PHOTOPROTON REACTIONS IN MEDIUM-WEIGHT NUCLEI

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ABSTRACT

The angular distributions, energy spectra and integrated yields of photoprotons were measured for nine medium-weight nuclei with atomic number between 28 and 50 in order to test the predictions of the Independent Particle Model of nuclear photodisintegration.

Thin metallic foils of natural isotopic abundance were irradiated with a bremsstrahlung spectrum of maximum energy 22 MeV. The photoprotons emitted by the targets were detected at angles of 45, 60, 90, 120 and 145 degrees with respect to the direction of the photon beam by means of silicon surface barrier detectors. By using pulse amplitude discrimination, pulse heights from the intense electron background were kept below the level of a four MeV proton for most of the target elements. The detection system was calibrated for energy with the alpha particles from a Ra\(^{226}\) source. Photoproton energy spectra were thus obtained at each of the five angles of observation for each of the nine target elements.

The results of the measurements are discussed in terms of the "resonance direct" interaction theory of Wilkinson. It is found that some extensions to this model are required to explain the experimental data.
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CHAPTER I

INTRODUCTION

In the study of nuclear structure much information has been obtained at high excitation energies by means of the nuclear photoeffect. One interesting aspect of the nuclear photoeffect is the study of the energy and angular distributions of the photonucleons. This thesis describes in some detail the experimental results on the yields and energy and angular distributions of photoprotons emitted by nine target elements between the closed proton shells at \( Z = 28 \) and \( Z = 50 \).

1.1 The Nuclear Photoeffect

The nuclear photoeffect or photodisintegration is the complete absorption by an atomic nucleus of an energetic photon whose energy, \( \hbar \omega \), is greater than the binding energy of a constituent nucleon or nucleon cluster and the subsequent de-excitation of the compound system by the emission of a particle or particle cluster. The photonuclear reaction can thus be thought of as a two-stage process (i) the formation of the compound system by photon absorption and (ii) the decay of the compound system into the products of the reaction.

The cross section \( \sigma(\omega) \) for the nuclear absorption of a photon of energy, \( \hbar \omega \), was first studied as a function of the photon energy by Baldwin and Klaiber(5) who found a strongly peaked curve now known as the giant resonance. Since that time many giant resonances have been delineated(73,76) and their systematics have been studied in terms of the following four parameters (i) the resonance energy, \( E_m \), (ii) the maximum cross section, \( \sigma_m \),
The resonance energy, or the photon energy at which the cross section is a maximum, is approximately constant at about 20 million electron volts (MeV) for $A \leq 40$, and decreases slowly with increasing mass number for $A > 40$ according to the relation (100),

$$E_m = 80 A^{-1/3} \text{ MeV.}$$

The maximum cross section varies approximately as $A^{3/2}$ (76); however, the value of the exponent of $A$ may not be very reliable due to the large errors inherent in the determination of the absolute cross sections. The width of the giant resonance varies between 4 and 10 MeV having its smallest values for nuclei containing magic numbers of neutrons or protons or both (24, 80).

The integrated cross section for photon absorption by a nucleus has a special significance since it indicates the electric dipole (E1) character of the giant resonance. The total integrated E1 absorption cross section has been calculated by Levinger and Bethe (60) using the E1 sum rule which applies to the sum of the squares of the transition matrix elements between the ground state and all excited states and yields the result

$$\sigma_{\text{int}} = 0.06 \frac{N^2}{A} \text{ MeV-barns.}$$

If exchange forces between the nucleons are taken into account, this equation must be modified. Levinger and Bethe (60) calculated that the integrated cross section would be increased by 30 percent if the neutron-proton potential was entirely exchange in character. The integrated cross
section then becomes

\[ \sigma_{\text{int}} = 0.06 \frac{N^2}{A} (1 + 0.8x) \text{ MeV-barns} \]

where \( x \) is the fraction of the potential which has exchange character.

Nathans and Halpern(76) found reasonable agreement between the \((\gamma, n)\) cross sections integrated to 25 MeV for medium and heavy nuclei and the sum rule expression for the integrated cross section. In light nuclei where competition from other particles, in particular protons, is relatively important the experimental values of the integrated \((\gamma, n)\) cross sections naturally fall below the predictions of the sum rule which predicts the total cross section for photon absorption. It can therefore be regarded as established that the giant resonance is El in character. This fact is further substantiated by the calculations of Blatt and Weisskopf(8) who found that the El transition probability predominates over the magnetic dipole (M1) transition probability by an order of magnitude and over the electric quadrupole (E2) transition probability by two orders of magnitude in the energy region of the giant resonance.

For some time the giant resonance phenomena have been explained on the basis of two models involving extreme assumptions about the inter-nucleon forces. In the collective model the assumption of strongly coupled nucleons is made and the normal modes of oscillation of the entire nucleus are of primary concern. The independent particle model (I.P.M.), however, assumes that each nucleon moves independently in a common potential which represents the average effect of interactions with all the other nucleons.

