier than helium in our universe. That motivation is only furthered by the fact that the only existing set of experimental data on this reaction below E_p = 200 keV [Hes57] is over 45 years old, and that data set (from E_p = 100-135 keV) provides only the total cross section via activation analysis. It contains no information on various modes of capture and disagrees by a factor of 2 with recent extrapolations of higher-energy data to low energies [Ang01].

The availability of more accurate measurement techniques (via precise observation of γ-rays from the ^{14}N(p,γ)^{15}O reaction) and the recent development of several new techniques of measurement at TUNL have led to the discovery of a new resonance in the ^{14}N(p,γ)^{15}O reaction at 127 keV. Because of the existence of this new resonance has large implications for astrophysics, which are detailed below and further explored in Chapter 6, it should be considered carefully in stellar models. Also, because the existence of a previously undiscovered resonance implies the discovery of a new energy level in ^{15}O, it provides a new piece of information to add to our understanding of nuclear structure.

1.2 Nuclear Astrophysics

The ^{14}N(p,γ)^{15}O reaction is the slowest reaction in the hydrogen-burning carbon-nitrogen-oxygen (CNO) cycle in most stars [Cla83]. The process by which this cycle converts four protons into a helium nucleus is shown in Figure 1.1. It consists of three radiative capture reactions (diagonal arrows in Figure 1.1), two beta decays (the vertical arrows), and a (p,α) reaction (the long return arrow) which converts the ^{15}N back into ^{12}C. The process is therefore catalytic, continuing until either the hydrogen is used up or weaker breakout reactions break the cycle. An example of such a reaction is ^{15}N(p,γ)^{16}O.

The reaction rates in a real stellar plasma also depend on the density and composition of the plasma in which they are taking place. The density, composition, and temperature of stellar material have great ranges within any particular star, depending on their location within the star. Obviously, the problem of a particular stellar model is beyond the scope of the present work. However, the impact of the current

Figure 1.1: Diagram of the CNO cycle. Diagonal arrows represent proton radiative capture reactions. Vertical arrows represent β+decays. The return arrow of the cycle (connecting ^{15}N to ^{12}C) represents the ^{15}N(p,α)^{1C}C reaction.

study of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction on the general reaction rate per particle pair (as a function of the plasma temperature) is a straightforward calculation, and will be presented in Chapter 6. When a detailed stellar model including the effects of this resonance is obtained, it may solve the mystery of the anomalous isotopic ratio of $^{15}\mathrm{N}/^{14}\mathrm{N}$ in micrometeorites [Her01a]. Such meteorites are believed to originate in Asymptotic Giant Branch (AGB) stars [Her01]. At the plasma temperatures typical of such stars, the contribution of the 127 keV resonance to the reaction rate is greatest (see Chapter 6).

1.2.1 Astrophyscal S-factor

The cross section for a radiative capture reaction at low energies is usually parameterized as an astrophysical S-factor, $S(E_{cm})$, which factors out two known terms

which vary rapidly as a function of the incident proton energy [Ang99]. The motivation for this parameterization is experimental. Stars last billions of years, and the nuclear reactions which power them proceed at an exceedingly slow rate due to the fact that an incident proton or alpha particle (generally the reacting particles because of their low charges) must tunnel through the Coulomb barrier in order to interact with the target nucleus via the strong force. In order to study these reactions in the laboratory, it is almost invariably necessary to extrapolate data at the energies where the reaction can be measured experimentally to lower energies characteristic of colliding particles in a stellar plasma (given by the Boltzmann distribution). The cross section and S-factor are related by:

$$\sigma(E_{cm}) = \frac{S(E_{cm})e^{-2\pi\eta}}{E_{cm}}. \tag{1.1}$$

The center of mass energy (E_{cm}) is used in Equation 1.1, and η (the derivation of η is described below) is the Sommerfeld parameter. The nuclear interaction is contained within S(E), and the 1/E term in Equation 1.1 is a geometrical factor of the interaction due to the wavelength of the incident proton. The exponential term is proportional to the probability that the incident proton will tunnel through the thick Coulomb barrier. This probability for transmission through the barrier is derived using the Wentzel-Kramers-Brillouin (WKB) approximation[Gri95]. A schematic of the application of the WKB approximation to the case of proton capture is shown in Figure 1.2.

The wavefunction of the incident proton, ψ, is given in the WKB approximation by:

$$\psi \approx \frac{C}{\sqrt{p(x)}} exp(\pm\frac{\imath}{h}\int p(x)dx), \tag{1.2}$$

where C is a complex normalization constant. $p(x)$ is the complex proton momentum

Figure 1.2: Diagram of the wavefunction (solid line) of an incident proton of energy E_p tunneling through the Coulomb potential barrier of a one-dimensional target nucleus in one dimension, superimposed on a plot of the interaction potential (dashed line). The dash-dotted line corresponds to the kinetic energy of the incident proton, E_p.

defined by $p(x) = \sqrt{2m_p(E_p - V(x)}$, where m_p is the proton mass, E_p is the proton energy, and $V(x)$ is the potential. The amplitude of the wavefunction for a proton tunneling through the Coulomb barrier of the target nucleus is given by solving the integral in Equation 1.2 through the Coulomb barrier. The probability, P, of the proton tunneling through the barrier is given by:

$$P \propto e^{-2\gamma}, \tag{1.3}$$

where γ is given by:

$$\gamma \approx K_1 \frac{Z}{\sqrt{E_p}} - K_2, \tag{1.4}$$

where Z is the charge of the target nucleus and the Ks are constants which will be discussed below. Equation 1.4 assumes that the radius of the nuclear surface is much less than the radius of classical turning for the proton in the Coulomb field. For a proton of $E_p = 127$ keV in the repulsive field of a nitrogen nucleus, the radius of classical turning is $\approx$70 fm. The nuclear radius, as given in Chapter 4, is < 5 fm, making the approximation relatively good.

The above steps (Equation 1.2-Equation 1.4) in the application of the WKB approximation are intended to help us derive the Sommerfeld parameter. The derivation is analogous to the derivation of Gamow's theory of alpha decay [Gri95]. K_2 is a constant which depends on the nuclear radius of the target nucleus. Its addition to the exponent will cause the energy-dependent part to be multiplied by a constant, which can be simply considered to be part of the measured astrophysical S-factor. The actual value of this constant is somewhat irrelevant, since the purpose of conversion to an S-factor is simply to remove known strong energy dependences from the cross section and hence make extrapolation of the cross section to lower energies more tractable. K_1, which is very relevant to the energy dependence of the cross section,

is given by [Gri95]:

$$K_1 = (\frac{e^2}{4\pi\varepsilon_0})\frac{\pi Z}{\hbar}\sqrt{\frac{m}{2E_p}} \qquad (1.5)$$
$$= 2\pi\eta = 0.9895Z(\mu/E_{cm})^{1/2}$$

In Equation 1.5, E_p is converted to E_{cm} because E_{cm} is typically used in calculations of thermonuclear reaction rates, and E_{cm} is given in units of MeV. μ is the reduced mass of the system in AMU. In more general reactions, the Z in Equation 1.5 is replaced by (Z_1Z_2), where Z_1 and Z_2 are the respective charges of the incident and target particles.

1.2.2 Energy levels in ^{15}O.

The existence of a new resonance must be considered in any nuclear theory which can predict nuclear energy levels. The exact form of the strong nuclear force remains one of the great unknowns in physics. Recent attempts to create realistic nucleon-nucleon potentials, typically employing many parameters which have been fit to nucleon-nucleon scattering data [Wir95], have still not been able to reproduce all experimental data for even systems of only three nucleons [Fri90]. Theorists are now working on including three-body forces, where three nucleons interact simultaneously in a manner different than two nucleon systems do, in models of three-body and four-body systems. For nuclei with A > 4, models such as the shell model [Kra88] are still the best hope for a description of a 15-body system such as the ^{15}O nucleus. Any model of the atomic nucleus must aim to reproduce as much of the observed behavior of nuclei as possible. Modern shell model computation using programs such as OXBASH have errors of ≈0.5 MeV in their predictions of the energy levels of light

nuclei [Bro00]. These shell-model programs also sometimes predict the existence energy levels which are not observed experimentally [Bro00, Wol90]. This is not always necessarily a failure of the models. Just as experiments are not complete without models to describe their data, nuclear models are not complete without sufficient experimental data to test their predictive powers. When a study of the feasibility of experimentally measuring the cross section of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction raised the possibility of the existence of a new energy level in $^{15}\mathrm{O}$, more importance was placed on the pursuit of a thorough experimental investigation of this reaction.

The energy level diagram for the residual nucleus in the reaction of interest, $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$, is shown in Figure 1.3 [Ajz91]. The direct capture (DC) picture of this reaction at low energies (below the $E_p = 278$ keV resonance, which corresponds to the 7.556 MeV state in $^{15}\mathrm{O}$) is that the incident proton is captured by the $^{14}\mathrm{N}$ nucleus by the emission of a γ-ray in a one-step process [Wel80]. The final state of the capture reaction may be either the ground state or one of the six bound states of $^{15}\mathrm{O}$. If the capture is to one of the subthreshold bound states, this state subsequently decays to the ground state by γ-ray emission.

Resonance contributions to the DC transition amplitude may be added to the DC model as discussed in Chapter 5. The contributions of resonances to the capture cross section generally take Breit-Wigner-like forms, causing the cross section to vary by orders of magnitude at energies near the resonance energies. The presence of a narrow and previously undiscovered resonance below E_p=278 keV would dramatically alter our understanding of the cross section for the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction in this energy range. The addition of a previously unobserved excited state in $^{15}\mathrm{O}$ would also represent a new energy level for theorists to reproduce with their models. This extra

Figure 1.3: Level diagram for ^{15}O.

demand on their models may eventually lead to a greater understanding of the models and of nuclear energy levels. In addition, a new experimental search for an analog state in ^{15}N could be motivated by such a discovery.

1.3 Development of the present experiment.

The experience of the low-energy radiative capture group at TUNL with frozen water targets [Sch95, Wul99] eventually led to the design of an advanced frozen target chamber. The new chamber was originally intended to be biased to a high negative voltage, in order to study the D(d,γ)^{3}He reaction at higher energies than were accessible via the ion source alone (80 keV). This was also done to maintain higher beam currents than existing acceleration systems at TUNL were capable of sustaining ($\approx$ 3μA). The chamber was also designed to be only 5 cm wide, with an outer shield of 1/4" thick plastic to separate detectors from the high voltage chamber. This allowed larger solid angles to be subtended by the detectors. The higher beam energies and solid angles attainable with this new target chamber led to test experiments involving radiative capture on frozen targets involving higher-Z nuclei, specifically carbon and nitrogen isotopes.

The development of the test experiment with the target of frozen, deuterated ammonia (ND_3) led to the development of the technique of measuring relative cross sections. An unexpected data peak from the test measurement and similar data from subsequent measurements, along with the techniques used to obtain and analyze the data, are detailed in the following sections. As an application test of the technique of measuring relative cross sections used in the present experiment, the ratio of cross sections for proton capture by ^{12}C and ^{13}C using a target of frozen benzene with a

known isotopic ratio of the carbon component was measured [Nes01]. The results and analysis of the carbon experiment are a straightforward variation on the primary experiment (using the nitrogen target), and are included for comparison to the present work as Appendix A.

1.4 Scope and goals of the experiment.

The original goals of this experiment were threefold. The first goal was to develop a technique of measuring an astrophysical S-factor with a biased target chamber by means of a relative cross section measurement. This first goal remains unchanged. The second goal was to measure the magnitude and slope of the astrophysical S-factor, and possibly the analyzing power (which will be described in Chapter 5), for the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction and to use this data to improve the extrapolation of the S-factor to low energies. This goal has been revised in light of the data we obtained, and became the new goal of providing a consistent and reasonable interpretation of the data. The key to this consistent picture is the existence of a narrow resonance at $E_p = 127$ keV. The evidence for the existence of this resonance shall be laid out in this dissertation, and the resonance parameters shall be quantified. The third goal of the present work is to understand the impact of the results on the astrophysically important $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction. The impact of our data on the rate of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction in a stellar plasma will be laid out in Chapter 6.

Chapter 2

Experimental Methods

The $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction has been studied at the Triangle Universities Nuclear Laboratory (TUNL) using proton beams which were stopped in thick targets of frozen ammonia. Beams of polarized and unpolarized protons with 80 keV of energy were created by the TUNL Atomic Beam Polarized Ion Source (ABPIS)[Cle90]. The energy of these protons was increased by means of a negatively biased target chamber. The γ-rays from this reaction were detected by large High Purity Germanium (HPGe) detectors. Also detected were γ-rays from a reference reaction, $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$, which is well studied in the 100 keV energy range [Sch95]. Signals from the HPGe detectors were processed by fast-electronic modules and sent to a computer for conversion into digital data.

2.1 Proton Beams

To create a proton beam, the TUNL Atomic Beam Polarized Ion Source (ABPIS) must first bleed hydrogen gas into the ion source through a very small hole in a copper nozzle. The neutral hydrogen atoms then drift from the nozzle into the Electron Cyclotron Resonance (ECR) ionizer, shown in Figure 2.1.

Figure 2.1: The ECR ionizer in the ABPIS. Electrons are accelerated by microwaves with a frequency equal to the cyclotron resonance frequency of the electrons in the magnetic field. Electrons colliding with gas in the ionizer create a source of plasma.

Electrons from nitrogen gas which is constantly fed into the ionizer are confined by a combination of a sextupole magnet and an external magnetic solenoid, so that the magnetic field increases in any direction they travel from the center. This configuration of magnetic field is referred to as a "minimum-B structure." The electrons are

excited by microwaves at the cyclotron frequency for electrons in the central magnetic field. Collision with the energetic electrons in the resulting nitrogen plasma ionizes the incoming hydrogen atoms.

These drifting ions are subsequently extracted and focused by electrostatic lenses into the cesium oven, where high-Z cesium atoms reduce the spreading of the beam due to the mutual repulsion of the ions (space-charge effect). The entire ion source is biased to +80 kV, and the positive ions are accelerated to ground potential by the electrostatic field inside an accelerator tube. This tube isolates the ion source from the ground-potential beamline. The exiting beam is focused by an electrostatic quadrupole at the end of the accelerator tube.

In the measurement utilizing the TUNL mini-tandem, which was done to investigate the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction with proton energies in excess of the $E_p = 278$ keV resonance, extra acceleration of the ions by the mini-tandem requires negative ions. These ions can be accelerated towards a biased stripping foil, stripped of their electrons, and re-accelerated back to ground potential. Negative ions are created with the ABPIS by increasing the temperature of the cesium oven, thereby raising the density of cesium atoms which the protons from the ECR ionizer pass through. The cesium acts as a charge-exchange gas, allowing ~ 10% of the protons to acquire two extra electrons before acceleration. In the case of negative ion production, the polarity of the ion source bias is reversed to accelerate the negative ions.

2.2 Beam Transport

Measurements with positive beam were carried out at the low-energy capture (LECAP) beam line, shown in Figure 2.2. The protons from the ABPIS were steered

60° to the right by a dipole inflection magnet, and continued down the LECAP line. Two electrostatic quadrupoles were used to change the beam focus in this beamline. Two magnetic steering elements were used to adjust the direction of the beam. A pair of adjustable collimation slits and a 1 cm diameter collimator defined the beam shape and direction.

Figure 2.2: Overhead view of the LECAP beam transport system.

In some of the present experiments, an electrostatic steerer in the LECAP beamline was connected to a KEPCO OPS 5000 fast operational-amplifier power supply.

The supply was switched on for 200-250 ms to steer the beam away from the target and into the collimator, and then off for 50-100 ms to allow the beam to reach the target. This cycling procedure had the effect of allowing a higher instantaneous beam current (about 30 μA) while maintaining the maximum average beam current on the ice targets which would allow our vacuum system to remain under 4 μTorr (about 8 μA). The exact cycle times for a particular measurement were determined by the achievable vacuum during the measurement. Data taken during the time when the beam was steered away from the target was sorted into a background spectrum, allowing us to acquire beam-on data and background data during the same measurement time period.

Surface contamination is of critical importance in any experiment involving low energy charged particle beams. Vacuum in the beamline was maintained by means of three turbo-molecular drag pumps (turbo pumps) backed by single-stage mechanical pumps. The vacuum was further enhanced by two liquid nitrogen filled cold traps on the beamline and the liquid nitrogen dewar used to cool the target. Vacuums below 1 μTorr were achieved with beam on a blank aluminum target. With beam on the ice targets the vacuum was typically maintained below 4 μTorr, with evaporating target material comprising the bulk of the rarefied gas present in the vacuum system.

2.3 Mini-tandem

Measurements with negative beam were carried out with the low-energy beam accelerator facility (LEBAF) shown in Figure 2.3. H^- ions from the ABPIS were steered right by 45° down the LEBAF line by a dipole inflection magnet. Two magnetic quadrupoles, one before and one after the mini-tandem accelerator, were used

to adjust the focus of the ion beam. Three magnetic steering elements were used to adjust the beam direction. Negative ions in the mini-tandem acceleration section of the beamline were accelerated in an accelerator tube towards a central chamber, which was maintained at 0-200 kV by a high voltage power supply. In the chamber, both electrons were stripped from the proton by a 2 μg/cm^2 thick carbon stripping foil. The protons continued through the chamber, and were then accelerated away from the chamber towards ground potential via a second accelerator tube. Due to incomplete charge stripping in the carbon foil, the beam contained some negative and neutral components, but only the positive component was bent 30° to the right by a dipole inflection magnet into the target chamber. The resulting protons can have energies up to nearly 480 keV.

The stripping foils for the mini-tandem were mounted on holder rings by floating carbon foils coated with collodion (a light plastic compound) away from their glass slides onto a water surface, then lifting them away on their holder rings. A protective coating of collodion evaporates within a few seconds of beam bombardment, leaving only the carbon foil. Vacuum in the beamline and chamber was maintained by three turbo pumps to levels on the 10^{-6} torr scale (always below 10^{-5} torr).

2.4 Frozen Targets and Target Chambers

2.4.1 High voltage chamber

An overhead view of the high voltage "frozen target" chamber used in the current work is shown in Figure 2.4. It was originally designed for studies of polarization observables in the ^{2}H(d, γ)^{4}He reaction. The chamber was designed to accommodate

Figure 2.3: Overhead view of the TUNL mini-tandem and LEBAF beam transport system.

NaI scintillating detectors, surrounded in large shields of plastic scintillators (for cosmic ray rejection) in a lead housing, for detection of the 23.8 MeV γ-rays at a large range of angles. It was also designed to get these detectors as close as possible to the target itself, to maximize the solid angle subtended by the detector.

Figure 2.4: Overhead view of the high voltage target chamber used for positive beam experiments. Also shown in the figure are the placements of HPGe detectors and their associated background-reduction shielding.

The chamber consisted of a 1 m long entrance beamline (to put the accelerator tube out of the way of detectors during measurements where the detectors were placed at angles greater than 90° relative to the beam) leading to a narrow chamber box. The surrounding ground potential detectors and lead shielding are insulated from the high voltage applied to the chamber by a 1/4" thick plastic layer. This plastic was also used to support the target chamber, which was suspended on the surrounding

plastic box by the liquid nitrogen dewar. An aluminum cold plate, on which the target material was deposited, extended downward from the liquid nitrogen dewar. This plate can be seen in a side view close-up of the chamber itself in Figure 2.5.

An exit port for charged particles such as protons from the $^2\mathrm{H}(\mathrm{d},\mathrm{p})^3\mathrm{H}$ reaction was located at a 30° angle below the entrance beamline, but was not used in the present experiment. Directly below the target plate was another exit port for pumping. A turbo pump was attached to this port via a 60 cm delrin (low vapor pressure plastic) tube, which served as electrical insulation between the pump and the high voltage chamber. The chamber was composed of stainless steel. The walls of the target chamber box were made 1/4" thick to prevent buckling under the force applied by atmospheric pressure.

The side of the box opposite the HPGe detector had a 3" hole, over which was placed a plastic cap with an O-ring seal between the cap and the outer wall of the chamber. Near the edge of this cap, an O-ring feedthrough for a tube to carry target material vapor was sealed in place with TorrSeal (low vapor pressure epoxy). The copper vapor feed tube could be rotated through this feedthrough, as shown in Figure 2.5, so that target vapors could be deposited on the surface. A metal plate attached to the vapor feed line also served as a tuning target. With the tuning target in place and insulated from the rest of the chamber by means of the plastic cap, it was possible to tune the proton beam so that it would strike the target plate and not the chamber walls. During data taking, the tuning target and vapor feed line were rotated so that they were out of the path of the beam. High voltage bias was supplied by a Glassman WR 100 regulated power supply or equivalent regulated supply with the current limitation set low to prevent supply overload in case of the formation of a short circuit to ground.

Figure 2.5: Side view of the high voltage target chamber used for the positive beam measurements. The target material vapor line and attached tuning target are shown in both the target making (solid lines) and data taking (dashed lines) positions.

2.4.2 Ground potential chamber

The ground potential chamber was also originally designed for study of another reaction, ^{2}H(p,γ)^{3}He. As shown in Figure 2.6, it consisted of a simpler version of the high voltage target chamber. This chamber was used only for the study of the 278 keV resonance with the mini-tandem. This chamber's dewar was insulated from the target chamber by an insulating plastic spacer ring with O-ring seals. This insulation, and the fact that the chamber sat at ground potential, made integration of the beam current possible. The positive terminal of a 90 V battery was connected to the target dewar to suppress the emission of secondary electrons from the frozen targets. The negative terminal was fed into a model 1000 Brookhaven beam charge integrator. The charge integrator produced a logic pulse corresponding to a set amount of integrated charge, and the pulses were counted by scaler modules in the computer interface to determine the total charge incident on the target.

