RADIATIVE CAPTURE REACTIONS

IN

LIGHT NUCLEI

by

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A Thesis submitted for the degree of Doctor of Philosophy
in the Australian National University.

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P R E F A C E.

This Thesis describes experiments in nuclear physics which were performed in the Research School of Physical Sciences at the Australian National University between March 1956 and December 1958. The work was done in the department of Nuclear Physics, under the supervision of Professor E.W. Titterton.

Some of these investigations have been performed in collaboration with other workers. The work described in Chapter II, which includes the design of a NaI(Tl) crystal spectrometer and the evaluation of various methods of γ-ray collimation was done individually by me. This spectrometer was then applied to the experiments reported in the chapters which follow. The Be $(p,\gamma)^{10}$ experiment in Chapter III was performed jointly with Dr. R.D. Edge. The experimental work was shared equally by us, but I was responsible for the determination of calibration pulse-height distributions and for the analysis of γ-ray spectra. The experiments with the 7.7 Mev cyclotron were performed in collaboration with Mr. A.H. Morton, a student in the department of Particle Physics. Mr. Morton's particular concern was with the operation of the cyclotron, while mine was with the detection of reaction γ-rays. The technique of extracting a stable mono-energetic proton beam was developed by Mr. Morton and Dr. W.I.B. Smith.
and is described elsewhere. The experimental work in Chapters IV, V, and VI was shared equally by Mr. Morton and myself. The experiments described in Appendices A and C have been performed independently by me.

Some of the work reported in this Thesis has been published as follows:


(ii) D.S. Gemmell, A.H. Morton and E.W. Titterton, Nuclear Physics, Vol. 10, No. 1 (1959). In press. "A Study of the Giant Resonance Regions of Be$^8$ and C$^{12}$ through the Inverse Reactions Li$^7$(p,$\gamma$)Be$^8$ and B$^9$ (p,$\gamma$)C$^{12}$." 


I am greatly indebted to my supervisor Professor E.W. Titterton who suggested the experiments on the reactions Be$^9$(p,$\gamma$)B$^{10}$...
and B_{11}^{11} (p,\gamma)C_{12}^{12} and who offered friendly advice and criticism during all aspects of my work. It is a pleasure to acknowledge the assistance given by Mr. N.F. Bowkett in operating the 1 Mev H.T. set, by Mr. R.V. Parkes in operating the cyclotron and by Mr. W.H. Owen and members of the main workshop staff in the construction of apparatus.

I wish to acknowledge gratefully the assistance given by my wife, Anne, who typed this thesis.

I am grateful to the Australian National University for the award of a Scholarship, during the tenure of which these studies were carried out.

No part of this Thesis has been submitted for a degree at any other University.

D.F. Gemmell
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CHAPTER I.

GENERAL INTRODUCTION.

1.1 Introduction.

The experimental investigations to be described in this thesis are concerned with the class of nuclear reactions in which electromagnetic radiation is emitted subsequent to the capture of a bombarding proton by a light nucleus \((A < 32)\). Radiative proton capture reactions form a useful method of studying the spectroscopy of nuclear levels and permit the examination of nuclear structure over a wide range of excitation energies. In the past, however, this approach has not been pursued to the same extent as other methods of nuclear spectroscopy such as the study of beta decay or reactions involving the emission of heavy particles. This has been mainly because of technical difficulties in the detection of \(\gamma\)-rays.

In this introductory chapter, a short review is given of some of the main topics relating to the experimental work described in the chapters to follow.

1.2 Gamma-ray Detection.

The wavelengths of nuclear \(\gamma\)-rays are too short\(^1\) to be

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\(^1\) For 1 Mev \(\gamma\)-rays, \(\lambda \approx 2 \times 10^{-11}\) cm.
measured by ordinary optical devices, and the somewhat indirect methods which must be used result in an energy determination which is quite inaccurate when compared with that achieved in atomic spectroscopy. The resolution obtainable is generally poor; of the order of a few percent. It is important to attain the best possible resolution in the study of capture reactions since the γ-ray decay scheme of highly excited nuclear levels is frequently complex, resulting in the production of many γ-rays of various energies and intensities which must be sorted out if significant information is to be obtained. Furthermore, the width of a nuclear level excited by a capture process is usually composed of partial widths mainly for particle emission and the γ-ray width is only a small fraction of the total. The probability of γ-ray emission is correspondingly lower, necessitating spectroscopic apparatus with a high detection efficiency.

To obtain a γ-ray detector with both high efficiency and good resolution poses a technical problem which, so far, has been only partially solved. In the following chapter a description is given of attempts to optimise both of these requirements using a large NaI(Tl) crystal as a scintillation spectrometer.

2. For mono-energetic γ-rays, the resolution is defined as the full width at half-height of the maximum intensity peak in the energy spectrum recorded experimentally, divided by the energy of this peak.
1.3 Classification of Transitions.

It is customary to classify γ-ray transitions as electric or magnetic multipoles according to the spin and parity changes involved in going from the initial to the final nuclear state. If a γ-ray carries away an angular momentum $L$ (measured in units of $\hbar$) it results from a $2^L$-pole transition. It is an electric $2^L$-pole (abbreviated as EL) if the parity change is given by $(-1)^L$ or a magnetic $2^L$-pole (ML) if the parity change is $(-1)^{L+1}$. This is shown in Table 1.1 for the first few multipoles.

<table>
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<td>L</td>
<td>Yes</td>
<td>EL</td>
<td>M2</td>
<td>E3</td>
<td>M4</td>
</tr>
<tr>
<td></td>
<td>No</td>
<td>ML</td>
<td>E2</td>
<td>M3</td>
<td>E4</td>
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1.4 Gamma-ray Selection Rules.

The requirement that angular momentum and parity must be conserved for the total system of nucleus plus γ-ray, results in selection rules imposed on the possible multipolarities allowed for a γ-ray transition between two states:

\[ |J_i - J_f| \leq L \leq J_i + J_f, \]  

where \( J_i \) and \( J_f \) are the angular momenta of the initial and final nuclear states and the parity change determines whether an electric or magnetic transition is involved. As a consequence of the fact that a \( \gamma \)-ray is a transverse electromagnetic vibration, it can be shown\(^4\) that all monopole radiative transitions are forbidden. In addition to these selection rules, there are others governing changes in the isotopic spin quantum number, \( T \) (discussed in § 1.6).

1.5 Gamma-ray Lifetimes.

The transition probability for \( \gamma \)-radiation between two nuclear states is sensitively dependent on the wave functions of the two states. Except perhaps for the case of the deuteron, these wave functions are not well known. It is possible, however, to obtain significant information concerning nuclear wave functions by a comparison of experimental \( \gamma \)-ray transition probabilities with theoretical values calculated on the basis of a specific nuclear model.

Formulae for the \( \gamma \)-decay lifetimes of nuclear states have been derived by Weisskopf\(^5\) from a single-particle model of the nucleus. This is a simplified version of the nuclear shell model in which a

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single proton is considered to move in a potential well which represents the average influence of all the other particles in the nucleus. The emission of a $L$-pole $\gamma$-ray is then attributed to this single proton changing its quantum state from one with orbital angular momentum $L$, and total angular momentum $J = L + \frac{1}{2}$, to one with zero orbital angular momentum. Using a value for the nuclear radius of $1.5A^{1/3} \times 10^{-13}$ cm, and on the assumption that this value is much smaller than the wavelength of the $\gamma$-ray emitted, Weisskopf's estimates of the radiative widths, $\Gamma_w$, (measured in eV) for the first few multipole orders are:

$$\Gamma_w^{(E1)} = 0.11A^{2/3}E_\gamma^3$$
$$\Gamma_w^{(M1)} = 0.021 E_\gamma^3$$
$$\Gamma_w^{(E2)} = 1.2 \times 10^{-7} A^{4/3}E_\gamma^5$$  \hspace{1cm} (1.2)
$$\Gamma_w^{(M2)} = 2.2 \times 10^{-8} A^{2/3}E_\gamma^5$$
$$\Gamma_w^{(E3)} = 8.7 \times 10^{-14} A^2E_\gamma^7,$$

where the $\gamma$-ray energy $E_\gamma$ is measured in Mev. These calculations were made under extremely simplified assumptions: e.g. the case where the radiating particle is a neutron was not included, and the effect of the recoil of the charged core of the nucleus was ignored. In spite of these and other simplifications, however, the results can be used to indicate the order of magnitude of known $\gamma$-ray widths. Consequently the expressions (1.2) for $\Gamma_w$ have come to be known as
Weisskopf units, and experimental γ-ray widths are conveniently measured in terms of these units. A γ-ray transition is said to have a width of $|M|^2$ Weisskopf units if $|M|^2 = \Gamma_\gamma/\Gamma_w$ where $\Gamma_\gamma$ is the experimentally determined radiative width and $\Gamma_w$ is the Weisskopf unit appropriate to the type (multipole and parity change) of transition.

It can be seen from expressions (1.2) that, other things being equal, the transition probability decreases rapidly with increasing multipole order and that for a given multipolarity, electric transitions are stronger than magnetic ones (subject, of course, to the operation of selection rules). It is to be expected, therefore, that most radiative transitions proceed primarily with emission of γ-rays of the lowest permissible multipole order ($L = |J_i - J_f|$ in most cases), or possibly with a mixture of the lowest two multipole orders (but of opposite class because of parity considerations). The transition probability for magnetic $2^L$-pole radiation is usually comparable with that for electric $2^{L+1}$-pole transitions. Thus, although it is rare to find a significant amount of $M2$ radiation in a mainly $E1$ transition, it is frequently found that in a mixed $M1$ and $E2$ transition the two multipolarities are comparable in strength. It sometimes happens, especially in heavy nuclei far from closed shells, that $E2$ widths are enhanced to many times the Weisskopf estimate and this is
interpreted as evidence for collective motion in the nucleus as described by Bohr and Mottleson.\cite{6,7}

The first classifications of known $\gamma$-ray transitions in terms of Weisskopf units were made by Goldhaber and Sunyar\cite{7} for isomeric nuclei where large spin changes are usually involved between the initial and final nuclear states, and by Wilkinson\cite{8} for the fast, energetic $E1$ and $M1$ transitions found in the light elements. These surveys, which have extended\cite{9,10,11}, indicate that the experimentally determined radiative widths conform in general to those expected on the basis of detailed shell model calculations, with a few exceptions such as the enhanced widths for $E2$ transitions mentioned above.

On the basis of experimental data, mainly from radiative capture reactions, Wilkinson\cite{10,11} has analysed over one hundred $\gamma$-

\begin{enumerate}
\item[8.] D.H. Wilkinson, Phil. Mag., \textbf{44}, 450, (1953).
\item[10.] D.H. Wilkinson, Phil. Mag., \textbf{1}, 127, (1956).
\end{enumerate}
ray transitions in elements with $A \lesssim 20$, in terms of Weisskopf units. Most of these were dipole transitions and it was found that the values of $|M|^2$ grouped about a mean, which showed that $E1$ transitions have a most probable width of about 0.032 Weisskopf units with a spread of a factor of seven either way. For $M1$ transitions the most probable width is 0.15 Weisskopf units with a spread of a factor of twenty either way. Wilkinson has shown further that these distributions are very similar to those obtained theoretically on the basis of the nuclear shell model with intermediate coupling. From the thirteen electric quadrupole transitions known in the light elements, it appears that $E2$ widths are frequently larger than the Weisskopf unit, a behaviour to be expected if some collective motion of the Bohr-Mottelson type exists superimposed on the general shell model structure.

1.6 **Isotopic Spin Selection Rules.**

The concept of isotopic (or isobaric) spin was developed by Wigner\textsuperscript{12} and is of great importance for nuclear reactions with light elements. On the assumption that nuclear forces are completely charge independent and that they depend only on space and spin co-

ordinates, it is possible to treat neutrons and protons as two alternative states of the "nucleon". To distinguish these two states, the nucleon is supposed to have as isotopic spin \( \mathbf{t} \), whose magnitude is \( \frac{1}{2} \), and which can have two possible orientations in "charge" space such that for a neutron \( t_z = \frac{1}{2} \) and for a proton \( t_z = -\frac{1}{2} \). This is analogous to the description of spin in ordinary space. The isotopic spins of all the nucleons in a nucleus combine according to the usual rules for the vector addition of angular momenta\(^{12,13}\) to give a total isotopic spin \( \mathbf{T} = \Sigma \mathbf{t} \), with \( T^2 = T (T+1) \), and a total z-component \( T_z = \Sigma t_z = \frac{1}{2}(N-Z) \). It is then possible to allot to a nuclear state a definite isotopic spin quantum number, \( T \), whose value is limited by the Pauli principle which can be generalised to include the isotopic spin quantum number and which requires that no two nucleons can have the same set of quantum numbers. This leads to the idea of an isotopic spin multiplet\(^{14}\). For a state of isotopic spin, \( T \), there are another \( 2T \) similar states which can be formed from this state by replacing some neutrons with protons and vice versa, without violating the Pauli principle. The \( (2T + 1) \) member states of an isotopic spin multiplet are found in members of an isobaric set of nuclei ex-


tending from $T_z = -T$ to $T_z = T$. These states will have the same spin, parity and energy. Several such sets of states are now established in the light nuclei\textsuperscript{13,15}. (To compare the experimentally determined energies of member states of an isotopic spin multiplet, it is necessary to estimate and allow for the perturbing effect of the Coulomb interaction and the n-p mass difference).

It follows from the assumption of charge independence in the nuclear Hamiltonian that the isotopic spin, $T$, of a state should be a good quantum number\textsuperscript{13}. For the light nuclei, Radicati\textsuperscript{16} and MacDonald\textsuperscript{17} have shown that in general the effects on nuclear wave functions of the Coulomb repulsion between protons and the n-p mass difference are so small that $T$ is expected to be a good quantum number in practice. The "goodness" of $T$ (after making small allowances for the Coulomb force and the n-p mass difference) can therefore be taken as a test of the assumption of charge independence, and this can be checked experimentally by testing the operation of certain isotopic spin selection rules.


If T is a good quantum number, it must be conserved in heavy particle reactions as was pointed out by Adair\textsuperscript{18}. Thus, a compound state with a given T value can only disintegrate into two particles in states with isotopic spins $T_1$ and $T_2$, if $T_1$ and $T_2$ combine vectorially to give T. That is

$$|T_1 - T_2| \leq T \leq T_1 + T_2 \ldots \ldots \ldots \ldots \ldots \text{(1.3)}$$

and conversely, collision of two particles with isotopic spins $T_1$ and $T_2$ can only form compound states with T between $|T_1 - T_2|$ and $T_1 + T_2$.

Electromagnetic transitions in nuclei are charge dependent and so it is not to be expected that isotopic spin will be conserved in the emission and absorption of $\gamma$-rays. Nevertheless, it has been shown by Radicati\textsuperscript{19} and Gell-Mann and Telegdi\textsuperscript{20} that definite selection rules do apply in such cases. The isotopic spin selection rules for radiative processes may be summarised as follows:

(a) For all multipolarities in all nuclei, $\Delta T = 0, \pm 1$, except in the case of E1 transitions in a self-conjugate ($T_z = 0$) nucleus.

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(b) For El transitions when $T_Z = 0, \Delta T = \pm 1$, only. Isotopic spin will be discussed further in Chapter III where an experiment is described in which $\gamma$-ray transitions are observed in apparent violation of this last selection rule.

1.7 The Giant Resonance.

One of the main interactions of $\gamma$-rays with nuclei is the well known "giant resonance" effect. Although this process is more familiar in connection with photo-disintegration experiments, it will be shown later that it is equally valid to consider it from the point of view of the inverse radiative capture reactions. The giant resonance shows up as a strong, broad maximum in the photo-disintegration cross-section in the region between 15 and 25 Mev excitation $^{21,22,23,24,25}$. The energy at the maximum decreases smoothly with increasing atomic number showing a dependence $^{22}$ close to $E_{\text{max}} \propto A^{-0.2}$. The width of the resonance is found to be substantially constant at about five

to six Mev and independent of atomic number except for significant reductions in the width for nuclei near closed shells. The integrated cross-section increases in an almost linear fashion with atomic number.

The first theoretical descriptions of the giant resonance absorption were given by Goldhaber and Teller who considered collective modes of dipole oscillation in the nucleus. In the model which they analyse in detail, γ-rays are considered to excite an ordered nuclear vibration in which the protons and the neutrons are considered as two interpenetrating clouds which oscillate back and forth through one another. Using reasonable values for the nuclear parameters involved, this model gives good agreement with the observed positions of giant resonance maxima and predicts an A dependence of $E_{\text{max}} \propto A^{-1/6}$, also in fair agreement with observations.

Because of the strength of the absorption, the giant resonance is generally ascribed to electric dipole radiation. Levinger and Bethe, without assuming any specific nuclear model, have derived an expression for the integrated cross-section for electric dipole absorption by a nucleus, based on the well-known sum-rules for the absorption of electromagnetic radiation. Their values, which

include an extra contribution from the neutron-proton exchange interaction, show agreement with the experimental ones suggesting that the assumption of dipole absorption is justified. The calculations indicate that quadrupole transitions can account for only about 6% of the experimentally observed cross-section. The integrated cross-section found on the Goldhaber-Teller model\textsuperscript{26} identically exhausts the non-exchange part of the dipole sum.

Other collective modes of oscillation have been examined\textsuperscript{28, 29, 30}, which differ in detail from the Goldhaber-Teller model described above but which give similar predictions for the position and strength of the giant resonance. However, in order to describe the subsequent de-excitation of the excited nucleus, it is further necessary to assume that the ordered dipole vibration following El absorption, is rapidly broken up through random nucleon collisions leading to the formation of a compound nucleus, which then decays by the emission of particles or quanta emerging in statistical competition with one another. Experimentally this description fits the energy distribution found for most of the photo-disintegration products (see


for example, references\(^{31,32,33}\)). However, from a study of the ratios of \((\gamma,p)\) to \((\gamma,n)\) yields from several elements, using \(17.6\) Mev \(\gamma\)-rays, it was shown by Hirzel and Waffler\(^{34}\), that there is also a yield of high energy protons which is too high to be explained by a compound nucleus theory. Observations of this anomalous emission of high energy photo-protons have been recorded in other experiments\(^{31,32,33}\), and it was proposed that this is a direct effect which does not proceed via a compound nucleus\(^{35,36}\).

Wilkinson\(^{37,38}\) has given a description of the giant resonance effect on the basis of the nuclear shell model. According to the model, the giant resonance can be pictured as composed of a number of individual levels, \(\mathcal{E}_n\), each of which is an eigen-state of the nuclear Hamiltonian. The only levels excited by \(\text{E}1\)-absorption

\begin{itemize}
  \item 34. O. Hirzel and H. Waffler, H.P.A., 20, 373, (1947).
  \item 35. P. Jensen, Naturwiss., 35, 190, (1948).
\end{itemize}
will be those whose wave functions contain appreciable amounts of an "ideal" state, $\Phi$, formed by operating on the ground-state wave function with the electric dipole operator, and these levels will have energies which cluster about the expectation value of the energy for $\Phi$. By using a velocity dependent potential (or, equivalently, an effective nuclear mass equal to about $1\over 2$ of its free value) in the calculations, the values of $E_{\text{max}}$ can be made to agree fairly well with those determined experimentally. The theory can also account for the observed variations in width of the giant resonance and for the strengths of the transitions. The anomalous emission of photoprotons from heavy nuclei is accounted for by a "resonance direct" mechanism by which a nucleon may be emitted from a single-particle state.