The first model suggested to explain the giant resonance was the collective model of Goldhaber and Teller(45). In this essentially classical...
model the neutrons and protons were treated as fluids which ran through each other, in bulk, in opposite directions. Various assumptions about the inter-nucleon forces were made which were equivalent to treating the neutrons and protons as either compressible nuclear fluids or incompressible and interpenetrating nuclear fluids. The compressible nuclear fluids model gave the correct dependence on $A$ of the resonance energy (100).

The collective model gives the correct order of magnitude for the integrated cross section; however, it is difficult to introduce into it any term corresponding to the exchange term. Since the collective model pictures the giant resonance as a single vibration there is no difficulty in explaining the sharpness of the resonance. The width of the resonance is interpreted as a frictional effect between the nucleons but cannot be accounted for on a quantitative basis.

The collective model led to the first prediction of the splitting of the giant resonance in strongly deformed nuclei. This prediction, which was based on the fact that there are two frequencies of oscillation along the major and minor axes of the ellipsoid, led to experimental investigation in the rare earth region (43), where excellent agreement was found with the independent calculations of Okamoto (80) and Danos (24). Again the compressible nuclear fluids model is favoured indicating some internal consistency in the use of this approximation.

The collective model does not attempt to describe the de-excitation process; however, absorption of the photon by the collective motion of the nucleus does imply the formation of a true compound nucleus state since the excitation energy is shared by all the nucleons. The decay of the compound nucleus is described by statistical theory (8) in which the absorption of the
photon is considered as a nuclear heating process which causes the evaporation of neutrons or other particles. The energy distribution of emitted neutrons has a peak at one or two MeV, corresponding to the nuclear temperature of the residual nucleus, with a Maxwellian tail at higher energies. Since the Coulomb barrier inhibits the emission of low-energy charged particles the proton energy spectrum peaks somewhat below the height of the barrier rather than at the nuclear temperature. The angular distribution of the emitted nucleons is expected to be isotropic(111) except that the high-energy particles, which leave the residual nucleus in its ground or one of its first few excited states, may have an $a + b\sin^2\theta$ distribution where $\theta$ is the angle between the direction of the incoming photon and the emitted nucleon(59). Since the $E_1-E_2$ interference term averages to zero, the statistical photoprotons should show no forward-backward asymmetry(9).

The first indication of the inadequacy of the statistical theory was the early work of Hirzel and Wüffler(51) who found that the relative yield of $(\gamma,p)$ to $(\gamma,n)$ reactions was a factor of 20 to 1000 too high in a number of nuclei with $A$ near 100. Diven and Almy(29) determined that the proton energy spectrum from silver consisted of a low-energy isotropic group in quantitative agreement with the statistical model as well as a very anisotropic high-energy tail whose yield was definitely in excess of the statistical predictions. The failure of the statistical model to account for the anomalous yield of high-energy photoprotons stimulated theoretical calculations on the basis of the I.P.M.

In the I.P.M. calculations of Wilkinson(109) a direct interaction is considered to take place between the absorbed photon and a nucleon, in which the nucleon is lifted into an excited single-particle state. On the basis of
this model Wilkinson was able to explain satisfactorily the giant resonance phenomena as well as to account quantitatively for the anomalously strong emission of high-energy protons.

When all the possible transitions which can be made by nucleons in the I.P.M. are considered, they have been shown to effectively exhaust the El sum rule(14). The main contribution to the El transition strength is from closed shells that lie near the Fermi surface. Although the El sum rule involves all the nucleons of the nucleus, the nucleons in low-lying states of the nucleus are excluded from El transitions by the Pauli principle since the levels to which these nucleons would make their El transitions are already full. These nucleons pass their contribution to the integrated cross section up to the levels near the Fermi surface by cancelling the negative contribution of the downward transitions from these full levels(110). The valence nucleons are relatively few in number and make little contribution to the integrated cross section.

Although a large number of El transitions is possible, Wilkinson(109), using a square well of infinite depth, showed that these transitions cluster in a way consistent with their representing the giant resonance. In the I.P.M. there are several sources of the width of the giant resonance. Important contributions come from the separation in energy of the fundamental single-particle transitions and the smearing out of the upper I.P.M. state due to the imaginary part of the optical model potential, which blurs the position of the single-particle states by dissolving them into a multitude of compound nucleus states. In non-closed shell nuclei the coupling of the valence nucleons to the excited core causes an additional broadening of the giant resonance.