A further consequence of the insulation ring was that the only part of the chamber which was electrically connected to the liquid nitrogen dewar, and which could be reached by the beam, was the target plate itself. This made a tuning target unnecessary, and allowed the use of a fixed vapor feed line. The chamber also contained a small collimator just in front of the target plate to ensure good beam collimation.

2.4.3 Target materials

Frozen target materials which were used fell into two basic classes: those which contained ^{14}N and those which did not. Ammonia was chosen both because its vapor pressure near 77 K (10^{-11} torr [Avs98]), the temperature of the liquid nitrogen coolant, was low enough to maintain vacuum. It was also chosen because it contains no carbon.

Figure 2.6: Side view of the target chamber used for mini-tandem experiments.

Compounds which contain appreciable carbon, such as melamine, were avoided as target materials. As protons stopping in the target break chemical bonds the residual carbon atoms in a carbon-compound target, such as melamine, combine to form a dead layer of carbon on the target surface. Small amounts of carbon from the breakdown of hydrocarbons (from the oil in the mechanical pumps and the grease on the O-rings) which become frozen to the target plate surface were still expected to be the principal target contaminant. Water vapor from air entering small leaks was another possible surface contaminant. Consequently, much time was spent finding and sealing all air leaks before the start of each experiment and heavily greased O-rings were avoided.

Because experiments in the high voltage chamber could not rely on beam current integration for normalization, deuterated ammonia (ND_3) with 98% D_3 from Cambridge Isotope Laboratories was used so that γ-rays from the well-understood $^2H(p,\gamma)^3He$ reaction could be used to extract the beam current incident on the targets. Because of the high cross section for the $^2H(p,\gamma)^3He$ reaction, 98 % D_3 was determined to be unnecessary after several experiments and the deuterated ammonia was mixed with an equal amount of normal ammonia (NH_3). This change in the isotopic mixture provided an additional check of the effect of neutron production by scattered deuterons in the target through the $^2H(d,n)^3He$ reaction.

Data were also taken with frozen heavy water (D_2O), which can withstand considerably higher beam currents because of its higher freezing point temperature. Further data were available from another experiment with an identical setup using a target of frozen benzene (C_6H_6) at very high beam currents ($\tilde{}$50 μA). The best control data were taken in the last measurement with ammonia enriched with 98 % ^{15}N ($^{15}NH_3$), with a 20% mixture of ammonia which was both deuterated and enriched in ^{15}N ($^{15}ND_3$, 98% enrichment in both ^{15}N and D).

2.5 Detectors

2.5.1 High Purity Ge (HPGe) detectors

In the present work photons were detected by large, high purity germanium (HPGe) detectors. The main detector element consists of a single cylindrical crystal of germanium with a central bore. This element is essentially a giant diode, and it is available in two types. In the p-type which were used in this work, the outer surface is coated with a thick (600 μm) layer of electron-donating lithium while the surface of the central hole is ion-implanted with a very thin (< 0.3 μm) layer of boron. In the n-type, typically smaller and used for lower-energy work, these contacts are reversed.

A HPGe detector converts the energy of a γ-ray into excited electron-hole pairs within the semiconductive Ge material, which are then separated by a bias voltage to their respective electrodes. This results in an amount of charge which is linearly proportional to the energy of the γ-ray to be collected by the preamplifier. The low energy required for the creation of an electron-hole pair (2.96 eV at 77K for Ge) creates many more such pairs than photons generated in a scintillator, and their collection efficiency is nearly 100%. This leads to much finer energy resolution in HPGe detectors (where typically FWHM <0.1% of E_γ) than in a typical NaI scintillation detector (where FWHM <5% of E_γ is considered excellent resolution).

Photons are converted into electron-hole pairs by three processes, described below, in any radiation detector. Each of these processes is associated with an attenuation length, $l_{att.}$, given by:

$$N_t = N_0 e^{-l/l_{att.}}, \tag{2.1}$$

where N_t is the number of γ-rays transmitted a depth l through the absorbing ma-

terial, and N_0 is the number of γ-rays incident on the absorbing material. The attenuation length, $l_{att.}$, depends on the process of absorption, the material absorbing the γ-rays, and the energy of the γ-rays. The energy dependence of the attenuation lengths for these three γ-ray interactions in Ge is well-known[Ort97], and can be seen from the plot of the inverse attenuation length in Figure 2.7.

The process which dominates at energies $<\tilde{}$ 1 MeV is the photoelectric effect, where the photon ionizes an atom and the photon energy is converted into kinetic energy for the electron and the energy required to ionize the Ge atom to which that electron is bound. This electron loses its kinetic energy in the crystal through the creation of a combination of electron-hole pairs and lattice phonons via scattering processes in the Ge crystal, the details of which are still not entirely understood [Ort97].

At intermediate energies ($\tilde{}$ 1-10 MeV), as shown in Figure 2.7, the dominant process of energy conversion becomes Compton scattering. An incident photon with energy E_γ scattering from an electron in the crystal at an angle θ will have an outgoing energy E' determined by conservation of energy and momentum as

$$E' = \frac{E_\gamma}{1 + \frac{E_\gamma}{m_e c^2}(1 - cos\theta)} \tag{2.2}$$

The scattered photon may be absorbed or scatter again, or it may escape the crystal undetected. If the scattered photon escapes, the recoiling Compton scattered electron leaves its kinetic energy in the crystal. This energy ranges from 0 to a maximum of $E_\gamma(1 - \frac{1}{1+\frac{2E_\gamma}{m_e c^2}})$ for photon scattering angles of 0-180°. In the measured spectrum, this results in a continuous distribution of energies up to a maximum for 180° scattering (the "Compton edge") just below the full-energy peak.

The third conversion process for γ-rays, pair production, is more important at

higher energies. It becomes the dominant process of conversion for γ-rays with energies above 10 MeV. In this process, photons with energies above 1022 keV may create an electron-positron pair with kinetic energy equivalent to E_γ-1022 keV. These two leptons lose their energy in the crystal, and then the positron annihilates with an atomic electron to create a pair of 511 keV γ-rays. These γ-rays may be detected in the crystal, resulting in a full-energy peak, but one or both of them may escape the crystal instead. This results in the formation of two "escape peaks" in the measured spectrum.

Figure 2.7: The inverse of the attenuation length for the three photon interaction processes in Ge and Si, given in units cm^{-1} as a function of the photon energy in MeV. The curves are labelled by their material (Ge or Si) and by one of: (PE) for the photoelectric effect, (CS) for compton scattering, or (PP) for pair production [Ort97].

HPGe detectors are delivered with an efficiency measured for the 1.332 MeV

γ-rays from a ^{60}Co source relative to the efficiency of a cylindrical 3"×3" NaI scintillator, the industrial standard γ-ray detector. The two HPGe detectors in the present experiment had efficiencies of 123% and 140%, respectively. The relevant absolute efficiency of each detector is not so simple, as it depends on the energy of the measured γ-rays and the geometry of the setup. The energy dependence of the efficiency of the 123% detector has been recently measured [Bru01] by observations of the ^{27}Al(p,γ)^{28}Si reaction. This reaction has a resonance at E_p = 992 keV and provides γ-rays with well-known relative intensities at energies from 1-10 MeV. The results of this measurement of relative efficiency agree well with Monte-Carlo simulations using the Electron Gamma Shower 4 (EGS4) code package [Sch97]. The 123% HPGe detector was used in all but one of the present measurements: the measurement at E_p=140 keV. The different efficiency of the 140% detector used for the E_p=140 keV measurement was corrected for by means of a calculation performed with a Monte Carlo code, the details of which will be discussed in Section 4.2.

2.5.2 NaI annulus

The NaI annulus is an example of a scintillation detector. Scintillation detectors work in a way similar to HPGe detectors in that they convert the energy of a photon into an electrical signal. The energy in a γ-ray is converted into the kinetic energy of energetic electrons by the same three processes as it is converted into electron-hole pairs in HPGe detectors. The electrons lose this kinetic energy to the energy required to excite the atomic electrons of the scintillator material. Photons emitted from the de-excitations are not re-absorbed by other atoms in a scintillator, but instead are free to travel between the reflective inner walls of the detector. A fraction of the photons are absorbed by a material with a low work function coating the face of a

photomultiplier tube (PMT) which is optically coupled to the scintillator using a light pipe and optical grease. The couplings is designed to have an index of refraction close to that of the scintillator material itself to minimize total internal reflection at the interfaces.

The photons which are absorbed by the PMT's inner-surface coating cause low-energy electrons to be ejected from the coating via the photoelectric effect. The coating is biased negatively with respect to the next stage of the PMT, called a dynode. The ejected electrons are accelerated by the electric field between the coating and the dynode. The impact of each electron with the dynode emits 4-5 electrons, which are accelerated towards the next dynode, which is at a more positive potential. This process, repeated through a succession of dynodes which are maintained about 100 volts apart, generates a large shower of electrons which are received at the detector's anode. The signal is patched to the control room for analysis by the circuit explained in Section 2.6.2.

2.5.3 ^{6}Li-glass detector

The ^{6}Li-glass detector is another example of a scintillation detector. This detector was used to determine the background level of thermal neutrons in the room via the ^{6}Li(n,^{3}H)^{4}He reaction in the lithium-doped glass of the scintillator. The Q-value for this reaction is 4.783 MeV, which is converted into the kinetic energy of the charged reaction products. The spectrum derived from this detector in the presence of thermal and/or epithermal neutrons is a low background with a peak corresponding to 4.783 MeV ^{4}He events.

The particular ^{6}Li-glass detector used was a 9.5 mm thick by 51 mm diameter

Bicron GS20 (6.6 % Li by weight, 95 % ^{6}Li enrichment) scintillator which is optically coupled to a PMT. To find the efficiency of the scintillator for thermal neutrons (E_{av}=0.025 eV) we may either calculate the probability of neutron absorption directly using the cross section for the ^{6}Li(n,^{3}H)^{4}He reaction presented in Figure 2.8 or we may more simply refer to the nomogram provided by Bicron in Figure 2.9. Shown in Figure 2.9 is an example of how to calculate the efficiency for a 2mm thick GS20 detector. Drawing a line on the nomogram between the detector material type and the glass thickness allows one to extrapolate the ^{6}Li area density. A line between the neutron energy and this extrapolated ^{6}Li area density can then be used to extrapolate the detector efficiency. As can be seen in Figure 2.9, a 2 mm thick slab of GS20 scintillator has an efficiency of 95%, or a transmission of 5%. The 9.5 mm thick slab (equivalent to 4.75 slabs, each 2 mm thick) used in the present work then has a transmission of approximately $(0.05)^{4.75} = 6.6 \times 10^{-7}$, giving an efficiency close enough to 100% to neglect this correction in analysis of the neutron flux data. We therefore assume that the detector has an efficiency of 100% for detecting thermal neutrons.

2.6 Electronics

2.6.1 Solid state detector electronics

In order to obtain digital data from the detectors, their output signals must first be amplified and shaped to a form which the analog-to-digital converter can handle. A separate trigger signal for each proper (larger than noise) input signal must also be generated in order to tell the computer that a signal has arrived. These signals must be processed with fast electronic modules, and carefully timed to arrive at the

Figure 2.8: Energy dependence of the ^{6}Li(n,^{3}H)^{4}He cross secion (bottom line) and total neutron absorption cross section (top line) for ^{6}Li. At thermal energies (2.5×10^{-8} eV), the cross section is dominated by the (n,^{3}H) reaction.

Figure 2.9: Nomogram for the neutron detection efficiency of ^{6}Li-glass detectors. Drawn on the nomogram is an example of how to extrapolate the efficiency of a particular type and thickness of ^{6}Li-glass detector (see text).

digitizing module together.

The electronics for processing signals from silicon and HPGe detectors are nearly identical. The processing electronics for the HPGe detectors used in this experiment are shown in Figure 2.10. The HPGe detector was biased with an ORTEC 659 5kV bias supply. The first of two preamplifier outputs from the HPGe were connected to an ORTEC 672 spectroscopy amplifier. The integrated signals were converted into output pulses with gaussian shapes and peak heights proportional to the integrated input signals. The integration time of the amplifier was chosen to be 6 μs and the gain was adjusted so that the full energy range of the acquired spectra was approximately 8.5 MeV, corresponding to an output pulse height of 10 Volts.

The second preamplifier output was connected to an ORTEC 454 timing filter amplifier (TFA), which created a short negative output pulse with accurate timing relative to the leading edges of the input pulse. The TFA integrated and differentiated the first 200 ns (this time is adjustable, and for scintillators is usually set shorter) of the input pulse to form an estimate of the amplitude for each input pulse, and generated an output pulse for each input pulse with an amplitude proportional to this estimate. The TFA output was read by an ORTEC 934 constant fraction discriminator (CFD). The CFD received input signals and created copies of those same signals, which were sent from the CFD through a 100 ns cable delay and back again. The CFD combined input signals with their delayed copies to create shaped signals with accurate zero-crossing timing. The zero crossings were detected by internal logic circuits which created three copies of a short NIM logic signal (a -1 V square pulse) if the TFA output was above the CFD threshold, which was set by an adjustable leading-edge discriminator attached to the input line. The CFD threshold was set to eliminate double-pulsing in the CFD. The detector threshold was adjusted by adjust-

ing the gain of the TFA so that the small, high-rate signals due to electronic noise were not amplified above the CFD threshold.

One of NIM outputs from the CFD was used to trigger a Phillips Systems 794 gate and delay generator (GDG). The width and delay of the GDG were adjusted so that its delayed NIM output signal could trigger a second GDG at the beginning of the output pulse from the spectroscopy amplifier. The width of the second GDG was adjusted so that the signal from its TTL output (a 2-5 V square pulse) encompassed the output of the spectroscopy amplifier. These signals were patched to an ORTEC 413 4-input peak-sensing analog to digital converter (ADC), which was activated by the TTL gate. The ADC converted the height of the gaussian pulse from the spectroscopy amplifier into a channel number and set a look-at-me (LAM) signal to inform the computer that it should record the information stored in the CAMAC modules. The other two outputs of the CFD were used as inputs to scalers and for cosmic ray rejection (see below). Details of the electronic interface between the ADC and the MicroVAX computer are described in Section 2.7.

2.6.2 Scintillation detector electronics

High-energy cosmic rays were a primary source of background in this experiment. Cosmic rays create showers of energetic electrons and γ-rays which generate continuous distributions of background counts in a measured spectrum even at high energies. Cosmic rays passing through lead shielding create their own showers and hence cannot be effectively removed with passive shielding. These radiation showers can be detected and removed, however, because they leave a signal in a detector placed close to (or around, if the detector is an annulus) the HPGe used for the measurement. The 511 keV γ-rays associated with escape peaks in the HPGe as well as Compton-

Figure 2.10: Output pulse analysis circuit for the HPGe detectors (see text).

scattered γ-rays may also leave signals in the annulus which are coincident with the signals in the HPGe. The energy signals of the NaI annulus were therefore analyzed to investigate the effects of using them to make software cuts for the HPGe signals.

For some measurements in this experiment a NaI(Tl) scintillator annulus was available to perform active rejection of cosmic ray signals; the circuit for this is shown in Figure 2.11. This rejection shield was originally designed for γ-ray polarimetry measurements via Compton scattering, and consists of an annular cylinder of NaI(Tl) scintillator divided into four segments. Photons from each segment of scintillator were detected in a photomultiplier tube powered with a Lecroy HV4032A high-voltage supply. The negative output pulse of each tube was sent through a Phillips 744 fan out to make two copies of the signal. One copy of each PMT's output signal was sent through a TFA and into a CFD to create NIM signals indicating a detected event, as with the HPGe electronics. The integration and differentiation of the TFA were set to 20 ns, and the CFD was given 20 ns of delay. The other copy was sent to a spectroscopy amplifier with an integration time of 1 μs. The output of the spectroscopy amplifier was delayed to coincide with the output of the HPGe detector's spectroscopy amplifier output pulse, and sent to an ADC for conversion to a digital signal. That ADC was gated by the same TTL pulses as the ADC for the HPGe detector.

Because of its low resolution, energy signals from the annulus were only of interest if they occurred in coincidence with signals from the HPGe detector. Copies of the four CFD outputs of the scintillator were combined with a Phillips 755 logic module, which was set to the minimum coincidence level so that any input NIM signal would generate a NIM output signal (thus creating a logical OR of the NIM input signals). The delayed output of this logic unit was sent to an ORTEC CO4020 fast coincidence

Figure 2.11: Electronic circuit for signals from the cosmic ray rejection annulus.

module along with a copy of the signal from the CFD of the HPGe. When a signal above their respective noise thresholds occurred simultaneously in both a segment of the annulus and the HPGe, the coincidence module generated a NIM signal. This signal was sent to a Bi Ra 2351 NIM-input module (the hit register), which records sorting bits which were to be read by the computer each time an event was processed in order to indicate the type of event being processed.

For the case of the ^{6}Li-glass detector (with a single PMT biased by the same LeCroy HV4032A supply that powers the NaI annulus PMTs), the negative output pulse of the PMT was simply put through an ORTEC 672 spectroscopy amplifier with an integration time of 1 μs. The output pulses from this amplifier were patched directly into a multi-channel analyzer (an ORTEC 413 ADC card in a PC running windows) which was operated with the program MAESTRO. This software allowed us to take and store a spectrum of one detector, along with its live time (run time with the total dead time from digitization of the pulses, which was less than 2%, subtracted). It also calculated the area under the peak of interest by means of a simple Gaussian fit, which was then divided by the live time to give an average number for the counts from the detector per second for 90,000 seconds. This was just over 24 hours, but sufficient because of the efficiency of the detector.

2.7 Data acquisition system

2.7.1 Computer Interface

The ADCs, hit register, and scalers were located in a CAMAC crate controlled by a Microprogrammed Branch Driver (MBD), as shown in Figure 2.10. The MBD

programming code was loaded through the MicroVAX 3200 computer with which it was interfaced. Whenever the MBD received an interrupt signal, a LAM from the LAM patch panel, it executed the data acquisition program (DAP) code which had been loaded into it. The MBD was capable of performing operations on 8 channels, but only one was used in the experiment. The DAP code indicated the modules to be read. Data from these channels were stored as a string of words in a memory buffer of the VAX. When that buffer was full the events stored in it were sorted by the TUNL data acquisition program, XSYS. XSYS referred to a specific set of instructions in an event language (EVL) code for sorting the recorded events into spectra. The data word read for the hit register was broken into individual bits to tell XSYS how to route digitized energy data from the HPGe and NaI annulus (when available) into spectra. The words were also spooled into large event files for later offline analysis. Acquired spectra were then displayed on the VAX as the data were being taken using the DISP package in XSYS.

2.7.2 Voltage control and beam pulsing

In addition to cosmic ray rejection bits from the shield, background was reduced by pulsing the beam in some experiments (see Section 2.2). This pulsing was controlled by an adjustable clock made from two GDG modules, shown in Figure 2.12, with one set for the beam-on time (50-100 ms) and one set for the beam-off time (200-250 ms). The TTL output of the beam-off module drove the external input of the power supply which powered the electrostatic steerer for the chopping, and the NIM output of the beam-on supply was sent to the hit register to indicate to XSYS that the events should be sorted into beam-on spectra. A short (5 ms) veto pulse for the steering time of the electrostatic steering system was created by a third GDG

which had as an input copies of the triggers of both GDG modules in the clock. The NIM output of this module was sent to the computer as a veto bit so that the beam-on and beam-off spectra generated would have the beam either fully on or off.

Figure 2.12: Beam pulsing and voltage ramping circuits.

In the single measurement in which the E_p=122, 124, 126, and 128 keV data points were taken, the target-bias voltage was changed automatically in 2 kV steps. The order of the chamber bias voltages during voltage stepping was 42→ 44→ 46→ 48→ 48→ 46→ 44→ 42 kV. That cycle was then constantly repeated from the beginning (42 kV) until the week-long measurement was over. The analog output of a TUNL-made digital ramp was passed through a voltage divider made of a 10-turn potentiometer and a 100 kΩ resistor in series, as shown in Figure 2.12. The resulting drive voltage

was sent through a Lecroy 588 linear fan-in to add an offset voltage and then sent out to the external input of the bias supply for the high voltage chamber. Separate TTL digital outputs from the TUNL digital ramp were passed through a Lecroy 688AL level adapter. The resulting NIM outputs indicating the voltage state of the chamber were patched to the hit register through 0.1 μF isolation capacitors (which prevented ground loops) to be used as sorting bits. The delay between bias changes was 60 seconds, approximately 200 times the estimated 300 ms second ramping time, so the contribution to the error from counts at incorrect voltages is estimated at $<0.5\%$.

Chapter 3

Results

The data collected by the procedures described in the previous section consist of a set of HPGe spectra. There are one or more accumulated spectra for each beam energy and target used. In addition, each accumulated spectrum has a corresponding background spectrum. Extracting information about the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction from this set of spectra involves a series of steps. The description of these steps in this chapter is intended to help form a consistent picture of the body of experimental work.

An example of the spectra taken with beam incident on a target of frozen ND_3 is displayed in Figure 3.1. The high resolution of the HPGe detectors can be seen in the forest of peaks at low energies, which represent the detection of monoenergetic γ-rays from the decay of naturally occurring unstable isotopes. These peaks correspond to transitions of known energies and provide data points for energy-calibrating the detection system. Gain shifts show up as shifts in the position of these sharp lines within the spectra, allowing us to check the stability of the gain of the detection

Figure 3.1: Example of the full, accumulated spectrum from a one-week measurement. The wide D(p,γ)^{3}He peak is clearly visible near 5.6 MeV, along with a host of low-energy lines from the decay of naturally occurring radioactive elements in the surrounding material.

system over time. This constant flux of calibration data allowed accurate corrections to the channel ranges used to determine the γ-ray yield for the two cases of gain shift mentioned below.