From the theoretical viewpoint, it is therefore of considerable interest to see to what extent the giant resonance can be regarded as composed of a clustering of individual levels. The level density in the giant resonance region is lower for light than for heavy nuclei, so that one would expect to observe any fine structure more readily in the light elements. Some experiments have been per-

formed which report fine structure in this region of the periodic table.

By irradiating photographic plates with bremsstrahlung, reactions such as \( \text{Li}^7(\gamma, t)\text{He}^4 \), \( \text{C}^{12}(\gamma, 3\alpha) \) and \( \text{O}^{16}(\gamma, 4\alpha) \) have been studied, and the results indicate the resonance absorption of \( \gamma \)-rays into narrow levels of the target nuclei. In reactions of this type, the energy of the photon absorbed can be determined uniquely from a knowledge of the energies of the reaction products which are usually stopped in the emulsion. A careful examination of the photo-neutron yield curves obtained with bremsstrahlung beams has been made by Katz et al. and by Penfold and Spicer. These workers find evidence in some light elements for discontinuities or

"breaks" in the slope of the yield curves which are interpreted as being due to the excitation of narrow levels. The "breaks" technique requires extremely fine control (± 5 kev) of the betatron since the difference between the observed integral yield curve with "breaks" and without "breaks" is small. Except in the region of the threshold, the differentiation of this yield curve requires taking small differences of large numbers, with the resultant disadvantage that any unavoidable fluctuations in the observed yield can produce relatively large fluctuations in the final result. Cohen at al.47 have used nuclear emulsions to examine the energies of photo-protons ejected from Be9, C12, and O16 by a bremsstrahlung beam. The first excited states of the residual nuclei are sufficiently well removed from the ground state to permit determination of the energies of the photons responsible for ejection of protons in a small energy range. After the application of corrections for the incident bremsstrahlung spectrum, for escape losses due to protons leaving the emulsion, and, at lower proton energies, for the estimated effect of transitions to excited states of the residual nuclei, their results show some evidence for fine structure in the giant resonance (γ,p) cross-sections in these elements.

Apart from experimental difficulties, most of the methods

using bremsstrahlung to determine precise $(\gamma,p)$ and $(\gamma,n)$ cross-sections in light nuclei suffer from the disadvantages inherent in the complexity of the bremsstrahlung spectrum, and in the ambiguity in determining the state to which transitions are made in the residual nucleus. Ideally, the way to measure a $(\gamma,p)$ cross-section in the giant resonance region would be to have a continuously variable, mono-energetic $\gamma$-ray source and a proton detector with resolution good enough to distinguish transitions to the ground state of the residual nucleus. So far, no such $\gamma$-ray source has been developed. However, there is an alternative method of tackling the problem.

It was pointed out by Bethe$^{48}$ in 1937 that the $(\gamma,p)$ and the $(p,\gamma)$ cross-sections for two mutually inverse reactions are related by detailed balancing. This means that the shape of the giant resonance may be determined equally well by bombarding what was, in the photo-disintegration reaction, the residual nucleus with continuously variable mono-energetic protons and using a $\gamma$-ray detector with resolution good enough to distinguish transitions to the ground state of what previously was the target nucleus. In general this would require a proton source with energy continuously variable between

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about 5 and 15 Mev, and it is only recently that such a source has been realised with the development of the tandem-style Van de Graaff generator. In Chapters IV and V of this thesis, however, experiments are described demonstrating the feasibility of measuring \((p,\gamma)\) yield curves using, as the source of protons, a fixed-energy cyclotron in conjunction with stopping foils and, as the \(\gamma\)-ray detector, the NaI(Tl) scintillation spectrometer described in Chapter II. This technique has been employed to measure the shapes of the giant resonances in Be\(^8\) and C\(^{12}\) by means of the inverse reactions Li\(^7(p,\gamma)Be\(^8\) and B\(^{11}(p,\gamma)C\(^{12}\). In Chapter VI similar experiments are described relating to regions of nuclear excitation just below the giant resonance in several other nuclei.
CHAPTER II.

THE DETECTION OF HIGH-ENERGY GAMMA RAYS WITH A LARGE NaI(Tl) CRYSTAL.

2.1 Introduction.

The energy determination for mono-energetic \( \gamma \)-rays is usually accomplished by one of two possible methods. In the first single interactions of the photons with matter are analysed and in the second the result of multiple interactions is examined. Examples of spectrometers based on the first method are the bent-crystal spectrometer\(^1\), the magnetic pair spectrometer\(^2\) and detectors using photonuclear reactions (reviewed by Bishop and Wilson\(^3\)). These spectrometers are capable of good energy resolution (< 5%), but their detection efficiencies are extremely low (generally only a small fraction of a percent) and to use them to measure the low intensity \( \gamma \)-rays characteristic of capture reactions would usually require intolerably long experimental times.

The scintillation counter is a spectrometer of the second type in which a \( \gamma \)-ray photon interacts many times with the detector

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and is frequently totally absorbed. The introduction of the single NaI(Tl) crystal detector\(^4\)\(^5\) made possible a spectrometer capable of both good resolution and high detection efficiency. For large crystals the detection efficiency can be made to approach 100\%.

Most of the \(\gamma\)-ray measurements described in the succeeding chapters were made using a large thallium-activated sodium iodide crystal, 4 in. long by 5 in. diameter. The following sections describe the behaviour of this crystal as a spectrometer.

2.2 Theoretically expected behaviour of the crystal.

A \(\gamma\)-ray quantum incident upon the crystal can interact with it in one of three main ways. At low \(\gamma\)-ray energies (below 250 kev) the photo-electric effect gives the major contribution to the total absorption cross-section, while for intermediate energies (0.25 to 7 Mev) the Compton effect predominates. For \(\gamma\)-ray energies above 1.02 Mev pair production is possible and at energies above 7 Mev in NaI(Tl) this becomes the main interaction process. Each process liberates electrons in the crystal and the resulting scintillations are amplified by the photo-multiplier as electrical pulses which can be

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sorted into a pulse-height spectrum.

Foote and Koch\(^6\) have carried out tests on a NaI(Tl) spectrometer 8 in. long by 5 in. diameter. They find that for incident \(\gamma\)-ray energies up to about 3 Mev the resolution obtained in the pulse-distribution is determined mainly by statistical fluctuations in the number of the photo-electrons liberated from the photo-cathode. For higher energies the resolution gets slightly worse due to energy losses from the crystal. Since, in the experiments to be described, the crystal was only used to examine \(\gamma\)-rays with energies greater than 3 Mev, it is of interest to examine briefly the causes of energy loss from the crystal. At these incident \(\gamma\)-ray energies the primary interaction of a photon with the crystal will usually be either a Compton scattering or a pair production event. A fraction of the incident energy may then be lost by any of the following processes:

(a) The scattered quantum following a Compton event may escape.

(b) Following pair production the electron and positron will come to rest in the crystal and the positron will annihilate producing two 0.511 Mev quanta, one or both of which may escape.

(c) For high energy \(\gamma\)-rays pair production releases fast electrons

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and positrons in the crystal and these lose part of their kinetic energy by the emission of bremsstrahlung radiation, some of which may escape.

(d) For high energy $\gamma$-rays the escape of electrons or positrons through the walls of the crystal becomes significant.

(e) Some portion of the energy of the secondary quanta mentioned in (a), (b) and (c) may escape following further interactions in the crystal.

For a large crystal the effect of processes (a), (c), (d) and (e) is to produce a low-energy "tail" on the pulse-height distribution for mono-energetic $\gamma$-radiation. The escape of one or both annihilation quanta, (b), gives rise to two satellite peaks in the distribution situated at 0.51 and 1.02 Mev below the total absorption peak.

The response function of a NaI(Tl) crystal spectrometer could, in principle, be calculated analytically from the theory of gamma-ray absorption and of energy loss of electrons.

7. i.e. the function relating the observed pulse-height distribution to the true $\gamma$-ray spectrum.

Figure 2.1 (a). Monte Carlo calculation of the pulse-height spectrum expected for 4.45 Mev γ-rays (Berger and Doggett)

(b). Spectrum of 4.45 Mev γ-rays obtained from the $^{11}_B(p,\gamma)^{12}_C$ reaction. (½ inch diameter axial collimation).
In practice, however, the multiplicity of the processes involved makes this approach impracticable, although an approximate analytical calculation has been carried out by Maeder et al.\(^8\) for small crystals. Instead, Monte Carlo calculations of the response function have been made. In these, the fate of a single quantum incident on a crystal is followed using random sampling in conjunction with known interaction probabilities to determine the sequence of interactions. When this is done for many incident quanta, a theoretical "pulse-height distribution" is obtained, which is an analogue of the physical one. Such a calculation was first done by Campbell and Boyle\(^9\) for small crystals up to 2.8 in. long by 1.6 in. diameter and for γ-ray energies of 6, 12 and 18 Mev.

Berger and Doggett\(^10\) have used an automatic computer to make a Monte Carlo calculation of the response functions for a variety of cylindrical crystals ranging in size from small (½ in. long by ½ in. diameter) to very large (9 in. long by 5 in. diameter) for radiation incident with several energies up to 4.45 Mev. Their spectrum for 4.45 Mev γ-rays, using a 4 in. long by 5 in. diameter NaI(Tl) crystal, is shown in Figure 2.1(a). It was derived by tracing the histories of 5,000 quanta incident in a finely collimated beam along the crystal axis. The peaks marked "A" and "B" in Figure 2.1(a) are the peaks

corresponding to the escape from the crystal of one and two positron annihilation quanta respectively. As is to be expected for a crystal of this size, the main peak in the spectrum is the total absorption peak. Another consequence of the size of the crystal is that the double escape peak, B, is much less intense than the single escape peak, A. This may be explained by assuming that if a positron comes to rest in such a position in the crystal that one annihilation quantum escapes, then the chances are high that the other quantum will be captured since it is emitted in the reverse direction and probably towards the bulk of the crystal.

2.3 Crystal Mounting.

The 4 in. long by 5 in. diameter crystal which was used in most of the experiments was supplied by Harshaw and is encapsulated in a machined aluminium container with 0.040 in. wall thickness. The backcap of the assembly is spun aluminium 0.020 in. thick and the glass window at the photo-tube end is 3/8 in. thick. The reflecting medium surrounding the crystal at the side walls is packed aluminium oxide 1/8 in. thick. At the backcap end a thin beryllium copper spring exerts a pressure around the periphery of the crystal, which holds it against the glass window. There is a silicone interface maintaining

11. The Harshaw Chemical Company, Cleveland, Ohio, U.S.A.
Figure 2.2. The assembled $\gamma$-ray spectrometer.
the optical coupling between the crystal and glass window and the reflector at the back of the crystal is aluminium oxide sprayed on a 0.005 in. aluminium foil sheet.

A special mounting for the crystal assembly and phototube was designed to give maximum flexibility in using the complete spectrometer. The arrangement which was used most frequently is shown schematically in Figure 2.2. The crystal was screwed, by means of its aluminium mounting ring, on to a steel plate $\frac{1}{2} \times 9 \times 9$ in. A $5\frac{1}{2}$ in. diameter hole in the centre of this plate permitted contact of the phototube with the glass window of the crystal assembly, and optical contact between the two was achieved using a high viscosity silicone grease. The phototube was screened from stray magnetic fields by a mu-metal shield and a steel cover $\frac{1}{16}$ in. thick. This cover also served as a light-tight shield and was a sliding fit over a 6 in. diameter ring screwed to the $\frac{1}{2} \times 9 \times 9$ in. plate. Light seals were effected using two thin neoprene gaskets in conjunction with the steel plate and, at the phototube base, by a neoprene ring which was clamped to the steel cover by a brass ring and fixed to the phototube base with black "Scotch" tape. This assembly could then be attached, by means of screws through the plate, to a rectangular steel block, $5\frac{1}{2} \times 9 \times 9$ in., which had been bored out to take the crystal and its aluminium mounting ring. The steel block and plate provided shield-
ing of the crystal against background radiation, and the overall dimensions of this assembly were 9 x 9 x 6 inches. These dimensions were chosen to facilitate using the steel block in conjunction with the 9 x 3 x 3 in. lead bricks in the laboratory.

Three stepped holes labelled (a), (b), and (c) in Figure 2.2, were bored in the block to allow a γ-ray beam to be collimated onto the crystal in one of three directions. The three methods of collimation used were:

(a) Axial collimation, where the γ-rays are collimated along the axis of the crystal cylinder.

(b) Side collimation, where the γ-rays are collimated along a diameter in the centre of the crystal cylinder.

(c) Diagonal collimation, where the γ-rays are collimated along a diagonal of the crystal cylinder.

The diameter of a collimating hole could be changed by inserting a suitably machined steel plug, and when not in use, collimating holes were plugged with solid pieces of steel. Collimating holes up to 1 in. diameter could be used with axial and side collimation and up to \( \frac{1}{2} \) in. diameter with diagonal collimation. Several lead shields having collimating holes of various sizes were made up for use with the steel block. When no collimation was required, the combination of crystal and phototube could be removed from the steel block by
unscrewing and removing the steel plate.

The phototube employed was selected from a group of six Dumont 6364 tubes on the basis of resolution and stability of gain at high counting rates. The chain of bleeder resistors for the phototube together with a cathode-follower output stage, was contained in a small head amplifier attached to the phototube socket. Positive going pulses from the last dynode passed through the cathode-follower into a Higinbotham non-overload amplifier and then into a Hutchinson-Scarrott multi-channel analyser.

2.4 Performance of the Crystal.

(a) Comparison with Theory.

Figure 2.1(b) shows a pulse-height spectrum of 4.45 Mev \( \gamma \)-rays obtained from the \( ^{11}\text{B}(p, \gamma)^{12}\text{C} \) reaction at the 160 kev resonance. The low energy tail produced by the 11.7 Mev \( \gamma \)-rays arising from the same reaction has been estimated and subtracted off to give this spectrum. Gamma rays from the boron target were collimated into a beam along the crystal axis by a \( \frac{1}{2} \) in. diameter hole in a 3 in. thick lead wall. The experimental spectrum, Figure 2.1(b), is there-

13. G.W. Hutchinson and G.G. Scarrott, Phil. Mag., 42, 792, (1951).
fore directly comparable with the theoretical one, Figure 2.1(a), since the Monte Carlo calculation was made assuming fine axial collimation. Although the theoretical result is in the form of a rather rough histogram, the gross features of the distributions may be compared.

The general shapes of the two spectra are similar except that the experimental one has a much more pronounced low energy tail which rises steeply at the lowest energies. Most of this tail (certainly that part above the dashed line in Figure 2.1(b) ) is probably attributable to low energy electrons and scattered quanta arising from interactions of the primary γ-ray beam with the walls of the collimating channel. This effect is not considered in the calculation where a "clean" collimation is assumed.

The experimentally determined photo-fraction calculated for the spectrum of Figure 2.1(b), excluding the low-energy pulses above the dashed line, is 0.27, which is significantly lower than the theoretical value of 0.49. The discrepancy could be due to the fact that the calculation does not take into account energy losses due to electron escape from the crystal. Also, at high γ-ray energies,

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14. That is, the number of pulses in the full-energy peak expressed as a fraction of the total number of pulses.
Figure 2.3. Spectra of 4.43 and 11.7 Mev γ-rays at the 160 kev resonance in the $^9\text{Be}(p,\gamma)^{12}\text{C}$ reaction. (1 in. diameter side collimation).
Figure 2.4. Spectrum of $\gamma$-rays from the 340 kev resonance in the $^{19}_F(p,\alpha,\gamma)^{16}O$ reaction. (1 in. diameter side collimation)
Figure 2.5. Spectrum of $\gamma$-rays from the 550 kev resonance in the $^{13}(p,\gamma)^{14}$ reaction. (1 in. diameter side collimation)
the approximations used by Berger and Doggett for the escape of secondary bremsstrahlung and annihilation quanta become less reliable.

2.4(b) Spectra at Various Gamma-ray Energies.

Figures 2.3 to 2.6 show a variety of pulse-height distributions obtained using γ-rays from some proton capture reactions. Each of these spectra was recorded using side collimation through a lead wall with a collimating hole 3 in. long by 1 in. diameter.

The distributions from the 4.4 and 11.7 Mev γ-rays arising from the 160 kev resonance in B\(^{11}\)(p,γ)C\(^{12}\) are shown in Figures 2.3(a) and (b). The estimated low-energy tail of the 11.7 Mev radiation is shown in Figure 2.3(a) as a broken line. The distribution produced by the 6.1 Mev γ-rays from the 340 kev resonance in the F\(^{19}\)(p,α,γ)C\(^{16}\) reaction is shown in Figure 2.4. The less intense 7.1 Mev radiation is also resolved. Figure 2.5 shows the spectrum of 8.1 Mev γ-rays arising from the 550 kev resonance in C\(^{13}\)(p,γ)N\(^{14}\). The small peak at about 4 Mev in the spectrum is due to cascade γ-rays proceeding through the 3.95 Mev level in N\(^{14}\).

It is evident from these spectra that, for higher energy γ-rays, the single escape peak becomes larger than the full-energy peak, and at energies of greater than about 12 Mev it is to be expected that the two peaks will not be resolved from each other. This
Figure 2.6. Resonance and non-resonance γ-ray spectra from the $^{7}\text{Li}(p,\gamma)^{8}\text{Be}$ reaction. (1 in. diameter side collimation)
effect is illustrated in Figure 2.6 which shows two spectra obtained using the Li$^7$(p,$\gamma$)Be$^8$ reaction at resonant and non-resonant proton energies. The lithium metal target was 80 kev thick. The two $\gamma$-ray lines from this reaction are a sharp 17.6 Mev line produced by transitions to the ground-state of Be$^8$, and a broad 14.8 Mev component produced by transitions to the broad 2.9 Mev state in Be$^8$ ($\Gamma \approx 2$ Mev). At the 441 kev resonance, transitions to the ground state predominate while at higher proton energies the main transitions are to the 2.9 Mev state. The broad nature of the 14.8 Mev component is clearly evident from the spectra.

2.4 (c) The Effect of Direction of Collimation.

With large cylindrical crystals it has generally been the practice to collimate (when necessary) the incident $\gamma$-radiation into a narrow beam along the crystal axis$^{6,15}$. Tests were carried out, however, which show that the best results with a 4 in. long by 5 in. diameter crystal are not necessarily obtained using axial collimation.

With a 3 in. long by $\frac{1}{2}$ in. diameter collimating hole in

Figure 2.7. Gamma-ray spectra recorded at the 340 kev resonance in the $^{19}(p,α,γ)^{16}$ reaction for three different directions of collimation.
a lead wall, spectra were recorded at various $\gamma$-ray energies between 4.4 and 17.6 Mev using, in turn, each of the three directions of collimation described in §2.3. At each $\gamma$-ray energy the result was the same: side collimation gave the best resolution, diagonal collimation was nearly as good and axial collimation was distinctly worse. This is illustrated in Figure 2.7 where the high-energy region of the $\gamma$-ray spectrum from the 340 kev resonance in the $^{19}\text{(p,}\alpha,\gamma)^{16}$ reaction is examined using the three directions of collimation. It will be seen that the proportion of annihilation quanta which escape from the crystal is significantly greater in the case of axial collimation than in the case of either of the other two methods. Further, it is apparent that the width of the full-energy peak is smaller for side collimation than it is for either diagonal or axial collimation.