The principal shortcoming of the Wilkinson single-particle model lies in the explanation of the position of the maximum energy of dipole photon

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absorption. When the I.P.M. level positions are calculated using a reasonable nuclear radius and well depth, the E1 transitions cluster at an energy which is about half as large as the experimentally observed giant resonance energy. This discrepancy is reconciled by the introduction of a velocity-dependent potential into the shell model, which gives rise to a reduced mass for the nucleon moving inside the nucleus of about half its normal mass\((109,83)\). In accordance with this effective nucleon mass, the energy spacing between unoccupied levels of the nucleus would be expanded by a factor of two; however, more recent experimental evidence based on other reactions\((20,21)\) indicates the contrary. Recently, however, detailed calculations for the photodisintegration of some double closed-shell nuclei including \(^{12}\text{C}(4,107)\), \(^{16}\text{O}(3,12)\), \(^{40}\text{Ca}(12,3)\), \(^{90}\text{Zr}(96)\) and \(^{208}\text{Pb}(4)\) with the inclusion of particle-hole interactions, neglected in the Wilkinson single-particle model, have shown that the shell model is capable of furnishing a comprehensive description of the photodisintegration of light nuclei in the giant resonance region as well as reconciling the location of the giant dipole resonance in medium and heavy nuclei.

In addition to giving a satisfactory account of the giant resonance phenomena, the I.P.M. of Wilkinson is also capable of describing the de-excitation of the nucleus. The nucleon, with which the absorbed photon interacted, is elevated to an excited single-particle state from which it may either be directly emitted without sharing its energy with the other nucleons or be dissolved into a compound nucleus state which subsequently decays by a statistical process. The "resonance direct" emission of nucleons leads to both a much higher proton to neutron yield ratio and a greater percentage of high-energy nucleons than a purely statistical process so that the I.P.M. gives better agreement with the empirical results on photonucleon emission than the
statistical process alone.

The angular distribution of direct photonucleons in the I.P.M. is given by \( a + b \sin^2 \theta \) where the \( b/a \) ratios depend on the angular momentum of the nucleon in the nucleus for the various El transitions. The interference of \( E2 \) absorption gives rise to a forward-backward asymmetry in which the angular distribution peaks forward of 90 degrees. When this asymmetry is present the angular distribution can be fitted to the expression

\[
Y(\theta) = a + b \sin^2 \theta (1 + p \cos \theta)^2.
\]  

Again, the direct interaction of the I.P.M. gives better agreement with experiment than the statistical model alone, since both the highly anisotropic angular distribution of the high-energy component of the photonucleon spectra and the forward-backward asymmetry observed in the photoproton angular distributions can be explained.

In the remainder of this thesis the theoretical basis of the I.P.M. of photonuclear reactions will be outlined in some detail and the results of this experimental work on photoproton reactions will be discussed in terms of the predictions of the I.P.M.
2.1 Absorption of Electromagnetic Radiation

2.1.1 The interaction Hamiltonian, electric dipole transition probability and the concept of effective charge

The theory of photon absorption by an atomic nucleus is based on time-dependent perturbation theory for the interaction between the electromagnetic field and the nucleons. The electromagnetic field interacts with both the charge of the nucleon and its magnetic moment.

The Hamiltonian for a non-relativistic charged particle in the presence of an electromagnetic field with vector potential $\vec{A}$ is (10)

$$H = \left\{ \vec{p} - \left( \frac{e}{c} \right) \vec{A} \right\}^2 \frac{2m}{2mc} - \frac{e\hbar}{2mc} \mu \vec{S} \cdot (\nabla \times \vec{A}) + V(\vec{r})$$

where $\vec{p}$ is the momentum of the particle,

$m$ is the mass of the particle,

$e$ is the charge of the particle,

$c$ is the velocity of the electromagnetic wave,

$\mu = \frac{e\hbar}{2mc}$ is the magnetic moment of the particle,

$\vec{S}$ is the spin of the particle and

$V(\vec{r})$ is the non-electromagnetic potential acting on the particle within the nucleus. The Hamiltonian may be expanded and written in the form

$$H = H_0 + H'$$
where $H_0$ is the Hamiltonian of the particle in the absence of the electromagnetic field and

$H'$ is the interaction Hamiltonian.

If terms in $A^2$ are neglected the interaction Hamiltonian is given by

$$H' = \frac{e\vec{A} \cdot \vec{p}}{mc} - \frac{e\hbar}{2mc} \mu_0 (\nabla \times \vec{A}).$$

Since $H'$ is much smaller than $H_0$ the effect of the radiation interaction through the operator properties of the vector potential $\vec{A}$ may be calculated by time-dependent perturbation theory. The vector potential $\vec{A}$ is first expanded into a series of electric and magnetic waves by the multipole expansion(8) in which the electric and magnetic multipoles are classified according to their angular momentum (multipole order) and parity.

The vector potential may also be written in the form(59)

$$\vec{A} = A_0 \tilde{u} e^{i(\omega t - k \cdot \vec{r})}$$

where $\tilde{u}$ is the unit vector along the photon polarization direction,

$\omega$ is the frequency of the electromagnetic wave,

$k$ is the photon wave number,

$\vec{r}$ is the vector position of the multipole field and

$A_0$ is the amplitude of the vector potential.