Shown in Figure 3.2 is an example of the spectra taken with beam incident on a target containing ^{14}N. Peaks from the reaction of interest and background lines are visible in the spectrum, including the strong doublet peak from thermal neutron capture on ^{56}Fe at E$_\gamma$=7.631 and 7.645 MeV [Ven99]. The energy and width of this doublet provided checks of the energy calibration and a measurement of the FWHM of the HPGe detectors near the energy of interest. Background in the primary data peak due to the ^{70}Ge(n,γ)^{71}Ge reaction shows up in the spectrum at 7.4158 MeV. The fact that peaks from individual transitions in ^{71}Ge are not observed is due to the fact that individual transitions between excited states of ^{71}Ge are suppressed by multiple γ-ray absorption in the detector crystal, while the peak to total calorimetry

Figure 3.2: Example of a section of Figure 3.1 from 7.0-7.8 MeV, showing the peak of interest at 7.414 MeV, as well as the double-peak around 7.64 MeV from thermal neutron capture by ^{56}Fe, as well as a weaker peak from thermal neutron capture in the surrounding lead (^{207}Pb) at 7.36777 MeV.

of the reaction is enhanced by the very same process. This appears in the data as an enhanced peak at the Q-value of the reaction, which for ^{70}Ge(n,γ)^{71}Ge is 7.4158 MeV [Lon81]. The effect of total calorimetry also suppresses the appearance of transitions in ^{74}Ge from ^{73}Ge(n,γ)^{74}Ge, which has a higher capture cross section and would otherwise be observed in the spectrum. A peak at the Q-value for that reaction would be enhanced; however, that peak is at is 10.196 MeV [Lon81]. That energy is above the energy at which the spectroscopy amplifier saturates ($\approx$9 MeV).

3.1 ^{14}N(p,γ)^{15}O yields

The first step is to extract the yield of γ-rays from the ^{14}N(p, γ)^{15}O reaction, the reaction of interest. This is a relatively simple procedure. It consists of subtracting the number of counts within a given energy range (2$\times$ the FWHM of the detector for 7.414 MeV γ-rays, and centered at 7.414 MeV) of the normalized, accumulated background spectrum taken with the beam off from the number of counts in the same energy range of the accumulated spectrum taken with the beam on. It is critical that the energy range be the same in both spectra. In the bulk of cases this was not a problem, the gain of the detection system remained stable to well within the resolution of the detector throughout the measurement of the background and the beam-on data. For the case of the benzene target (used as a control) and for one of the two measurements with ^{14}N-ammonia at E$_p$=140 keV, corrections to the channel range used had to be made for small but detectable gain shifts in order for the channel range to correspond to the same energy range in the beam-on spectrum and the background spectrum.

A channel range of 2$\times$ the FWHM of the detector, centered around 7.414 MeV,

is chosen for both beam-on and background. A linear background is fit to the data and subtracted. Energy ranges on either side of the peak of interest are chosen in the flat regions where we observe no background peaks which can be seen in Figure 3.2. The same energy range is used to fit both the beam-on and background spectra. The number of background counts after subtraction of the flat background (the net background) from the energy range of interest is normalized to the size of the 7.631+7.645 MeV ^{56}Fe(n,γ)^{57}Fe doublet (also with a flat background subtracted). This peak is chosen because both it and the background peak which lies beneath the peak of interest are thermal neutron capture peaks. Therefore, the size of these peaks should be linearly proportional to the flux of thermal neutrons in the experimental area, and any change in that flux will cancel in the normalization.

The number of background counts in that channel range is also checked by normalizing the net number of background counts to the number of counts in the 1.461 MeV background peak due to the decay of naturally occurring ^{40}K. This peak is very large, and its size is proportional to the time for which the detector accumulates data minus its total dead time (the live time). The normalization of the net number of background counts was checked a second way by normalizing it to the time read by the scaler modules (including the correction of about 0.5% for the measured computer dead time), and both cases returned an identical result to that obtained with normalization to the ^{56}Fe neutron capture doublet. The computer dead time correction factor was calculated by dividing the number of detector triggers out of the fast CFD modules, which correspond to digitizeable events, by the number of events which are actually digitized and stored by the MBD. This dead time correction can be slightly different between beam-on and beam-off data taking because of the extra dead time introduced by the flux of γ-rays from the D(p,γ)^{3}He reaction.

Figure 3.3: Spectrum taken with a neutron source and the usual lead shield, but no steel target chamber present, showing the dramatic reduction on the 7.64 MeV doublet from the ^{56}Fe(n,γ)^{57}Fe reaction.

The fact that three different normalization techniques lead to the same result suggests that any change in the cosmic muon flux between beam-on and background measurements was small. More importantly, it implies that our beam is not itself a source of neutrons. As can be seen in Figure 3.3, when the steel target chamber is removed the thermal neutron capture doublet near 7.64 MeV is dramatically reduced (to 15% of its former value relative to the ^{70}Ge(n,γ)^{71}Ge peak). The spectrum was taken with an AmBe neutron source in combination with an 8 inch thick layer of thermalizing wax.

It is possible that scattered deuterons in the target could be a source of neutrons through a secondary D(d,n)^{3}He reaction. This reaction has a cross section of 12 mb at the approximate energy of any deuterons elastically scattered by the proton beam[Bro90]. This generation of extra thermal neutrons with our beam would alter

the results derived by normalization to the strong ^{56}Fe capture doublet. This effect was never observed in our data, and changes in the deuterium concentration of the target (by mixing it with normal ammonia) did not alter the results. Neutrons from any conceivable beam-related process would also have led to a constant increase in thermal neutron flux at higher proton energies. As will be shown later in our final yield curve, this behavior is inconsistent with our results.

Figure 3.4: Spectrum taken with a pulser at the test input at each of the indicated attenuations for 100 seconds.

The linearity of the system (important for our linear energy calibration) was checked by connecting a pulser to the test input of the HPGe preamplifier and changing the gain of the pulser by factors of 1/2, 1/5, and 1/10. A 100 second measurement was taken at each setting. The centroid channel numbers of the tall pulser peaks in the resulting spectrum, shown in Figure 3.4, are fit by a linear function with deviations below the FWHM of the peaks. This demonstrates that any nonlinearities in the detection system are below the detector resolution and therefore not important

in the results. Differences in the peak heights in Figure 3.4 are due to the differences in the number of channels of amplifier resolution for the various amplitudes of pulser input.

In a few cases the detector resolution was negatively affected by electrical noise from the nearby high voltage target chamber. For measurements where the resolution was negatively affected, a window corresponding to the FWHM with the target chamber biased was used for both the beam-on and background data. This has the effect of including more background counts in the measurement and making the error on the **net** counts slightly larger.

Figure 3.5: Composite spectrum of all beam-on measurements (solid line) with $E_p \geq$ 130 keV and ^{14}N-ammonia targets, shown with a composite of all their corresponding background measurements (dotted line) in an energy range near the peak of interest.

A composite of all accumulated spectra for measurements with $E_p \geq 130$ keV and ^{14}N-ammonia targets is shown in Figure 3.5. The solid line represents beam-on data while the dashed line represents a composite of corresponding background

spectra. Each background spectrum in this composite was normalized to the live time of its corresponding beam-on accumulated spectrum by the prominent 1.461 MeV background line, as previously described. This normalization makes the area of the beam-on and background peaks from thermal neutron capture by ^{56}Fe the same to within statistical uncertainty even in the composite spectrum (which has much higher statistics than any individual spectrum). There is a clear enhancement of nearly 400 counts in the peak of interest, which overlaps a peak due to thermal neutron capture by ^{70}Ge. This excess is 5 standard deviations above the background measurement, indicating the presence of a proton-induced reaction leading to an excited state of 7.414 MeV in ^{15}O. This result will be discussed in more detail in Chapter 4.

3.2 $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ yields

To map the excess yield of 7.414 MeV γ-rays as a function of energy, the count excess from each individual measurement (derived as described above) must be properly normalized. This requires knowledge of the number of protons incident on the target. The high bias on the target chamber makes direct integration of the current impossible, however. For this experiment, the number of incident protons in a particular measurement is derived from a simultaneous measurement of the $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ reaction. This reaction is well-studied at low energies using the same ion source and γ-ray detection system used to detect γ-rays from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction.

Shown in Figure 3.6 is the full γ-ray spectrum from the $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ reaction. Visible in the figure are the full-energy photopeak, the first and second escape peaks, and the Compton edge. Even though the escape peaks represent additional data, it can be seen that the full-energy photopeak alone contains enough statistics to

Figure 3.6: γ-ray spectrum from the ^{2}H(p, γ)^{3}He reaction.

accurately derive the number of protons incident on the target.

Displayed in Figure 3.7 is an example plot of the full-energy photopeak from the ^{2}H(p, γ)^{3}He reaction in an accumulated spectrum. The energy of the incident protons in this measurement was 124 keV. The energy of a γ-ray from the ^{2}H(p, γ)^{3}He reaction is equal to the center of mass energy of the system added to the Q-value of the reaction. In a "thick" target, one which stops the beam entirely, the center of mass energy of the system varies and an energy distribution of γ-rays is emitted from the target. The cross section and stopping power also vary strongly as functions of the proton energy in the target, resulting in a characteristic shape to the distribution of emitted γ-rays. This distribution of emitted γ-rays is given by:

$$Y(E_{cm}) = \frac{\sigma(E_{cm})}{Stp(E_{cm})} N_p(E_{cm}) \epsilon d\Omega \tag{3.1}$$

where Stp(E_{cm}) is the stopping power of the target material, $\sigma(E_{cm})$ is the reaction

Figure 3.7: Section of the γ-ray spectrum representing only the full-energy photopeak from the ^{2}H(p,γ)^{3}He reaction (points with statistical error bars) with a deconvolution fit (solid line).

cross section, $N_p(E_{cm})$ is the number of incident protons (given by a step function which goes to zero above the incident proton energy), ϵ is the γ-ray detector efficiency, and $d\Omega$ is the fraction of the total solid angle subtended by the detector. The astrophysical S-factor for the ^{2}H(p,γ)^{3}He reaction may be used to replace $\sigma(E_{cm})$ by means of Equation 1.1, and is parameterized by the linear function S(E)=S_0+S_1E. For the ^{2}H(p,γ)^{3}He reaction, $S_0 = 0.166 \pm 0.005$ eV b and $S_1 = 0.0071 \pm 0.0004$ b [Sch97].

The stopping powers ($STP(E)$) for many materials are relatively well measured and understood [And77]. Shown in Figure 3.8 is a plot of the stopping power for protons in ammonia as a function of the proton energy in the lab frame. Because we are dealing with frozen ammonia, there is some concern that the stopping powers, which are typically measured for vapors, might have an error of 5-10%. As will be

shown in Chapter 4, however, the absolute scale of the stopping power will cancel in the final analysis and only the form of the stopping power will remain. This is a benefit of the use of a relative cross section measurement. As can be seen in Figure 3.7, the form of our stopping power (these forms are very similar for most materials) produces an accurate fit to the data, so we expect the contribution to the total error of our measurement from uncertainty about error in the stopping power to be negligible.

Figure 3.8: The stopping power for protons in ammonia as a function of the proton energy in the lab frame [And77].

The fraction of the total solid angle is usually simple to approximate from the geometry of the setup. In the present experiment, however, the target chamber was designed so that the HPGe could be placed within 5 cm of the target. The detector crystal dimensions are larger than the separation of the target and the detector crystal's front face. Typically, a detailed simulation using the electron gamma shower (EGS4) simulation package[Nel85] would be required to determine the product of efficiency and solid angle, as well as to account for the effects of a finite target geometry.

However, in this case the $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ reaction is being used to normalize the yield from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction in the same detector, from the same target, and using the same proton beam. As will be shown in Chapter 4, the solid angle terms in the analysis cancel and only a correction for the relative efficiency for detection of the γ-rays of 5.6 and 7.4 MeV is required. The mid-plane of the cylindrical Ge crystal typically subtends 1.01 Steradians out of the total 4π Steradians of solid angle from the center of the target, and this value is used in the present calculations for $\mathrm{d}\Omega$. This value, like the stopping power of ammonia, is just a stand-in which will cancel in the final analysis. Therefore, any error in the measurement of the solid angle should also contribute a negligible error to the total error of the measurement when compared to statistical error.

Figure 3.9: Absolute efficiency data for the 123% efficiency HPGe as a function of γ-ray energy [Sch97, Bru01], and simulations of the detector's efficiency as a function of energy using three different Monte-Carlo simulation packages.

The absolute efficiency for a source at 25 cm from the front face of the 123% HPGe

Figure 3.10: Spectrum from the ^{27}Al(p,γ)^{28}Si reaction at E_p=997 keV, showing the many γ-ray energies emitted by the residual nuclei [Bru01]. Peaks which are marked "X" are transitions to the ground state of ^{28}Si. Other peaks are due to transitions between excited states in ^{28}Si.

has been previously measured at TUNL using a variety of radioactive sources [Sch97]. The efficiency curve is displayed in Figure 3.9. The efficiency of the 123% detector has also been re-measured since it was put in a new cryostat by using γ-rays from the ^{27}Al(p,γ)^{28}Si reaction at the $E_p = 992$ keV resonance with the TUNL FN tandem accelerator [Bru01]. In the ^{27}Al(p,γ)^{28}Si measurement, a beam of approximately 500 nA of $E_p = 997$ keV protons were incident on a 1 mm thick Al target with a Cu backing. The 123% HPGe was placed at a distance of 14.5 cm and an angle of 45° from the target. This measurement provides data extending from 1.7 MeV to 10 MeV. A spectrum from this recent measurement is shown in Figure 3.10, and data from the efficiency measurement at several higher energies is shown with the data of Schmidt in Figure 3.9. The results are in good agreement with the EGS4 Monte-Carlo simulation from the earlier work, which is displayed in Figure 3.9 along with recent Monte-Carlo simulations using the GEANT and MCNP programming packages. Because of their agreement with [Bru01] in the energy range of interest for this experiment (5.6-7.4 MeV), the Monte-Carlo calculation of [Sch97] is adopted here. Simulations using EGS4 show that the overall efficiency per unit solid angle is slightly lower when the detector face is 5 cm from the target compared to the efficiency at the 25 cm distance used in the measurement. However, this absolute scale in the efficiency will cancel later in the data analysis. The simulations also indicate that the ratios of efficiencies at 5.6 MeV and 7.4 MeV are identical to within 0.3% for an identical detector at the two distances from the γ-ray source.

The yield function of Equation 3.1 has a discontinuity at the incident proton energy, but is smeared by the detector's nearly Gaussian response function into the smooth distribution shown in Figure 3.7. The response function, shown in Figure 3.11, is a fit to observed peaks due to γ-rays from ^{66}Ga sources [Sch95]. This re-

Figure 3.11: The HPGe response function [Sch97] for the 123% detector.

sponse function, given in Equation 3.2 is basically a skewed Gaussian form. The non-Gaussian contributions are given by functions built from the complimentary error function (erfc) and fit to the measured line shapes, as can be seen in Equation 3.2 [Sch95].

$$Resp(E-E_\gamma,\sigma) = \frac{0.75}{\sqrt{2\pi}\sigma}e^{-(E-E_\gamma)^2/2\sigma^2} + \frac{0.22}{3.5e^{-\sigma^2/6}}erfc(\frac{(E-E_\gamma)}{\sqrt{2}\sigma} + \frac{\sigma}{1.72\sqrt{2}}) + 0.01erfc(\frac{(E-E_\gamma)}{\sqrt{2}\sigma}) \tag{3.2}$$

$$Y(E_{obs}) = Y(E_{cm}) \star Resp(E-(E_{cm}+Q)) \tag{3.3}$$

Equation 3.3 was applied to Equation 3.1 to create convolved yield curves like the one shown in Figure 3.7. The widths (σ) for the fits performed were taken from

observations of the resolution of background lines in the spectra corresponding to nearly mono-energetic γ-rays. The widths of sharp background lines at low ($\approx$1 MeV) and high ($\approx$7.5 MeV) energies were fit by a linear function for each measurement, and the width during a particular measurement at 5.6 MeV for the ^{2}H(p,γ)^{3}He reaction were taken from this fit. The known astrophysical S-factor and stopping power for ammonia were put into the equations, and the total incident proton charge was varied as the only fit parameter to minimize the value of χ^2. The resulting value for the incident proton charge was used to normalize the excess yield to give the excess counts in the detector per Coulomb of incident beam.

3.3 Neutron yields

The ^{6}Li-glass detector described in Chapter 2 counted for three days. It was placed in the same type of lead shield and same location as the HPGe detector during the measurement of the ^{14}N(p,γ)^{15}O reaction. All accelerators were turned off at the time. The data acquisition and analysis program MAESTRO returned a background-subtracted yield of 14,300 $\pm$ 431 counts in the ^{6}Li(n,α)^{3}H peak with a total live time of 93,000 seconds in the measurement. The surface area of the detector is 55.2 cm^2, resulting in a flux of 28 $\pm$ 1 n/m^2/s for the 100% efficient detector.

3.4 Minitandem yields

The spectra from the minitandem experiment are slightly different from those taken at the LECAL line, due to the higher proton energies. The spectrum of γ-rays from the ^{2}H(p,γ)^{3}He reaction is shown in Figure 3.12. Also shown are a con-

volution fit (solid line) and a convolution fit which neglects both the escape peak associated the 6.176 MeV γ-ray due to the ^{14}N(p,γ)^{15}O reaction and a 35% contamination layer in the top 10 keV of target material (dashed line). This contamination will be shown in Chapter 4 to be inconsequential to the data analysis. The escape peak associated the strong 6.176 MeV γ-ray from the ^{14}N(p,γ)^{15}O reaction at 5.665 MeV (see below) is labelled in Figure 3.12. The size and width of that escape peak were calculated using the known full-energy to escape peak ratio at this energy and the observed size of the full-energy (6.176 MeV) peak. The incident proton charge determined from the deconvolution fit after the addition of the contamination layer and escape peak to the fit curve was 0.0309 Coulombs, which agrees to within 3% with the 0.0300 Coulombs measured by current integration.

The γ-rays from the ^{14}N(p,γ)^{15}O reaction displayed in Figure 3.13 involve capture to both the ground and subthreshold excited states of ^{15}O at 6.176 and 6.793 MeV, which can be seen in Figure 1.3. These excited states undergo a subsequent decay, almost always to the ground state, which can be observed by the HPGe detector, resulting in peaks at 6.176 and 6.793 MeV in addition to the peak from direct capture to the ground state near 7.55 MeV (the excited state energy associated with the resonance) where the cross section for capture is maximum. These three peaks also result in the aforementioned escape peaks in the data at 0.511 and 1.022 MeV below the peak energies, though usually the second escape peak (and E$_\gamma$-1.022 MeV) is very small. Both unlabelled peaks in Figure 3.13 are escape peaks.

Figure 3.12: The full energy photopeak from the $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ reaction with 0.03 Coulombs of incident 300 keV protons (points). Also shown is a deconvolution fit to the data (solid line), and a deconvolution fit which neglects both surface contamination and the escape peak associated with proton capture to the 6176 keV excited state of $^{15}\mathrm{O}$ (dotted line).

Figure 3.13: Logarithmic plot of the spectrum of γ-rays from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction with 0.03 Coulombs of incident 300 keV protons. Specific transitions are labelled, unlabelled peaks are all first escape peaks of labelled peaks.

Chapter 4

Analysis of the Data

4.1 The yield curve.

The excess yield of 7.414 MeV γ-rays from each measurement, as described in Section 3.1, is normalized to the number of Coulombs of incident protons extracted from the deconvolution fit to the D(p,γ)^{3}He data for that measurement. The resulting set of data (one point per measurement) is plotted in Figure 4.1.

The points representing data taken with ammonia targets containing ^{14}N (triangles) show a step function structure consistent with a narrow resonance at $E_R = 127 \pm 2$ keV. As protons with $E_p > E_R$ slow to a stop in the target, they pass through the resonance energy, where they have a much higher probability of emitting γ-rays. A step function (dotted line) is fit to these data points to determine the thick target yield. Data points representing targets not containing ^{14}N (labeled in the plot) show no excess yield. Particularly interesting are the circled data points, which were taken in a single measurement by stepping the voltage of the target chamber in 2 kV steps,

Figure 4.1: The excess yield (counts per Coulomb) of 7.414 MeV γ-rays as a function of incident proton energy for various targets. Circled data points were taken in a single measurement (see text), and labeled data points represent data taken for the indicated control targets (those with no ^{14}N content). Triangles represent yields from ^{14}N-ammonia targets, which are fit with a step function (dotted line).

as described in Section 2.7.2. This created many cycles of steps over the course of the week-long measurement, with data from the HPGe detector sorted into separate spectra for each voltage. This allowed us to extract γ-ray yields from the sets of data taken at each voltage as described in Chapter 3, as if the data sets had been taken in separate measurements at each voltage. Even more revealing are two points at E_p = 160 keV. The data point for the ^{15}N-enriched ammonia target showed an excess yield consistent with zero (15 $\pm$ 18 counts per Coulomb). When ^{14}N ammonia was substituted for the last half of the measurement under identical conditions, the excess yield was again consistent with the existence of the resonance (210 $\pm$ 40 counts per Coulomb).

The error bars in this plot represent a combination of statistical errors and systematic errors. Statistical errors arise from the counting statistics in the peak of interest, the normalizations, and the background subtraction. Sources of systematic error considered were the detector efficiency ratio for γ-rays at 5.6 and 7.4 MeV, the error in the γ-ray attenuation coefficient ratios at 5.6 MeV and 7.4 MeV for the target chamber walls and high voltage shielding, and the error in the abundance of deuterium. The overwhelming error contributions in any single measurement are the statistical errors from the counting statistics associated with the 7.414 MeV γ-rays.