In an attempt to explain these effects, at least qualitatively, it is convenient to distinguish between two causes of energy loss from the crystal: these are (i) the escape of 0.511 Mev annihilation quanta and (ii) the escape of energy as a result of any of the other processes listed in §2.2. Escape of type (i) gives rise to two satellite peaks which may or may not be resolved in the pulse-height distribution, while, at these incident $\gamma$-ray energies, escape of type (ii) produces an asymmetrical broadening of
both the full-energy peak and the satellite peak. If, in Figure 2.7, we take the relative heights of the full-energy peak and the single-escape peak as an index of the resolution, then both types (i) and (ii) of escape are included since a broadening of the full-energy peak will also give a contribution under the single escape peak, thereby increasing the apparent height of this peak.

Let us further make the assumption that the electrons and positrons created in the crystal have zero range, an approximation which is expected to be reasonable for the γ-ray energies considered.  

This means that the interactions of primary photons with the crystal may be regarded as equivalent to a linear source of secondary photons (annihilation quanta, bremsstrahlung and quanta from Compton scattering events) along the direction of collimation of the incident beam, and with a source strength decreasing in an approximately exponential fashion with depth in the crystal. We now consider the relative merits of different directions of collimation in terms of the escape from the crystal of some of these secondary quanta.

The escape of bremsstrahlung and Compton photons will give rise to energy losses of the type (ii) referred to previously.

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16. e.g. 4 Mev electrons have a range of 0.86 cm. in NaI(Tl) and this is small compared with the crystal dimensions.
These quanta are emitted preferentially towards the forward direction. Therefore, to reduce this type of energy loss for a crystal whose dimensions of length and breadth are comparable, one would expect to obtain the best results by using a type of collimation which gave a depth in the crystal (measured along the direction of the γ-ray beam) larger than the width (measured at right angles to the beam). This would account, in part, for the improvement obtained in going from axial collimation (depth, d = 4 in.) to side (d = 5 in.) or diagonal (d = 6.4 in.) collimation.

It can be seen from Figure 2.7 that the breadth of the full-energy peak is smaller for side than for diagonal collimation in spite of the greater depth available with the latter method. This could be due to a corresponding decrease in the "effective" width of the crystal and possibly due to uneven light collection. Side collimation has the advantage that the scintillations occur at points in the crystal which are roughly equi-distant from the cathode of the photo-tube making light collection uniformly efficient.

Energy losses of the type (i) occur when positron annihilation quanta escape. The two 0.511 Mev quanta are usually

created only after the positron has come to rest so that they are emitted isotropically and in opposite directions. Of the total number of primary interactions between the incident γ-rays and the crystal, the fraction which contribute to the escape peak will be determined mainly by the fraction of the interactions occurring near a face of the crystal. To a first approximation, we may consider an annihilation quantum as escaping if it is formed in the crystal, at a point whose distance from any face of the crystal is less than a mean free path (for 0.511 Mev quantum in NaI(Tl), this is 1.21 in.\textsuperscript{18}). If we regard the sources of annihilation quanta as distributed exponentially along the line of collimation through the crystal, then it is evident that the fraction of these sources located less than 1.21 in. from a crystal face is smaller for side collimation than for axial collimation. For diagonal collimation, the fraction seems to be about the same as that for side collimation (Figure 2.7). The advantage of the increased depth obtained with diagonal collimation is probably off-set by the fact that the γ-rays enter and leave the crystal through corners (rather than through flat, or nearly flat, surfaces as in the other methods of collimation) with a consequent increase in the escape of annihilation quanta.

\textsuperscript{18} G.R. White, NBS Report No.1003 (1952).
Figure 2.8. Gamma-ray spectra recorded at the 340 kev resonance in the $^{19}$F$(p,\alpha,\gamma)^{16}$O reaction showing the effect of collimating diameter.
2.4(d) The Effect of Collimating Diameter.

Using side collimation through a 3 in. lead wall, spectra were recorded at \( \gamma \)-ray energies between 4.4 and 17.6 Mev to examine the effect of variation in diameter of the collimating channel. The results were similar throughout the range of \( \gamma \)-ray energies and are shown in Figure 2.8 for the case of fluorine \( \gamma \)-rays. With no collimation at all the energy escape from the crystal is much higher than with collimation: indeed for uncollimated 6.1 Mev \( \gamma \)-rays the escape peak is more intense than the full-energy peak. By collimating the \( \gamma \)-rays into a beam, it is ensured that most of the secondary radiation is emitted from points not far from the line of collimation in the crystal, and the chances of escape are correspondingly lower. In going from 1 in. to \( \frac{1}{2} \) in. diameter collimation an improvement in resolution is obtained, but further reduction in the diameter of the hole (1/4 and 1/8 in. diameter collimations were tried) does not reduce the proportion of energy escape appreciably. In fact, for collimating diameters of less than \( \frac{1}{2} \) in. there is a disadvantage in that the low-energy "tail" in the spectrum begins to grow prohibitively large relative to the full-energy peak. This effect is probably due to an increase in the fraction of low-energy quanta and electrons coming from the walls of the collimator (since the wall area
of the collimating hole is increased relative to its cross-sectional area at smaller diameters).

Even if a "clean" and infinitely fine collimation could be obtained in practice, one would not expect to get resolution which was significantly better than that obtained with $\frac{1}{2}$ in. diameter collimation. The reason for this is that, although the secondary electrons released in the crystal travel mainly forwards, § 2.4(c)) there is still an appreciable side-ways scatter and even for an infinitely fine collimation this scatter is of the order of $\frac{1}{2}$ in. diameter.

2.4(e) Cosmic Rays.

The passage of cosmic rays (mainly relativistic $\mu$-mesons) through the crystal gives rise to background pulses whose heights are proportional to the energy losses of the incident particles. For $\mu$-mesons which pass right through the crystal, the energy loss is roughly proportional to the path length of the meson in the crystal. This effect has been studied by Hudson and Hofstadter\textsuperscript{19} who find that the most probable energy loss of $\mu$-mesons in NaI is 12.5 Mev per in. of path length, and that for a given path length, the

Figure 2.9. Spectrum of cosmic-rays.
energy straggling is unsymmetrical having a high-energy "tail". Thus \( \mu \)-mesons travelling down the crystal axis would be expected to lose about 50 Mev, and those going along a diameter would lose 62.5 Mev. The longest possible path in the crystal is a diagonal one (6.4 in.) which corresponds to 80 Mev energy loss.

Figure 2.9 shows a cosmic ray spectrum recorded at an elevation of 2,000 feet above sea level over a period of one hour. The energy scale extends out to 160 Mev and has been calibrated using 17.6 Mev \( \gamma \)-rays from the 441 kev resonance in the \( \text{Li}^7(p,\gamma)\text{Be}^8 \) reaction. The linearity with energy of the crystal and its associated electronics has been checked up to 17.6 Mev and after taking precautions against possible saturation effects in the photo-tube and the pulse amplifier, it was assumed that the energy scale was linear to 160 Mev. The spectrum of Figure 2.9 was recorded under two feet of concrete, and with the axis of the crystal vertical. It shows a prominent peak at 60 Mev superimposed on a pulse distribution which steadily decreases with energy.

The position of the peak in the spectrum corresponds well with that to be expected from \( \mu \)-mesons with large path-lengths in the crystal. The exact shape and position of this peak will be a
function of the shape and orientation of the crystal cylinder, the angular distribution of the incident $\mu$-mesons, and the straggling of the energy losses in the crystal. The large number of smaller pulses is probably due to $\mu$-mesons with short path lengths. At energies greater than 100 Mev in the spectrum it is unlikely that the pulses are due to the passage of $\mu$-mesons. These pulses extending out to at least 160 Mev are probably caused by stars or nuclear interactions initiated in the crystal by the N component of the cosmic rays. Standil\textsuperscript{20} has observed such pulses extending out to over 1,000 Mev in NaI.

CHAPTER III.

ISOTOPIC SPIN IMPURITIES IN LIGHT NUCLEI.

3.1 Introduction.

Assuming the complete charge independence of nuclear forces and neglecting the effect of the Coulomb interaction between nucleons, the isotopic spin, $T$, of a nuclear state is expected to be a good quantum number. There is so far no direct evidence against the assumption of charge independence; the results of nucleon-nucleon scattering experiments and the well-known correspondence of levels in mirror nuclei both support the assumption. A further test of this hypothesis can be made by checking the isotopic spin selection rules (§ 1.6) by experiment. In this chapter, an experiment is described which investigates an apparently flagrant violation of the selection rule $\Delta T = \pm 1$, for $E1$ transitions in a self-conjugate nucleus.

3.2 Theoretical Considerations.

The Coulomb interaction can be treated as a perturbation in an otherwise charge-independent system (other charge-dependent electromagnetic interactions are negligible). This, in effect, makes $T$ a "nearly good" quantum number and results in a partial relaxation of the isotopic spin selection rules. Each nuclear state is no longer
a pure T state, but is a mixture of states with differing isotopic spins. Following the treatment given by Radicati\textsuperscript{1}, this may be written, in first order perturbation theory,

\[ \Phi = \xi(T) + \sum_{T'} \alpha_{T'}(T') \xi(T') \]

where \( T \) is the isotopic spin of the nuclear state expected on the basis of complete charge independence and \( \xi(T) \) is the corresponding wave function. \( \xi \) is the actual wave function and \( \xi(T') \) are states with the same spin and parity as \( \xi(T) \) but with differing isotopic spins, \( T' \). As a result of the Coulomb interaction, these states are mixed in with amplitudes \( \alpha_{T}(T') \). If \( E_T \) and \( E_{T'} \) are the expectation values of the energies of the states \( \xi(T) \) and \( \xi(T') \) respectively, then

\[ \alpha_{T}(T') = \frac{H_{TT'}^C}{E_T - E_{T'}} \]

where \( H_{TT'}^C \) is the matrix element of the Coulomb interaction between the states \( \xi(T) \) and \( \xi(T') \). If a particular transition from a state whose isotopic spin is mainly T, is forbidden by the selection rules, but is allowed from a state of isotopic spin \( T' \), then the transition will occur with an intensity \( \alpha_{T}^2(T') \) times that expected if there were no isotopic spin selection rules. For \( \gamma \)-ray and \( \alpha \)-particle emission, the intensity expected without these selection rules can be estimated

from the most probable values of \(|M|^2\) and of \(\Theta^2\) (the reduced \(\alpha\)-particle width in single-particle units \(\hbar^2/mR\)). These are given by Wilkinson\(^2\) for elements with \(A \leq 20\). Alternatively \(\alpha_T^2(T^i)\) may be estimated from a knowledge of the branching ratio between a forbidden and an allowed transition.

Explicit shell model calculations of the values to be expected for \(H_{TT}^0\) in the ground states of some light nuclei have been made by Radicati\(^1\) and in more detail by MacDonald\(^3\). Small admixtures to the states of mainly \(T = 0\) were found, corresponding to values of \(\alpha_T^2(T^i)\) ranging from \(1.3 \times 10^{-5}\) (\(\text{He}^4\)), to \(1.0 \times 10^{-3}\) (\(\text{B}^{10}\)) and \(6.7 \times 10^{-3}\) (\(\text{O}^{16}\)). The impurities in states up to a few Mev excitation are not expected to be very different from these estimates for the ground states\(^3\), but at higher excitations (above 10 Mev, say) the values of \(\alpha_T^2(T^i)\) will grow larger because of the decrease in separation of states with differing \(T\) values.

Even if there were no Coulomb interaction and \(T\) were an absolutely good quantum number, El transitions with \(\Delta T = 0\) in \(T_Z = 0\)


nuclei would not be completely forbidden. Gell-Mann and Telegdi\textsuperscript{5} have shown that such El transitions would still occur with an intensity appropriate to a magnetic quadrupole transition. That is, they would be reduced by a factor of order $2 \times 10^{-7} \frac{E_y^2}{E}$ where $E_y$ is the $\gamma$-ray energy in Mev (see Equations 1.2). Even for $\gamma$-ray energies as high as 10 Mev this factor is still only $2 \times 10^{-5}$. Consequently it is to be expected that the presence of isotopic spin impurities in states will remain the chief cause of failure of selection rules.

Wilkinson\textsuperscript{2} has demonstrated that the isotopic spin selection rules for El transitions are certainly effective. He has shown that for thirteen established forbidden transitions, the values of $|M|^2$ form a group which is unmistakably lower (a factor of 100 or so) than that for the bulk of the allowed El transitions.

3.3 The 6.89 Mev Level in $^{10}$B.

One of the largest isotopic spin impurities reported in a light nucleus is found\textsuperscript{6} in the 6.89 Mev level in $^{10}$B. This broad

Figure 3.1. Relevant Energy levels of $^{10}$B.
level, of width 140 kev, is observed as the 330 kev resonance in the reactions Be + p and can decay by the emission of deuterons, α-particles, protons or γ-radiation. It is thought to be a 1− state formed by the capture of s-wave protons. The isotopic spin of the level has been taken to be mainly T = 0 because of the intense El γ-ray transition to the T = 1 state at 1.74 Mev (see Figure 3.1) and also because of the alternative modes of decay by deuteron or α-particle emission to the ground states of Be8 and Li6 respectively. The reduced deuteron width, in particular, is quite large (θd = 0.18). If the T = 0 assignment is correct, however, it is difficult to interpret the strong transitions to the 1+, T = 0 states at 0.72 and 2.15 Mev which are highly forbidden by the isotopic spin selection rules for electric dipole transitions in self-conjugate nuclei. These transitions imply a T = 1 impurity in the 6.89 Mev level of approximately 20% intensity, which is distinctly the largest impurity yet found in a light nucleus.

An alternative explanation which would avoid the necessity


8. On the basis of forbidden α-particle transitions, it has recently been suggested that the 17.22 and 13.06 Mev states in C12 and O16 respectively, may contain similarly large impurities. At excitations this high, however, such impurities are not unexpected.
for such a large isotopic spin impurity is that, instead of the transitions occurring from a single state of $^9\text{Be}$ at 6.89 Mev, they occur from two closely adjacent levels, one of which has mainly $T = 0$, and the other mainly $T = 1$. It was to test this possible explanation that the following experiments on the $^9\text{Be}(p,\gamma)^{10}\text{Be}$ reaction were performed.

3.4 Experimental Method.

It should be possible to distinguish two such levels by measuring the excitation functions for the individual high-energy $\gamma$-rays as a function of proton energy. Transitions from the $T = 1$ state would be resonant at a proton energy differing from the resonance energy for transitions from $T = 0$ state. Also, since the emission of deuterons or $\alpha$-particles from the $T = 1$ state would be discouraged by isotopic spin selection rules, it might be expected that the total widths would not be the same for the two levels.

To determine these excitation functions, accurate spectra of the high-energy $\gamma$-radiation from $^{10}\text{Be}$ were recorded at 10 kev intervals for proton energies between 220 and 440 kev. Each spectrum was then analysed to find the contributions from its constituent $\gamma$-rays and yield curves for these $\gamma$-rays were then plotted.
The thin beryllium targets were prepared by evaporation in vacuo of beryllium metal onto thin copper backings. A sensitive balance was used to weigh the target backings before and after evaporation and then the target thickness was determined from the data of Warshaw. After some trials, targets could be made 10 kev thick at 300 kev proton energy and these were used in the experiments. Some difficulty was experienced in obtaining beryllium which was free of fluorine since the metal is frequently prepared from the fluoride. Beryllium from several different sources was tried and that which was finally used in the experiments had a fluorine contamination which was low enough to permit analysis of spectra at proton energies of 340 and 350 kev, in spite of the large number of 6.1 Mev γ-rays from the strong $^{19}$F$(p,γ^{0})^{16}$ resonance at 340 kev.

The protons were accelerated in a 1.2 Mev Cockcroft-Walton H.T. set, and after deflection by a 90° analysing magnet, passed through a 3/8 in. diameter collimating hole to the target. A cylindrical electrode in front of the target was held at a negative potential to prevent the escape of secondary electrons from the beryllium, and the number of protons incident on the beryllium was monitored with a beam-current integrator. To reduce deposition of carbon, the target was held at a temperature of 90°C by passing

hot water over the target backing, and a liquid air trap was inserted in the target tube to remove oil vapour.

3.5 Detection of Gamma Rays.

The total cross-section for the \( \text{Be}^9(p,\gamma)\text{B}^{10} \) reaction is small (approximately 20\( \mu \)b at the 330 kev resonance) and the \( \gamma \)-ray spectrum is complex. The 4 in. long by 5 in. diameter NaI(Tl) crystal spectrometer described in the previous chapter is therefore well suited as the \( \gamma \)-ray detector by virtue of its high sensitivity and resolution.

The combination of crystal and phototube was placed with its axis vertical and with the curved face of the crystal 6 in. from the beryllium target in line with the incoming protons. Gamma radiation from the target was collimated by a 1 in. diameter hole in a lead wall 3 in. thick, and by a similar hole in the large steel block in which the crystal was seated and which served to shield the crystal from scattered and background radiation. Pulses from the phototube passed through a non-overload amplifier into a 70-channel Sunvic pulse-height analyser.

Irradiation of up to six hours at beam currents of 20\( \mu \)A were required at each of these proton energies in order to obtain spectra sufficiently accurate for analysis. Figure 3.2(a) shows a...
Figure 3.2 (a). Spectrum of high-energy $\gamma$-rays from the reaction $\text{Be}^9(p,\gamma)\text{B}^{10}$ at a proton energy of 300 kev.

(b). Spectrum of $\gamma$-rays from the 340 kev resonance in the reaction $\text{F}^{19}(p,\alpha,\gamma)\text{O}^{16}$. 
typical pulse-height spectrum at a bombarding proton energy of 300 kev. Only γ-rays with energies above 4 Mev were recorded. The full-energy peaks and peaks corresponding to the escape from the crystal of one 511 kev positron annihilation quantum are marked.

For comparison, the pulse-height distribution produced with the same experimental geometry by the 6.14 Mev γ-rays arising from the 340 kev resonance in the $^{19}$F$(p,\alpha,\gamma)^{16}$ reaction is shown in Figure 3.2(b). The lower intensity 7.12 Mev radiation is also resolved. The size of the escape peak A could be further reduced and the resolution improved by collimation of the γ-rays into a narrower beam (see §2.4(d)). However, in this experiment a 1 in. diameter collimator gave a satisfactory compromise between resolution and counting rate. The shoulder, marked B in Figure 3.2(b), on the 6.14 Mev spectrum is due to the peak corresponding to escape of both annihilation quanta and is greatly reduced relative to the single-escape peak A because of the large dimensions of the crystal.

3.6 Analysis.

For the analysis of spectra it was necessary to know the individual pulse-height distributions produced by γ-rays of 4.7, 5.2, 6.2 and 6.9 Mev energy. These were derived by interpolation between the distributions produced by a number of calibration γ-rays of known energy. By substituting a $^{13}$C target for the beryllium one, the pulse-height distribution produced by 8.1 Mev γ-rays arising from
Figure 3.3. Yield curves of high-energy γ-rays. The probable errors are shown at a proton energy of 370 kev.
the 550 kev resonance in the $^{13}\text{(p,}\gamma)^{14}$ reaction was measured. Similarly, the 340 kev resonance in the $^{19}\text{(p,}\alpha,\gamma)^{0}\text{O}$ reaction gave the distribution for 6.1 Mev radiation. Pulse-height distributions were also measured using the reaction $^{11}\text{(p,}\gamma)^{12}$ at the 160 kev resonance; this gave two calibration $\gamma$-rays of 4.4 and 11.7 Mev energy. During the measurement of each calibration distribution, the geometry of the experiment was kept exactly the same as that for the beryllium runs, in order to avoid any errors due to changes in resolution.