The interaction is then classed in terms of multipoles by expanding $e^{-i k \cdot \vec{r}}$ in a Taylor series about $r = 0$ so that

$$e^{-i k \cdot \vec{r}} = 1 - i k \cdot \vec{r} - \frac{1}{2}(k \cdot \vec{r})^2 \ldots$$
The El interaction is obtained by using only the first term in the expansion(8). Further multipoles are obtained by a suitable grouping of terms in successive higher order approximations. The next term, \(-iE_r\), in the expansion of the exponential gives both the E2 and the M1 interactions.

The transition probability from the ground state \(o\) to an excited state \(n\), \(|c_n|^2\), is proportional to \(|H'^{on}|^2\) where \(H'^{on}\) is the matrix element of the interaction Hamiltonian between the states \(o\) and \(n\) and is given by

\[
H'^{on} = \int \psi_n^* H' \psi_o \, dt
\]

where \(\psi_o\) and \(\psi_n\) are the wave functions of the ground and excited states respectively.

Blatt and Weisskopf(8) have performed the appropriate multipole expansion of the vector potential \(A\) and determined that the probability of excitation of a nucleus from its ground state \(o\) to an excited state \(n\) by the absorption of a photon of appropriate energy and multipolarity, \(L\), is

\[
|c_n|^2 = S(w_{on}) \sum_{L=1}^{\infty} \sum_{M=1}^{L} \frac{k^{2L-1}}{L[(2L+1)!]} \frac{M}{Q_L + Q'_M} \left[ M + M' \right] \left[ M + M' \right] 2.7
\]

where \(S(w_{on})\) is the flux of photons with energy \(\hbar w_{on}\).

\(M\) is the component of the angular momentum, \(L\), of a quantum of multipole radiation along a reference axis,

\(Q_L^M\) is the multipole moment for the transition due to the interaction of the electric multipole with the charge of the nucleon,

\(Q'_M\) is the multipole moment for the transition due to the interaction of the electric multipole with the magnetic moment of the nucleon,
\( M_L \) is the multipole moment for the transition due to the interaction of the magnetic multipole with the charge of the nucleon and 
\( M_L' \) is the multipole moment for the transition due to the interaction of the magnetic multipole with the magnetic moment of the nucleon.

If these matrix elements are evaluated, the contribution of the \( M_l \) interaction is found to be about 10\% of the \( El \) interaction and the \( E2 \) interaction is about 1\% of the \( El \) interaction for the absorption of photons in the region of the giant resonance. It has further been shown that the \( El \) interaction is predominantly with the charge of the nucleon. Thus the transition probability reduces to

\[
|c_{n_{El}}|^2 = S(\omega_{on}) \frac{8\pi^3}{3} \hbar^2 \left[ \left| Q_1 \right|^2 + \left| Q_{-1} \right|^2 \right]. \tag{2.8}
\]

The matrix element \( Q_1 \) is given by

\[
Q_1(o,n) = e \sum_{i=1}^{Z} \int \bar{r}_i \, Y_{1}(\theta_i, \phi_i) \, \psi_n^* \, \psi_o \, d\tau \tag{2.9}
\]

where the summation is taken over all the protons in the nucleus. The sum over the protons can be written as a sum over all the nucleons of the nucleus by introducing the concept of effective charges \( e^* \) for \( El \) radiation. If a proton is given a displacement \( \bar{r}_i \), the nucleus of mass \((A-1)\) and charge \((Z-1)\) recoils with a displacement \(-\bar{r}_i/(A-1)\) and the net dipole moment \( \bar{D}_i \) is

\[
\bar{D}_i = e(\bar{r}_i - (Z-1)/(A-1) \, \bar{r}_i) = \frac{N}{A-1} \, e\bar{r}_i. \tag{2.10}
\]
The separation between the centres of mass, \( \vec{d}_1 \), is then

\[
\vec{d}_1 = \vec{r}_1 + \frac{1}{(A-1)} \vec{r}_1 = \frac{A}{A-1} \vec{r}_1
\]

so that the dipole moment becomes

\[
\vec{D}_1 = \frac{N}{A} e \vec{d}_1
\]

when the transformation to relative coordinates is made. The effective charge for protons is then

\[
e' = \frac{N}{A} e
\]

In a similar manner the effective charge for neutrons can be shown to be

\[
e' = \frac{Z}{A} e
\]

The matrix element \( Q^M_1 \) then becomes

\[
Q^M_1 = \sum_{i=1}^{A} e_i \int \vec{r}_1 \, Y^M_1 (\theta_1, \phi_1) \, \psi_n^* \psi_o \, d\tau
\]
2.1.2 Cross section for compound nucleus formation, dipole sum rule and integrated electric dipole cross section

If the assumption is made that the source emits photons within a certain energy interval, $\Delta E$, which is large compared to the level spacing $W_1(E_n)$ of the E1 levels in the compound nucleus at the excitation energy $E_n$, then many states are excited and the total excitation probability is

$$\left| C_{nE1} \right|^2 = \frac{\sum_{E_n} \omega_{\text{on}} \Delta E}{\sum_{E_n} \omega_{\text{on}}} \left| C_{nE1} \right|^2 = \left\langle \left| C_{nE1} \right|^2 \right\rangle_{\text{av.}} \frac{\Delta E}{W_1(E_n)} \tag{2.16}$$

where $\frac{1}{W_1(E_n)}$ is the density of final states of the compound nucleus at the excitation energy $E_n$.