4.2 The efficiency of the 140% HPGe detector.

The measurement at E_p=140 keV was performed with a 140% efficient HPGe detector. The relative efficiency of each detector at 5.6 and 7.4 MeV (energies of particular interest in the data analysis) were calculated with a Monte Carlo simulation using the EGS4 code system[Nel85]. The physical dimensions of each detector's Ge

crystal, which were used to create the input file for the simulation, were obtained from the specification sheets provided with the detectors by ORTEC. The input file also assumed a point source 6.35 cm from the front face of the detectors. EGS4 works by calculating the track of a γ-ray from the point source. The code first determines whether the photon, generated with a random direction, passes through the detector. If it intersects the detector, EGS4 determines the coordinates at which its flight path enters and leaves the crystal. Breaking up the line between these two points into smaller sections, it checks the probability of each process of absorption mentioned above to determine if the γ-ray is absorbed, where it is absorbed, and by which process. If the process is Compton scattering or pair production (which is always associated with two 511 keV γ-rays from the annihilation of the positron), the new γ-rays created by the interaction are also tracked through the detector to determine if they are absorbed. The amount of energy deposited in the Ge crystal by the Compton scattered electron is determined from the angle the Monte Carlo code assigns (the angle is assigned randomly, but the probabilities are weighted to reproduce the known angular distributions of Compton scattered γ-rays) to the new γ-ray using Equation 2.2.

The deposited energies are sorted into bins of adjustable width; for the present work a width of 100 keV has been used. As can be seen in Figure 4.2, all counts in the full-energy peak of the HPGe spectrum are in the top 100 keV bin. The number of 5.6 MeV γ-rays will be used to normalize the number of 7.4 MeV γ-rays which are observed. As can be seen in Figure 4.2, the ratios of the counts in the 7.4 MeV γ-ray photopeak to the counts in the 5.6 MeV photopeak are very similar for these two detectors. Therefore the ratios of the efficiency at 7.4 MeV to that at 5.6 MeV are very similar for the two sizes of HPGe, with the 140% HPGe's efficiency ratio

Figure 4.2: Output of the Monte Carlo code (EGS4) used to simulate the detector efficiencies. Shown are the counts in 100 keV wide bins from point sources which generate 300,000 7.4 MeV and 5.6 MeV γ-rays situated 6.35 cm from the face of the 123% and 140% HPGe detectors. The simulation uses the geometric measurements of the detectors specified in the data sheets, which were provided by ORTEC along with the detectors.

only 1.005 times that of the 123% HPGe. The ratios of 7.4 MeV γ-rays to 5.6 MeV γ-rays observed in the 140 keV measurement and in the measurement with the frozen benzene target, both made with the 140% HPGe, were corrected for this factor prior to their inclusion in Figure 4.1.

4.3 Extraction of resonance parameters.

The resonant cross section for a radiative capture reaction is generally approximated by the Breit-Wigner expression [Ang99]:

$$\sigma(E) = \frac{\pi}{\kappa^2}\xi\frac{\Gamma_p(E)\Gamma_\gamma(E)}{(E-E_R)^2+\Gamma(E)^2/4}, \tag{4.1}$$

where $\Gamma_p(E)$ and $\Gamma_\gamma(E)$ are the partial widths of the entrance and exit channels of the reaction (the incident proton and the outgoing γ-ray), respectively. $\Gamma(E)$ is the

total width of the resonance, given by the sum of $\Gamma_p(E)$ and $\Gamma_\gamma(E)$. κ is the wave number of the incident proton in the center of mass frame, and ξ is the statistical factor given by:

$$\xi = \frac{2J+1}{(2I_1+1)(2I_2+1)} \tag{4.2}$$

where J is the angular momentum of the excited state, and I_1 and I_2 are the spins of the incident and target particles. For the $^{14}\mathrm{N(p,\gamma)^{15}O}$ reaction, $\xi = (2J+1)/6$.

The spectroscopic factor for this state (which will be derived later) is very small, which in turn results in a very small resonance width ($\Gamma(E)$). This is mentioned here since it allows us to simplify Equation 4.1 by neglecting the energy dependence of the proton and γ-ray partial widths. The usual energy dependence of the γ-ray partial width is given by [Lan58]:

$$\Gamma_\gamma(E) = \Gamma_\gamma(\frac{E-E_f}{E_R-E_f})^{2\lambda+1}, \tag{4.3}$$

where λ is the multipolarity of the radiation and E_f is the energy of the final state (for our purposes, the ground state, so $E_f = 0$). The fraction in Equation 4.3 is obviously very close to 1 for energies E over the width of the resonance when $\Gamma_\gamma \ll E_R$, since $E_R \pm \Gamma \approx E_R$. So for any energy E close enough to E_R for the resonance to contribute strongly to the capture cross section, the width is approximately constant.

The energy dependence of the proton partial width is given by:

$$\Gamma_p(E) = 2\gamma^2 P_l(E,r) = \Gamma_p(E_R)\frac{P_l(E,r)}{P_l(E_R,r)}, \tag{4.4}$$

where $P_l(E,r)$ is the penetration function for the proton at proton energy E and for the radius, r, of the target nucleus. l is the angular momentum of the captured proton, and r is defined as:

$$r = r_0(A_1^{1/3} + A_2^{1/3}). \tag{4.5}$$

The value of r_0 in Equation 4.5 is taken to be the NACRE value of 1.4 fm, which is close to the standard value of 1.25 fm. A_1 and A_2 are the masses of the incident and target nuclei, respectively [Ang99]. $P_l(E,r)$ is a rapidly varying function of E. Again, however, in the energy range where the resonance strongly contributes to the capture cross section we have $E_R - \Gamma < E < E_R + \Gamma \approx E_R$. Therefore $P_l(E,r) \approx P_l(E_R,r)$, and the proton partial width will also be approximately constant. Energy dependence in the proton partial width will not affect our calculations of its magnitude at the resonance energy, and the assumption that the resonance width is much less than the resonance energy will be justified below.

4.3.1 The resonance strength.

The "thick target" yield from a resonance is the yield obtained when the energy loss of particles in the target causes the energy range of the particles in the target to extend from several resonance widths (Γ) above the resonance energy (E_R) to several widths below the resonance energy. This yield is given by the integrating the differential yield over the resonance energy [Fow48]:

$$y_{tot} = \int_{E_R-\Delta E}^{E_R+\Delta E} \frac{\sigma(E)}{STP(E)} dE. \tag{4.6}$$

In Equation 4.6, the condition of a "thick target" implies that the energy spread around the resonance energy (ΔE) is large. This allows us to replace the limits of the integral with $\pm\infty$. When the Breit-Wigner form of Equation 4.1 is put in the above integral, along with the assumption that the resonance is narrow enough to justify the assumption that the stopping power is constant ($STP(E) \approx STP(E_R)$), the solution may be reduced to [Fow48]:

$$y_{tot} = \frac{\pi}{2}\sigma_R\Gamma/STP(E_R) = \frac{4\pi^2}{2\kappa^2 STP(E_R)}\omega\gamma, \tag{4.7}$$

where $\omega\gamma$, the resonance strength, is equal to $\omega\frac{\Gamma_\gamma\Gamma_p}{\Gamma}$. ξ is again defined by Equation 4.2. This means that we may write the magnitude dy of the step function in the yield curve of Figure 4.1 for a detector subtending a fraction of the total solid angle of $d\Omega/4\pi$ and with an efficiency ϵ as:

$$dy = \frac{\pi\epsilon d\Omega}{2\kappa^2 STP(E_R)}\omega\gamma. \tag{4.8}$$

In Equation 4.8, respective values of ϵ, $d\Omega$, κ, $STP(E_R)$, and dy are 0.066, 1.01 Steradians, 7.55×10^{11} cm^{-1}, 2.12×10^{-14}eV cm^2/target nucleus, and 3.84×10^{-17}counts/incident proton. These values give a resonance strength for this new resonance of $\omega\gamma$=(4.5 $\pm$ 0.9) $\times$ 10^{-6} eV, approximately $1/3000^{th}$ the strength of the narrow resonance at E_p=278 keV.

4.3.2 The proton partial width.

The proton partial width, Γ_p, may be estimated in a relatively model-independent way with the assumption that $\Gamma_p \ll \Gamma_\gamma$. A model-dependent estimate, to follow, will justify this assumption. Under our assumption, $\Gamma \approx \Gamma_\gamma$, which reduces $\omega\gamma$ to $\omega\Gamma_p$. For this case we find that $\Gamma_p = 2.8 \times 10^{-5}/(2\mathrm{J}+1)$ eV. The maximum proton partial width is therefore 1.4×10^{-5} eV for J = 1/2.

Γ_p may be estimated according to the simple model that $\Gamma_p \approx 2\gamma^2 P_l$, where γ^2 is the reduced width of the state and P_l is the proton penetrability at 127 keV. γ^2 is given by $S_l\gamma_W^2$, where γ_W^2 is the Wigner limit for the reduced width of the excited state given by $\gamma_W^2 = 3\hbar^2/2\mu r^2 = 2.93$ MeV [Ang99]. μ and r are the reduced mass of the system in amu and the nuclear radius defined by Equation 4.5, respectively. Values of J^π are assumed for the 7.414 MeV state and minimum values of l are determined for proton capture to an excited state with that total angular momentum. Values

for $P_l(127\text{ keV})$ are determined by a TUNL-written FORTRAN code, which solves the Schrödinger equation for the incident proton in the Coulomb potential of the target nucleus. These values for P_l represent the probability that the proton will tunnel through the Coulomb barrier. Values of Γ_p/S_l for the first few possible J^π of the 7.414 MeV state are contained in Table 4.1. S_l (as determined in a following subsection) has not yet been factored in.

J^π	l	$P_l(127\text{ keV})$	Γ_p/S_l
$1/2^+$	0	1.4×10^{-6}	8.2 eV
$3/2^+$	0	1.4×10^{-6}	8.2 eV
$1/2^-$	1	3.0×10^{-7}	1.8 eV
$3/2^-$	1	3.0×10^{-7}	1.8 eV
$5/2^-$	1	3.0×10^{-7}	1.8 eV
$5/2^+$	2	3.7×10^{-9}	0.00059 eV

Table 4.1: Possible values of J^π for the 7.414 MeV excited state in ^{15}O, the associated angular momentum (l) for proton capture to such a state, the penetration function (P_l) for the necessary angular momentum at $E_p = 127$ keV, and estimates of Γ_p for the transitions.

4.3.3 The γ-ray partial width.

An estimate of the γ-ray partial width, Γ_γ, may be obtained from the Weisskopf estimates[Kra88] for a particular multipolarity of radiation (L). The estimates for the γ-decay constants of the electric multipole transitions of energy E_γ are generally given by:

$$\lambda_E(E_\gamma, L) = \frac{8\pi(L+1)}{L[(2L+1)!!]^2}\frac{e^2}{4\pi\epsilon_0\hbar c}\left(\frac{E_\gamma}{\hbar c}\right)^{2L+1}\left(\frac{3}{L+3}\right)^2 c r^{2L}, \tag{4.9}$$

with r defined as in Equation 4.5. The units of λ as defined above are s^{-1} for γ-ray energies in MeV. Weisskopf estimates of the magnetic multipole transition probability

are similarly given by:

$$\lambda_M(E_\gamma, L) = \frac{8\pi(L+1)}{L[(2L+1)!!]^2}\frac{e^2}{4\pi\epsilon_0\hbar c}(\mu_p - \frac{1}{L+1})^2(\frac{\hbar}{m_p c})^2 \times(\frac{E_\gamma}{\hbar c})^{2L+1}(\frac{3}{L+2})^2 c r^{2L-2}. \tag{4.10}$$

The probability of an electromagnetic transition depends on the transition matrix element of the appropriate multipole operator, m(σ,L), where σ denotes electric or magnetic transitions, and L is the angular momentum of the radiation. The transition matrix element is of the form:

$$m_{fi} = \int \psi_f^* m(\sigma, L)\psi_i dv, \tag{4.11}$$

where the integral is carried out over the volume of the nucleus emitting the γ-ray. The ground state is assumed to be a pure single-particle shell model state in making the Weisskopf estimate, with a spectroscopic factor of 1 [Sch86]. Since it is a transition to the ground state that has been observed in the present work, ψ_f is the ground state of ^{15}O. Any estimate of the transition probability based on the assumption that the state which contributes a resonance amplitude to the capture cross section is a single-particle state in ^{15}O should be multiplied by the spectroscopic factor, S_l, of the excited state. $\sqrt{S_l}$ is the fraction of the single-particle shell model wavefunction in the state. This will have the effect of reducing estimates of Γ_γ by a factor of S_l. The estimates of Γ_γ given by $\Gamma_\gamma = \hbar/\lambda$, where λ is given by either Equation 4.9 or Equation 4.10, may be found in Table 4.2. As in the estimate of the proton partial width, S_l has not yet been factored in. The actual γ-ray width of the state may be found by multiplying the estimates in Table 4.2 by the associated value of S_l for the state, assuming a particular value of J$^\pi$.

We will primarily focus on E1 and M1 radiation from this state for two reasons. First, the 7.414 MeV state would be more likely to decay to one of the intermediate

J^π	Radiation	Γ_γ/S_l
$1/2^+$	E1	163 eV
$3/2^+$	E1	163 eV
$1/2^-$	M1	15 eV
$3/2^-$	M1	15 eV
$5/2^-$	E2	0.040 eV
$5/2^+$	M2	0.0031 eV

Table 4.2: Possible values of J^π for the 7.414 MeV excited state in ^{15}O, the primary associated electromagnetic transitions to the $J^\pi = 1/2^-$ ground state, and estimates of Γ_γ for the transitions.

excited states in ^{15}O by E1 or M1 radiation than to decay to the ground state if such a decay involved a higher multipolarity. Secondly, the two shell model calculations which predict a "missing" state in the A=15 system both predict a $1/2^-$ state, which would decay to the ground state via an M1 transition.

4.3.4 The spectroscopic factor.

As can be seen in Table 4.2 and Table 4.1, values of Γ_p are considerably smaller than corresponding values of Γ_γ except for the case of $J^\pi = 5/2^+$. This fact justifies the model independent estimate of Γ_p. For the case for $J^\pi = 5/2^+$, $\Gamma_\gamma \ll \Gamma_p$, and the model independent estimate of Γ_p changes into a model independent estimate of Γ_γ. Factoring out S_l from both widths in either situation gives:

$$\omega\gamma = \xi S_l \frac{\Gamma_{\gamma,est}\Gamma_{p,est}}{\Gamma_{\gamma,est} + \Gamma_{p,est}}, \tag{4.12}$$

where $\Gamma_{\gamma,est}$ and $\Gamma_{p,est}$ are the pure single-particle estimates ($S_l = 1$) for the partial widths from Table 4.2 and Table 4.1. Assuming a value of J^π for the resonance, we may insert the corresponding values of $\Gamma_{\gamma,est}$ and $\Gamma_{p,est}$ from the tables into Equation 4.12. We may also insert the experimentally determined value of the resonance strength, $\omega\gamma$

$= 4.5 \times 10^{-6}$ eV (from Equation 4.8), and the appropriate value of ξ from Equation 4.2, leaving S_l as unknown to solve Equation 4.12 for. The spectroscopic factors for the first few possible values of J^π are presented in Table 4.3.

J^π	l	Radiation	S_l
$1/2^+$	0	E1	1.7×10^{-6}
$3/2^+$	0	E1	8.6×10^{-7}
$1/2^-$	1	M1	8.6×10^{-6}
$3/2^-$	1	M1	4.3×10^{-6}
$5/2^-$	1	E2	0.0011
$5/2^+$	2	M2	0.67

Table 4.3: Possible values of J^π for the 7.414 MeV excited state in ^{15}O, the minimum associated angular momentum (l) for proton capture to such a state, the strongest γ-decay mode for decay to the $1/2^-$ ground state of ^{15}O, and a calculation of S_l for the state.

4.4 Background yield simulation.

A quantitative understanding of the size of the background peak from the ^{70}Ge(n,γ)^{71}Ge reaction may be derived from a detailed simulation of the situation. There are many γ-ray transitions to be considered in the simulation. The relative strength of individual γ-ray transitions from thermal neutron capture are well documented, but a complete map of the cascades does not exist. Inspection of the γ-ray level branching ratios [Mar99] shows a common decay scheme for ^{71}Ge. The decay from 7.416 MeV to the ground state proceeds primarily through cascades involving three γ-rays. The typical energy ranges of these three γ-rays are 6-7 MeV, 0.8-1.2 MeV, and 200-500 keV. About 68% of all γ-rays emitted from neutron capture are γ-rays from 20 transitions near these energies, which are generally grouped in such 3-γ cascades. Other similar cascade schemes can be seen in cascades of lower probability, making a simu-

lation of the probability of capturing three such γ-rays a good approximation to use in estimating the efficiency of total calorimetry for the ^{70}Ge(n,γ)^{71}Ge reaction.

EGS4 was used to simulate the full-energy photopeak efficiencies of the 123% HPGe detector. The average efficiency was computed for capturing a 6.2 MeV γ-ray, a 0.9 MeV γ-ray, and a 300 keV γ-ray emitted at random directions from a point source inside the detector. The weighted average energy of the transitions from the 20 strongest transitions mentioned above which fell into the three distinct energy ranges (also mentioned above) were used to determine the three γ-ray energies used in the simulation, which conveniently add up to 7.4 MeV. The full-energy photopeak efficiencies for the three γ-ray energies were multiplied to form the total efficiency of calorimetry (capture of all three γ-rays) at that point. This was repeated for many random coordinates in the crystal and the results for the total efficiency were averaged over all the points to give an average probability of total calorimetry (resulting in a count in the 7.416 MeV background peak) of 0.00457 per thermal neutron capture by a ^{70}Ge nucleus. While that number appears small, it is 4 times higher than the probability of the single strongest transition and approximately 40 times higher than the probability of observing any single transition above 5.6 MeV, where the background might permit such a transition to be seen. This explains the lack of individual transitions in ^{71}Ge in our data.

The total expected background yield per hour in the background peak is then given by:

$$Y_{bkg}/hr = f_{nat}\rho N_A dA\sigma\phi(3600s)P_{abs}, \tag{4.13}$$

where f_{nat} is the fraction of ^{70}Ge in natural Ge, ρ is the density of Ge in mol/cm^3, N_A is Avogadro's number, ϕ is the measured thermal neutron flux of 28 thermal neutrons/m^2/s, d is the average distance a neutron will traverse in the crystal (7.95

cm), $A = 70$ cm^2 is the surface area, and $P_{abs} = 0.00457$ is the simulated probability that all γ-rays in the cascade will be absorbed in the crystal. The calculated background yield per hour is thus $Y_{bkg}/hr = 0.84$ counts/hour.

Dividing the counts in the 7.416 MeV background peak from the composite spectrum (before normalizing to time) by the time of the measurement gives an average measured background count rate of 0.816 ± 0.044 counts/hour, which compares very well with the simulated value. The fact that the rate derived from the simulation agrees to within the statistical error of the data indicates that it is possible to quantify the size of the background peak, and hence to rule out a significant source of background other than the ^{70}Ge(n,γ)^{71}Ge reaction.

4.5 The 278 keV resonance

The thick target yield of the 278 keV resonance is somewhat complicated by the contamination of the target surface. From the deconvolution fit to the data, we may reconstruct a profile of the contaminant level. The profile from the fit provides us with the fraction of contamination as a function of proton energy in the target, and can be used to correct the yield of γ-rays as a function of beam energy. The corrected yield may then be compared to the observed yield of γ-rays from the measurement at 300 keV. As will be shown below, the contamination turns out to have a negligible effect on the resonant yield of the ^{14}N(p, γ)^{15}O reacion at $E_p = 300$ keV.

Plotted in Figure 4.3 is a graph of the fractional contamination given by $N_{exp}(E)/N_{obs}(E)$ - 1 (solid line). $N_{exp}(E)$ is the expected number of counts from the deconvolution fit shown in Figure 3.12 at center of mass energy, E, and N_{obs} is the number of counts in the data at that energy. This data has been smoothed, but is essentially a statistical

Figure 4.3: Contamination profile of the target for the minitandem experiment as a function of beam energy (solid line) and a fit to the contamination layer (dashed line).

fluctuation around zero with a surface layer of contamination. The thin contaminant layer is due to vacuum hydrocarbons which freeze at the surface. This should result in a 100% surface contamination in a thin layer, but this is smeared into a nearly Gaussian shape by the detector response function (FWHM=0.2%=12 keV). This fit to the contaminant fraction will allow us to examine the importance of the contamination to the γ-ray yield of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction at 300 keV.

A fit to the previous world data set[Sch86, Pix57, Hes57, Heb63] for the astrophysical S-factor of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction at energies near the E_p = 278 keV resonance (and extending to E_p = 0 keV) is shown in Figure 4.4. The fit was constructed with a polynomial background and a Breit-Wigner cross section for the E_p = 278 keV resonance. The γ-ray yield from a proton of a given kinetic energy, E,

Figure 4.4: Fit to the low energy S-factor world data set.

stopping in the target can be obtained by evaluating:

$$Y(E) = 3\frac{\epsilon d\Omega}{4\pi}\frac{q}{1.602\times 10^{-19}}\int_{E_{init}}^{0}\frac{S(E)}{STP(E)}(e^{-2\pi\eta})/E_{cm}dE, \qquad (4.14)$$

where E_{cm}=14E/15 for the ^{14}N(p,γ)^{15}O reaction. The factor of 3 in Equation 4.14 accounts for the 3 deuterium nuclei in each target molecule. This results in a yield curve as a function of initial proton energy shown in Figure 4.5.

As can be seen in Figure 4.5, the bulk of the yield of the ^{14}N(p,γ)^{15}O reaction (97.9%) will come from protons whose energies are below E_{cm} = 285 keV in the lab frame, due to the presence of the E_p = 278 keV resonance. In the center of mass frame for the ^{2}H(p,γ)^{3}He reaction, this corresponds to E_{cm} = 190 keV. As Figure 4.3 shows, this is the lower energy cutoff of the contaminant layer. Because the maximum contamination level is 35%, we estimate the thickness of the contamination on the target surface to be only 4 keV. The effect of this dead layer on the yield of the ^{14}N(p,γ)^{15}O reaction is below 1% and may be safely neglected. Summing over the three major decay γ-ray peaks shown in Figure 3.13, the measurement yielded 3476 $\pm$ 60 counts from the ^{14}N(p,γ)^{15}O reaction. This is 6% higher than the expected

Figure 4.5: Plot of the integrated yield of γ-rays from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction due to protons of energy E_{init} stopping in a thick target of ND_3.