The analysis of each $^{9}\text{(p,}\gamma)^{10}$ spectrum was performed by successively subtracting the distributions due to the 6.9, 6.2, 5.2 and 4.7 Mev $\gamma$-rays. The resulting yield curves for the individual $\gamma$-rays are shown in Figure 3.3, where the errors are calculated on the assumption that the distributions due to individual $\gamma$-rays are known exactly. The high values on the 6.2 Mev curve at proton energies of 340 and 350 kev (shown on the dashed lines in Figure 3.3) are due to the 340 kev resonance in the $^{19}\text{(p,}\alpha,\gamma)^{0}\text{O}$ reaction. The result of a run on a blank copper target showed that the contaminating fluorine was located in the beryllium. Since the fluorine resonance is narrow $^{7}(\Gamma = 2.9$ kev) and the fluorine (being distributed evenly throughout the beryllium) formed a target 10 kev
Figure 3.4. Excitation function of the total radiation resonant at 330 kev, after allowing for barrier penetrability and proton wavelength. The broken line represents a single level dispersion formula with a resonance energy of 310 kev and a width of 135 kev.
thick, the $\gamma$-ray spectra were only affected at two proton energies. The yield curve for 6.2 Mev $\gamma$-rays was corrected for this effect by drawing a smooth curve through the points on either side of the fluorine resonance.

It will be seen from Figure 3.3 that, excluding the 6.9 Mev radiation, all the $\gamma$-rays are resonant at the same energy to within ten kilovolts, and the yield curves have approximately the same width. This indicates that the explanation of two close states is untenable, and that the transitions to the $T = 1$ and to the $T = 0$ states do, in fact, come from the same state of $B^{10}$.

From the yield curves shown in Figure 3.3, it is possible to derive the excitation function of the total resonant radiation. After using Coulomb wave functions according to Christy and Latter\textsuperscript{10} to allow for the effect of barrier penetrability for s-wave protons, the curve shown in Figure 3.4 was obtained. This could be fitted with a Breit-Wigner one-level dispersion formula with resonance energy 310 kev and a width of 135 kev.

The ground state transition, which was well resolved in all of the spectra, is clearly non-resonant over the range of proton en-

\textsuperscript{10} R.F. Christy and R. Latter, Rev. Mod. Phys., 20, 185, (1948).
ergies examined. The steadily increasing yield suggests that the 6.9 Mev \( \gamma \)-rays arise from the tail of the 993 kev resonance which is known to emit \( E_1 \) radiation predominantly to the ground state of \( B^\alpha \). However, the yield of 6.9 Mev \( \gamma \)-rays at 330 kev was calculated to be five times that expected from the tail of the higher resonance. This discrepancy prompted a separate study of the ground state transition over the range of proton energies between 300 and 1060 kev and the results of this are given in Appendix A.

3.7 Other Measurements.

(a) Angular Dependence of Gamma Rays.

At proton energies of 260, 330 and 400 kev, the high-energy \( \gamma \)-ray spectra at 0° and 90° to the incoming proton beam were closely compared. It was not possible to distinguish between the shapes of the spectra, and the total yields of high-energy \( \gamma \)-rays at 0° and 90° did not differ by more than 2%.

(b) Low Energy Gamma Rays.

Spectra of the low-energy \( \gamma \)-rays accompanying this reaction were obtained at several bombarding energies, and each spectrum was checked to ensure that the low-energy \( \gamma \)-rays fitted into

Figure 3.5. Spectrum of low-energy $\gamma$-rays from the reaction $\text{Be}^9(p,\gamma)\text{B}^{10}$ at a proton energy of 310 kev. The $\gamma$-ray energies corresponding to the various peaks are indicated.
the expected decay scheme of $\text{B}^{10}$ in conjunction with the high-energy radiation (see Figure 3.1). This fit was found to be in agreement with the high-energy data to within $20\%$, except for the presence of a 1.7 Mev $\gamma$-ray, which is discussed in the next section. Figure 3.5 shows a typical spectrum of low-energy $\gamma$-rays recorded at a proton energy of 310 kev with a 2 in. long by 1.5 in. diameter NaI(Tl) crystal. The photo-peak from the 0.41 Mev $\gamma$-ray is obscured in the distribution by the Compton edge due to the 0.72 Mev $\gamma$-ray.

(c). The 5.11 and 5.16 Mev Levels.

It has been suggested by Clegg $^{12}$ that a 1.7 Mev $\gamma$-ray transition occurs from the 6.89 Mev level in $\text{B}^{10}$ with an intensity $20\%$ of that for the transition to the 1.74 Mev level. A 1.7 Mev $\gamma$-ray was also observed in the present work (Figure 3.5) and a coincidence experiment was performed at a bombarding proton energy of 310 kev to investigate possible cascades involving it and other $\gamma$-radiation.

A NaI(Tl) crystal 2 in. long by 1.5 in. diameter was used to examine the low-energy $\gamma$-rays and the spectrum from this crystal was observed in coincidence with pulses from the big crystal

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$^{12}$ A.B. Clegg, Phil. Mag., 1, 1116, (1956).
which were larger than the pulse height corresponding to 3.5 Mev. The small crystal was placed at 90° to the proton beam while the big crystal remained at 0°. In order to obtain a workable coincidence counting rate, the big crystal was used without any collimation and with its curved face only 2 in. from the target. It was also necessary to keep the small crystal 6 in. from the target to avoid a spurious 1.7 Mev summation peak from the cascade γ-rays of 1.02 and 0.72 Mev occurring in coincidence with the intense transition to the 1.74 Mev state. From a comparison of the gated and ungated spectra it was estimated that less than 5% of the 1.7 Mev γ-rays are in coincidence with any γ-radiation of energy greater than 3.5 Mev.

It is unlikely that the 1.7 Mev γ-ray is a transition from the 1.74 Mev level to the ground state of B\textsuperscript{10} as this would imply a magnetic octupole transition and furthermore one would expect it to be observed in coincidence with 5.2 Mev γ-rays. Clegg\textsuperscript{12} has measured the energy of the γ-ray carefully and his results indicate that it arises from a transition from the 6.89 Mev level in B\textsuperscript{10} to the 2\textsuperscript{−}, T = 0 level at 5.11 Mev (Γ = 2 kev), and not to the neighbouring 2\textsuperscript{+}, T = 1 level at 5.16 Mev (Γ < 1 kev). Since the 5.11 Mev level has T = 0, it should decay mainly by alpha-particle emission to the ground state of Li\textsuperscript{6}; the only γ-ray transitions would be a weak M2 trans-
ition to the 1.74 Mev level and comparably weak isotopic-spin forbidden El transitions to the other low lying levels. As a result of the bias level in the big crystal, the 1.7 Mev γ-rays could only be observed in coincidence with transitions from the 5.11 Mev level to the 0.72 Mev level or to the ground state. The absence of such coincidences supports the suggestion that the 5.11 Mev level (if it is the level involved) decays principally by α-particle emission.

If the transitions proceed through the 5.11 Mev level, as indicated by Clegg, then this implies that the 1.7 Mev γ-ray is an abnormally strong M1 transition (|M1|^2 = 6.2). Recent measurements by Meyerhof\(^{13}\) suggest that it is in fact mainly the 5.16 Mev and not the 5.11 Mev level which is involved, implying a strong El transition (|M1|^2 = 0.39) for the 1.7 Mev γ-ray. The 5.16 Mev level has \(T = 1\) and the partial widths for γ-ray and α-particle emission are thought to be comparable\(^{14}\). It can decay by allowed E2 or M1 transitions to all of the low lying levels. Assuming that, of these transitions, 50% go the 0.72 Mev level and the ground state, then the present results indicate that, the 5.16 Mev level (if it is involved), has a γ-ray

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Figure 3.6. Ratios of yields of high-energy γ-rays to that of the 5.15 Mev γ-ray, as a function of proton energy.
width which is less than 10% of the total width.

3.8 Discussion of Results.

The position and width of the 330 kev resonance determined from this experiment are in agreement with those found previously\(^7\) and the isotropy of the high-energy \(\gamma\)-radiation, together with the fact that the \(\gamma\)-ray branching ratios are independent of angle (§ 3.7(a)), support the suggestion that this is an \(s\)-wave resonance.

It is interesting to note that the \(\gamma\)-rays resonant at 330 kev do not bear quite the same ratios to one another over the resonance (see Figure 3.6). This effect could be caused by direct transitions mixing with the resonant radiation. Evidence for the presence of direct transitions to the ground state of \(^{10}\)Be is given in Appendix A and this suggests the possibility that there may be direct transitions to some of the excited states as well.

If we assume that a single-level Breit-Wigner formula is applicable to the \(^9\)Be\(^{(p,\gamma)}\)\(^{10}\)Be reaction at the 330 kev resonance, then the peak cross-section for the emission of 5.2 Mev \(\gamma\)-rays to the 1.74 Mev level is given by the equation:

\[
\delta_{\text{max}}(5.2) = \frac{2J+1}{(2s+1)(2j+1)} \cdot \frac{\pi X^2 \Gamma_p(5.2) \Gamma_p}{\Gamma^2/4},
\]

\[\text{………………(3.3)}\]
where $\Gamma_\gamma(5.2)$ is the partial width for emission of 5.2 Mev $\gamma$-rays, $\Gamma_p$ is the partial proton width (given by Meyerhof and Chase\textsuperscript{15} as 40 kev), $\Gamma$ is the total width of the 6.89 Mev level (equal to 130 kev\textsuperscript{7}), and $\lambda$ is the proton wavelength (equal to 8.8 x 10\textsuperscript{-13} cm.). Each of these values is quoted in the centre-of-mass system. The angular momentum of the 6.89 Mev level in B\textsuperscript{10} is $J = 1$, that of the ground state of Be\textsuperscript{9} is $j = 3/2$ and the proton spin is $\frac{1}{2}$. This gives a statistical factor of 3/8 in Eq.(3.3).

Using the value determined by Carlson and Nelson\textsuperscript{16,17} of $\sigma_{\text{max.}}(5.2) = 12 \pm 4 \times 10^{-30} \text{ cm}^2$, Eq.(3.3) gives a value of $\Gamma_\gamma(5.2) = 1.4 \pm 0.5$ eV. The partial widths, $\Gamma_\gamma(6.2)$ and $\Gamma_\gamma(4.7)$, for the emission of $\gamma$-rays to the states at 0.72 and 2.15 Mev respectively can be derived from the intensities of these $\gamma$-rays relative to the 5.2 Mev $\gamma$-ray. The values of $|M|^2$ obtained for the two transitions forbidden by the isotopic spin selection rules, together with the value of $|M|^2$ for the allowed transition, can be used to estimate the


\textsuperscript{17} This figure agreed well with that determined later (Appendix A) by a calibration using the known cross-section at the 993 kev resonance\textsuperscript{11}. 
isotopic spin admixture in the 6.89 Mev level, from the relation:

\[
\frac{\alpha_0^2(1)}{1 - \alpha_0^2(1)} = \frac{|M|^2}{|M|^2} \text{ allowed} - \frac{|M|^2}{|M|^2} \text{ forbidden.} \quad \cdots (3.4)
\]

This data is summarised in Table 3.1.

**TABLE 3.1.**

| To State in B10at | E \(_\gamma\) (Mev) | \(\gamma\text{-Intensity at 330kev}\) | Partial \(\Gamma_{\gamma}(\text{eV})\) | Weisskopf Unit \(\Gamma_w(\text{eV})\) | \(|M|^2\) | \(\alpha_0^2(1)\) | \(\alpha_0(1)\) |
|------------------|-----------------|-----------------------------------|-----------------|---------------------------------|------|---------|---------|
| 0.72             | 6.2             | 55                                | \(0.76 \pm 0.3\) | 120                             | 0.0063 | 24\%   | 0.49    |
| 1.74             | 5.2             | 100                               | \(1.4 \pm 0.5\) | 69                              | 0.020  | 22\%   | 0.47    |
| 2.15             | 4.7             | 22                                | \(0.30 \pm 0.1\) | 54                              | 0.0056 | 22\%   | 0.47    |

The mean of the two estimates of \(\alpha_0^2(1)\) gives a value of 23\% for the intensity of the \(T = 1\) admixture in the 6.89 Mev level.

This figure is in agreement with the 20\% estimated by Wilkinson and Clegg. An "impurity" of this size implies either an abnormally large matrix element \(H_{01}^0\) or, alternatively, the presence of a contaminating \(1^-, T = 1\) state very close to the 6.89 Mev level (see Eq. (3.2)).

The large proton reduced width of the 6.89 Mev state (\(g_P^2 = 0.3\) for s-wave protons) indicates that this state belongs predomini-
stantly to the configuration (1s^4lp^5)2s with the Be^9 ground state as unique parent. It is expected that there will be a group of four states in Be^{10} with this configuration, corresponding to the coupling of an s-wave nucleon to the 1s^4lp^5 partial configuration. These four states will be (1^-, T = 0), (1^-, T = 1), (2^-, T = 0) and (2^-, T = 1). In Be^{10}, a 1^-, T = 1 level at 5.96 Mev and a 2^-, T = 1 level at 6.26 Mev are known, together with a state of unknown spin and parity at 6.18 Mev. It seems likely that the 7.48 Mev level in Be^{10} is the T_z = 0 analogue of the 6.26 Mev level in Be^{10} and consequently it is to be expected that a 1^-, T = 1 level (the T_z = 0 analogue of the 5.96 Mev level in Be^{10}) will also be found in the region of 7 Mev excitation in Be^{10}.

Wilkinson and Clegg have suggested that large Coulomb matrix elements may occur between states of the configuration (1s^4lp^5)2s. Shell model calculations using an oscillator potential well give small values for these matrix elements, but much larger values (0.4 Mev) can be obtained if a well of finite depth is used. Even if we consider a value of H_{01}^C for the 6.89 Mev level as large as 0.5 Mev, it follows from Eq. (3.2) (assuming that the perturbation treatment is still applicable) that in order to explain the value

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20. F.C. Barker, Phil. Mag., 2, 286, (1957); and private communication (1958).
of \( a_0(l) \), the contaminating \( l^- \), \( T = 1 \) state must lie between 5.9 and 7.9 Mev in \( ^{10}B \). For smaller values of \( H \), these limits will also be correspondingly smaller. So far, however, this \( l^- \), \( T = 1 \) state has not been located.

The region of excitation in \( ^{10}B \) between 6.89 and 8 Mev has been carefully examined without revealing the presence of any other \( l^- \) level. Indications of possible states of unknown spin and parity at 7.01 and 7.20 Mev have been observed in the reactions \( ^{9}Be(p,d)^{8}Be \) and \( ^{9}Be(p,\alpha)Li^{6} \). These states are uncertain, however, and the absence of \( \gamma \)-rays to lower lying states makes it improbable that either of these is \( l^- \), \( T = 1 \). The results of the present \( ^{9}Be(p,\gamma)^{10}B \) experiment rule out the possibility of a \( l^- \), \( T = 1 \) state within \( \pm \) 100 kev of the 6.89 Mev level. It therefore seems likely that the level we seek lies below 6.89 Mev. It has not been observed in the \( ^{9}Be(p,\gamma)^{10}B \) reaction, but it could be one of four levels of unknown spin and parity between 6.1 and 6.8 Mev which have been established from the \( ^{9}Be(d,n)^{10}B \) reaction. It is to be expected that the \( l^- \), \( T = 1 \) state has an isotopic spin admixture (in this case of \( T = 0 \)) which is similar in magnitude to that found in the 6.89 Mev level. Consequently, if it were one of the levels between 6.1 and 6.8 Mev it is surprising that it has not been observed in the reaction \( Li^{6}(\alpha,\gamma)^{10}B \) which has been used to study levels in \( ^{10}B \) up
to 6.74 Mev excitation. The possibility of the contaminating state lying outside the range 6 to 8 Mev excitation cannot be excluded. However, the $1^-, T = 1$ level at 5.96 Mev in Be$^{10}$ must have an analogue somewhere in the level structure of Be$^{10}$ and one would expect it to be in the neighbourhood of 7 Mev excitation. The exact position of this level poses a problem which, so far, remains unsolved.

CHAPTER IV.

USE OF A 7.7 MEV CYCLOTRON TO STUDY RADIATIVE CAPTURE REACTIONS.

4.1 Introduction.

In Chapters IV and V, experiments are described in which the yields of high energy γ-rays from capture reactions are measured as functions of proton energies in the range of 4 to 7.7 Mev using as proton source, the 7.7 Mev 30-in. cyclotron recently completed in Canberra. This cyclotron is a high beam-current machine designed for use as injector for the 10 Bev proton synchrotron under construction in the laboratory. For the present purposes, it was modified to provide small extracted beams of high energy homogeneity and stability. The proton energy was reduced in small steps by inserting aluminium foils into the beam and capture γ-radiation from the target was detected by means of the large sodium iodide crystal spectrometer described in Chapter II.

There are many experimental difficulties associated with the use of a fixed-energy cyclotron for this type of experiment but until recently, variable-energy sources of protons in this energy range were not available. The main problem in performing the experiments lay in the high level of background radiation associated
Figure 4.1. General experimental arrangement.
with the cyclotron, which is notoriously a "dirty" machine. Several preliminary experiments were done to determine the best way of overcoming this, and in the sections which follow a description is given of the experimental arrangement which was finally adopted, together with an analysis of the background problem and details of measurements which were made to determine the specifications of the proton beam incident on the target.

4.2 Experimental Method.

The general experimental arrangement is shown in Figure 4.1. After leaving the ion source, the protons in their first, second and third orbits were required to pass through correctly located beam-defining slits of width 1, $\frac{1}{2}$ and $\frac{1}{2}$ mm. respectively. This ensured that only those protons leaving the ion source during a few electrical degrees about a particular dee voltage phase angle were accepted for acceleration. (This method is described in detail by Morton and Smith\textsuperscript{1}). Radial movement of the second slit served to restrict the circulating beam current to about 10μA, excess ions being stopped before acquiring sufficient energy to produce background radiation. An internal probe at 180° from the oscillator

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sampled the beam at the final orbit and the current to this was adjusted to 5 μA. The current to the probe was very sensitive to magnetic field fluctuations and, by stabilising it at 5 μA, machine conditions could be held constant over long periods of time. The remaining beam current of 5 μA was completely extracted and, after undergoing a momentum analysis in the fringing magnetic field, the extracted beam was limited to 1 μA by a system of narrow slits 2 mm. high and 2.5 mm. wide just outside the machine. These slits were backed with tantalum in order to reduce the neutron emission from them. The final 1 μA was then focussed onto the target by a quadrupole magnetic lens system composed of four sections.

Proton energy variation was achieved by inserting 1.64 mg.cm.\(^{-2}\) aluminium foils at the position indicated on Figure 4.1. These foils were mounted on two movable racks, each of which could be placed in any one of 7 different positions. Any number of foils between 0 and 6 could be inserted by suitable positioning of the first rack, and between 0 and 42 foils, in steps of 7 foils, by positioning of the second rack. Combining the positions of the two racks it was possible to increase in steps of 1 foil from 0 to 48 foils. By this means the proton energy could be reduced from 7.7 to 3.5 Mev in 48 steps; the steps were not equal in energy, but
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sampled the beam at the final orbit and the current to this was adjusted to 5 \( \mu \text{A} \). The current to the probe was very sensitive to magnetic field fluctuations and, by stabilising it at 5 \( \mu \text{A} \), machine conditions could be held constant over long periods of time. The remaining beam current of 5 \( \mu \text{A} \) was completely extracted and, after undergoing a momentum analysis in the fringing magnetic field, the extracted beam was limited to 1 \( \mu \text{A} \) by a system of narrow slits 2mm. high and 2.5mm. wide just outside the machine. These slits were backed with tantalum in order to reduce the neutron emission from them. The final 1 \( \mu \text{A} \) was then focussed onto the target by a quadrupole magnetic lens system composed of four sections.