The cross section for the formation of the excited compound nucleus is then obtained by dividing the total transition probability by the flux of incident photons

$$S(\omega) d\omega = S(\omega) \frac{\Delta E}{h} \tag{2.17}$$

The cross section for the creation of the compound nucleus by the absorption of E1 radiation is then

$$C_{E1}(E) = \frac{\alpha E^3}{3} k \frac{\left\langle \left| q_1^1 \right|^2 + \left| q_1^{-1} \right|^2 \right\rangle_{\text{av.}}}{W_1(E_n)} \tag{2.18}$$

This cross section does not show any fine structure in the giant resonance but rather represents an average over many individual transitions. Since the
energy interval, $\Delta E$, is small compared to the width of the giant resonance, however, the latter can be resolved.

The total contribution of the El transitions to the photoeffect or to the integrated El cross section can be evaluated by the use of certain relations, called sum rules, which apply to the sum of the squares of the transition matrix elements between the ground state $o$ of the nucleus and all the excited states $n$. The simplest sum rule is the one for El matrix elements which is given for $M=0$ by the equation (3)

$$Q_{1}^{0}(o,n) = \left( \frac{3}{4\pi} \right)^{1/2} \sum_{i=1}^{A} e_{i}^{*} \int \psi_{n}^{*} \tilde{z}_{i} \psi_{o} d\tau$$  

where $e_{i}^{*} \tilde{z}_{i}$ is the dipole operator of the nucleon, $i$.

In order to derive the sum rule the following commutation relation is used (3)

$$\sum_{n} \left[ \tilde{z}_{i}(o,n) p_{i',z}(n,o) - p_{i',z}(o,n) \tilde{z}_{i}(n,o) \right] = i\hbar \delta_{i'i}$$  

where $\tilde{z}_{i}(o,n)$ is the matrix element of the $z$ coordinate of the nucleon, $i$, for a transition between nuclear states $o$ and $n$ and $p_{i',z}(o,n)$ is the matrix element of the $z$ component of the momentum of the nucleon, $i'$, between quantum states $o$ and $n$ of the nucleus.

In the Heisenberg representation the following matrix relation holds

$$p_{i',z}(o,n) = m \tilde{z}_{i}(o,n)$$  

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Since each matrix element oscillates with the exponential $e^{im\omega_{on}t}$, equation 2.21 may be written in the form

$$p \cdot z(o,n) = \frac{i m}{\hbar} (E_o - E_n) z_i(o,n) .$$  \hspace{1cm} 2.22

Equation 2.20 then becomes

$$\frac{i m}{\hbar} \sum_n \left[ z_i(o,n)z_i(n,o) \left\{ E_n - E_o \right\} - \left\{ E_o - E_n \right\} z_i(o,n)z_i(n,o) \right] = i m \delta_{ii} .$$  \hspace{1cm} 2.23

If $i = i'$ then this expression reduces to the form

$$\sum_n (E_n - E_o) z_i(o,n) z_i(n,o) = \frac{n^2}{2m} .$$  \hspace{1cm} 2.24

Since $z_i(n,o)$ is Hermitian $z_i(n,o) = z_i^*(o,n)$ so that

$$\frac{2m}{\hbar^2} \sum_n (E_n - E_o) |z_i(o,n)|^2 = 1 .$$  \hspace{1cm} 2.25

Combining equations 2.19 and 2.25 and substituting the effective charges for neutrons and protons leads to the El sum rule

$$\frac{2m}{\hbar^2} \sum_n (E_n - E_o) \left| q_1^0(o,n) \right|^2 = \frac{2}{4\pi} \sum_{i=1}^A \left( e_i' \right)^2$$

$$= \frac{3e^2}{4\pi} \left[ \frac{Z(N)}{A} \right]^2 + N\left( \frac{Z}{A} \right)^2 \right]$$
The corresponding sums for $|q_1^{0}(o,n)|^2$ and $|q_1^{-1}(o,n)|^2$ give the same result(3). Blatt and Weisskopf(3) have also shown that (Equation 7.15, Chapter XII)

$$\sum_n |q_1^{0}(o,n)|^2 = \int_0^\infty \langle |q_1^{0}(o,n)|^2 \rangle \frac{dE_n}{W_1(E_n)}. \tag{2.27}$$