3289 counts shown in Figure 4.5, which is within the estimated error of the effective solid angle (6%) for the detector. If the charge determined from the deconvolution fit to the $^{2}\mathrm{H}(\mathrm{p},\gamma)^{3}\mathrm{He}$ peak (0.0309 C) is used, rather than the charge measured by current integration, the difference shrinks to 3%. The level of accuracy achieved for the E_p = 278 keV resonance cannot be achieved for the E_p=127 keV resonance, but indicates the power of the simultaneous measurement and the validity of the data analysis.

4.6 Comparison to previous work.

4.6.1 Previous $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ studies.

One previous study of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction which could have previously revealed the presence of the 7.414 MeV state in ^{15}O was done with a solid target of titanium nitride (TiN)[Hes57] . Numerous studies have been done near the E_p

= 278 keV resonance, but all these studies stop at a lower energy limit of E_p = 200 keV[Sch86][Pix57][Heb63] . The experiment, in 1957, was performed at E_p = 100 - 135 keV in 5 keV steps, and did not report seeing the effects of the 127 keV resonance. Aside from the obvious age of the data set, there are several reasons why this experiment may not have detected this resonance. The primary candidate for this arises from the possibility of a downward shift in the energy of the protons actually hitting the ^{14}N.

There are several sources of error in the measurement of the beam energy. The first source of error in the beam energy lies in the beam transport. The voltage of the ion source was measured with a precision voltage divider and a potentiometer, which gives the average voltage. A 6 kV peak-to-peak ripple in the bias supply of the ion source results in a deviation about the true average voltage, creating the possibility that protons with up to 3 kV less energy were tuned to the target. This opens up the possibility that the beams incident on the target for the measurements at E_p = 130 and 135 keV were already below the stated energy.

The second source of error can be found in the target itself. Two contaminant layers add to shift the beam energy downward from whatever energy is incident on the target surface. First, the incident protons could encounter a layer of carbon on the surface of the target. Furthermore, the 25 mA of beam stated in the paper corresponds to 3 kW of beam heating at proton energies above the resonance energy. This heat is deposited in a few micrograms of the TiN target, heating it. This can cause the breakdown of hydrocarbons in the vacuum system at the surface (depositing a carbon layer), or possibly depleting the ^{14}N in the target surface.

In their paper, Hester and Lamb present a composite of the decay curves (counts in their detector as a function of time after each bombardment) from all measurements

Figure 4.6: Data from the composite decay curve (β decay rate as a function of time) of Hester and Lamb, with two fits. The data represents the sum of data sets at all measured energies. The first fit (solid line) includes contributions from both the decay of ^{15}O and ^{13}N. The second fit (dashed line) assumes a pure ^{15}O source.

in their experiment [Hes57]. Shown in Figure 4.6 are the data (obtained from a scan of Figure 2. in [Hes57]) of the "composite decay curve", along with two fits. The solid line represents a fit including residual ^{15}O ($t_{1/2}$ = 122 sec) from the ^{14}N(p, γ)^{15}O reaction and ^{13}N ($t_{1/2}$ = 10 min) from the ^{12}C(p,γ)^{15}O reaction, which would occur if a carbon layer was present. The dashed curve is a fit which assumes (as Hester and Lamb do in their data analysis) only ^{15}O as a source of counts in the detector. The exact value of χ^2/DOF for solid curve in Figure 4.6 depends on the scale of the composite decay curve (not included in the figure from which the data were scanned), while for the dashed curve χ^2/DOF is 7 times this value independent of this scale. Clearly room exists for a significant surface contamination of the target, but details about the use of individual targets in the experiment are unavailable. Therefore, a calculation of the thickness of such a contaminant layer for any single measurement is impossible. The "composite decay curve" only gives information about the average

carbon contamination for all the measurements.

The second issue involves questions about the behavior of a water-cooled TiN target surface under 3 kW of heating. Experience at TUNL with TiH targets with beam currents 3 orders of magnitude lower than those of Hester and Lamb [Hes57] suggests that water cooling is not sufficient to maintain the surface temperature of such a target. Without exact knowledge of the thickness of their titanium "sheets" it is not possible to calculate the exact surface temperature of their targets. An estimation can be made from the assumptions of a standard thin Ti sheet (0.063 inches thick) at the stated beam spot diameter (0.75 inches), perfect contact with cooling water at 300 K, and a small cooling component from blackbody radiation at the surface. Titanium is a poor conductor of heat, and even this thin layer would reach steady-state surface temperatures of 1150 K. Under such a proton bombardment (4.5 incident protons per molecule of TiN in the stopping layer in one 360 second measurement), a surface layer depleted in nitrogen may have formed. This layer would have the effect of further reducing the energy of the incident protons before they reached the TiN target layer. The possibility of thin contaminant layers is of particular importance because the stopping power peaks near 100 keV for protons in almost all materials, and therefore such layers have the greatest stopping effect on protons of these energies.

One further detriment to the sensitivity of the previous $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ experiment to a resonance at $E_p = 127$ keV is that it relies on a plastic scintillator to count the $\beta+$ decays of the residual ^{15}O nuclei in the target. The accuracy of this measurement relies on the retention of residual ^{15}O in the surface of the target material. ^{15}O formed in the thin surface layer that stops the beam may have escaped the target at the temperatures calculated above, or the activated surface layer may have been sputtered away by the large number of protons.

For all these reasons, we believe our measurement is more reliable than the previous measurement of the $^{14}N(p,\gamma)^{15}O$ reaction at these energies. It is also worth noting that the resulting number arrived at by Hester and Lamb in this experiment is 2.5 times the currently accepted value at 100 keV, based on extrapolations and models.

4.6.2 Stripping reactions.

Experiments involving stripping reactions, such as a neutron time-of-flight (TOF) measurement of the $^{14}N(d,n)^{15}O^*$ reaction, are also sensitive to the presence of energy levels in this energy range. An example of such a TOF measurement is presented in Figure 4.7 [Bom71]. The TOF spectrum in Figure 4.7 was obtained using a pulsed beam of 6 MeV deuterons on a target of 210 $\mu g/cm^2$ thick layer of melamine (a carbon-nitrogen compound). Neutrons from the $^{14}N(d,n)^{15}O^*$ reaction were observed by a neutron detector 7.36 m from the target and at an angle of 37.5° relative to the beam. The population of excited states in ^{15}O results in neurons with energies corresponding to the state created. A Time to Amplitude Converter (TAC) converts the time between a "start" signal from the target (created by the incident pulse of neutrons) and a "stop" signal from the neutron detector into an output pulse. The amplitude of the output pulse is proportional to the time between the start and stop pulses. The peaks in the spectrum of these pulses, shown in Figure 4.7, correspond to energy levels in the residual ^{15}O nucleus. The size of these peaks is proportional to the transition matrix element for populating that state, and hence is proportional to its spectroscopic factor.

A background peak (origin unknown) at 7.4 MeV is marked to show the general level of sensitivity of these experiments. From the spectroscopic factors derived above,

we calculate that the labeled peak is 250 times the size we would expect of a peak from the 7.414 MeV excited state in ^{15}O. Even more sensitive experiments using the ^{14}N(^{3}H,d)^{15}O* reaction in a charged particle spectrometer have a sensitivity only down to a spectroscopic factor of 10^{-5} for s-wave capture (and worse for higher-l partial waves)[Cha01]. Therefore, we believe stripping reactions have not uncovered the existence of this newly discovered resonance because its spectroscopic factor is simply too small. This is not an unheard of occurrence in nuclear physics, the TUNL high resolution laboratory is currently mapping dozens of such unknown levels in heavier nuclei, but the discovery of a new energy level is somewhat unusual in such a light nucleus where there are few energy levels and those levels are generally separated by at least hundreds of keV.

Figure 4.7: Data from a TOF measurement of the ^{14}N(d,n)^{15}O* reaction[Bom71]. A peak (believed to be background) at 7.4 MeV is labeled.

Chapter 5

Direct Plus Resonance Model Calculations

In this chapter, the direct-capture-plus-resonances (DCPR) model is applied to the data [Wel80]. This model is used not so much to fit the data itself (although the potentials involved are adjusted to reproduce the observed data) as to provide some predictions of angular distributions of γ-rays and analyzing powers from the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction at the resonance energy. The DCPR model calculation can also provide a quantitative idea of how this new resonance affects capture to excited states in ^{15}O. These calculations should prove useful for anyone planning future experiments on this new excited state of ^{15}O.

5.1 The direct capture model.

The direct capture (DC) model describes a nuclear reaction in terms of a transition matrix element between the initial (scattering) and final (bound) states, involving a suitable Hamiltonian H. This process is a direct transition between the free-particle (scattering state) and the bound state in a single step, as described in the introduction. No intermediate, or compound, nucleus is formed. The timescale for such an interaction is exceedingly short, $\approx 10^{-22}$ seconds. A detailed description of this model may be found in [Rol73, Pot77, Wel80].

The DC model is based on first-order perturbation theory, where the Hamiltonian is written as:

$$H = H_0 + H', \tag{5.1}$$

in which H_0 contains the kinetic and potential energy terms of the projectile and target particles while H' describes the perturbative part of the Hamiltonian represented by the electromagnetic interaction operator. The interaction between the initial state ($|\chi>$) and the final state ($|\psi_b>$) is written as the transition matrix element T_{fi}, and is given by:

$$T_{fi} = < \psi_b | H | \chi > . \tag{5.2}$$

Fermi's Golden Rule allows us to immediately write the transition rate for such an interaction as:

$$W_{fi} = \frac{2\pi}{\hbar} |T_{fi}|^2 \rho_f, \tag{5.3}$$

where ρ_f is the density of final states accessible via the interaction H'. For capture to a final state of angular momentum J via the emission of a photon, the density of

states is given by:

$$\rho_f = V^{-1} \sum_{m_J, C} \delta(E_i - E_f - E_\gamma), \tag{5.4}$$

where m_J represents the magnetic substates of the final state, C represents the possible polarization states of the photon, and V is an arbitrary normalization volume. The δ-function simply insures that the γ-ray has an energy equal to the difference between the energies of the system before and after the reaction (conservation of energy). The cross section may be expressed in terms of the transition rate as:

$$\frac{d\sigma}{d\Omega} = \frac{W_{fi}}{\Phi_i} \tag{5.5}$$

where Φ_i is the flux of incident particles. The incident flux can be written as $\Phi_i = n_i v_i$, where n_i is the volume density of incident particles and v_i is their velocity. With the same normalization volume V for Φ_i as for ρ_f, $\Phi_i = v_i/V$. We may then write the differential cross section as the average of cross sections over all possible magnetic substates of the projectile and target (m_p and m_t, respectively):

$$\frac{d\sigma}{d\Omega} = \frac{2\pi}{\hbar v_i} \frac{1}{(2j_p+1)(2j_t+1)} \sum_{m_p, m_t, m_J, C} |T_{fi}|^2 \tag{5.6}$$

We may add resonance terms to the DC transition amplitude to form the DCPR transition amplitude. The total radial part of the transition amplitude in the DCPR model is generally written as:

$$T_{fi} = < \psi_b | H'_L | \chi > + \sum_n \frac{< \psi_b | F_n | \chi >}{(E - E_{R_n}) + i\Gamma_n/2}, \tag{5.7}$$

where H'_L is the radial part of the electromagnetic interaction operator given in Equation 5.8, F_n is a strength parameter for the n^{th} resonance, and E_{R_n} is the energy of this resonance. The strength parameters of the resonances are fit to data, and we add

as many resonances as necessary to fit the data. The radial matrix elements represented by Equation 5.7 are integrated numerically by the FORTRAN code HIKARI. Details of the physics considered in this calculation are described below, and results of the calculation are described in Section 5.2.

5.1.1 Electromagnetic interaction Hamiltonian.

Evaluation of Equation 5.6 requires that we specify the interaction operator. For the electromagnetic interaction is generally given by [Jac99]:

$$H'_L = -\frac{\epsilon_L}{c} \int \vec{J}(\vec{r}) \cdot \vec{A}(\vec{r}) d\vec{r} \tag{5.8}$$

where $\vec{J}$ is the nuclear charge current density due to motion of the protons and $\vec{A}$ is the vector potential of the electromagnetic field. Here we also include a scaling factor, the recoil effective charge (ϵ_L) for the appropriate multipolarity of radiation for which we wish to create an interaction operator (H'_L). The effective charge, ϵ_L, is given by $(A_2^L Z_1 + (-1)^L Z_2 A_1^L)/(A_1 + A_2)^L$, where the A's are the masses of the projectile (subscripted 1) and target (subscripted 2) nuclei, the Z's are their respective charges, and L is the multipolarity of the radiation [Wel80]. This effective charge scales the strength of the direct capture matrix element, accounting for the effective charge of the incident nucleon due to the recoiling nucleus.

The multipole expansion of the radiation field allows this interaction operator to be enormously simplified because we may retain only the dominant terms. The terms considered here will be only the electric and magnetic dipole transitions. Because of the low energy of this capture reaction and the prohibitive angular momentum barrier, we shall only consider capture by s-wave (allowed values of $J^\pi = 1/2^+, 3/2^+$) and p-wave (allowed $J^\pi = 1/2^-, 3/2^-, 5/2^-$) protons. The resonant contribution to the

electromagnetic transition could then consist of E1, M1, E2, M2, and M3 radiation. Either E1, M1, or E2 contributions would be the lowest-order multipole in any of these cases, and the lowest order multipole possible is generally the dominant one. We therefore restrict ourselves to dipole (E1 and M1) terms only for direct capture.

For *electric* transitions, we may use Siegert's theorem [Sie37] to replace $\vec{J}(\vec{r})$ with a charge density $\rho(\vec{r})$. Neglecting meson exchange currents, and invoking the long-wavelength approximation for the low energy ($\lambda \approx 168$ fm for a 7.4 MeV photon) gives an E1 transition operator whose radial part is given simply by r. Higher-order multipole operators (E2, and E3) have been derived but their contributions to the cross section are not considered in this simple picture. It has been shown that the total cross section for electric radiation of multipolarity L from a spin zero target in this picture [Wel80] may be finally reduced to:

$$\sigma_T^L = 4\pi\epsilon_L^2 (\frac{197}{137}) \frac{k_f}{E_{cm} k_i} \frac{(2j_f+1)}{(2j_i+1)} B_L^2 \sum_{l_i, j_i} |T^L_{l,j;l_i,j_i}|^2, \tag{5.9}$$

$$B_L^2 = \frac{L+1}{(2L+1)L} \frac{k_f^{2L}}{[(2L-1)!!]^2}, \tag{5.10}$$

$$T^L_{l,j;l_i,j_i} = i^{(l_i-l-L)} C(jLj_i, \frac{1}{2}0\frac{1}{2}) \sqrt{S_l} < u_b(r)|r^L|\chi_{l_i j_i}(r) > . \tag{5.11}$$

In Equation 5.9-Equation 5.11, k_i and k_f are the incident and outgoing wave numbers, l_i and j_i are the orbital and total angular momentum quantum numbers of the incident nucleon, l and j are the orbital and total angular momentum quantum numbers for the bound particle, and $C(jLj_i, \frac{1}{2}0\frac{1}{2})$ is the appropriate Clebsch-Gordan coefficient. In Equation 5.11, $|u_b(r) >$ and $|\chi_{l_i j_i}(r) >$ are the radial parts of the bound and scattering state wavefunctions, respectively. The cross section for a target with angular momentum j_t, a single particle in the final state with total angular momentum j_{spf}, and a total final state angular momentum j_f is multiplied by a factor of $\frac{2j_f+1}{(2j_t+1)(2j_{spf}+1)}$ [Wel82]. The problem is therefore reduced to a set of multiplicative

factors and the radial integral part of Equation 5.11.

For M1 transitions, the single-particle M1 operator for *proton* capture, O_{M1}, which is given in [Wel82] reduces to:

$$O_{M1} = \mu_0(g_t\vec{J}_t + \mu\vec{\sigma} + \alpha\vec{L}), \tag{5.12}$$

where:

$$\mu_0 = \frac{e\hbar}{2m_pc}, \tag{5.13}$$

in which m_p is the proton mass, and:

$$\alpha = \frac{Z + A^2}{A(1 + A)}, \tag{5.14}$$

where A and Z are the mass and charge of the target nucleus, respectively. In Equation 5.12, g_t and $\vec{J}_t$ are the gyromagnetic ratio and angular momentum of the target nucleus, respectively. The matrix element of $g_t\vec{J}_t$ vanishes in the direct capture model, because that operator contains only the internal coordinates of the core, which have been integrated out in this model. $\vec{L}$ is the operator for the relative orbital angular momentum of the target nucleus and projectile, μ is the magnetic moment of the bound single particle, and $\vec{\sigma}$ is the Pauli spin matrix. It has been shown [Wel82] that the total cross section for direct M1, $\sigma_T(M1)$, is evaluated to be:

$$\sigma_T(M1) = \frac{2j_f + 1}{(2j_i + 1)(2j_{spf} + 1)}\frac{2\pi}{3}\frac{k_f^3}{E_{cm}k_i}\frac{e^2(\hbar c)^2}{m^2c^4}(2\mu - \alpha)^2\frac{l_i(l_i + 1)}{2l_i + 1}|T(M1)|^2, \tag{5.15}$$

where:

$$T(M1) = \int u_b(r)\chi(r)dr. \tag{5.16}$$

In Equation 5.15, j_{spf} is the total angular momentum of the bound particle in the final state, j_f is the total angular momentum of the residual nucleus, and other quantities are defined in the same way as those in Equation 5.9.

5.1.2 Wavefunctions.

In order to calculate the cross section, a bound (final) state wavefunction and a scattering (initial) state wavefunction are also required. The partial waves of the wavefunction of the incident proton *before* being distorted in the optical model potential are given by the partial wave expansion of the incident plane wave [Gri95]:

$$e^{ikz} = \sum_{l=0}^{\infty} i^l (2l+1) j_l(kr) P_l(cos\theta), \tag{5.17}$$

where the $j_l(kr)$ are spherical Bessel functions and $P_l(cos\theta)$ are the Legendre polynomials. We use the optical model potential to generate the distorted partial waves of the scattering state wavefunction. The potential, V_{OM}, was taken to have the form:

$$\begin{aligned} V(r) = -V_0 f(a_0, r_0, r) + i4a_d V_{im} \frac{d}{dr} f(a_{im}, r_{im}, r) \\ +(\frac{\hbar}{m_\pi c})^2 \frac{V_{SO}}{r} \frac{d}{dr} f(a_{SO}, r_{SO}, r) \vec{L} \cdot \vec{S} + V_{Coul}(r, r_C). \end{aligned} \tag{5.18}$$

In Equation 5.18 V_0 specifies the strength of the real, central nuclear potential. V_{im} specifies the strength of the imaginary nuclear potential, which represents absorption into other channels and is taken to have a derivative Woods-Saxon form (surface peaked). V_{SO} specifies the strength of the spin-orbit interaction. r_C represents the Coulomb radius parameter of the nucleus. The function $f(r_x, a_x, r)$ is the Woods-Saxon form factor, which is given by:

$$f(r_x, a_x, r) = \frac{1}{1 + exp(\frac{r - r_x A^{1/3}}{a_x})}. \tag{5.19}$$

The imaginary part of the potential accounts for the absorptive (inelastic) process, which we expect to be small at these energies because few inelastic channels (capture to subthreshold states) are open. The solution to the radial Schrödinger equation for these partial waves gives the radial part of the incident proton wavefunction. The

parameters of the potentials defined by Equation 5.18, which are used to used to generate the distorted waves and final state wavefunction, are given in Table 5.1.

	Distorting potential	Bound State Potential
V_0	50 MeV	46 MeV
r_0	1.4 fm	1.4 fm
a_0	0.65 fm	0.65 fm
V_{im}	5 MeV	N/A
r_{im}	1.4 fm	N/A
a_{im}	0.65 fm	N/A
V_{SO}	1.5 MeV	1.5 MeV
r_{SO}	1.4 fm	1.4 fm
a_{SO}	0.65 fm	0.65 fm
r_C	1.2 fm	1.2 fm

Table 5.1: Optical model and bound-state potential parameters used by HIKARI.

The radial part of the bound state wavefunction $|u_b>$, is formed by the solution to the radial Schrödinger equation for a bound state in a potential defined by Equation 5.18 for the bound projectile, and its magnitude is scaled by the fraction of single-particle wavefunction that the bound projectile actually has (S_l, the spectroscopic factor). The depth of the potential well given in Table 5.1 was calculated by adjusting the well depth to reproduce the binding energy of the bound proton (the Q-value of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction), and the form of the wavefunction is allowed to have only two nodes (one at $r = 0$ and one at $r = \infty$). It thereby represents the lowest energy $l = 1$ single-particle proton state possible in ^{15}O. The radial part of the bound-state wavefunction, calculated by HIKARI, is plotted in Figure 5.1. For capture to excited states, this potential is adjusted to fit the binding energy of the proton in the excited state instead of the ground state. This binding energy is given by the Q-value of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction minus the energy of the excited state itself.

Figure 5.1: Radial part bound-state wavefunction for the ground state of ^{15}O.

5.1.3 Angular distributions.

In order to investigate the angular distribution of the cross section for capture to the ground state, we expand the cross section in terms of Legendre polynomials (P_k) as:

$$\sigma(\theta) = A_0(1 + \sum_{k=1}^{\infty} Q_k a_k P_k(cos\theta)), \tag{5.20}$$

where expressions for the amplitude A_0 and a_k coefficients are derived in terms of the reduced transition matrix elements, $T^L_{l_j;j_i j_i}$, by [Sey83]. These expressions are used by HIKARI to calculate the a_k coefficients, which are then used to calculate the angular distribution of γ-rays from the target according to Equation 5.20. The Q_k's are experimental finite geometry correction coefficients (defined by experimental apparatus). For these calculations, all Q_k's are set equal to 1.