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Figure 4.2. Side view of target chamber.
a single foil corresponded to 68 kev at 7.7 Mev, and to 125 kev at 3.5 Mev.

To minimise the background radiation produced in the neighbourhood of the crystal, it was undesirable to stop the beam in the target or target backing. Furthermore for good resolution in the yield curves, thin targets were essential and consequently the targets used in these experiments were either in the form of thin foils or were deposited on thin foil backings. When 3mg.cm. -2 gold foil was used as target backing the amount of background attributable to the gold was found to be negligible. The target chamber is shown in more detail in Figure 4.2. The targets were hung in the beam and stretched by a small weight suspended from them. The small amount of power (of the order of 1/10 Watt) dissipated in such thin targets made water cooling for them unnecessary. The beam leaving the target passed into a 2 ft. long tubular "beam"-catcher" lined with lead which stopped the protons at a sufficient distance from the detector so that background radiation effects were negligible. The beam current falling in the catcher was measured and integrated.

Gamma rays from the target were detected at 90° to the proton beam by the 4 in. long by 5 in. diameter NaI(Tl) crystal
seated inside its steel block. The crystal was mounted with its axis vertical and located 19\(\frac{1}{2}\) in. from the target centre. One inch diameter side collimation was adopted (§ 2.3). The steel block was surrounded on all sides, above and below, by a 6-in. thickness of lead shielding through which a 1-in. diameter collimating hole was drilled so that \(\gamma\)-rays from the target could pass along a diameter of the crystal. To reduce the neutron flux through the crystal, a 0.010 in. cadmium sheath was wrapped around the lead shield enclosing the steel block, and 8-in. thick layers of a cast mixture of 50\% paraffin wax and 50\% borax were placed on three sides of the assembly as shown in Figure 4.1. Gamma rays from the target had to be detected through the neutron shield, which transmitted 70\% of 20 Mev \(\gamma\)-rays.

4.3 The Background Problem.

The two main sources of background radiation in the range of \(\gamma\)-ray energies examined with the crystal were (a) the cosmic-ray \(\mu\)-mesons and (b) the proton beam from the cyclotron.

Traversals of the crystal by natural \(\mu\)-mesons produced a spectrum of pulses whose height depended on the pathlengths of the mesons in the crystal (§ 2.4(e)). In the energy range of interest
(15 to 25 Mev) the distribution of cosmic-ray pulses was smooth, showing a slight decrease towards higher energies. The rate of accumulation of these pulses was found to be constant so that a knowledge of the time taken to record a spectrum enabled the cosmic-ray contribution to be estimated and subsequently subtracted.

It was desirable to reduce the number of points in the experimental arrangement at which the proton beam was either stopped or slowed down since these points yielded neutrons through \((p,n)\) reactions and also \(\gamma\)-rays through \((p,\gamma),(p,p'\gamma),(n,\gamma)\) and \((n,n',\gamma)\) reactions. Wherever such places could not be avoided (e.g. at the 180° probe, the external slit, the stopping foils, the target and the beam catcher) the materials used were chosen, where possible, to keep background radiation to a minimum. Inside the concrete block-house several large blocks of paraffin were placed in such a way as to shield the crystal against neutrons arising from the 180° probe and from the external slits.

The quadrupole focussing magnets were of great assistance in solving the problem of background radiation. Early runs were made without these lenses, and since the divergence of the extracted beam was about \(\frac{1}{2}^\circ\), further slits were required just outside and just inside the concrete shielding wall to limit the size of the beam.
spot at the target. A much bigger internal beam had to be accelerated in order to get enough beam at the target and the additional slits, being close to the crystal were a troublesome source of background. The introduction of focussing magnets made these slits unnecessary and allowed a higher beam current at the target. This in turn meant that thinner targets could be used with a resulting improvement in the energy resolution of the experiment.

The biggest contribution to the background came from the radiative capture of neutrons in the crystal and its surroundings. The radiative capture cross-section for fast neutrons in iodine is high, being 29 mb at a neutron energy of 5.9 Mev\(^2\). The most energetic neutrons detected by the crystal were produced mainly in the target material itself (e.g. for 7.7 Mev protons, the reactions \(\text{Li}^7(p,n)\text{Be}^7\) and \(\text{B}^{11}(p,n)\text{C}^{11}\) produce neutrons with maximum energies of 4.5 and 4.0 Mev respectively\(^3\)). The Q-value for the reaction \(\text{I}^{127}(n,\gamma)\text{I}^{128}\) is 7.0 Mev\(^4\) so that capture of fast neutrons in the crystal produce \(\gamma\)-rays with energies extending up to about 12 Mev. In addition to observing such background pulses in the recorded \(\gamma\)-ray spectra,

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further evidence of radiative neutron capture in the crystal was found by running the cyclotron for a short time and then, with the beam off and the gain of the pulse amplifier turned up, examining the spectrum resulting from radioactivity induced in the crystal. The spectrum of 2 MeV $\beta$-particles from $^{128}$I with their characteristic half-life of 25 minutes was readily observed.

Background pulses in the crystal were sometimes observed as a steeply falling edge in the spectrum extending up to energies as high as 15 Mev (e.g. see Figure 5.1). From an examination of the position of this edge for differing beam currents, it was established that these pulses were due to the "pile-up" of lower-energy pulses. The occurrence of this "pile-up" did not appear to affect the resolution obtained for the capture $\gamma$-rays being studied, provided that the "pile-up" edge was kept below the region of interest in the spectrum.

The aluminium stopping foils contributed primarily through the reactions:

$$\text{Al}^{27}(p,n)\text{Si}^{27} \quad Q = -5.61 \text{ Mev.}$$

$$\text{Al}^{27}(p,\gamma)\text{Si}^{28} \quad Q = +11.6 \text{ Mev.}$$

$$\text{Al}^{27}(p,p')\text{Al}^{27}(\gamma)\text{Al}^{27}.$$  

At 7.7 Mev proton energy the first of these reactions produces
neutrons with a maximum energy of 1.8 Mev leading to neutron capture pulses in the crystal with maximum energies of 9 Mev. Furthermore this reaction has a threshold at a bombarding energy of 5.8 Mev which means that the foil system acts as an increasing source of neutrons as more foils are inserted into the beam until a proton energy of 5.8 Mev is reached after which insertion of further foils does not increase the neutron yield. The second and third reactions listed on the previous page, take place at effectively all proton energies and consequently the number of $\gamma$-rays from these sources increases until the beam is completely stopped in the foils. In addition, some of the $\gamma$-ray energies lie in the region of interest in the crystal spectrum, (since the maximum $\gamma$-ray energy from the $\text{Al}^{27}(p,\gamma)\text{Si}^{28}$ reaction is 19.0 Mev). To provide further shielding against these $\gamma$-rays, an extra 3-in. thickness of lead was interposed between the stopping foils and the crystal just outside the foil chamber (Figure 4.1). It was then found that, with the beam currents used in these experiments, the only significant background in the region 15 to 25 Mev was due to cosmic rays.

The main difficulty in connection with the background was to be able to keep the extracted beam current high enough so that the yield of proton capture $\gamma$-rays from the target was large
relative to the cosmic-ray background and, at the same time, to prevent background pulses from other sources from "piling-up" and obscuring the regions of interest in the γ-ray spectra. With the shielding arrangements shown in Figure 4.1, this difficulty was to a large extent overcome, and it was possible to record the spectra of γ-rays from (p,γ) reactions in the target with a counting efficiency of 1/7,100 for 20 Mev γ-rays. With beam currents of the order of 1 μA, the cosmic-ray background level was usually less than 10% of the peak intensity of the capture γ-rays of interest. At proton energies near 4 Mev the background "pile-up" pulses began to spread into the region being examined in the spectrum. This effect, when combined with the reduced energies of the proton capture γ-rays, set a lower limit of about 4 Mev proton-energy which could be usefully reached in these experiments.

4.4 Measurements On The Proton Beam.

4.4(a) Proton Energy.

Calibration of the scale of proton energy relative to the number of aluminium foils inserted into the beam depended on a precise knowledge of the maximum proton energy and of the thickness of the aluminium foils. From measurements of the magnetic field, R.F. voltage, and other specifications of the cyclotron, the
energy of the extracted beam was calculated to be 7.72 Mev. This agreed with the value of 7.70 ± .050 Mev obtained by measuring the range of the beam in aluminium and using the range-energy relations given by Aron et al., for protons in aluminium. Both of these values agreed with that obtained by the magnetic deflection experiments described in §4.4(c).

The aluminium foil used in the racks was selected for its uniformity in thickness. Many pieces of foil of differing sizes but cut from the same batch, were weighed and all gave the result that the thickness was within the limits of 1.637 ± 0.004 mg.cm.² This result was again confirmed when the experiments were completed and small pieces cut from foils on the rack were weighed. The proton energy was then derived as a function of number of foils using the data of Aron et al.

As well as introducing a source of unwanted background radiation, the method of varying the proton energy by the foils system had the unavoidable disadvantages that the beam was spread by scattering and the proton energy resolution worsened by straggling. These effects are described in the following sections.

Figure 4.3. A Dyeline print showing the beam cross-section with no stopping foils inserted. The dark lines are shadows cast by steel cross-wires.
4.4(b) Scattering by the Foils.

Preliminary tests had shown that, as a result of Rutherford scattering in the aluminium, the stopping foils had a slight defocussing effect on the beam. This effect became significant with the insertion of many foils and consequently to prevent the size of the beam spot from becoming too large at low proton energies the target chamber was designed with the foil racks only 5 in. up-beam from the target. This necessitated careful shielding of the crystal against radiation from the foils.

The magnitude of the scattering effect was measured experimentally by removing the beam catcher and fixing in its place a 5 mg.cm.\(^{-2}\) by 2 1/4 in. diameter aluminium window on the far end of the target chamber. After passing through this window the beam struck a 2 1/2 in. square piece of "Dyeline" photo-sensitive copying paper placed about 1/4 in. outside the window. After a short exposure to the beam and a brushover with developer, the paper showed the exact position and shape of the beam. A set of thin steel cross wires inserted between the window and the paper cast a shadow which defined the centre of the beam tube. This technique was also found to be a speedy aid in lining up the beam tube and focussing magnets. Figure 4.3 is a typical Dyeline print.
With the extracted beam properly focussed and with no stopping foils inserted, the cross-section of the beam in the region of the target was a rectangle 0.3 by 0.2 in. Scattering in the aluminium window and in the air gap between the window and paper did not blur the picture appreciably because of the short distances involved. When a few foils were inserted the edges of the rectangle blurred and with the introduction of more foils the beam produced on the Dyeline paper, a diffuse circular spot whose diameter increased with increasing numbers of foils. These measurements showed that the maximum beam spot diameter at the target was 0.5 in. at a proton energy of 3.9 Mev. This was considered satisfactory as the hole in the lead collimator for the crystal was of 1 in. diameter, which meant that the geometry, and hence the resolution for the crystal, was essentially constant over the range of proton energy used.

4.4(c). Energy Resolution of the Proton Beam.

The energies of the protons incident on the target were spread about a mean value as a result of (a) the inherent energy spread of the extracted beam and (b) straggling in the aluminium stopping foils, arising from the statistical nature of the slowing-down process. The energy spread due to both of these effects was measured experimentally by the following method.
Figure 4.4. Experimental arrangement used for measurements with the 20° deflecting magnet.
The beam catcher was removed from the end of the target chamber and was replaced by a 20° deflecting magnet. The target was replaced by a strip of lead sheet 0.010 in. thick with a vertical slit 0.050 in. wide in its centre. The beam passing through the slit was deflected by the magnet and a picture of the beam cross-section was recorded by the Dyeline paper technique described in the previous section. The experimental arrangement is shown schematically in Figure 4.4. The distances from the slit to the magnet and from the magnet to the Dyeline paper were arranged so that the beam was focussed on the paper to form an image of the slit.

The magnet current was adjusted and stabilized so that, without stopping foils in the beam, the centre of the slit image coincided with the centre of the cross-wires at the window. The magnet had been calibrated previously using a flux-meter and, from a determination of the magnetic field required to deflect the beam through 20°, the maximum proton energy was again calculated to be 7.7 Mev.

The breadth of this image was then measured and found to be larger than that expected for an image produced by a mono-energetic beam of protons. Assuming that all of the spread in the image was the result of a spread in proton energies it was possible to place an
Figure 4.5 (a). Crude straggling curve determined from measurement of image spread as a function of the number of stopping foils inserted into the beam. The circles are the experimental points and the full line is the theoretical value given by Bohr for the straggling.

(b). The theoretical straggling curve of Bohr together with the 6 values determined experimentally.
upper limit on the energy spread of the 7.7 Mev extracted beam at 35 kev. This agreed well with the \( \frac{1}{2} \% \) energy spread predicted from a knowledge of the positions and widths of the beam-defining slits inside the cyclotron.

The effect of straggling was examined by introducing stopping foils one at a time and measuring the spread in the image as a function of the displacement of the image centre relative to the cross-wires. This result could then be translated into a plot of total energy spread versus mean beam energy. These measurements were made up to six foil thicknesses only, since scattering in the foils necessitated exposure times at lower energies which were prohibitively long. When the initial spread of 35 kev was subtracted from these measurements, a crude estimate of the straggling curve was obtained (Figure 4.5(a)).

The straggling for charged particles which are losing energy in traversing a stopping medium has been calculated by Livingstone and Bethe, who find that the rate of change with distance of the mean square fluctuation in the energy, \( E \), is given by

\[
\frac{\mathrm{d}}{\mathrm{d}x} \left[ \langle E^2 \rangle_{\text{av}} - \langle E_{\text{av}} \rangle^2 \right] = 4n_e e^2 Z^2 N \left[ Z' + \sum_n k_n l_n e_n / m^2 \log(2m^2/l_n) \right],
\]

where \( Z' \) is the total number of "effective" electrons defined as the number of electrons in the atom, \( Z \), excluding those in the inner

\[6. \text{ M.S. Livingstone and H.A. Bethe, Rev. Mod. Phys., 9, Section 16, (1937).}\]
shells for which $I_n > 2mv^2$. The sum goes over the shells which are not excluded. $Z_n$ is the number of electrons in the nth shell, $I_n$ their average excitation energy, and the $k_n$ certain constants between $2/3$ and $4/3$. $x$ is the amount of stopping material traversed, $N$ is the number of atoms per cm$^3$ in the material, and $ze$, $m$ and $v$ are the charge, mass and velocity of the particles being slowed down. For high energies the sum over $n$ may be neglected, and with $Z'$ replaced by $Z$, the formula given by Bohr on the basis of classical mechanics is obtained:

$$\frac{d}{dx} \left[ (E^2)_{av} - (E_{av})^2 \right] = 4\pi e^2 Z^2 N x \quad \ldots \ldots \ldots (4.2)$$

This gives for the straggling:

$$\sqrt{\Delta E^2} = \left( 4\pi e^2 Z^2 N x \right)^{1/2} \quad \ldots \ldots \ldots (4.3)$$

In the case of 7.7 Mev protons in aluminium, the Bohr formula is expected to be correct to within 10% down to an energy of 4 Mev - the energy range covered in the present experiments. At still lower energies the straggling would become appreciably larger than Bohr's value due to the growth of the summation term in Eq.(4.1).

The theoretical straggling curve of Bohr is shown on

Figure 4.5(b) together with the straggling measurements resulting from magnetic analysis. The agreement is reasonable up to 6 foil thicknesses and the Bohr curve was assumed to hold for proton energies down to 4 Mev.


The energy stability of the extracted proton beam was checked experimentally by two methods. In the first, changes in deflection of the full-energy beam by the 20° magnet were observed as a function of time, and in the second, the yield of γ-rays from the $^{27}\text{Al}(p,\gamma)^{28}\text{Si}$ reaction was monitored at a proton energy of 7.50 Mev (3 foils). There is a steeply rising portion of the $^{27}\text{Al}(p,\gamma)^{28}\text{Si}$ yield curve centred about 7.50 Mev where the γ-ray yield is doubled for an increase in proton energy of only 150 kev (see Figure 6.5). Consequently any fluctuation in the proton energy produces a marked change in γ-ray yield. Both of these measurements indicated that after an initial "warm-up" period of half an hour the extracted beam was stable to within +0 and -30 kev over many hours of operation.
5.1 Introduction.

It was pointed out in §1.7, that it is of considerable interest to determine to what extent the giant resonance can be regarded as composed of a clustering of individual nuclear levels. Some experiments which have been performed using photo-nuclear reactions, report fine structure in the \((\gamma,p)\) and \((\gamma,n)\) cross-sections in the giant resonance region for a few light nuclei. However, these experiments suffer from certain disadvantages inherent in the use of bremsstrahlung radiation as the \(\gamma\)-ray source (§1.7) and for this reason it was considered of value to perform the experiments, described in this chapter, using the inverse, radiative capture reactions.

Bethe has demonstrated that the cross-sections for two mutually inverse $(\gamma,p)$ and $(p,\gamma)$ reactions are related by detailed balancing, provided that the same proton and $\gamma$-ray wavelengths (in the centre-of-mass system) are referred to in both processes. Generally the only cases of practical importance will be where the residual nuclei are left in their ground states. The cross-sections for the reactions $^{12}_6((\gamma,p)B^{11}$ and $^{11}_6(p,\gamma)^{12}$ have been related in this way by Mann and Titterton. They made a measurement of $\sigma_{\gamma p}$ at $17.6$ Mev $\gamma$-ray energy relative to the known cross-section for $^{12}_6(\gamma,3\alpha)$. Their value of $\sigma_{\gamma p} = 1.19 \pm 0.21$ mb. is in good agreement with the value of $1.09 \pm 0.16$ mb. calculated from the inverse reaction by detailed balance. Wright et al have also related their $^{14}_6(\gamma,p)^{13}$ results by detailed balance to those known for $^{13}_6(p,\gamma)^{14}$.

Although the method of radiative proton capture has the advantages that the energy of excitation is known precisely and that good resolution is obtainable with a mono-energetic beam of protons, it has not been exploited previously in an investigation of the giant

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7. Detailed balance is discussed in more detail in Appendix B.
resonance. The reason is that, until recently, variable-energy sources of protons in the range of 5 to 15 Mev required to cover the region were unavailable. In the experiments reported in the following sections, the Canberra 7.7 Mev cyclotron was used as a source of protons of good energy homogeneity. The energy was varied using stopping foils as described in the previous chapter.

Two reactions have sufficiently high Q-values to allow the giant resonance region to be covered by (p,γ) studies in this energy range; they are Li$^7(p,γ)$Be$^8 + 17.2$ Mev and B$^11(p,γ)$C$^{12} + 16.0$ Mev. The maximum proton energy of 7.7 Mev enabled nuclear excitations of 24.0 Mev and 23.0 Mev to be reached in Be$^8$ and C$^{12}$ respectively. The ground states of both Be$^8$ and C$^{12}$ are 0$^+$ and since most of the emitted γ-radiation will be electric dipole (§1.7) it is expected that the levels which contribute most significantly to the giant resonance will be 1$^-$ (and T = 1, if isotopic spin is still a good quantum number at such high excitation). The low level density associated with such light nuclei together with the fact that only 1$^-$ levels are expected to be significant suggests that Be$^8$ and C$^{12}$ are most suitable nuclei in which to search for fine structure in the giant resonance. The approximate position of the giant resonance in C$^{12}$ is known$^{10}$, but Be$^8$ is an unstable nucleus and cannot be

investigated by photo-disintegration techniques; it was hoped to
establish the position and properties of the giant resonance for it.