The integrated El cross section for the formation of the compound nucleus is then found by combining equations 2.18, 2.26 and 2.27 to be given by the equation

$$\sigma_{\text{El int}} = \frac{2\pi^2 e^2}{mc} \frac{NZ}{A} \text{MeV-cm}^2. \tag{2.28}$$

In the case of compound nucleus formation the integrated El cross section is the sum of A individual nucleon cross sections of approximately equal magnitude since each nucleon is contributing in proportion to the square of its effective charge. In the shell model treatment, which follows, the Pauli principle excludes El absorption by certain nucleons because the levels to which they would make their transitions are already full. Thus each tightly bound nucleon makes its contribution by passing its dipole strength up to those nucleons which are allowed El absorption so that their El strength is considerably enhanced(57).
2.1.3 Independent particle model calculations of the photon absorption probabilities and integrated electric dipole cross sections

The probability of absorption of a radiation of energy $h\omega_{on}$ was given in equation 2.8. If the expression 2.15 for the matrix element $Q_1^M$ is substituted into this equation the transition probability from the ground state $o$ to an excited state $n$ then becomes

$$|c_n^E|^2 = S(\omega_{on}) \frac{8\pi^3}{3} \frac{(E_n-E_o)}{\hbar c} \sum_{M=1}^A \sum_{i=1}^A \left[ (e_i^o)^2 \left| \int \bar{r}_i Y_1^M (\theta_i, \phi_i) \psi_n^* \psi_o \, d\tau \right|^2 \right]$$  \hspace{1cm} (2.29)

The cross section for the $E_1$ transition between the states $o$ and $n$ is defined as the transition probability per unit flux so that

$$(\sigma_{E_1})_{on} \, dE = \frac{8\pi^3}{3} \frac{(E_n-E_o)}{\hbar c} \sum_{M=1}^A \sum_{i=1}^A \left[ (e_i^o)^2 \left| \int \bar{r}_i Y_1^M (\theta_i, \phi_i) \psi_n^* \psi_o \, d\tau \right|^2 \right] \hspace{1cm} (2.30)$$

The integrated $E_1$ cross section is then

$$\sigma_{E_1, \text{int}} = \frac{8\pi^3}{3\hbar c} \sum_{i=1}^A \sum_{n}^A \sum_{M=1}^A \left[ (E_n-E_o) \left| \int \bar{r}_i Y_1^M (\theta_i, \phi_i) \psi_n^* \psi_o \, d\tau \right|^2 \right] \hspace{1cm} (2.31)$$

In order to evaluate the $E_1$ transition probabilities and integrated cross section on the basis of the I.P.M. the single-particle wave functions of the ground and excited states must be determined. In the I.P.M. the nucleons are considered to move independently of each other in a constant potential which represents the average effect of all interactions with the other nucleons. For ease of calculation the most common choices of this
potential are the square well potential and the harmonic oscillator potential. The single-particle wave functions are then the eigenfunctions of the Schrödinger equation employing such a potential.

For a central potential the eigenfunctions of the wave equation can be separated into a radial part, $R_{nl}(r)$, and a second part which includes the angular dependence and spin functions, $Y^m_{jm}(\theta, \phi, s)$, so that the wave function for the $i$th nucleon assumes the form

$$\psi_{nljm}(r, \theta, \phi, s) = R_{nl}(r) Y^m_{jm}(\theta, \phi, s)$$

where $n$ is the number of radial nodes of the wave function,

- $l$ is the orbital angular momentum of the nucleon,
- $s$ is the spin angular momentum of the nucleon,
- $j$ is the vector sum of the orbital and spin angular momenta of the nucleon and
- $m$ is the component of $j$ on a reference axis.

The $Y^m_{jm}$ represent eigenfunctions of the total angular momentum $j$ and the component of the total angular momentum on a reference axis. These eigenfunctions are formed by combining the orbital angular wave functions, $Y^m_{lm}$, with the spin wave functions, $\chi_{ms}$, in the form

$$Y^m_{jm} = \sum_{m_\ell} \sum_{m_s} C_{l\ell s}(j, m; m_\ell, m_s) Y^m_{\ell m_\ell}(\theta, \phi) \chi_{ms}$$

where $C_{l\ell s}(j, m; m_\ell, m_s)$ is the appropriate Clebsch Gordon coefficient.

The solution of the Schrödinger equation leads to a set of energy eigenvalues, $E_{nljm}$. In a spherically symmetric potential field the levels
are \((2j + 1)\)-fold degenerate since all states with the same \(n\ell j\) have the same energy, \(E_{n\ell j}\).