If measuredangular distributions are to be compared to these calculations, the Q_k coefficients will have to be factored in. The Q_k coefficients for detection of a γ-ray which includes the entire HPGe response function are given by Equation 5.7

in [Sch95], but here we are interested in only the full-energy peak of the response function. For this case the Q_k coefficients are obtained by running a Monte-Carlo simulation in which γ-rays are emitted from a source at angles β_i relative to the HPGe detector's symmetry axis. If N γ-rays in the simulation deposit their full energy in the HPGe detector, the Q_k coefficients are given by [Bru94]:

$$Q_k = \frac{1}{N}\sum_{i=1}^{N} P_k(cos(\beta_i)), \tag{5.21}$$

where the P_k functions are the Legendre polynomials as in Equation 5.20.

For a HPGe detector with dimensions identical to those of the HPGe detector used in the present experiment, and situated 5 cm from a source of 7.4 MeV γ-rays, the first five Q_k coefficients for the full-energy peak have been calculated using Equation 5.21 and EGS4. These coefficients are listed in Table 5.2.

Q_0	Q_1	Q_2	Q_3	Q_4
1	0.95	0.87	0.75	0.61

Table 5.2: First five values of Q_k coefficients for the HPGe detector geometry used in the present experiment.

Beams of polarized protons are often used in radiative capture reactions to investigate the J^π of states. This is typically done by observing the difference in the number of detected photons from a reaction for the case where the incident protons are polarized to be spin-up (N_+) versus the number detected for spin-down (N_-). "Spin-up" is defined by the Madison convention to be parallel to the Y-axis, which is defined by $\vec{k}_{in} \times \vec{k}_{out}$, as shown in Figure 5.2, and "spin-down" is anti-parallel to the Y-axis. The angle of the detector, θ, in the reaction plane is also shown in Figure 5.2. This allows us to define the analyzing power A_y:

$$A_y(\theta) = \frac{N_+ - N_-}{P_-N_+ + P_+N_-}, \tag{5.22}$$

Figure 5.2: The Madison convention, defining axes for polarized protons. The photon polarization vector convention is not shown here.

where $P_{\pm}$ are the polarizations of the spin-up and spin-down proton beams, respectively. The product of the analyzing power and the cross section may also be expanded in terms of the first associated Legendre polynomials, $P_k^1(cos\theta)$. That expansion is given by:

$$\sigma(\theta)A_y(\theta) = A_0(\sum_{k=1}^{\infty} Q_k b_k P_k^1(cos\theta)), \tag{5.23}$$

where expressions for the b_k coefficients are also defined by Seyler and Weller [Sey83], and are used by HIKARI to calculate the angular distribution of the analyzing power, $A_y(\theta)$.

For purposes of notation we may specify the ^{14}N(p,γ)^{15}O reaction in terms of the angular momenta of its component parts as:

$$a(x, L)c, \tag{5.24}$$

and the orbital angular momentum, l, of the incident proton. We adopt the channel-spin coupling scheme [Wel80]:

$$\vec{x} + \vec{a} = \vec{s}, \qquad \vec{l} + \vec{s} = \vec{J}, \qquad \vec{L} + \vec{c} = \vec{J} \tag{5.25}$$

The transition matrix elements (TMEs) are labeled by the quantum numbers $^{2s+1}l_j$ in this representation, and each is associated with certain multipolarities (here, only E1 or M1) of the outgoing γ-ray.

5.1.4 Allowed transitions.

The appropriate TMEs to be considered for the ^{14}N(p,γ)^{15}O reaction are determined by examination of the angular momenta and parities of the constituent particles. The incident proton has a spin-parity of $1/2^+$, while the ground state of

^{14}N is 1^+. For capture to the ground state ($1/2^-$) of ^{15}O, the γ-ray will carry away 1 unit of angular momentum, since we only consider E1 and M1 radiation here. E1 radiation has odd parity, so capture to the ground state (odd parity) via an E1 transition must be of even parity (l=0,2,4,...). M1 radiation has even parity, so capture to the ground state via an M1 transition must have odd parity (l=1,3,5,...). Considering only capture to the ground state (which we observe) via E1 or M1 (as appropriate) radiation, we have 6 contributing TMEs. For p-wave capture, J = s$\pm$1,0 and s = {1/2,3/2}. The four TMEs for p-wave capture and an M1 (no J = 5/2, L=2 transition) transition to the ground state are thus $^2p_{1/2}$, $^4p_{1/2}$, $^2p_{3/2}$, and $^4p_{3/2}$ TMEs. For s-wave capture, J = s and s = {1/2,3/2}. Therefore the two contributing TMEs for s-wave capture and E1 radiation are $^2s_{1/2}$ and $^4s_{3/2}$.

5.2 Numerical calculations.

The radial integral for $T^L_{l,j;l_i,j_i}$ is evaluated numerically by HIKARI. It is used in Equation 5.9 to solve for the cross section for electric transitions, and in Equation 5.15 for magnetic transitions. The complex TMEs are then used to calculate a_k, and b_k coefficients using the expressions of [Sey79, Wel80] (see Appendix B). From these coefficients, angular distributions of cross sections and analyzing powers are calculated using Equation 5.20 and Equation 5.23.

5.2.1 S-factors from HIKARI.

In the previous work of [Sch86], Schröder's extrapolation of the S-factor for the ^{14}N(p,γ)^{15}O reaction to low energies indicates that proton capture via the ^{14}N(p,γ)^{15}O reaction at energies below 200 keV is dominated by capture to the 6.176 MeV and 6.793

Figure 5.3: Level diagram of the important levels and transitions in the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction below E_p=200 keV.

MeV excited states. These states, and their J^π values, are shown in Figure 5.3. The J^π of the new resonance is unknown, but based on theoretical results (Chapter 6) we will assign it a value of $J^\pi = 1/2^-$ for the purpose of performing S-factor calculations. Its width is such that it will not significantly affect the calculation at energies more than a few keV away from the resonance energy. Although capture to the excited states at 6.793 and 6.176 MeV is included here for completeness in the calculation of the total S-factor at off-resonance energies, we have not observed capture to these states which could be attributed to the new resonance. Therefore, we will only consider capture to the ground state in calculations of the angular distributions of γ-rays (below).

Figure 5.4: S-factor for capture to the ground state (long dashes), 6.176 MeV excited state (dot-dashed line), 6.793 MeV excited state (short dashes), and their sum (solid line) as a function of proton energy from E_p = 0-0.4 MeV.

The HIKARI calculations were run with both the ground state and the two subthreshold states at 6.176 and 6.793 MeV in Figure 5.3 as final states over a range of incident proton energies from 0-350 keV. For the ground state calculation, the ground

state of ^{15}O was considered to be a pure $p_{1/2}$ single particle state with a spectroscopic factor of 1. Capture to the 6.176 MeV excited state was also done using a $p_{1/2}$ single particle state, but with a spectroscopic factor of 0.1 [Bom71, Sch86]. Capture to the 6.793 MeV state was done twice, because it is considered to be an admixture of $s_{1/2}$ and $d_{3/2}$ single particle states with spectroscopic factors of 0.49 and 0.25, respectively [Bom71]. The results of the two calculations for capture to the 6.793 MeV state were added to obtain the total cross section for capture to that final state. The results of the capture calculations to each of these three final states are shown in Figure 5.4.

The quantum numbers and spectroscopic factors of the final states in the HIKARI input file were adjusted according to the known level assignments for these states. While other bound states were examined as final states, only these three were found to contribute significantly to the cross section. Our calculation indicates much the same behavior as the extrapolation of [Sch86] above 200 keV, but a total off-resonance S-factor below 200 keV which is in better agreement with the recent extrapolation of [Ang01], which ignores both the data of [Hes57]. Recent studies of the 6.793 MeV state [Cha01, Ang01] indicate that it has a width substantially lower than previously believed, and therefore its high-energy tail should not strongly influence the capture cross section. Because of this, we do not include it as a subthreshold resonance in our calculation of the capture cross sections below 200 keV in the ^{14}N(p, γ)^{15}O reaction. The only resonances included in the present calculation are those at E_p = 127 and 278 keV. The resulting total S-factor at zero energy which we calculate for the ^{14}N(p, γ)^{15}O reaction is 1.5 keV b, compared with a value of 1.8 keV b by [Ang01]. The value of the S-factor at zero energy was previously given by [Sch86] as 3.2 keV b. The strength of the E_p = 127 keV resonance was adjusted (by adjusting a strength factor of that resonance in the HIKARI input file) to create an area under the reso-

nance curve consistent with the step size in our data according to the formulae laid out in Chapter 4.

Figure 5.5: The total S-factor from Figure 5.4, along with data of various authors on the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction.

The total S-factor data from previous experimental studies of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction are plotted against the total S-factor from the HIKARI calculation in Figure 5.5. The calculation is in good agreement with the world data set at most energies, except for the data of Hester and Lamb. This is possibly due to the lack of inclusion of carbon as a significant contaminant in their analysis of their data, as mentioned earlier. Such a contamination could have caused the inclusion of extra counts in the data due to the protons incident on the surface layer of carbon. The $^{12}C(\mathrm{p},\gamma)^{13}\mathrm{N}$ reaction has a much higher cross section than the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction, and extra counts from the β-decay of $^{13}\mathrm{N}$ would lead to an S-factor which was artificially high. A surface layer of carbon would also explain the lack of observation of a resonance at 127 keV due to a downward energy shift in the incident protons before they reached the nitrogen beneath the carbon layer.

5.2.2 Angular distributions from HIKARI.

The angular distributions of the cross section and analyzing power were calculated as a function of the incident proton energy. In the case of capture to the ground state of ^{15}O, only the resonances at E_p = 127 and 278 keV were included. The angular distributions were integrated from E_p = 0-140 keV, and this integration was weighted by the total cross section at each energy and the inverse of the stopping power of ammonia at that energy. These calculations represent the angular distributions which would be observed for an experimental setup identical to the current one, after accounting for the Q_k coefficients.

The effect of the resonance on the angular distribution will depend on its J^π. Using expressions from Appendix B, we may write down a_1, a_2, b_1, and b_2 coefficients

of Equation 5.20 and Equation 5.23 in terms of their contributing TMEs:

$$\begin{aligned}
1.0 = {} & 2.0^2 s_{1/2}^2 + 4.0^4 s_{3/2}^2 + 2.0^2 p_{1/2}^2 + 2.0^4 p_{1/2}^2 \\
& + 4.0^2 p_{3/2}^2 + 4.0^4 p_{3/2} \qquad (normalization) \\
a_1 = {} & -4.0^2 s_{1/2}\ {}^2p_{1/2} cos(\phi_{2s_{1/2}} - \phi_{2p_{1/2}}) + 4.0^2 s_{1/2}\ {}^2p_{3/2} cos(\phi_{^2s_{1/2}} - \phi_{^2p_{3/2}}) \\
& - 2.828^4 s_{3/2}\ {}^4p_{1/2} cos(\phi_{^4s_{3/2}} - \phi_{^4p_{1/2}}) - 8.944^4 s_{3/2}\ {}^4p_{3/2} cos(\phi_{^4s_{3/2}} - \phi_{^4p_{3/2}}) \\
a_2 = {} & -4.0^2 p_{1/2}\ {}^2p_{3/2} cos(\phi_{^2p_{1/2}} - \phi_{^2p_{3/2}}) + 1.265^4 p_{1/2}\ {}^4p_{3/2} cos(\phi_{^4p_{1/2}} - \phi_{^4p_{3/2}}) \\
& - 2.0^2 p_{3/2}^2 + 1.6^4 p_{3/2}^2 \\
b_1 = {} & 1.33^2 s_{1/2}\ {}^2p_{1/2} sin(\phi_{^2s_{1/2}} - \phi_{^2p_{1/2}}) + 1.886^2 s_{1/2}\ {}^4p_{1/2} sin(\phi_{^2s_{1/2}} - \phi_{^4p_{1/2}}) \\
& + 0.667^2 s_{1/2}\ {}^2p_{3/2} sin(\phi_{^2s_{1/2}} - \phi_{^2p_{3/2}}) - 2.981^2 s_{1/2}\ {}^4p_{3/2} sin(\phi_{^2s_{1/2}} - \phi_{^4p_{3/2}}) \\
& + 1.33^4 s_{3/2}\ {}^2p_{1/2} sin(\phi_{^4s_{3/2}} - \phi_{^2p_{1/2}}) - 2.357^4 s_{3/2}\ {}^4p_{1/2} sin(\phi_{^4s_{3/2}} - \phi_{^4p_{1/2}}) \\
& + 6.667^4 s_{3/2}\ {}^2p_{3/2} sin(\phi_{^4s_{3/2}} - \phi_{^2p_{3/2}}) - 2.981^4 s_{3/2}\ {}^4p_{3/2} sin(\phi_{^4s_{3/2}} - \phi_{^4p_{3/2}}) \\
b_2 = {} & -0.667^2 p_{1/2}\ {}^2p_{3/2} sin(\phi_{^2p_{1/2}} - \phi_{^2p_{3/2}}) + 1.193^2 p_{1/2}\ {}^4p_{3/2} sin(\phi_{^2p_{1/2}} - \phi_{^4p_{3/2}}) \\
& - 0.943^4 p_{1/2}\ {}^2p_{3/2} sin(\phi_{^4p_{1/2}} - \phi_{^2p_{3/2}}) - 0.211^4 p_{1/2}\ {}^4p_{3/2} sin(\phi_{^4p_{1/2}} - \phi_{^4p_{3/2}}) \\
& + 1.784^2 p_{3/2}\ {}^4p_{3/2} sin(\phi_{^2p_{3/2}} - \phi_{^4p_{3/2}})
\end{aligned} \tag{5.26}$$

Here the $\phi_{^{2s+1}l_J}$ terms are the phases of the TMEs, while the amplitudes of the TMEs are denoted by ${}^{2s+1}l_J$. If the new resonance decays to the ground state via M1 radiation, then the $l = 1$ TMEs will dominate the expressions for the a_k and b_k coefficients listed in Equation 5.26. If it decays via E1 radiation, then the $l = 0$ TMEs will be dominant.

The shape of the angular distribution of cross section will be largely (or completely, for the case of *pure* dipole radiation) determined by the first three Legendre polynomials in Equation 5.20. The assumptions of s-wave and p-wave capture via the emission of E1 or M1 radiation reduce the possible distributions of cross section and

Figure 5.6: The first three Legendre polynomials $P_{0-2}(\cos\theta)$, as a function of θ.

analyzing power to four cases. We shall also consider a fifth case, where the existence of a resonance at 127 keV is neglected (the 278 keV resonance is still included in the model for this case), to show the effect of the existence of the new resonance.

In the case of no resonance, the angular distribution of the cross section will have direct s-wave (E1) and p-wave (M1) contributions to the capture amplitude, as well as a resonant s-wave (E1) contribution from the tail of the $J^\pi = 1/2^+$ resonance at 278 keV. The E1 contributions should be somewhat dominant over the direct M1 component, giving $l = 0$ TMEs in Equation 5.26 which are relatively large compared to the $l = 1$ TMEs. Therefore, the angular distribution of cross section given by Equation 5.20 should be dominated by A_0 (isotropic) with some a_1 (cosine-like, as can be seen in the shape of $P_1(cos\theta)$ in Figure 5.6) and almost nothing from a_2 ($sin^2\theta$-like). If we now include the resonance as a $J^\pi = 1/2^-$ or $3/2^-$ state in ^{15}O, then p-wave capture with an M1 decay to the ground state will dominate. In this case, the $l = 1$ TMEs will almost entirely dominate the expressions for the a_k and b_k coefficients, and a_1 will be tiny. a_2 and A_0 both contain terms with $l = 1$ TMEs which

and no relative phases. Therefore they should be somewhat comparable. Hence, we would expect some isotropic component mixed with some component with a valley (for a_2 positive) or a peak (for a_2 negative) at 90°. For $J^\pi = \{1/2^+,3/2^+\}$, s-wave capture via E1 radiation will dominate. We would then expect both a_1 and a_2 to be small compared to A_0, giving a nearly pure isotropic distribution of γ-rays.

The angular distributions of the cross section without the resonance, and for the four most probable values of J^π are plotted in Figure 5.7. The angular distributions of cross section and analyzing power for the cases where the value of J^π=$\{1/2^-,3/2^-\}$ are indistinguishable, and as such are plotted together. HIKARI calculated the cross section, which is normalized to A_0, at 18 different angles for each proton energy in the input file. The plots of the calculations represent integrations of these angular distributions at proton energies from $E_p = 0$ - 140 keV in the target, where the weights assigned in the integration were the cross section and the inverse of the stopping power (which gives the target areal density). Because of the rapid variation in the cross section, the angular distributions were calculated at many energies in the resonance region. These contributions dominate the integration of the angular distribution for the three cases where the resonance is included. Therefore, these angular distributions represent what would be observed in the detector in a real experiment, when the Q_k coefficients have been included.

When the 127 keV resonance is not included, as in the curve in Figure 5.7 labeled "no resonance", the capture cross section is nearly isotropic. The interference of s and p TMEs (see Equation 5.26) give rise to a small a_1 coefficient which adds the cosine shape to the isotropic (A_0) component, as expected. For the cases of a $J^\pi = \{1/2^-,3/2-\}$ resonance, the angular distribution of the cross section shown in Figure 5.7 has an isotropic component and a central peak. This indicates a negative a_2

Figure 5.7: Angular distributions of the cross section, integrated from 0-140 keV.

coefficient, and the general shape is again as expected. Such a large asymmetry (approximately a factor of 2) between 0° and 90° opens the possibility of determining the parity of the new state in a direct measurement. The two cases of s-wave capture (J^π = $\{1/2^+,3/2^+\}$) are isotropic, as shown in Figure 5.7. In any case, the measurement of the angular distribution of the cross section for the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction would not definitely indicate the J^π of the new resonance.

Figure 5.8: The first two non-zero associated Legendre polynomials $\mathrm{P}^1_{1,2}(\cos\theta)$, as a function of θ.

Even in the case of an isotropic angular distribution, a sufficiently precise measurement of the angular distribution of the analyzing power may reveal the nature of the new state in ^{15}O. The shape of the angular distribution of the analyzing power for γ-rays from the new resonance should also be deducible from the TMEs and the expansion of Equation 5.23. For pure dipole radiation there are only two terms, which are determined by b_1 and b_2. For *pure* E1 radiation (pure s-wave capture), A_y will be identically zero, while pure M1 radiation (pure p-wave capture) can only give rise to a finite b_2 coefficient. A finite b_2 coefficient will add a sinusoidal component to the angular distribution of the analyzing power, as can be seen from Figure 5.8. A finite

b_1 coefficient requires both M1 and E1 radiation. The general form of the contribution to the analyzing power for each of these coefficients, is given in Equation 5.23. The first two non-zero associated Legendre polynomials are plotted in Figure 5.8. An angular distribution of the analyzing power would have to be measured and a TME fit performed in order to extract the amplitudes and phases of the E1 and M1 TMEs.

The integrated angular distributions of the analyzing powers below $E_p = 140$ keV for the four cases considered are plotted in Figure 5.9. The integrated distributions use the analyzing powers from the HIKARI output file, which are calculated at 18 angles from 0° to 170°, as with the differential cross sections shown in Figure 5.7. For the case with no resonance, the presence of both E1 and M1 TMEs with definite phase differences creates a finite b_1 coefficient, as can be seen from Equation 5.26. This interference of E1 and M1 terms creates a significant 90° analyzing power, as can be seen in Figure 5.9 (see also Figure 5.8 and Equation 5.23). For the M1 cases, the angular distributions of analyzing power are (as was the case for the angular distribution of cross section) indistinguishable, and nearly 0 for all angles. For the cases of $J^\pi = 1/2^+,3/2^+$, E1 decay dominates and the angular distributions of analyzing power are small, but have slightly different shapes. HIKARI calculates small analyzing powers in all cases where the resonance is included. This indicates very small relative phases between the contributing TMEs. As can be seen in Equation 5.26, the b_k coefficients depend on the sine of the relative phase, therefore very small relative phases will suppress the b_k coefficients.

Figure 5.9: The angular distributions of the analyzing power, integrated from 0-140 keV.

Chapter 6

Conclusions

The strengths and energies of resonances in the ^{14}N(p, γ)^{15}O reaction at proton energies of a few hundred keV are of central importance in understanding the energy production and evolution of stars, and the set of data from the present experiment has revealed the existence of previously unknown resonance at $E_p = 127 \pm 2$ keV. Later in this chapter we will use the strength of the resonance, $\omega\gamma$, to draw conclusions about the astrophysical significance of this new resonance, which turns out to dominate the rate of the ^{14}N(p, γ)^{15}O reaction in stellar plasmas with a T_9 (temperature in billions of Kelvin) near 0.11. From the step size in the yield curve it was possible to extract the parameters of the resonance, namely the strength ($\omega\gamma$ = 4.5 $\pm 10^{-6}$ eV), width ($\Gamma \approx 10^{-4}$ eV) and the spectroscopic factor (on the order of 10^{-5}-10^{-6}). Such a level in ^{15}O was not observed in the prior work of [Hes57], leading us to re-examine their data and question the level of carbon buildup on the target surfaces used in that work. The exclusion of the data of Hester and Lamb will allow us to conclude (in Section 6.2) that the rate of the ^{14}N(p, γ)^{15}O reaction in stars near temperature of

the sun is about half what was previously believed. The existence of a new resonance also implies that we should re-examine shell model calculations of the energy levels of the mass-15 system, in order to see if it has an impact on our current understanding of the structure of the mass-15 system.