5.2 Preparation of Targets.

5.2(a) Lithium.

Attempts were made to prepare a target by evaporation
of lithium metal onto a 3 mg/cm$^2$ gold foil backing. However, it was
found that the lithium rapidly diffused into the gold resulting in
an effectively thicker and non-uniform target.

The lithium targets which were finally used were
thin, uniform foils of natural lithium metal which were made initially
by rolling a small piece of the metal, bathed continuously in par-
affin oil, between two sheets of brass. The thinnest obtainable by
this technique were about 0.003 in., this limitation being set by a
tendency for the lithium to work-harden and to adhere to the brass.
Later targets were made as thin as 0.001 in. by rolling the metal in
a kerosène bath between two thin polythene sheets. These foils were
transferred to a bath of dry xylene where they were clamped in a
target holder. The whole target assembly was then rapidly transferred
to the target chamber which was immediately pumped out. By observing
these precautions it was possible to get a target into position in
the cyclotron without any visible oxidation, hydrolysis or other
chemical action taking place on the lithium.
5.2(b) **Boron.**

Early runs were made using B₂O₃ targets evaporated onto a 3 mg.cm⁻² gold foil backing. However, in order to improve the yield from the target without losing resolution due to increased target thickness, it was desirable to make a pure boron target. After some experimentation carbon crucibles were made which were satisfactory for the evaporation in vacuo of pure amorphous boron (melting point 2,300°C). The greatest difficulty in making these crucibles which were carved out of graphite blocks, lay in achieving a proper gradation in the electrical resistance so that excessive heat was not conducted away to the water-cooled copper electrodes at each end. The targets used for most of the runs were layed on the gold foil as 1.25 in. diameter discs of thickness about 0.8 mg.cm⁻². Chemical tests on the material which was evaporated onto a copper plate holding the foil in the evaporator indicated that only about 50% of the target was boron, the remainder being carbon. Comparison between yields from the boron and B₂O₃ targets confirmed this figure for the fraction of boron. The presence of carbon in the target was attributed to some of the molten boron combining with carbon in the crucible to form boron carbide which melts at 2,350°C.

The possible presence of small amounts of contaminant
Figure 5.1. $^{11}B(p,\gamma)^{12}C$ spectrum at $E_p = 7.7$ Mev. The dashed line represents the background level at this energy.
elements in the lithium and boron targets did not raise any difficulties in these experiments. This was because all of the likely contaminants had low Q-values for (p,γ) reactions and could not possibly give rise to γ-rays in the high energy region under study. The same remarks apply to the isotopes Li⁶ and B¹⁰ which were present in the targets in their natural isotopic abundances of 7.5% and 18.8% respectively.

5.3 Measurements.

The experimental arrangement was as described in Chapter IV. The proton energy was varied between 7.7 Mev and 4 Mev by insertion of stopping foils. At each proton energy, for a given number of micro-coulombs of integrated beam current, a pulse-height spectrum of γ-rays was recorded on a kick-sorter which was biased to avoid counting the large number of pulses below 15 Mev. Each spectrum took about 15 minutes to record.

In Figure 5.1 a typical B⁷⁷ (p,γ)C¹² spectrum at a bombarding energy of 7.7 Mev is shown with the background level marked in. Transitions to the ground state and the 4.4 Mev first excited state of C¹² are clearly resolved at 23.0 and 18.6 Mev γ-ray energies. The steeply ascending counting rate in channels below 10 is due to
Figure 5.2. $^{7}\text{Li}(p,\gamma)^{8}\text{Be}$ spectrum at $E_p = 7.7$ Mev.
the background effects discussed in §4.3. As the bombarding proton energy was reduced, both peaks in the spectrum moved downwards in energy. This effect, combined with the slight upward movement of the "pile-up" edge as more stopping foils were inserted, tended to obscure the low-energy peak (first excited state transition) at lower proton energies. The peak due to the ground-state transition was resolved at all energies and by counting the number of pulses under this peak on each spectrum recorded, the yield of ground-state γ-rays was determined as a function of bombarding proton energy.

To obtain good statistics several long runs were carried out. In each run the cyclotron (after an initial warm-up period of half an hour) was operated continuously for approximately 20 hours. During this time, spectra were recorded starting at full energy, slowly working down to 4 Mev and then returning to full energy. In all, a total of seventeen such runs was made with targets of boric oxide and boron of various thicknesses, but the final result comprises data taken in six runs on a 40 kev thick elemental boron target and is shown in Figure 5.3(a).

Similar experiments were carried out with lithium targets. In this case the pulse-height spectrum at maximum energy was as shown in Figure 5.2 and the main transition was to the broad 2.9 Mev level of Be⁶; transitions to the ground state were less intense and showed up only as a small peak on the side of the broad one. Yield
Figure 5.3 (a). Yield curve for $^{11}\text{B}(p,\gamma)^{12}\text{C}$ ground-state $\gamma$-rays.

(b). Two determinations of the $^{7}\text{Li}(p,\gamma)^{8}\text{Be}$ yield curve.
curves were determined as before, but the number of runs over the available proton energy range was limited to six as the cyclotron had to be dismantled and moved into its position as injector for the proton synchrotron. The yield curves were derived by plotting the total number of pulses under the peaks corresponding to transitions to both the 2.9 Mev and the ground states of Be$^8$.

In preliminary experiments with lithium targets it was found that the neutron production from them was appreciably greater than that from the boron targets. This was attributed mainly to a higher cross-section for the $(p,n)$ reaction in lithium, and partly to $(\alpha,n)$ reactions resulting from the $(p,\alpha)$ reactions occurring in the target. As a result of this enhanced background, it was necessary to use smaller beam currents to avoid distortion of the spectrum by "pile-up". The yield curves regarded to be of the most significance (i.e. obtained at the lowest beam currents) are the two shown in Figure 5.3(b).

5.4 Calibration of the NaI(Tl) Crystal.

At the completion of experiments on the cyclotron, the target chamber, crystal and its shielding were transferred in toto to the 1 Mev Cockcroft-Walton accelerator in the High Tension laboratory.
The resolution and calibration of the crystal for high-energy $\gamma$-rays were then determined under geometrical conditions identical with those previously used on the cyclotron. A thick (350 kev) layer of lithium metal evaporated onto a copper backing was used as the target. At the $441$ kev resonance the thick target yield from the $\text{Li}^7(p,\gamma)\text{Be}^8$ reaction was taken as $1.90 \times 10^{-8}$ $\gamma$-rays per proton. The resolution of the crystal at a $\gamma$-ray energy of $17.6$ Mev was estimated at $15\%$.

It was essential to prevent chemical deterioration of the lithium target during transfer from the evaporator to the target chamber. To avoid this, after deposition of the target, the evaporator and the target chamber were filled with dry argon and the target was transferred under a jet of argon. The target chamber was then evacuated. This technique permitted a bright mirror-like deposit of lithium metal to be maintained throughout.

Using this calibration the peak cross-section for the emission of ground-state $\gamma$-rays in the $\text{B}^{11}(p,\gamma)\text{C}^{12}$ reaction was found to be:

$$\frac{d\sigma}{d\Omega(90^\circ)} = (2.7 \pm 0.5) \times 10^{-28} \text{cm}^2/\text{4}\text{sterad.}$$

Similarly, the maximum value of the cross-section for the $\text{Li}^7(p,\gamma)\text{Be}^8$

reaction was found to be:

\[
d \sigma / d \Omega (90^\circ) = (1.4 \pm 0.4) \times 10^{-28} \text{ cm}^2 / 4\pi \text{ sterad.}
\]

where both transitions to the ground state and the 2.9 Mev state are included. Most of the uncertainty in these cross-sections is due to the difficulty in determining the fraction of the total spectrum which is represented by counts under the peaks.

5.5 Results.

5.5(a) \( B^{11}(p,\gamma)C^{12} \).

The yield curve for ground-state \( \gamma \)-radiation emitted in the \( B^{11}(p,\gamma)C^{12} \) reaction, shown in Figure 5.3(a), reveals the position of the maximum of the giant resonance at a bombarding proton energy of \( (7.2 \pm 0.1) \text{ Mev.} \) The curve also exhibits small peaks at proton energies of 5.95 and 6.75 Mev (corresponding to excitations in \( C^{12} \) of 21.4 and 22.1 Mev respectively), but it is otherwise smooth within the resolution of the experiment (indicated on Figure 5.3(a) for various proton energies). No significance is attached to the low point at 7.3 Mev proton energy because it did not occur reproducibly and it is within the statistical errors of the experiment.

It was noticed that as the proton energy was varied, the ratio of the intensity of the ground-state \( \gamma \)-ray transition to
that of the first excited state transition did not appear to change appreciably. The transition to the 4.4 Mev level in $^{12}C$ merged into the "pile-up edge" for bombarding proton energies below about 5.5 Mev. However, throughout the range 5.5 to 7.7 Mev the ratio of ground-state to first excited state transitions was essentially constant and was roughly 2:1.

5.5(b) $^{7}\text{Li}(p,\gamma)^{8}\text{Be}$.

The two yield curves obtained at the lowest target currents are shown in Figure 5.3(b). One run with a target 145 kev thick and a beam current of 0.35 $\mu$A gave a yield curve which rises from lower proton energies and flattens off at about 6.0 Mev proton bombarding energy. The other run, made with a thicker target (280 kev) and a much lower beam current of 0.05 $\mu$A, gave a yield curve which displays a broad resonance with a maximum at a bombarding energy of 5.8 Mev. The experimental evidence suggested that during the first run (at 0.35 $\mu$A), the pile-up of background radiation may have been distorting the spectra, and for this reason it is believed that the second yield curve taken at the lower proton beam current should be accepted. The experimental points recorded in this run are shown in Figure 5.5. There is no significant evidence for structure in the resonance, but the statistical accuracy and the experimental resolution (which ranged between 300 and 350 kev) are not sufficient to
rule out the possibility of small peaks of the type observed in the
$^{11}_B (p,\gamma)^{12}_C$ reaction.

The ratio of ground-state to first excited state transitions
appeared to be essentially constant over the range of proton energy
used. This ratio was approximately 1:5.

5.5(c) Transitions to Higher Excited States of $^{8}_Be$.

Evidence for and against the existence of excited
states in $^{8}_Be$ between 3 and 10 Mev excitation, has been produced by
many authors (see Appendix C). Inall and Boyle report the reaction
$^{7}_Li (p,\gamma)^{8}_Be (\alpha)^4_He$ proceeding via levels in $^{8}_Be$ at 4.1, 5.3 and 7.5
Mev. In the $^{7}_Li (p,\gamma)^{8}_Be$ experiment using the cyclotron a search was
made for possible $\gamma$-rays proceeding through these states. None were
found, although many spectra were recorded with good statistical
accuracy at several bombarding proton energies. It was estimated
that if 10% of the total $\gamma$-ray transitions had involved states at
4.1 or 5.3 Mev with widths less than $\frac{1}{2}$ Mev, these would have been
observed. Lower energy $\gamma$-rays going to a state at 7.5 Mev probably
would not have been resolved from the background.

5.6 Comparison with Other Results.

5.6(a) $^{11}\text{B} (p,\gamma)^{12}\text{C}$. 

Application of the principle of detailed balancing (see Appendix B) to the maximum cross-section for emission of ground-state $\gamma$-rays in the $^{11}\text{B} (p,\gamma)^{12}\text{C}$ reaction ($\S$ 5.4) yields a value of

$$\frac{d\sigma}{d\Omega}(90^\circ) = (24 \pm 4) \text{ mb} / 4\pi \text{ sterad}.$$ 

for the peak cross-section for the emission of protons to the ground state of $^{11}\text{B}$ in the inverse $^{12}\text{C} (\gamma,p)^{11}\text{B}$ reaction, the peak being located at a $\gamma$-ray energy of $E_{\text{max}} = (22.55 \pm 0.1) \text{ Mev}$. 

Halpern and Mann $^{10}$ in photo-disintegration experiments found a maximum cross-section of $(34 \pm 8) \text{ mb} / 4\pi \text{ sterad}$ at the same angle. Their value, however, is dependent on large corrections for target absorption which are sensitive to the energy distribution of the emitted protons. This distribution has since been measured more accurately by Cohen et al. $^5$, who find a peak cross-section at $90^\circ$ of $21\text{ mb} / 4\pi \text{ sterad}$ for the emission of ground-state protons.

The previous best estimates of the position ($E_{\text{max}}$) and width ($\Gamma$) of the giant resonance in $^{12}\text{C} (\gamma,p)^{11}\text{B}$ are given by Halpern and Mann $^{10}$ ($E_{\text{max}} = 21.5 \text{ Mev}, \Gamma = 1.7 \text{ Mev}$) and by Haslam...
Figure 5.4. Present $^{11}$B$(p,γ)^{12}$ yield curve compared with the results of some previous experiments.
et al. \( (E_{\text{max}} = 21.2 \text{ Mev}, \Gamma = 2 \text{ Mev}) \). The present results indicate that these values for \( E_{\text{max}} \) are about 1 Mev too low. If we assume that the giant resonance is symmetrical about the peak, then the yield curve of Figure 5.3(a) has a width of 3.6 Mev in proton bombarding energy. This corresponds to a width in the \( ^{12}\text{C}(\gamma,p)^{11}\text{B} \) reaction of \( \Gamma = 3.3 \text{ Mev} \), a value significantly larger than those previously estimated. Within the limitations of their statistical accuracy, the results of Cohen et al. show fair agreement with the present results as far as the position and the width of the giant resonance are concerned. (see Figure 5.4)

The yield curve of Figure 5.3(a) is shown again in Figure 5.4 where it is compared with the results of some other experiments. The results of Cohen et al. were obtained by irradiating a thin poly­thene foil with 24 Mev bremsstrahlung and using photographic plates to examine the photo-protons emitted. The resultant energy spectrum of photo-protons was corrected for the escape of protons from the emulsion and multiplied by a bremsstrahlung correction factor. On the assumption that the residual \( ^{11}\text{B} \) nucleus is always left in its ground state, the \( ^{12}\text{C}(\gamma,p)^{11}\text{B} \) cross-section curve shown in Figure 5.4 was obtained. Their curve gives some indication of fine structure in the

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giant resonance region, but the statistical accuracy is not good and
the possibility of proton emission to excited states of $^{11}$ complicates the interpretation of the curve below 22 Mev photon energy.
Between 22 and 24 Mev the curve becomes uncertain because of poorer statistics combined with the need for larger bremsstrahlung corrections. However, the histogram does show a fairly definite group at 21.5 Mev which could be identified with the peak observed in this experiment at 21.4 Mev.

The $^{11}$B($p,n$)$^{11}$C yield curve of Blaser et al. is also shown on Figure 5 and although one would not necessarily expect a close agreement between the $(p,n)$ and the $(p,\gamma)$ results, their curve does show two peaks at excitations in $^{12}$C of 21.3 and 21.8 Mev.

Moreover, the present $^{11}$B($p,\gamma$)$^{12}$C results are strikingly similar to those reported by Barber at the Washington photonuclear conference (May, 1958) relating to the reaction $^{12}$C$(e,e'p)$B. The proton energy distribution measured in this experiment showed a peak at 6 Mev, which corresponds to the giant resonance peak at 22.5 Mev excitation in $^{12}$C. With the resolution available (about 170 kev) fine structure was not observed, but below the main peak alternate

plateaux and steeply rising portions occurred. A plateau at 21.4 Mev was observed with good reproducibility, and additional plateaux at 21.8 and 20.9 Mev were suggested by the data.

On the other hand, the present results are not compatible with the most recent data presented by Katz et al. for the "breaks" in the $^{12}$C ($\gamma$,n)$^{11}$ activation curve. In the region of $^{12}$C excitation between 21 and 23 Mev the Saskatoon work suggests that most of the absorption occurs into sharply defined narrow states (with widths less than 50 kev) at 21.08, 21.22, 21.58, 22.02 and 22.88 Mev, the last three being very strong. The vertical lines in Figure 5.4 represent the relative integrated cross-sections under the levels corresponding to these "breaks". Over this region the resolution of the present experiment varies from 60 to 85 kev and is good enough to have indicated such structure. Further evidence against this interpretation of the "breaks" data has been reported by Wolff et al. who have measured the total absorption cross-section in $^{12}$C from 20.3 to 20.8 Mev using monochromatic $\gamma$-rays from the $^3T^3$p($\gamma$,He$^4$) reaction. They found that no resonances of width less than 70 kev are present with more than 1/2 to 1/5 of the integrated cross-sections reported by Katz; but state that their results are not inconsistent with the

Figure 5.5. The most reliable Li$^7$(p,$\gamma$)Be$^8$ yield curve data. The full line represents the data of Bair et al., normalised at $E_p = 4.8$ Mev.
broader structure reported by Cohen et al. 5.

The broken curve in Figure 5.4 is the 90° yield curve for ground state γ-rays from the B (p,γ)C reaction obtained by Bair et al. 17 These workers did not measure absolute cross-sections, and their yield curve is shown normalised to the present result at a proton energy of 4.5 Mev.

5.6(b) Li(7(p,γ)Be 8.

Figure 5.5 shows the fit of the present data to the yield curve determined by Bair et al. 18 for γ-rays of energy greater than 10 Mev. The two results are normalised at a proton energy of 4.8 Mev.

Be is an unstable nucleus and therefore the present results cannot be compared with any photo-disintegration data for Be(γ,p)Li 7. However, the yield curve of Figure 5.5 indicates that the inverse reaction, if it could be observed would have a broad giant resonance maximum at about $E_{\text{max}} = 22.3$ Mev and an approximate width of $\Gamma = 5$ Mev.

Wilkinson 19 has suggested that the smoothly increasing

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trend of the $^{18}\text{Li}^{7}(p,\gamma)^{8}\text{Be}$ excitation function at proton energies above about 2.5 MeV, was mainly due to a process inverse to the giant resonance in $^{8}\text{Be}(\gamma,p)^{7}\text{Li}$. The peak cross-section determined in this experiment (§ 5.4) for the emission of $\gamma$-rays to both the ground and first excited states of $^{8}\text{Be}$ in the $^{18}\text{Li}^{7}(p,\gamma)^{8}\text{Be}$ reaction is:

$$\frac{d\sigma}{d\Omega}(90^\circ) = (1.4 \pm 0.4) \times 10^{-28} \text{ cm}^2 / \text{4\pi sterad.},$$

which is about half the value calculated by Wilkinson on the basis of a postulated giant resonance in $^{8}\text{Be}$ with parameters $E_{\text{max}} = 22.2$ MeV, $\Gamma = 2.3$ MeV and an integrated cross-section of 30 MeV - mb. If the experimental width of $\Gamma = 5$ MeV were used in the calculation together with the same integrated cross-section of 30 MeV - mb (this figure is not likely to be greatly in error) then the predicted peak cross-section agrees well with that determined in this experiment.