The substitution of equation 2.52 for the single-particle wave function leads to an expression for the El transition probability between the levels \(n\ell j, j\) and \(n'\ell ' j', j'\) for the \(i\)th nucleon of the form

\[
|c_{qq'}^{\text{El}}|^2 = S(w_{qq'}) \frac{8\pi^2\varepsilon^2}{\hbar^2 c} \times \\
\sum_{M=1}^{2J+1} \left| (E_{q'j} - E_{qj}) \left( \int_{j}^{m} r^{\ell+1} \Omega^{m}_{j} r^{\ell} \Omega^{m}_{j'} d\Omega \right)^2 \right|
\]

where \(q, q'\) denote the quantum numbers \(n\ell j\) and \(n'\ell ' j'\) respectively of the initial and final states. The total El transition probability between these levels is then

\[
|c_{qq'}^{\text{El}}|^2 = \sum_{i} |c_{qq'}^{i}|^2 = (2J + 1) \left| c_{qq'}^{i} \right|^2
\]

where the factor \((2J + 1)\) is the number of nucleons in the initial filled level. Wilkinson(109) has written this transition probability in terms of the square of the radial overlap integral, \(D\), measured in units of the square of the nuclear radius, \(R^2\), where

\[
DR^2 = \left| \int_{R_{n\ell j}}^{\infty} \frac{R_{n'\ell '} r^2}{R_{n\ell j}} d\rho \right|^2
\]

The angular part of the matrix element is given by
\[ S_{l_j} = \left| \frac{1}{2} \int \int Y^m_{j'} \cos \theta^m_j \, d\Omega \right|^2. \]  

If it is recalled that \( |q_J^{0}(o,n)|^2 = |q_J^{1}(o,n)|^2 = |q_J^{-1}(o,n)|^2 \), it can be seen from the relation (14)

\[ \cos \theta = \left( \frac{4\pi}{3} \right)^{1/2} Y^0_1(\theta, \psi). \]

that the angular part of equation 2.34 may be written in the form

\[ \left| \frac{1}{2} \int \int Y^m_{j'} Y^m_1 \, d\Omega \right|^2 = \frac{3}{4\pi} S_{l_j}. \]

The total probability of absorption of El radiation between the \( n, l, j \) and \( n', l', j' \) levels of the nucleus is then

\[ \left| c_{El} \right|^2 = \sum_{qq'} s(\omega_{qq'}) \frac{4\pi e^2}{\hbar^2 c} \left[ (E_q - E_{q'}) S_{l_j} \lambda_{n'l'} \, R^2 \right]. \]

where \( S_{l_j} \) is called the enhancement factor.

The enhancement factors for the three possible dipole transitions from each level \( l, j \) are given in Table 2.1 below.
Table 2.1

<table>
<thead>
<tr>
<th>Level</th>
<th>Transition</th>
<th>Enhancement Factor</th>
</tr>
</thead>
<tbody>
<tr>
<td>( j = \ell + 1/2 )</td>
<td>( \ell + 1/2 \rightarrow \ell + 1 + 1/2 )</td>
<td>( \frac{(\ell + 1)(\ell + 2)}{2\ell + 3} )</td>
</tr>
<tr>
<td>( j = \ell + 1/2 )</td>
<td>( \ell + 1/2 \rightarrow \ell + 1 - 1/2 )</td>
<td>( \frac{\ell + 1}{(2\ell + 3)(2\ell + 1)} )</td>
</tr>
<tr>
<td></td>
<td>( \ell + 1/2 \rightarrow \ell - 1 + 1/2 )</td>
<td>( \frac{\ell(\ell + 1)}{2\ell + 1} )</td>
</tr>
<tr>
<td></td>
<td>( \ell - 1/2 \rightarrow \ell + 1 - 1/2 )</td>
<td>( \frac{\ell(\ell + 1)}{2\ell + 1} )</td>
</tr>
<tr>
<td>( j = \ell - 1/2 )</td>
<td>( \ell - 1/2 \rightarrow \ell - 1 + 1/2 )</td>
<td>( \frac{\ell}{(2\ell + 1)(2\ell - 1)} )</td>
</tr>
<tr>
<td></td>
<td>( \ell - 1/2 \rightarrow \ell - 1 - 1/2 )</td>
<td>( \frac{\ell(\ell - 1)}{2\ell - 1} )</td>
</tr>
</tbody>
</table>

The relative importance of the various El transitions is largely determined by the squares of the radial overlap integrals which are given in Table 2.2 for an infinite square well\((109)\).
Table 2.2

D is the square of the radial overlap integral in units of the square of the nuclear radius.