6.1 Nuclear model calculations.

Some energy level calculations can be found which are in agreement with the existence of this excited state. The treatment of Wolters, et al. [Wol90] predicts an "extra" state at 8.4 MeV in ^{15}N with J$^{\pi}$=1/2^{-}. The calculation uses a shell model code with a basis of harmonic oscillator-like bound states in a model space built from $2\hbar\omega$ basis states. An analog to this state in ^{15}O at 7.414 MeV in ^{15}O would constitute a pair of analog states with an energy shift similar to that of the 8.313 MeV state in ^{15}N and its analog at 7.556 MeV ^{15}O, which can be seen in the isobar diagram shown in Figure 6.1[Ajz91]. A state in ^{15}N, which could correspond to the analog of the new state in ^{15}O, is not reported in the current literature. Any such analog state may have a similar spectroscopic factor, because corresponding E1 and M1 transitions in conjugate nuclei generally have the same strength [War69], and therefore such a state in ^{15}N would be difficult to detect. Also, an OXBASH (OXford Buenos-Aries SHell model code) calculation [Bro00] indicates a missing $2\hbar\omega$ state in ^{15}O with J$^{\pi}$=1/2^{-} at 7.468 $\pm$ 0.5 MeV. That calculation is well within error of the energy of the newly-measured state, and implies that several higher-energy excites states (with questionable spin-parity assignments) be assigned new values of J^{π}. Given the existence of this new excited state, it might be prudent to experimentally examine the spin-parity assignments of these states, for which no current experimental data

is available. Unfortunately, further details of the new state (strength, width, or spectroscopic factor) are not available from the calculation, so that a quantitative comparison between the present measurement and these theoretical predictions is not possible.

Figure 6.1: The isobar diagram for A=15.

6.2 Reaction rates.

Given the values of the astrophysical S-factor as a function of energy, the calculation of the stellar reaction rate per particle pair is a straightforward procedure. The

problem requires the evaluation of an integral of the product of the energy-dependant cross section and the Boltzmann distribution of energies found in the plasma. The Maxwellian-averaged reaction rate is given by [Ang99]:

$$N_A < \sigma v >= N_A \frac{(8/\pi)^{1/2}}{\mu^{1/2}(k_B T)^{3/2}} \int_0^\infty \sigma E e^{-E/k_B T} dE, \tag{6.1}$$

where N_A is Avogadro's number, μ is the reduced mass of the system, k_B is Boltzmann's constant, T is the temperature in Kelvin, and σv is the product of the cross section with the relative velocity of the particles. Using tables of cross sections from NACRE [Ang99] compilation, it was a relatively simple procedure to create a numerical integral in *Mathematica*^{tm} to evaluate Equation 6.1. Using the HIKARI extrapolation of the total S-factor to low energies (shown in Figure 5.4) and adopting the compilation of data presented by the NACRE collaboration [Ang99] at high energies, we can evaluate the reaction rate for the ^{14}N(p, γ)^{15}O reaction by performing the above integral numerically in *Mathematica*^{tm}.

For the impact of the resonance alone on the reaction rate in a plasma, only the value of $\omega\gamma$ and the resonance energy are important[Ang99], and these value are the most model-independent resonance parameters which are possible to derive from the data. The contribution of the new resonance to the reaction rate in a plasma may be obtained by inserting the Breit-Wigner cross section of Equation 4.1 into the reaction rate integral given by Equation 6.1. With temperature in T_9, the contribution of a narrow resonance (narrow enough so that the Boltzmann distribution is approximately linear across the resonance width) reduces to [Ang99]:

$$N_A < \sigma v >_R = \frac{1.539 \times 10^{11} (\omega\gamma)}{A^{3/2} T_9^{3/2}} e^{-11.605 E_R / T_9} \tag{6.2}$$

where the resonance energy in the center of mass frame (E_R = 118.5 keV) and the resonance strength ($\omega\gamma$) are expressed in MeV, the reaction rate $N_A<\sigma v>_R$ has units

of $cm^3mol^{-1}s^{-1}$, and A is the reduced mass in amu. In Equation 6.2 we neglect the insignificant "tail" contribution to the rate for this narrow resonance. The additional contribution to the reaction rate given by Equation 6.2 may be simply added to the rate given by performing the numerical integral of Equation 6.1, where in performing that integral we use a flat extrapolation of the S-factor across the resonance energy. The method of adding a contribution for the new resonance given by Equation 6.2 to the rate derived by numerical integration of Equation 6.1 for an S-factor which ignores the new resonance gives results identical to those for brute-force numerical integration of the full S-factor for several test temperatures across a range of plasma temperatures from T_9=0.008-10. This is the range of temperatures presented in the NACRE tables. Using Equation 6.2 is much faster computationally, however, due to the computational time necessary to integrate over such a narrow resonance.

It should be noted, in light of the angular distribution for the case where the new resonance is M1 (consistent with the shell model calculations), that the measurement of $\omega\gamma$ at 90° (and hence the spectroscopic factor) could be approximately 35% too high. This remains to be observed, but will have definite consequences for astrophysics. Our calculations will proceed for the case of an isotropic distribution, since there is insufficient evidence to conclude otherwise. Equation 6.2 allows one to recompute the rate for the case of an $\omega\gamma$ of only 3 μeV, for the case of an M1 resonance. Which rates are adopted will depend on which result is borne out by a study of the angular distribution of cross section or by stellar model predictions which are more consistent with observational astronomy. Such a determination is beyond the scope of the present work. It should also be noted that the error in the strength of the resonance is approximately 20%, and so the error in the contribution of this resonance to the reaction rate (linearly proportional to $\omega\gamma$) is also 20%.

The results of the reaction rate calculation are shown as the ratio of the calculated rate to the rates published in the NACRE compilation in Figure 6.2, along with a calculation neglecting the resonance in Figure 6.3. Near T_9=0.11, where the resonance energy is close to the peak burning energy in the Gamow window, the calculated rate is 35 times the rate published by the NACRE collaboration [Ang99]. In Figure 6.3 we can see that the low temperature rate of the ^{14}N(p, γ)^{15}O reaction is smaller than the corresponding NACRE rate when a strong contribution to the cross section from the 6.793 MeV state is not included, in accord with recent measurements of the width of this state [Ang01, Cha01]. Calculation of the rates with the inclusion of a strong contribution to the total S-factor from this resonance has been performed, resulting in the ratio of calculated rates to the NACRE rates of $1 \pm 5\%$ across the energy range from T_9=0.008-10.

Figure 6.2: The ratio of calculated reaction rate to the NACRE rates as a function of T_9.

In very hot stars, the rate of the CNO cycle may be limited by the β-decay of ^{13}N ($\tau_{1/2}$ = 598 s). However, most of a very hot star's original hydrogen has already

Figure 6.3: The ratio of the calculated reaction rate to the NACRE rates as a function of T_9 in the absence of both a 7.414 MeV state and contribution from the subthreshold resonance at 6.793 MeV.

been used up [Her01], possibly making the $^{14}N(p,\gamma)^{15}O$ reaction the slowest (and therefore controlling) reaction in the CNO cycle in spite of the limitation imposed by the β-decay process. Also, in most stars which are ≈30-1000% heavier than the Sun, β-limiting is not a factor [Cla83]. Therefore, the discovery of this new resonance has an impact on the rate of energy generation in the entire CNO cycle, as well as the isotopic abundances produced in this cycle in its steady state.

For completeness, we should fit the reaction rate with an analytical approximation which may be used in computational models of stellar evolution. Rather than create an arbitrary function to fit the rapidly varying reaction rate (which appears to be customary in nuclear astrophysics), we use the method of rational interpolation to create a standardized fit [Nel00]. The simple tables of coefficients used in such a fit should make computation much easier. This approximation is given by Rate=$10^{(P_1/P_2)}$, where P_1/P_2 is a rational polynomial fit to a table of the $\log_{10}$ of the

reaction rates. P_1 and P_2 are simple polynomials of the form $C_0 + C_1x + C_2x^2$+..., and x is given by $\log_{10}(T_9)$.

The polynomials were kept to the minimum number of terms necessary to achieve 10% agreement with the calculation and eliminate singularities (where P_2=0) over the entire range of T_9 for which the fit is valid. In this particular case, an unusually high number of coefficients are needed to accomodate the variations in the reaction rate. These variations are due to the presence of the two very narrow, low-energy resonances in the $^{14}N(p,\gamma)^{15}O$ reaction at $E_p = 127$ and 278 keV. The typical number of rational interpolation coefficients needed to describe the rate of an reaction in a plasma over this range of temperatures is 10-20 [Nel00], considerably less than the 34 coefficients needed to describe the $^{14}N(p,\gamma)^{15}O$ reaction. The range for which the fit is good is T_9 = 0.008 - 10, as in the present compilation of rates [Ang99]. For comparison, the sun has a T_9 near 0.012; a T_9 of 10 is only accessible to violent objects such as the accretion disks around black holes or the shockwave of a supernovae [Bec82]. The present data format should make the creation of stellar models easier and make the input data tables simpler to update when new data on the $^{14}N(p,\gamma)^{15}O$ reaction at truly low energies ($E_p < 80$ keV) becomes available.

$P1(C_0$-$C_5)$	$P1(C_6$-$C_{11})$	$P1(C_{12}$-$C_{16})$	$P2(C_0$-$C_5)$	$P2(C_6$-$C_{11})$	$P2(C_{12}$-$C_{16})$
2.0673084	-175.41455	-24.686854	1.00000000	-32.080265	4.0493300
0.42994493	58.369550	94.557645	-0.51043720	41.845299	15.389956
-13.635733	367.82842	40.872354	-4.5897260	83.449985	6.4139192
7.2245650	182.39821	0.43977404	3.5908415	5.0465969	1.0872110
43.461514	-255.42462	-1.6855844	8.0965536	-70.696962	0.07148332
-63.415757	-296.24728		-24.157371	-46.660287	

Table 6.1: Coefficients for the polynomials in the rational approximation to the reaction rate.

6.3 Future directions for study.

Currently, an experiment at the TUNL Low Energy Nuclear Astrophysics (LENA) lab is planned to verify the findings of the present work [Cha01]. This measurement will involve a 5 mA proton beam with up to 200 keV of energy. Unlike the experiment of Hester and Lamb, the targets (TiN and TaN) will be tested using a 1 MV Van de Graff accelerator. Protons from this accelerator will observe the target surface via observation of higher-energy resonances in the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction. If verified, work should be done in two main areas. First, the experiment will have to move back from LENA to TUNL to measure the angular distributions of the analyzing power and γ-rays. Only TUNL has a polarized ion source which can produce enough polarized protons to make such a measurement. Statistics will make the measurement very challenging, and very difficult to do simultaneously because the large HPGe detectors would have to be pulled back from the target to fit in at those angles. That would reduce counting statistics. Certainly every effort should be made to reduce neutron background in the detectors. The resonance should also be observed again with the improvement of heavy thermal neutron shielding to reduce background. Complete elimination of the background peak alone could improve statistical error by about 30%.

If the resonance is M1, as the shell-model calculations suggest, then a look at Figure 5.7 tells us that placing detectors at 0° and 90° would be a prudent first step. If the angular distribution of γ-rays proves to be isotropic, then the resonance may be E1. Measurement of the analyzing power at 45° and 90° could bear this out, but getting enough counts to achieve the statistical error necessary for a definite measurement is complicated by the small solid angles subtended by the detectors at

the distances which they would have to be pulled back to in order to achieve the 45° and 90° angles. An analyzing power measurement would also be extremely difficult to achieve, given the small analyzing powers shown in Figure 5.9. It could, perhaps, be accomplished in a series of long measurements which are done in exactly the same manner. Such measurements could be averaged to reduce the error to an acceptable level.

Because the resonance dominates the reaction rate in plasmas with T_9 near 0.11, further work to improve the statistical accuracy of the present measurement will become necessary as models of stellar evolution improve. The new evaluation of the reaction rates for the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction as a function of T_9 should also be carefully considered in such models, to see if they add to or lessen the agreement of these models with observational astronomy. We have created the analytical approximation given in Table 6.1 for exactly this purpose.

Chapter 7

Appendix A

7.1 Acknowledgements for the appendix.

The paper included here as Appendix A was co-authored by C. D. Nesaraja, A. P. Tonchev, C. R. Brune, J. H. Kelley, R. M. Prior, K. Sabourov, D. R. Tilley, S. O. Nelson, and H. R. Weller and published in Physical Review C [Nes01]. The design of the experiment flowed from the simultaneous measurement technique developed for the measurement of the $^{14}\mathrm{N}(\mathrm{p},\gamma)^{15}\mathrm{O}$ reaction, and as a straightforward variation on that experiment it is included here.

The ratio of S factors for (p,γ)reactions on ^{12}C and ^{13}C at astrophysically relevant energies

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Abstract

We present measurements of the ^{12}C(p,γ_0)^{13}N /^{13}C(p,γ_0)^{14}N S factor ratio at E_p=160 keV. The overall result of 0.33$\pm$ 0.03 is in good agreement with the value of this ratio obtained from previous experimental values of the S factors for the ^{12}C(p,γ_0)^{13}N and the ^{13}C(p,γ_0)^{14}N reactions. A value for the total S factor ratio of S[^{12}C(p,γ_{total})]/ S[^{13}C(p,γ_{total})] = 0.24$\pm$ 0.03 is computed by correcting the ^{13}C(p,γ_0) S factor for transitions to the other intermediate states. The direct + resonance model calculation of this ratio and the impact on the CNO cycle are also explored.

7.2 Introduction:

The ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions are the first members in the CNO cycle and therefore determining their S factors is of interest at astrophysical energies. The CNO cycle is one of the main series of reactions in hydrogen burning. The CNO cycle requires the presence of C, N and O isotopes which can be represented in a sequence of reactions as shown in Figure 7.1:

In stars more massive than the sun, hydrogen-burning proceeds predominantly via the operation of the CNO cycle. An example of the hydrogen burning that has

predicted the changes of chemical composition of the observable surface layers of the stars is the abundance ratio $^{12}C/^{13}C$. The $^{12}C/^{13}C$ ratio is an important measure of stellar evolution [Rol88] and nucleosynthesis [Kol90].The changes in the ratios of $^{12}C/^{13}C$ in stars happen as they evolve from the main sequence to the first ascent giant branch. As stars leave the main sequence and advance up the red giant branch, the convection zone grows and penetrates to greater depths where it begins to dredge up material that has been hot enough for the CNO cycle to convert to N. This is when the primordial ^{12}C is converted into ^{13}C and ^{14}N by the reactions $^{12}C(p,\gamma)^{13}N$ $(\beta+)^{13}C$ and $^{13}C(p,\gamma)^{14}N$, hence reducing the $^{12}C/^{13}C$ ratio. During the late Asymptotic Giant Branch (AGB) phase, the stars suffer thermal instabilities in the Helium shell where partial Helium burning occurs causing the $^{12}C/^{13}C$ ratio to increase.

There have been extensive theoretical fits made by stellar atmospheric models on the $^{12}C/^{13}C$ ratio in some selected carbon rich stars [Lam86][Ohn00]. These selected carbon rich stars are believed to be burning on the CNO cycle by means of the reactions $^{12}C(p,\gamma)^{13}N(\beta+)^{13}C(p,\gamma)^{14}N$. At equilibrium in the CNO cycle, the $^{12}C/^{13}C$ abundance ratios should be equal to the inverse ratio of their cross sections for proton capture. Therefore, under conditions of equilibrium, the abundance ratios for the $^{12}C/^{13}C$ can be derived from the ratio of the cross sections or S factors. Thus an accurate estimate of the $^{12}C/^{13}C$ -ratio should increase our understanding of the processes that lead to the formation of carbon stars [For97][Wal98]. The $^{12}C/^{13}C$ ratio is also an important tracer of the past star formation rate and stellar mass function [Pra96][Gre97].

The $^{12}C(p,\gamma)^{13}N$ (Q= 1.943 MeV) and the $^{13}C(p,\gamma)^{14}N$ (Q=7.551 MeV) reactions have been studied over a wide range of beam energies extending down to around 80 keV [Kin94][Sea51][Woo52][Heb60][Hes61][Bai50][Hal50][Lam57][You63][Rol74]. In the

most recent results, King *et al.* [Kin94] measured the cross section and S factor for the ^{13}C(p,γ)^{14}N reaction in the energy range E$_p$(lab)= 120–950 keV. Enriched ^{13}C evaporated on a Ta backing was used in King's experiment, and the resulting gamma ray yield was taken from the primary and secondary transitions. Prior to King's work, other investigators [Sea51][Woo52][Heb60][Hes61] had measured the cross section to as low as E$_p$(lab) ≈100 keV. The data from Hebbard [Heb60] and Hester [Hes61] were limited to capture cross sections to the ^{14}N ground state. The experimental S factor [Kin94] was summed from primary and secondary transitions to the ground state with an uncertainty of ∼12%. The result was in agreement with the results from the work of Hebbard [Heb60] and Hester [Hes61] for most of the proton energies except in the energy range E$_p$(c.m.) = 140-170 keV.

The experimental S factor for the ^{12}C(p,γ)^{13}N reaction was previously measured [Bai50][Hal50][Lam57][You63][Rol74] in the energy (lab) range from 72–541 keV. Most of the experimental data are available from the work done by Rolf and Azuma [Rol74]. There is a general agreement in the energy range E$_p$= 200–450 keV. However some discrepancies exist below this range [Bai50][Hal50][Lam57]. The discrepancies of the existing data that lie near the energy presently being studied [E$_p$(lab)=160 keV] show that accurate data are needed to enable an extrapolation of the S factors down to the 20-30 keV region.

The direct + resonance model calculations which will be presented are intended to provide insight into the elementary physics considerations which determine the ratio of the cross sections in the (p,γ) reactions on ^{12}C relative to ^{13}C .

7.3 Experimental Method:

Measurements of (p,γ) reactions yields for capture to the ground state of ^{13}N and ^{14}N were carried out at the Triangle Universities Nuclear Laboratory (TUNL). Figure 7.2 shows the experimental setup used in the current experiment. In this section detailed descriptions of the beam, targets and γ-ray detector are presented.

7.3.1 Proton Beam

An unpolarized beam from the TUNL Atomic Beam Polarized Ion Source (ABPIS) was used. The maximum beam energy from the ABPIS is 80 keV. To increase the beam energy to 160 keV, a -80-kV bias was applied to the target chamber, which was attached to the beam line with a multi-section acceleration tube. A measurement of the beam current integration was not possible due to the strongly biased target.

Stopping the beam completely in the target created a range of incident proton beam energies between E_p= 160 keV and E_p = 0 keV. Consequently, this range of incident beam energies creates a corresponding range of outgoing gamma rays that leads to a broadening of the full energy spectra. These gamma rays range from energies of 7.699-7.551 MeV for the ^{13}C(p,γ)^{14}N reaction and 2.091 - 1.943 MeV for the ^{12}C(p,γ)^{13}N reaction, respectively. The effective proton energy is defined to be the diminished proton beam energy which gives rise to one half of the observed yield. This was evaluated using known stopping powers [And77] and the assumption of a constant S factor below E_p =160 keV. It was found that E_p(eff) = (143$\pm$3) keV for both reactions.

7.3.2 Target

The experiments were performed by stopping a 60 μA proton beam in a thick frozen benzene target. The target consisted of a mixture of $^{13}C_6H_6$ and $^{12}C_6H_6$. The $^{13}C_6H_6$ was purchased from the Cambridge Isotope Laboratories and had a ^{13}C isotopic enrichment of 99%. The method used by the company to determine the content of the ^{13}C - labeled benzene was gas chromatography and mass spectrometry with an estimated error of $\pm$0.5 %. The masses of the mixed liquids were measured within 0.1% by a microbalance and the masses were used to determine the $^{12}C/^{13}C$ target composition ratio. The background for the lower energy gammas from $^{12}C(p,\gamma)$ was higher than for $^{13}C(p,\gamma)$. The concentration of ^{12}C used was 3.02, 3.77 and 3.57 times that of ^{13}C in three different runs to increase its statistics, in order to compensate for this higher background.

In this experiment, liquid benzene was allowed to flow into the target chamber. Since the vapor pressure of benzene at room temperature is $\sim$95 Torr, the benzene immediately vaporized due to the low pressure inside the target chamber. The vapor that touched the cooled surface (77 K) of the aluminum disk immediately froze and turned into a benzene ice-target layer. The low beam-line vacuum pressure ($\sim 10^{-6}$ torr) and the high intensity 60 μA proton beam assisted in the slow vaporizing of the outer layer of the benzene ice target, thus continuously exposing a new layer of target. Deposition of the benzene target for a duration of 30 seconds was sufficient to produce a thick target that stops the proton beam completely. In addition new targets were made every 4 to 8 hours.

7.3.3 Gamma Detection and Signal Processing

The outgoing gamma rays from the (p,γ) reactions were measured with a 123% efficient (relative to a 3" $\times$ 3" NaI detector) coaxial HPGe detector. The raw signals from the HPGe detector are initially fed into the preamplifier and then fanned out to a spectroscopy amplifier and a timing filter (TFA). The linear signal was shaped by the spectroscopy amplifier before it was fed into the ADC. The TFA and constant fraction discriminator (CFD) amplified and discriminated the timing signal which was subsequently sent through a gate and delay generator. The resultant proper signal (shape and width) eventually gated the ADC. All signals were sent to a CAMAC module which was read through a computer operating with the TUNL XSYS system.

The energy calibration of the detector was obtained using the energies of known background lines of ^{228}Th (2.614 MeV), ^{40}K (1.461 MeV) and the ^{56}Fe thermal neutron capture doublet peak at (7.647 MeV and 7.631 MeV). The typical energy calibration was about 0.9 channel keV in an 8000 channel spectrum. The background peak from the environmental ^{228}Th radio nuclide with an energy of 2.614 MeV is another line that contributed to the difficulty of determining an accurate yield for the ^{12}C(p,γ)^{13}N reaction.