5.7 Discussion.

The results of these experiments show that there is certainly structure in the giant resonance observed in the reaction $^{11}\text{B}(p,\gamma)^{12}\text{C}$ (and therefore in the reaction $^{12}(\gamma,p)^{11}\text{B}$). The structure appears to be of a broad nature and quite unlike that suggested by the "breaks" experiments. The discrepancy between the two results warrants further investigation since, from charge symmetry consider-
ations, one would expect that any fine structure observed in the giant resonance in the $^{12}\gamma p^{11}$ reaction would be similar to that observed in the $^{12}\gamma n^{11}$ reaction.

The present yield curve (Figure 5.3(a)) suggests that if only a few $l^-$ states (with $T = 1$ if isotopic spin is still a good quantum number at these excitations) contribute to the giant resonance in $^{12}C$, then these states must be broad, having widths of the order of $\frac{1}{2}$ Mev or more. In this region of high excitation one would expect many broad levels, and it is possible that, in this experiment, we pick out some of the $l^-, T = 1$ states, of which there may be only three or four with large matrix elements for the emission of electric dipole radiation to the ground state (§ 1.7). An alternative possibility is that there is a large number of such $l^-, T = 1$ levels participating in the giant resonance. If the separations of these levels were less than their widths, or (if the levels are narrow) less than the resolution obtained in this experiment, then the observed yield curve would represent the envelope of such levels.

The fact that a single broad $l^-, T = 1$ level ($\Gamma = 1.3$ Mev) is known to exist at 17.2 Mev excitation in $^{12}C$, is perhaps an

indication that only a few broad levels participate. However, to settle the question definitely, it would be desirable to repeat these experiments with an energy resolution which is an order of magnitude better than that obtained in this work with the cyclotron.

The γ-ray branching ratios for the two reactions $\text{Li}^7(p,\gamma)\text{Be}^8$ and $\text{B}^{11}(p,\gamma)\text{C}^{12}$ differ markedly. In the $\text{Li}^7(p,\gamma)\text{Be}^8$ reaction γ-ray transitions favour the first excited state by a factor of 5 over the ground state and in the $\text{B}^{11}(p,\gamma)\text{C}^{12}$ reaction they favour the ground state by a factor of 2 over the first excited state. The reason for the difference is not easy to see on the basis of simple ideas on γ-ray widths. In both cases we are dealing with p-shell nuclei; both $\text{Li}^7$ and $\text{B}^{11}$ have 3/2⁻ ground states and both $\text{Be}^8$ and $\text{C}^{12}$ are α-particle nuclei with 0⁺ ground states and 2⁺ first excited states. Other things being equal, one would expect, transitions from a 1⁻, $T = 1$ level to the first excited state to be a factor of 5 stronger than those to the ground state because of the higher statistical weight attached to the 2⁺ level. However, because of the increased transition probability associated with the higher energy of the ground-state γ-rays (see Eq.1.2) this factor would be reduced by $(E_1/E_0)^3$ where $E_0$ and $E_1$ are the energies of the transitions to the two states, assuming dipole radiation. For $\text{Be}^8$ and $\text{C}^{12}$ this
still leaves a factor of about 3 in favour of the first excited state transition.

Capture of s-wave protons by Li$^7$ or B$^{11}$ can lead only to 1$^-$ or 2$^-$ states in Be$^8$ and C$^{12}$. Only the 1$^-$, $T = 1$ states can decay by El radiation to the 0$^+$ ground state, but both 1$^-$ and 2$^-$ levels with $T = 1$ decay to the 2$^+$ first excited state. The capture of d-wave protons, whose contribution will be smaller because of the centrifugal barrier, can lead only to 1$^-$, 2$^-$, 3$^-$ and 4$^-$ states. Of these, only the 1$^-$, $T = 1$ states can decay by El emission to the ground state, but the 1$^-$, 2$^-$ and 3$^-$ states with $T = 1$ can decay to the 2$^+$ state. Such considerations indicate that a factor of 5 in favour of the transition to the first excited state, as observed in the Li$^7(p,\gamma)$Be$^8$ reaction, is understandable in these terms, and that for some reason, El $\gamma$-ray transitions to the 4.4 Mev state of C$^{12}$ are inhibited by a factor of about 10.

It is somewhat surprising that the $\gamma$-ray branching ratios do not appear to change with proton energy. One would expect that in going from a 1$^-$, $T = 1$ to a 2$^-$ or 3$^-$, $T = 1$ intermediate state, say, these ratios would be affected quite significantly. The fact that no appreciable changes in relative intensities are seen suggests that either (a) there are only a few intermediate states involved
each having the same spin, parity and isotopic spin quantum numbers; probably $1^-$, $T = 1$, since there is always an appreciable fraction of ground-state radiation present, or (b) the intermediate states are so dense that each spectrum recorded is an average over many levels and this average does not change its character appreciably with proton energy. In order to distinguish between these possibilities it would be desirable to perform the experiments again with much better energy resolution as mentioned earlier, and if possible also measure the angular distributions of the γ-rays emitted.
CHAPTER VI.

REGIONS OF HIGH NUCLEAR EXCITATION IN OTHER ALPHA-PARTICLE NUCLEI.

6.1 Introduction.

The experimental methods described in the previous chapter for the reactions Li\(^7\)\((p,\gamma)\)Be\(^8\) and B\(^{11}\)\((p,\gamma)\)C\(^{12}\) were also employed to study regions of nuclear excitation lying between 14 and 20 Mev in Ne\(^{20}\), Mg\(^{24}\), Si\(^{28}\), and S\(^{32}\).

The reactions used were:

\[
\begin{align*}
\text{F}^{19}\((p,\gamma)\)\text{Ne}^{20} & \quad + \quad 12.9 \text{ Mev} \\
\text{Na}^{23}\((p,\gamma)\)\text{Mg}^{24} & \quad + \quad 11.7 \text{ Mev} \\
\text{Al}^{27}\((p,\gamma)\)\text{Si}^{28} & \quad + \quad 11.6 \text{ Mev} \\
\text{P}^{31}\((p,\gamma)\)\text{S}^{32} & \quad + \quad 8.8 \text{ Mev}.
\end{align*}
\]

In the resultant α-particle nuclei, the high binding energy of a proton, together with the 7.7 Mev available from the Canberra cyclotron, enabled maximum nuclear excitations of 20.2, 19.1, 19.0, and 16.3 Mev respectively, to be reached.

Previous to the present experiments none of these \((p,\gamma)\) reactions had been studied at bombarding energies above 2.5 Mev\(^{1,2}\).

The $^{19}\text{F} + p$ reaction has been examined up to 5.3 Mev by Willard et al.\(^3\) but their yield curves include all $\gamma$-rays with energies above 3 Mev, which they considered to arise mainly from the $^{19}\text{F}(p,\alpha,\gamma)^{16}\text{O}$ reaction.

It was the aim of the experiments to be described in the following sections to measure the yield curves for the high energy $\gamma$-rays from the reactions listed on the previous page using protons with energies between 4 and 7.7 Mev.

6.2 Experimental Method.

Except for targets the general experimental arrangement was identical with that used previously. The same shielding against background radiation was again found adequate in these experiments. The target elements were all mono-isotopic, and again small traces of contaminant elements in the target material were found to be no problem.

Yield curves were determined by recording $\gamma$-ray spectra at each proton energy for a given number of protons, and by measuring the number of counts under the peaks corresponding to $\gamma$-ray transitions to the first excited state and to the ground state, of the final nucleus. Each spectrum took about 20 minutes to record

Figure 6.1. Pulse-height spectra produced by $\gamma$-rays from (a) the $^{19}_p(p,\gamma)Ne^{20}$ reaction at $E_p = 7.63$ Mev, and (b) the $^{23}_p(p,\gamma)Mg^{24}$ reaction at $E_p = 6.40$ Mev. The arrows indicate transitions to the ground and first excited states of $Ne^{20}$ and $Mg^{24}$. The dotted curves are the background counts with the targets withdrawn from the beam.
with beam currents of 0.4 μA for sodium, aluminium and phosphorus, and below 0.08 μA for fluorine.

6.3 Results.

6.3(a) $^{19}_p(p,\gamma)^{20}$.

The fluorine targets were poly-tetra-fluorethylene ("Fluon") foils of thickness 0.98 mg.cm$^{-2}$. These foils were found to suffer progressive damage produced by the passage of the proton beam; they became brittle and, to prevent their crumbling in the region of beam impact, the beam current was kept below 0.08 μA. After each spectrum was recorded the foil was raised so that the succeeding measurement was made on an unused section of the foil. The targets were sufficiently uniform and the running conditions constant enough, for this method to give reproducible results.

A γ-ray spectrum taken at a proton energy of 7.63 Mev is shown in Figure 6.1(a). Throughout the entire range of bombarding proton energies the γ-ray transitions were observed to go mainly to the first excited state of $^{20}$Ne at 1.6 Mev above the ground state. The smaller fraction going to the ground state was observed as a "shoulder" on the pulse height spectrum, and the ratio of intensities of the two transitions did not appear to change significantly with proton energy.
Figure 6.2. Yield curve for $^{19}_{\text{F}}(p,\gamma)\text{Ne}^{20}$. 
In determining the yield curve the counts under the combined ground-state plus first excited state peaks were recorded as a function of proton energy. Two determinations of this yield curve were made using different target foils; they agreed well and the average is shown in Figure 6.2. Pronounced broad resonances occur with maxima at proton energies of 5.5, 6.85 and probably 7.65 Mev. The target thickness and the total energy spread of the proton beam are shown marked on Figure 6.2 at various proton energies. Using the same cross-section calibration as described in the previous chapter, the $90^\circ$ cross-section for the sum of the two transitions at 7.7 Mev proton energy was determined to be:

$$\frac{d\sigma}{d\Omega}(90^\circ) = (8.5 \pm 2) \times 10^{-28} \text{ cm}^2/\text{4}\pi \text{ sterad.}$$

6.3(b) $\text{Na}^{23}(p,\gamma)\text{Mg}^{24}$.

Thin targets of sodium were made by rolling the metal in a kerosene bath using the technique developed for lithium described previously in §5.2(a). It was found much easier to roll sodium targets than lithium because sodium did not work-harden nearly so easily. Sodium foils 0.002 in. thick were used in two determinations of the yield curve over the full extent of the energy range, but two additional runs were made from 7.0 to 7.7 Mev to confirm the existence of a small peak at a proton energy of 7.45 Mev. Gamma-ray
Figure 6.3. Yield curve for Na$^{23}(p,\gamma)$Mg$^{24}$. 
transitions were predominantly to the first excited state of the residual nucleus — in this case the 1.4 Mev level in Mg$^{24}$. Only a small fraction of the transitions occurred to the ground state of Mg$^{24}$ as illustrated in the γ-ray spectrum of Figure 6.1(b) recorded at a proton energy of 6.40 Mev.

The yield curve of Figure 6.3 represents the average of all runs taking the total counts under the combined ground state plus first excited state peaks in the pulse height spectra. In addition to the small peak at 7.45 Mev, a broad resonance occurs with a maximum at a proton energy of 5.95 Mev, and there is an indication of a small peak at 4.6 Mev. However the rise in the background count together with the reduced energy and intensity of the γ-rays made it difficult to determine a yield curve with precision in the region of low proton energy. The cross-section at 7.7 Mev for the sum of both transitions was:

$$\frac{d\sigma}{d\Omega}(90^\circ) = (1.4 \pm 0.4) \times 10^{-28} \text{cm}^2/\mu\text{radian} \text{ sterad.}$$

6.3(c) Al$^{27}(p,\gamma)$Si$^{28}$.

An aluminium foil 1.64 mg.cm$^{-2}$ thick was used as the target for this experiment and a γ-ray spectrum recorded at 6.87 Mev proton energy is shown in Figure 6.4(a). Transitions at all proton
Figure 6.4. Pulse-height spectra produced by γ-rays from (a) the $^{27}\text{Al}(p,\gamma)^{28}\text{Si}$ reaction at $E_p = 6.87$ Mev, and (b) the $^{31}\text{P}(p,\gamma)^{32}\text{S}$ reaction at $E_p = 7.50$ Mev. The arrows indicate transitions to the ground and first excited states of $^{28}\text{Si}$ and $^{32}\text{S}$. The dotted curves are the background counts with the targets withdrawn from the beam.
Figure 6.5. Yield curve for $\text{Al}^{27}(p,\gamma)\text{Si}^{28}$. 
energies were mainly to the first excited state at 1.8 Mev in Si$^{28}$, with a smaller fraction going to the ground state. Figure 6.5 shows the yield curve derived using the sum of both transitions and represents the average of three separate determinations of the yield curve. Peaks in the curve appear at proton energies of 6.63, 6.87 and 7.20 Mev, and possibly at 6.1 and 7.7 Mev. The cross-section at 7.7 Mev for the sum of the two transitions was:

$$\frac{d\sigma}{dn}(90^\circ) = (1.9 \pm 0.5) \times 10^{-28} \text{ cm}^2/4\pi \text{ sterad}.$$ 

6.3(d) $P_3^1(p,\gamma)S_3^{22}$.

To prepare targets, some amorphous red phosphorus was first shaken into suspension in ethyl alcohol and passed through a fine filter paper many times. A few crystals of sugar were then dissolved in a portion of the resulting colloidal suspension. A drop of this suspension was placed on a clean 3 mg.cm$^{-2}$ gold foil backing and allowed to spread into a disc of about 1 in. diameter. After the alcohol had evaporated there remained a very thin uniform layer of red phosphorus which adhered well to the gold backing.

Three different targets, with thicknesses measured by weighing to lie between 6 and 8 mg.cm$^{-2}$, were used in three separate determinations of the yield curve. Figure 6.4(b) shows a pulse-
Figure 6.6. Yield curve for $\text{P}^{31}(p,\gamma)\text{S}^{32}$ ground-state $\gamma$-rays.
height spectrum recorded at a proton energy of 7.50 Mev. Throughout the entire range of bombarding proton energies, transitions to the ground state and first excited state of $^3\text{He}$ at 2.24 Mev occurred with about equal probability; the yield curve derived for the ground state transition only, is shown in Figure 6.6. Peaks appear at incident proton energies of 7.50, 7.15, 6.65 and probably 5.6 Mev. The cross-section at 7.7 Mev for ground state transitions only was:

$$\frac{d\sigma}{d\Omega}(90^\circ) = (0.5 \pm 0.2) \times 10^{-28} \text{ cm}^2/4\pi \text{ sterad.}$$

To complete the systematics for alpha-particle nuclei up to and including $^3\text{He}$, an attempt was made to observe the reactions:

$$^3\text{He}(p,\gamma)^4\text{He} + 19.8 \text{ Mev}$$

and $$^1\text{H}(p,\gamma)^1\text{H} + 12.1 \text{ Mev}$$

The tritium target used was adsorbed to a thickness of 100 \(\mu\)g/cm\(^2\) on a thin layer of zirconium which was deposited on a 0.010 in. platinum backing, and the $^1\text{H}$ target was deposited to a thickness of 30 \(\mu\)g/cm\(^2\) on a 0.010 in. tantalum backing. These targets were clamped to a water-cooled copper holder which held them in the proton beam.

In the \(\gamma\)-ray spectra recorded, no peaks could be discerned above the cosmic-ray background level at energies correspond-
ing to transitions to the ground states of He$^4$ and O$^{16}$, the reason being that the targets were not thick enough. Runs at several proton energies between 5.0 and 7.7 Mev were made with beams of up to 2 µA. Unfortunately higher beam currents could not be used because of the danger of over-heating and outgassing the targets. It was possible, however, to place upper limits of $4 \times 10^{-28}$ cm$^2$/4π sterad and $1.2 \times 10^{-27}$ cm$^2$/4π sterad, on the cross-sections, $d\sigma/d\Omega(90^\circ)$, for emission of ground-state γ-rays in the reactions H$^3(p,\gamma)$He$^4$ and N$^{15}(p,\gamma)$O$^{16}$, respectively.

6.4 Discussion of Results.

There is little previous information available on energy levels in the regions of nuclear excitation studied in these experiments. In the cases of F$^{19}(p,\gamma)$Ne$^{20}$, Na$^{23}(p,\gamma)$Mg$^{24}$ and Al$^{27}(p,\gamma)$Si$^{28}$, the levels can be compared with those found from previous measurements carried out with protons of energies up to 6.6 Mev on the reactions F$^{19}(p,n)$Ne$^{19}$, Na$^{23}(p,n)$Mg$^{23}$ and Al$^{27}(p,n)$Si$^{27}$.

The F$^{19}(p,n)$Ne$^{19}$ reaction has been studied by Blaser et al.$^4$ using a 6.6 Mev cyclotron and examining the activity induced in stacked foils, and by Marion et al.$^5$ using a 6 Mev Van de Graaff generator.

and a BF$_3$ counter. Both (p,n) yield curves show a broad peak at a
proton energy of about 5.4 Mev which could correspond with the broad
peak found at 5.50 Mev in the present $F^{19}(p,\gamma)Ne^{20}$ yield curve.
Blaser et al. also determined a yield curve for the $Na^{23}(p,n)Mg^{23}$ re-
action, and their broad peak at 6.00 Mev proton energy shows agree-
ment with that found in the $Na^{23}(p,\gamma)Mg^{24}$ yield curve at 5.95 Mev.

The peak at 6.1 Mev suggested in the $Al^{27}(p,\gamma)Si^{28}$ result
may be that observed by Blaser et al. at 6.17 Mev in their study of
the $Al^{27}(p,n)Si^{27}$ reaction.

Again, it could not be ascertained within the limits of
the experimental resolution, whether the peaks in the present (p,\gamma)
yield curves are composed of a small number of individual broad levels
or a large number of closely spaced levels (c.f. § 5.7). The regions
of excitation covered in these nuclei are just below the giant reson-
ance. However, the cross-sections are still large suggesting that
most of the \(\gamma\)-radiation emitted is electric dipole in character. As
in the cases of $Li^7(p,\gamma)Be^8$ and $B^{11}(p,\gamma)C^{12}$, the \(\gamma\)-ray branching ratios
were again found to be essentially constant with proton energy.

Table 6.1 summarises the values estimated for the relative intensities, $N_1/N_0$, of first excited state to ground state trans-
itions in each of the nuclei studied with the cyclotron.
TABLE 6.1.

<table>
<thead>
<tr>
<th>Target Nucleus and spin</th>
<th>Li$^7$</th>
<th>B$^{11}$</th>
<th>F$^{19}$</th>
<th>Na$^{23}$</th>
<th>Al$^{27}$</th>
<th>P$^{31}$</th>
</tr>
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<tbody>
<tr>
<td>Final Nucleus</td>
<td>Be$^8$</td>
<td>C$^{12}$</td>
<td>Ne$^{20}$</td>
<td>Mg$^{24}$</td>
<td>Si$^{28}$</td>
<td>S$^{32}$</td>
</tr>
<tr>
<td>Max $E_{\gamma_1}$ (Mev)</td>
<td>21.1</td>
<td>18.6</td>
<td>18.6</td>
<td>17.7</td>
<td>17.2</td>
<td>14.1</td>
</tr>
<tr>
<td>Max $E_{\gamma_0}$ (Mev)</td>
<td>24.0</td>
<td>23.0</td>
<td>20.2</td>
<td>19.1</td>
<td>19.0</td>
<td>16.3</td>
</tr>
<tr>
<td>$N_1/N_0$</td>
<td>5</td>
<td>1/2</td>
<td>2</td>
<td>4</td>
<td>2</td>
<td>1</td>
</tr>
</tbody>
</table>

The values of $E_{\gamma_1}$ and $E_{\gamma_0}$ are the energies of the $\gamma$-rays to the first excited state and ground state respectively measured at the maximum proton energy of 7.7 Mev.

Since each of the four nuclei, F$^{19}$, Na$^{23}$, Al$^{27}$ and P$^{31}$ have ground states whose parity is even, the reactions studied with these nuclei will involve the capture mainly of p-wave protons. The structure in the yield curves for these nuclei appears to be more pronounced and narrower in character than that for the two p-shell nuclei studied. This may result in part from the narrower proton widths associated with p-wave capture (as a result of the centrifugal barrier) compared with the mainly s-wave capture by Li$^7$ and B$^{11}$. By the same token, one would expect the structure to become more pronounced with increasing atomic number, because of the increasing Coulomb barrier. This suggestion also receives some support from the observed yield curves.
The experiments performed with the cyclotron have shown surprisingly that there is considerable structure in the (p,γ) yield curves in these regions of high nuclear excitation. The type of structure varies with different nuclei, from the small subsidiary peaks observed in the B^{11}(p,γ)C^{12} reaction to the large fluctuations in yield observed in some of the other reactions. However, to determine in more detail what types of level are involved and whether many levels or just a few are contributing to these yield curves, these experiments should be repeated with a greatly improved energy resolution and if possible the angular distributions of the γ-rays should be measured.

The recent development of the tandem-style Van de Graaff generator as a continuously variable, mono-energetic proton source should make such measurements possible in the near future. The relatively small amount of background radiation associated with these accelerators will permit much thinner targets to be used, and this, together with the small spread in proton energy, should allow the determination of yield curves with good energy resolution over the whole range of the giant resonance in many nuclei. To obtain the ultimate resolution with a tandem generator, it will be necessary to
use targets of the order of 1 kev thick. Since the maximum beam current available with these machines so far is only about 5 μA, it seems likely that a high resolution experiment will be hampered by the low yield of capture γ-rays. If these γ-rays are to be detected with a large NaI(Tl) crystal then the cosmic ray background could prove to be a limiting factor in the experiment. This problem could possibly be solved by surrounding the crystal with a tank containing a liquid scintillator which could be used in anti-coincidence with the crystal.
APPENDIX A.

THE GROUND-STATE GAMMA RAY IN THE Be$_9^9$(p,$\gamma$)B$_{10}^{10}$ REACTION.

It was shown in Chapter III that the $\gamma$-rays emitted to the ground state of B$_{10}^{10}$ in the reaction Be$_9^9$(p,$\gamma$)B$_{10}^{10}$ are not resonant at the 330 kev resonance. However, assuming the value of $12 \pm 4 \mu$b.
given by Carlson and Nelson for the maximum cross-section for emission of 5.2 Mev $\gamma$-rays at this resonance, the data shown in Figure 3.3 indicate a value of $3 \pm 1 \mu$b. for the cross-section at 330 kev for emission of ground-state $\gamma$-rays. This is roughly five times too high to be explained as arising from the low-energy tail of the 993 kev resonance, which is known to emit $\gamma$-rays mainly to the ground state of B$_{10}^{10}$. To investigate this discrepancy further, a yield curve of ground-state $\gamma$-rays was measured in the range of proton energies between 300 and 1060 kev.

Gamma rays were detected at $0^\circ$ to the proton beam and 1 in. diameter side collimation was used. The beryllium target was 16.5 $\mu$g.cm$^{-2}$ thick, which corresponds to a stopping power of 6.2 kev at 330 kev proton energy. The excitation function obtained with

Figure A.1 Excitation function for ground-state γ-rays from the reaction $^{9}\text{Be}(p,\gamma)^{10}\text{B}$.
this target is shown in Figure A.1. The level parameters given by Hornyak and Coor were used to calibrate the cross-section scale and to derive a curve (shown in the figure) which represents a single-level Breit-Wigner formula, allowing for the variation of partial widths with proton energy.

The results show that while the Breit-Wigner curve gives a good fit to the experimental points in the region of the resonance peak at 993 keV, at lower proton energies the fit is not good. The cross-section at 330 keV was again found to be about 3 µb., and there appears to be a significant non-resonant contribution in this region of proton energy. If we assume that this contribution is due to a direct reaction or to the tail of a distant resonance, then it might be expected that its excitation function will follow a curve of the form $P_{1}^{\frac{1}{2}}/E$, where $P_{1}$ is the barrier penetrability for protons of orbital angular momentum 1 and with bombarding energy $E$. The full line and the dashed line in Figure A.1 are the best fits to the experimental points which can be obtained by adding non-resonant s-wave and p-wave excitation functions of this form to the Breit-Wigner curve. The factor $P_{1}^{\frac{1}{2}}$ has been taken from the data of Christy and Latter for s- and p-wave protons incident on Be. The best fit

appears to be obtained with an s-wave contribution although the spread in the data is not small enough to exclude the possibility of p-wave. One would expect s-waves to give the main contribution because of the absence of a centrifugal barrier and because they lead to the emission of electric dipole radiation to the ground state of $^{10}\text{B}$. 

The non-resonant radiation in this case may imply a direct reaction involving a one-stage process without formation of a compound nucleus. Such a process has been postulated to explain successfully the non-resonant $\gamma$-ray transitions which are known to occur in the $^7\text{Li}(p,\gamma)^8\text{Be}$ reaction at similar proton energies. If direct $\gamma$-ray transitions are present in the $^9\text{Be}(p,\gamma)^{10}\text{B}$ case, then it is possible that they occur to the excited states of $^{10}\text{B}$ as well as to the ground state. When added (possibly with interference) to the resonant transitions to these states, these direct effects may explain the fact that the transitions to the 0.72, 1.74 and 2.15 Mev levels do not bear quite the same ratios to one another over the width of the 330 kev resonance (Figure 3.6).

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APPENDIX B.

THE RELATIONSHIP BETWEEN THE CROSS-SECTIONS FOR THE INVERSE REACTIONS $B^{11}(p,\gamma)C^{12}$ and $C^{12}(\gamma,p)B^{11}$.

The cross-sections for two mutually inverse $(p,\gamma)$ and $(\gamma,p)$ processes are related by detailed balancing as was pointed out by Bethe in 1937, and it was just these inverse reactions $B^{11}(p,\gamma)C^{12}$ and $C^{12}(\gamma,p)B^{11}$ which he visualised as an example.

If, in the centre-of-mass system, $\sigma_{\gamma p}$ is the cross-section for the reaction $(Z + 1)^{A+1}(\gamma,p)Z^{A}$, where the nucleus $Z^{A}$ is left in its ground state, and $\sigma_{\gamma' p}$ is the cross-section for the inverse process, $Z^{A}(p,\gamma)(Z + 1)^{A+1}$, where the nucleus $(Z + 1)^{A+1}$ is left in its ground state, and if the same $\gamma$-ray and proton wavelengths, $\lambda_{\gamma}$ and $\lambda_{p}$ (in the C.O.M. system), are referred to in both cases, then the following relation holds:

$$\sigma_{\gamma p} = \sigma_{\gamma' p} \left( \frac{2j + 1}{2j' + 1} \right)^2 \left( \frac{\lambda_{\gamma}}{\lambda_{p}} \right)^2, \quad \ldots \ldots \ldots (B.1)$$

where $j'$ is the spin of the nucleus $(Z + 1)^{A+1}$; $j$ that of the nucleus $Z^{A}$; $s$ that of the proton ($\frac{1}{2}$); and $(2s' + 1)$ is the effective statistical weight for unpolarised electromagnetic radiation $(2s' + 1 = 2$)

since there are two directions of polarisation). Eq. (B.1) also holds for differential cross-sections provided that they refer to the same angle of orientation.

The maximum cross-section for the emission of ground-state $\gamma$-rays in the reaction $^{11}\text{B}(p,\gamma)^{12}\text{C}$ is $d\sigma_{p\gamma}/d\Omega(90^\circ) = (2.7 \pm 0.5) \times 10^{-28}$ cm$^2$/4$\pi$ sterad. (§ 5.4) and it occurs at a bombarding proton energy of 7.2 Mev ($\bar{\chi}_p = 1.85 \times 10^{-13}$ cm.) corresponding to a $\gamma$-ray energy of 22.55 Mev ($\bar{\chi}_\gamma = 8.73 \times 10^{-13}$ cm.). Inserting these values into Eq. (B.1) and setting $j^\prime = 0$ for $^{12}\text{C}$ and $j = 3/2$ for $^{11}\text{B}$, we find that the implied maximum cross-section for the emission of ground-state protons in the inverse reaction, $^{12}\text{C}(\gamma,p)^{11}\text{B}$, also measured at 90$^\circ$, is $d\sigma_{\gamma\gamma}/d\Omega(90^\circ) = (24 \pm 4)$ mb. /4$\pi$ sterad.
APPENDIX C.

THE LOW-LYING LEVEL STRUCTURE OF Be$^8$.

There has been much disagreement in recent years concerning the level structure in Be$^8$ in the region between 0 and 14 MeV excitation. Several levels have been reported by various observers, but so far, the only ones which seem to be present with any degree of certainty are the narrow 0$^+$ ground state ($\Gamma \approx 5$ eV), the broad 2$^+$ level at 2.9 MeV ($\Gamma \approx 1.2$ MeV) and a broad 4$^+$ level at 11.6 MeV ($\Gamma \approx 6.7$ MeV). These levels correspond well with the only three states expected in this region of excitation on the basis of either a simple shell model or an α-particle model of Be$^8$.

Nevertheless the evidence for other even states in Be$^8$ at 4.1, 5.3 and 7.5 MeV is considerable. Most of this evidence comes from studies of the reactions Li$^7$(d,n)Be$^8$, B$^{11}$(γ,t)Be$^8$, B$^{10}$(γ,d)Be$^8$, Li$^7$(p,γ)Be$^8$ (α)He$^4$ and from photo-disintegration experiments using γ-rays from the 441 kev resonance in the Li$^7$(p,γ)Be$^8$ reaction. Other evidence for some of these levels has also been

found in $\alpha - \alpha$ scattering experiments⁵ and in the $^9\text{Be}(p,d)^8\text{Be}$ reaction.

On the other hand, the results of some other experiments⁷⁻¹⁶ using such reactions as $^7\text{Li}^7(d,n)^8\text{Be}$, $^7\text{Li}^7(p,\gamma)^8\text{Be}^*\alpha^4\text{He}$, $\alpha - \alpha$ scattering, $^8\text{B}^8\text{B}^8(p,\alpha)^7\text{Be}$, $^8\text{B}^8\text{B}^8(d,\alpha)^8\text{Be}$, $^8\text{Li}^8(\beta^-)^8\text{Be}$, $^8\text{Be}^{(\beta^+)^8\text{Be}}$, $^6\text{Li}^6(\text{He}^3,p)^8\text{Be}$ and $^8\text{Be}^{(d,t)^8\text{Be}}$, do not show evidence for the additional levels.

The conflicting results of different workers appear to be due mainly to limited experimental resolution and to the difficulty

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7. 5.5(c) of this Thesis.
in obtaining sufficient statistical accuracy. Any states in this region, apart from the ground and 2.9 Mev states, are involved only weakly (if at all) in most of the reactions listed above. Furthermore, it is likely that any such levels (being even and highly unstable to α-particle emission) are broad, which would make them difficult to observe experimentally.

In an attempt to shed further light on the existence of these states, an experiment is at present in progress in this laboratory. The reaction being studied is Li$^7$(p,γ)Be$^8$ (α)He$^4$ at the 441 kev resonance and although the work is as yet incomplete, it is considered worthwhile to report briefly on the progress so far.

This reaction has been examined previously by Inall and Boyle, who measured the integrated spectrum of α-particles in coincidence with γ-rays. The α-particle ranges were measured with absorbing foils and the γ-rays were detected with a NaI(Tl) crystal 1 cm. thick by 2 cm. square. After differentiation their spectrum showed evidence for levels with widths less than $\frac{1}{2}$ Mev in Be$^8$ at 4.09, 5.31 and 7.51 Mev. The corresponding γ-rays occur with intensities of 2%, 2% and $\frac{1}{2}$% respectively of the total intensity. La Vier et al. have measured the α-particle spectrum using magnetic analysis (but without γ-ray

No evidence for the extra levels was discovered although minimum observable intensities were less than the intensities found by Inall and Boyle. More recently Meyer and Staub using this reaction, have examined the region of excitation in Be between 0 and 4 Mev using magnetic analysis of the α-particles in coincidence with γ-rays and the region between 4 and 9 Mev using magnetic analysis only. They found no levels between 4 and 9 Mev present with intensities of more than 2% of the total number of excited Be nuclei.

It is felt that the region of excitation in Be between 4 and 10 Mev warrants further study with the Li⁷(p,γ)Be⁸(α)He⁴ reaction. The technique of Inall and Boyle suffers from the disadvantage that it produces an integrated spectrum in which small statistical fluctuations can produce a relatively large effect in the differential spectrum. The other methods which measure the differential spectrum without γ-ray coincidences are hampered by the recoil broadening of the emitted α-particles (approx. ± 0.2 Mev). This effect is removed if the direction of emission of the γ-ray is defined by a coincidence requirement.

In the present experiment using this reaction the differential spectrum of α-particles is examined in coincidence with γ-radiation. In order to obtain a high counting efficiency for the γ-rays from such weakly excited levels, the large NaI(Tl) crystal de-
scribed in Chapter II is used as the $\gamma$-ray detector. The $\alpha$-particles are detected with a CsI crystal which has to be thin ($\sim 0.005$ in.) to avoid large numbers of background pulses due to the passage of high energy electrons. A fast coincidence unit of the type described by Garwin has been built having a resolving time of 100 μsec.

The first experiments were done bombarding a thin lithium target with 450kev protons and with a 1.64 mg.cm$^{-2}$ aluminium foil interposed between the target and the $\alpha$-particle detector which was situated at 90° to the proton beam. The foil was necessary to screen the photo-tube from light emitted from the target and also to prevent scattered protons from reaching the CsI crystal. The use of the foil however, meant that only $\alpha$-particles with energy greater than 2.2 Mev could reach the crystal. Alpha particles which penetrate the foil reach the crystal with reduced energy and this, coupled with the fact that the scintillation response of CsI is reduced for low-energy $\alpha$-particles, means that poorer resolution is obtained. The reduction in the $\alpha$-energies introduced by the foil also prevents the peak due to $\alpha$-particles from the 2.9 Mev level from being observed in the pulse-height spectrum. There is another disadvantage in that the intense 8.8 Mev $\alpha$-particles from the Li$^7$(p,$\alpha$)He$^4$ reaction are detected.

Figure C.1 Plan view of the target chamber and α-particle detector used to study the reaction $\text{Li}^7(p,\gamma)\text{Be}^{8}\text{(α)He}^4$. The calculated trajectories of α-particles of various energies are shown. Scale: $3/8'' = 1''$. 
by the crystal and a portion of the large number of these gives rise to background pulses in the region of interest between 0 and 5 Mev α-particle energy.

An alternative method is therefore being tried at present in which no stopping foil is required. The experimental arrangement is shown in Figure C.1. The α-particles are detected at 90° to the proton beam through a 1 mm. slit. The particles passing through the slit are influenced by a magnetic field of 5.4 kilo-gauss so that only those α-particles with energies between 1 and 5 Mev reach the \( \frac{1}{2} \times 1 \times 0.005 \text{ in. CsI crystal detector.} \) Scattered protons from the target follow the same trajectory as \( \frac{1}{2} \) Mev α-particles (shown in Figure C.1) and do not reach the crystal. Similarly the 8.8 Mev α-particles are not bent sufficiently to reach the crystal. The geometry is arranged so that when the slit is correctly positioned, light from the target cannot reach the photo-tube directly and the inside of the analysing chamber is painted black to stop reflected light from reaching the phototube. The perspex window is present to allow alignment onto the target of the γ-ray detector, which is at 90° to the beam and opposite the α-particle detector; this window is covered with a thin light-tight foil when the experiment is in progress.

The magnetic field is produced by six permanent magnets
Figure C.2  Spectrum of particles recorded using the apparatus shown in Figure C.1. The target was natural lithium metal.
of the magnetron type. To obtain the required magnetic field strength the gap width has been reduced from $1\frac{1}{2}$ in. to $5/8$ in. and the area of the field has been halved by the introduction of soft-iron pole-pieces. To protect the photo-tube from the stray field of these magnets, a 7 in. long by $1\frac{1}{2}$ in. diameter perspex light pipe was introduced and a soft-iron shield surrounds the front end of the photo-tube. The CsI crystal is cemented to the light pipe with "Araldite" and is prepared by water polishing down to a thickness of 0.005 in. The resolution obtained is found to depend critically on the finish given to the surface of the crystal and for this reason the final polishing is done with glycerine with a trace of water added to it.

Figure C.2 shows an ungated spectrum recorded with this arrangement. The target was natural lithium metal and was 60 kev thick at the bombarding proton energy of 450 kev. The slit was deliberately adjusted to allow a small fraction of the 8.3 Mev $\alpha$-particles to strike the crystal. (This gives a useful energy calibration point). The two peaks at lower energies are due to the $\text{Li}^6(p,\alpha)\text{He}^3$ reaction. Although $\text{Li}^6$ has an isotopic abundance of only 7.5%, the $(p,\alpha)$ reaction is very intense. The $\alpha$-particle and the $\text{He}^3$ are emitted in this process with nearly the same momentum, and since their charges are the same, they follow the same trajectory in the magnetic field (i.e. that of 1.85 Mev $\alpha$-particles). The increase in counts below
about \( \frac{1}{2} \) Mev in the spectrum is due to doubly scattered protons and to electrons passing through the crystal.

The \((p,\alpha)\) reactions occurring in Li\(^6\) and Li\(^7\) are not resonant at 441 kev so that the optimum target thickness to use in this experiment is 12 kev, the width of the 441 kev resonance in the Li\(^7\)(p,\(\gamma\))Be\(^8\) reaction. Preliminary spectra have been recorded in coincidence with \(\gamma\)-rays, using natural lithium targets 10 kev thick at 441 kev proton energy. The \(\alpha\)-particles from the broad 2.9 Mev level can be observed clearly in these spectra provided that the beam current is kept at a low value (less than \(\frac{1}{2} \) \(\mu\)A) so that the random coincidences due to the products of the Li\(^6\)(p,\(\alpha\))He\(^3\) reaction are not too numerous. However, to obtain statistical accuracy for regions of higher excitation in Be\(^8\), beam currents of 10 \(\mu\)A or more are required. With beam currents of this order and a coincidence resolving time of 100 \(\mu\) sec., it is found that most of the pulses in the coincidence \(\alpha\)-spectra corresponding to excitations in Be\(^8\) of up to about 6 Mev, are due to random coincidences with the products of the Li\(^6\)(p,\(\alpha\))He\(^3\) reaction.

These random coincidences may be reduced by shortening the resolving time. However, there is a danger then of discriminating against the lower-energy \(\alpha\)-particles because of their longer flight
times. The best method of eliminating the effect of the $\text{Li}^6(p,\alpha)\text{He}^3$ reaction is to use a separated $\text{Li}^7$ target. A technique has been developed similar to that suggested by Ajzenberg \(^{20}\) for the electrodeposition of $\text{Li}^7$ metal from a solution of the chloride in dry pyridine. Thin uniform targets of easily controlled thickness can now be prepared by this method and it is hoped to obtain significant information using these targets in the near future.

19. The flight time of a 1 Mev $\alpha$-particle from target to crystal in the apparatus of Figure C.1 is 34 μ sec.