A measurement was also conducted with a ^{13}C enriched foil to observe if any ^{12}C buildup may have occurred on the target. No significant ^{12}C buildup was observed from the runs. The benzene target stability and possible contaminant buildup on the benzene targets were also gauged over the course of the experiment by monitoring the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N count rates for each new target run in the HPGe detector. No significant change in the count rates was observed for the duration of the experiment.

The relative efficiencies of the HPGe detector for the two gamma ray energies were obtained by measuring the photopeak efficiency for the E_γ ranging from 1.77 to 10.7 MeV produced by means of the ^{27}Al(p,γ)^{28}Si reaction at E_p = 997 keV. Utilizing the E_p = 992 keV resonance from the ^{27}Al(p,γ)^{28}Si reaction for the efficiency calibration has two main advantages: the gamma rays arising from ^{27}Al(p,γ)^{28}Si reaction have a wide energy range from 1.77 to 10.7 MeV and the intensities are well studied [Ant77]. In order to perform this measurement a 997 keV proton beam from the TUNL FN tandem accelerator, was stopped in a 1 mm Al target on a 0.8 mm thick Cu backing. The gamma rays produced from the resonance just below the proton energy, were detected with the same HPGe detector, positioned at $\theta = 45^\circ$ and a distance of 14.5 cm from the end of the aluminum target. Three hours of collecting data with an average beam current of $\sim$500 nA allowed us to obtain statistical uncertainties in the photopeak area of 0.5 to 3% for the strongest branches. The resulting experimental data agree well with the previous measurements of Schmid *et al.* [Sch96] at the low energy interval of 0.8 to 4 MeV. A plot of the data obtained from the ^{27}Al(p,γ)^{28}Si reaction is shown along with the low-energy source data of Schmid *et al.* in Figure 7.3. The solid line represents a fit to these data using the functional form given in Ref. [Sch96]. The results of these data and the fit give a ratio of the efficiencies for the 2.1 and 7.7 MeV γ-rays of $\epsilon_{12}/\epsilon_{13} = 2.88 \pm 0.09$ which is in very good agreement with Monte Carlo calculations [Ton00].

All measured gamma yields were corrected for attenuation through the stainless steel target chamber (6.35 mm thick) and Plexiglas (1.0 cm thick). At 0.2 to 2.4 MeV the experimental values for the attenuation coefficients has been obtained using a ^{226}Ra source. Measuring the relative intensities of the gamma lines of this source in the position of the target, with and without the chamber, allowed us to determine the

total attenuation coefficient at 2.1 MeV. At the higher energy (7.7 MeV) this information was obtained using the theoretical prediction from NIST (National Institute of Standards and Technology). The resulting attenuation through the stainless steel and Plexiglas were 16% and 2% for the γ-rays from the ^{13}C(p,γ)^{14}N reaction and 23% and 3% for the γ-rays from the ^{12}C(p,γ)^{13}N reaction, respectively. At 2.1 MeV the experimental value, measured with the help of ^{226}Ra source, gives the total attenuation coefficient of 28%. This is in good agreement with the theoretical one. The background was reduced by laying 10-20 cm of passive Pb shielding around the HPGe detector as shown in Figure 7.2. The method for determining the low-energy background for these measurements was based on the assumption that the background count rate is proportional to the data acquisition time and that an appropriately time-normalized "beam-off" spectrum could be subtracted from the target spectrum. The assumption was verified by noting that essentially all of the background was removed using this procedure. Both the background and the γ-ray spectra were corrected for ADC dead times.

7.4 Data Analysis:

The gamma-ray pulse height spectra from the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N are shown in Figure 7.4. The spectra were acquired at $\theta_{lab} = 90°$. The peaks correspond to gamma-ray energies of 2.091 MeV and 7.699 MeV for the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions, respectively. Background corrected yields of counts Y_{12} and Y_{13} were extracted for each peak.

The measured ground state yields were used to extract the S factor ratio of the ^{12}C(p,γ)^{13}N relative to the ^{13}C(p,γ)^{14}N reactions. The S factor is defined as

$$S = \sigma E_{c.m.} exp(2\pi\eta) \tag{7.1}$$

where η is the Sommerfeld parameter which is related to the center of mass energy in keV by the relationship:

$$\eta = \frac{1}{2\pi}(31.297) z_{target} z_{projectile} \sqrt{\frac{\mu}{E_{c.m}}}. \tag{7.2}$$

In both the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions, the direct capture to the ground state is expected to proceed primarily by an E1 transition. This can arise from s-wave and d-wave capture. Due to the low energy and angular momentum barrier, the s-wave capture is expected to dominate.

In the case of pure s-wave capture, the S factor can be assumed to be constant over a small energy interval at the energies of the present work. This is confirmed by direct capture-plus-resonance model calculations, which will be described below. For energies ranging from the beam energy E_b= 160 keV to 0 keV and assuming the cross sections are isotropic, the ratio of the total yields for the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions is described by the following equation:

$$\frac{Y_{12}}{Y_{13}} = \frac{\epsilon_{12} n_p N_{12} A_{12} \int_{E_b}^{0} \frac{\sigma_{12}(E)}{STP(E)} dE}{\epsilon_{13} n_p N_{13} A_{13} \int_{E_b}^{0} \frac{\sigma_{13}(E)}{STP(E)} dE} \ , \tag{7.3}$$

where, ϵ is the detector efficiency, n_p is the number of protons incident on the target the target, N is the atomic number density of carbon in the target, A is the attenuation correction, σ is the cross section and STP(E) is the stopping power. The ratio of the solid angles for the ^{12}C(p,γ)^{13}N relative to the ^{13}C(p,γ)^{14}N reaction is one since the same detector was used for both measurements. Using Eq.(1) and substituting it into Eq. (3) above, we see that under the assumption that the S factors are constant,

the ratio of the integrals in Eq. (3) is just a constant (which is $E_{c.m}[^{12}C(p,\gamma)]/E_{c.m}[^{13}C(p,\gamma)]= 0.994$ at any specific value of E(lab)) times the ratio of the S factors. A knowledge of the STP(E) is not needed since it is the same in both integrals and, since the complete energy dependent part of both integrands are identical when written in terms of E(lab), the integrals are identical and simply divide out. Note that the fact that the value of η is the same for both reactions was used in obtaining this result. The number of protons n_p are similarly identical and therefore cancel. The S factor ratio is now given by the expression :

$$\frac{S_{12C(p,\gamma)}}{S_{13C(p,\gamma)}} = 0.994\frac{\epsilon_{13}N_{13}A_{13}}{\epsilon_{12}N_{12}A_{12}}\frac{Y_{12}}{Y_{13}} \ . \tag{7.4}$$

Hence, the S factors ratios are deduced from the measured ratios of the yields with corrections for the γ-ray attenuation in the chamber, the ratio of the number densities of ^{13}C and ^{12}C in the target and the relative efficiency of the HPGe detctor for the two gamma-ray energies. The simultaneous measurement of the yields for the $^{12}C(p,\gamma)^{13}N$ to $^{13}C(p,\gamma)^{14}N$ has made it possible to eliminate systematic errors from the number of protons hitting the target, the stopping power of the target material and the solid angle of the detector.

Table Table 7.6 shows the results for the S factor ratios for three different target compositions of carbon ($^{12}C/^{13}C$) in benzene. Included in the table are the principal sources of uncertainties and their magnitudes in percentage. The total uncertainty for the S factor was obtained by combining the individual uncertainties in quadrature. An average of these three results gave us a value for the S factor ratio of 0.33 ± 0.03 for capture which goes only via a single transition from the continuum to the ground state. Experimental data for the $^{13}C(p,\gamma)^{14}N$ reaction [Kin94] have been used to estimate the contribution to S factor of this reaction from the transitions to excited

states. This factor at an effective proton energy of 143 keV was found to be 27 ± 5%. In the case of ^{12}C(p,γ)^{13}N, there are no transitions to excited states at these proton energies. This leads to a total S factor ratio of S[^{12}C(p,γ_{total})]/ S[^{13}C(p,γ_{total})] = 0.24± 0.03.

7.5 Comparison of the Experimental and Theoretical Results:

Our experimental S factor ratio of ^{12}C(γ_0)/^{13}C(γ_0) was 0.33 ± 0.03 at E$_p$(eff)=143 keV. This result was compared with the S factor ratio obtained from previous ^{12}C(γ_0)^{13}N and ^{13}C(γ_0)^{14}N S factor results. For the comparison, the ratio was deduced from Rolfs and Azuma's [Rol74] ^{12}C(γ_0) S factor results and King *et al.*'s ^{13}C(γ_0) S factor results [Kin94]. The ratio obtained was 0.34 ± 0.11. This result for the S factor ratio agrees very well with our result.

Direct Capture (DC) calculations have been performed in an attempt to understand the experimental ratio of the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N reactions. These calculations are intended to give a better understanding of the physics determining the cross section for proton capture to the ground states of ^{13}N and ^{14}N, respectively, over the energy range of 0 to 160 keV.

The DC model describes a reaction occurring between two particles, the projectile and the target, where both are assumed to be structureless. In this framework, radiative capture is considered a one-step process where no intermediate, or compound, nuclear state is formed. The DC reaction may be described as a transition from the initial (continuum) state to the final (bound) state mediated by the electro-

magnetic interaction. The bound final state in our calculation was described by the wave function for a bound proton in a Woods-Saxon well. In order to simplify our calculations the spectroscopic factors for both cases were set equal to 1.0. Because the current experiment was performed at the very low energy (E_p = 160 keV), there are no inelastic channels available. Therefore, the scattering potential is real. The potential used was taken to be the same Woods-Saxon potential which generated the bound state wavefunction. As will be seen from the discussion below, at these proton energies the particle interaction happens far away from the nuclear-nuclear interaction ($\sim$3 fm). It is therefore reasonable to expect that the Coulomb potential will be dominant term in the optical model. Table Table 7.6 shows the parameters set that were used.

In both the ^{12}C(p,γ)^{13}N and the ^{13}C(p,γ)^{14}N reactions the direct capture to the ground state proceeds primarily through E1 ($s_{1/2} \rightarrow p_{1/2}$ and E1 ($d_{3/2} \rightarrow p_{1/2}$) single-particle transitions. In both reactions the orbital angular momentum of the projectile is l=0 or l=2 and its laboratory energy is 160 keV. Moreover, the first resonances in ^{13}N and ^{14}N have energies of 457 and 558 keV, respectively. The widths of these relatively strong resonances are nearly the same: 39 and 40 keV, respectively. This means that at the very low energies, the ground state transitions will feel the strength of these resonances to a similar extent. Because of the similar influences of the resonances, Direct Capture calculations, which have already been successfully applied to the interpretation of the experimental results for low energy protons, may be able to explain the S factor ratio of these two reactions.

Our first goal is to reproduce the existing ground state data [Kin94][Rol74]. Including the lower lying E1 resonances allowed us to fit the data fairly well, as shown in Figure 7.5. The calculated total capture cross section ratio (0.32) for the ^{12}C(p,γ_0)^{13}N

to ^{13}C(p,γ_0)^{14}N the reactions is in fair agreement with the measured value of 0.33 $\pm$ 0.03. It should be noted here that the dotted line, which represents the DC plus the first resonance, is slightly offset from the S factor experimental data, probably as a result of our neglect of the second resonance [Rol73]. This offset is more pronounced at the high energy tail of the first resonance of the ^{12}C(p,γ)^{13}N and ^{13}C(p,γ)^{14}N. When all of the resonances are omitted, pure DC calculations fail to reproduce the absolute cross sections in both cases, but give a ratio of 0.32. This simplification gives us the opportunity to understand the physical origin of the observed ratio at low beam energies without the complications of interfering resonances.

The total cross section for direct E1 radiative proton capture is given by [Rol73]:

$$\sigma(E1) = 0.072\mu^{\frac{3}{2}} d^2 \frac{E_\gamma^3}{E_p^{\frac{3}{2}}} g(J)(l_i 010|l_f 0)^2 R(E1)^2 \tag{7.5}$$

where $\mu = m_p m_t/(m_p + m_t)$ is the reduced mass, d= $(Z_p/m_p - Z_t/m_t)$ is the effective charge correction factor, E_γ is the outgoing γ-ray energy, E_p is the projectile energy in MeV, g(J)= $(2J_f + 1)/((2j_p + 1)(2j_t + 1))$ is the spin factor and $(l_i 010|l_f 0)$ is the Clebsch-Gordan coefficient. The E1 radial integral R(E1) is given by :

$$R(E1) = \int_0^\infty u(r)\mathcal{O}_{E1}(r)\chi(r)r^2 dr \tag{7.6}$$

where u(r) is the continuum wave function, χ(r) the bound-state radial wave function and $\mathcal{O}_{E1}$ the E1 operator.

As can been seen from Eq. (5), two factors dominate the DC prediction of the capture ratio. First, the energy dependence of the cross section depends strongly on the E_γ^3 factor. Thus, this factor will greatly enhance the ^{13}C(p,γ)^{14}N reaction cross

section relative to the $^{12}C(p,\gamma)^{13}N$ reaction cross section by a factor of 50. This offset is due to the difference in the binding energies in both reactions. The Q-value for the $^{13}C(p,\gamma)$ reaction is 7.55 MeV instead of 1.94 MeV for the $^{12}C(p,\gamma)$ reaction. The contributions of other factors like reduced mass, effective charge, the spin factor and the Clebsch-Gordan coefficients to the total cross section are essentially the same. The only remaining factor which can compensate for the deficiency of the calculated $^{12}C(p,\gamma)/^{13}C(p,\gamma)$ cross section ratio (0.02) to the experimental one (0.33) is the radial integral. Thus, we would like to take a closer look at Eq. 6.

The radial integrands of Eq. 6 for the direct capture E1 transitions in $^{12}C(p,\gamma)$ and $^{13}C(p,\gamma)$ reactions are shown in Figure 7.6. As can been seen the major part of the process is dominated by contributions from regions far outside the nuclear radius. The maximum of the radial integrand for the $^{13}C(p,\gamma)^{14}N$ reaction occurs at 6.7 fm, while the maximum for the $^{12}C(p,\gamma)^{13}N$ reaction occurs at 13.2 fm.

The large Coulomb barrier greatly suppresses the wave function of the incoming low-energy proton at small distances such that the cross section is very dependent on the final state wave function at rather large radii (i.e., the tail of this wave function). The radial parts of the bound state wave functions are shown in Figure 7.7. As can been seen from this figure the radial part of the wave functions, formed by using Woods-Saxon potentials are almost identical at a small radii ($r < 4$ fm). However, as a result of the large binding energy, the radial wavefunction for ^{14}N is considerably smaller than that for ^{14}N at large radii, where the reaction occurs. It is this effect which leads to an enhancement of the radial integral of Eq.6 for the case of $^{12}C(p,\gamma)^{13}N$ relative to $^{13}C(p,\gamma)^{14}N$, offsetting the factor of 50 favoring the $^{12}C(p,\gamma)^{13}N$ reaction and producing a final ratio of 0.32 in this pure DC calculation.

7.6 Conclusion:

In summary, we have measured the S factor ratio of ^{12}C(p,γ_0)^{13}N to ^{13}C(p,γ_0)^{14}N at an effective proton energy of 143 keV. No intermediate transitions which lead to the ground state are included in our result. The method used enabled us to reduce systematic errors and led to a value of 0.33 with a 10 % error for the ratio of the S factor of the ^{12}C(p,γ)^{13}N reaction divided by that of the ^{13}C(p,γ)^{14}N reaction. Direct capture calculations explain the physical origin of the observed ratio at low beam energies, showing that the observed ratio results largely from the fact that the capture reaction occurs at relatively large distances (6 to 13 fm). Corrections to the ^{13}C(p,γ) S factor for transitions to other "intermediate" states allow us to compute an S factor ratio which includes intermediate transitions leading to the final state. Our result gives a value for the total S factor ratio of S[^{12}C(p,γ_{total})]/ S[^{13}C(p,γ_{total})] = 0.24$\pm$ 0.03. This ratio can be used to test equilibrium and non-equilibrium models of the CNO cycle and the predicted relative abundances of ^{12}C to ^{13}C by comparing the latter with the observed value.

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Figure 7.1: Sequence of reactions in the CNO cycle.

Figure 7.2: The physical setup of the current experiment.

Figure 7.3: Experimentally determined photopeak efficiency curve for the HPGe detector used in the present experiment. The open triangles were taken from [Sch96]. The open circles are the present measurements obtained using the $^{13}C(p,\gamma)^{14}N$ reaction. The solid line represents a fit to the experimental data using an exponential function having an argument which was quadratic in energy.

Figure 7.4: Gamma-ray pulse height spectra from the $^{12}C(p,\gamma)^{13}N$ and $^{13}C(p,\gamma)^{14}N$ reactions.

Target Composition $^{12}C/^{13}C$	**3.02/1**	**3.77/1**	**3.57/1**
Isotope	$^{12}C/^{13}C$	$^{12}C/^{13}C$	$^{12}C/^{13}C$
Statistical Uncertainty (%)	12.3/12.8	15.1/8.7	11.7/6.5
Attenuation Uncertainty (%)	3	3	3
Relative Uncertainty (%)	3	3	3
Isotopic Abundance (%)	0.05	0.05	0.05
Total (%)	13.0/13.1	15.4/9.7	12.4/7.8
Uncertainty for the S factor Ratio $^{12}C/^{13}C$ (%)	18.4	18.2	14.6
S factor $^{12}C(\gamma_0)/^{13}C(\gamma_0)$	**0.31±0.06**	**0.36±0.07**	**0.31±0.05**

Overall Result for the S factor ratio of $^{12}C(\gamma_0)/^{13}C(\gamma_0)$ at E_p(eff)=143keV: S factor = 0.33 ± 0.03

Table 7.1: Summary of the Statistical and Systematic Uncertainties in the S factor ratio.

Parameter	Reaction	
	$^{12}C(p,\gamma)^{13}N$	$^{13}C(p,\gamma)^{14}N$
V_R(MeV)	43.45	53.48
r_R (fm)	1.17	1.17
a_R (fm)	0.75	0.75
r_c (fm)	1.25	1.25

Table 7.2: Optical model parameters used to generate the scattering wave functions for the $^{12}C(p,\gamma)^{13}N$ and the $^{13}C(p,\gamma)^{14}N$ reactions. Proton energy in both reactions was set to 160 keV in order to calculate the optical model parameters that were used. The well depth was adjusted to fit the binding energies of the ground state of the final nucleus in each reaction.

Figure 7.5: S factor and cross section for the γ-ray transition to the ground state in ^{13}N and ^{14}N as a function of beam energy (E_p= 150-800 keV). The solid line represents only the DC model and the dotted line represents the DC plus the first E1 resonance in each case. The experimental data have been taken from [Kin94][Rol74].

Figure 7.6: The integrand of the E1 direct-capture matrix element leading to the ground state plotted as a function of the radius.

Figure 7.7: The total radial bound state wave functions for ^{13}N and ^{14}N used in Direct Capture calculations. DC has been performed assuming a spectroscopic factor of 1.0.

Chapter 8

Appendix B

8.1 Expressions for A_0, a_k, and b_k coefficients.

In the channel spin representation, we adopt the channel-spin coupling scheme [Wel80] of Chapter 5:

$$\vec{x} + \vec{a} = \vec{s}, \qquad \vec{l} + \vec{s} = \vec{J}, \qquad \vec{L} + \vec{c} = \vec{J} \tag{8.1}$$

The TMEs are, as in Chapter 5, labeled by the quantum numbers ${}^{2s+1}l_j$ in this representation. We denote reduced transition matrix elements T and T' as:

$$\begin{aligned} T &= < pL(c)J\pi||T||l(xa)sJ\pi > \\ T' &= < p'L'(c)J'\pi'||T||l'(xa)s'J'\pi' >, \end{aligned} \tag{8.2}$$

where p is the mode of radiation (electric or magnetic), and π is the parity of the state. We also define quantities $\hat{y} = (2y+1)^{1/2}$.

A_0, a_k, and b_k are then defined in terms of the above formalism by:

$$A_0 = (\frac{\lambda}{4\pi})^2 \frac{1}{\hat{x}^2\hat{a}^2} \sum_T \hat{J}^2 |T|^2 \tag{8.3}$$

$$a_k = (\frac{\lambda}{4\pi})^2 \frac{1}{A_0\hat{x}^2\hat{a}^2} \{\sum_{TT'} -1^{s-c+1}(1/2)(1+(-1)^{L+p+L'+p'+k})\hat{l}\hat{l}'\hat{L}\hat{L}'\hat{J}^2\hat{J}'^2 \tag{8.4}$$
$$\times(l0,l'0|k0)W(lJl'J';sk)(L1,L'-1|k0)W(LJL'J';ck)Re(TT'^*)\}$$

$$b_k = (\frac{\lambda}{4\pi})^2 \frac{1}{A_0\hat{x}^2\hat{a}^2} \{\frac{3\sqrt{x}\hat{x}\hat{k}}{((x+1)k(k+1))^{1/2}} \tag{8.5}$$
$$\times \sum_{TT'}(1/2)(1+(-1)^{L+p+L'+p'+k})\hat{s}\hat{s}'\hat{l}\hat{l}'\hat{L}\hat{L}'\hat{J}^2\hat{J}'^2(-1)^{a-x+c-J-s+l}$$
$$\times(l0,l'0|k0)W(xsxs';a1)(L1,L'-1|k0)$$
$$\times W(LJL'J';ck)X(lsJ;l's'J';k1k)Re(iTT'^*)\}$$

In Equation 8.3-Equation 8.5, W are the Racah coefficients, X are the 9-j symbols, and $(j_1m_1, j_2m_2|j_3m_3)$ are the Clebsch-Gordan coefficients. The summation is performed by the FORTRAN code LSCOEF over the contributing TMEs, resulting in the forms of a_0, a_1, b_1, and b_2 shown in Chapter 5. A_0 was calculated by HIKARI, but is not listed as it is not important in the shape of the angular distributions. For that only the normalized (to a_0, which is given in Chapter 5) a_ks and b_ks matter.

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