<table>
<thead>
<tr>
<th>Transition</th>
<th>( l ) = 0</th>
<th>1</th>
<th>2</th>
<th>3</th>
<th>4</th>
<th>5</th>
<th>6</th>
</tr>
</thead>
<tbody>
<tr>
<td>( n\ell \rightarrow 1\ell + 1 )</td>
<td>-</td>
<td>0.28</td>
<td>0.38</td>
<td>0.44</td>
<td>0.49</td>
<td>0.53</td>
<td>0.56</td>
</tr>
<tr>
<td>( 2\ell \rightarrow 2\ell + 1 )</td>
<td>-</td>
<td>0.23</td>
<td>0.28</td>
<td>0.33</td>
<td>0.37</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>( 3\ell \rightarrow 3\ell + 1 )</td>
<td>-</td>
<td>0.22</td>
<td>0.25</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>( 1\ell \rightarrow 2\ell - 1 )</td>
<td>-</td>
<td>0.092</td>
<td>0.065</td>
<td>0.050</td>
<td>0.039</td>
<td>0.036</td>
<td>-</td>
</tr>
<tr>
<td>( 1\ell \rightarrow 2\ell + 1 )</td>
<td>-</td>
<td>0.001</td>
<td>0.002</td>
<td>0.002</td>
<td>0.003</td>
<td>-</td>
<td>-</td>
</tr>
<tr>
<td>( 2\ell \rightarrow 3\ell - 1 )</td>
<td>-</td>
<td>0.12</td>
<td>0.09</td>
<td>-</td>
<td>-</td>
<td>-</td>
<td>-</td>
</tr>
</tbody>
</table>

These results indicate that the strongest overlap occurs for transitions in which the number of nodes of the radial wavefunction is small and does not change.

From these tables the relative strengths of the various El transitions based on the Wilkinson model may be estimated.

If equation 2.32 is substituted into equation 2.31 the integrated El absorption cross section based on the I.P.M. then becomes

\[
C_{\text{El int}} = \frac{8\pi^3 e^2}{3\hbar c} \sum_{qq'} \sum_{M=-1}^1 (E_{q'} - E_q) \times \left| \int \int \int \int R_{n'\ell',l',m'} R_{n\ell,l,m} d\Omega d\Omega' \right|^2
\]

2.41

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which may be reduced to the form

\[ \sum_{E_{\text{int}}} = \frac{2\pi^2 e^2 m}{mc} \sum_{qq'} \left[ \frac{2m(E_{q'} - E_q)}{\hbar^2} |S_{j'dn'l'_f} R|^2 \right] \]

according to equation 2.40.

This expression was evaluated for $\text{Pb}^{208}$ by Buskirk(14) who found good agreement with the El sum rule.

2.2 Decay of the Excited Nucleus

The I.P.M. of Wilkinson(109) was the first suggestion that both the "resonance direct" and the statistical emission were the result of the same absorption process. In this process a single nucleon is elevated from the initial single-particle state into a multitude of final states which are smeared out in energy due to interaction of the excited nucleon with the rest of the nucleus. The wave function of each of these states contains in its shell-model expansion a term corresponding to the I.P.M. state which was responsible for the absorption in the basic I.P.M. picture. Each of these states may then lead to "resonance direct" emission through the part of the wave function which represents its share of the single-particle motion or be statistically emitted through the remaining portion of the wave function corresponding to collective motion.

The absorption process can be more crudely pictured as the elevation of a nucleon from a closed shell to the appropriate I.P.M. state from which it may be directly emitted without interacting with the other nucleons in
what is termed a "resonance direct" interaction, or alternatively it may interact with the rest of the nucleus to form an ordinary compound nucleus state which subsequently decays by statistical emission of a nucleon.

It is then interesting to enquire into the relative probability of these two decay mechanisms. The escape width for a proton in a "resonance direct" interaction is given by Blatt and Weisskopf(8) (Chapter VIII, Equation 7.12) as

\[ \frac{\Gamma_{\text{esc}}}{\text{esc}} = T_{\ell} \frac{W}{2\pi} \]

where \( T_{\ell} \) is the transmission coefficient through the potential (Coulomb and centrifugal) barrier of the nucleus for a proton of angular momentum, \( \ell \), and

\( W \) is the energy level spacing.

From equation 2.43 it can be seen that the probability of a proton being directly emitted is proportional to the transmission coefficient. The transmission coefficients for s, p, d, f and g protons were calculated for energies up to 16 MeV by Mitchell(71) for three nuclei within the Z-range of this experiment. These calculations were based on the Coulomb wave functions \( F_\ell \) and \( G_\ell \) of Sharp et al.(91) using the formula

\[ T_{\ell} = \frac{4\pi}{A_{\ell}^2 \nu^2 + 2\nu + B_{\ell}^2} \]

where

\[ A_{\ell}^2 = F_\ell^2 + G_\ell^2, \]

\[ B_{\ell}^2 = F'_\ell^2 + G'_\ell^2, \]

\[ \nu = \frac{K}{K} \]

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k is the wave number of the nucleon outside the nucleus and 

K is the wave number of the nucleon within the nucleus.

The results for Ni, Nb and Sn are reproduced in Figure 2.1. The proton transmission coefficients for other nuclei within this range can be found by interpolation.

The transmission coefficient for a proton has also been given by the following approximate expression of Blatt and Weisskopf(8)(Chapter VIII, Equation 5.6),

\[ T = \frac{4kP}{K} \quad 2.45 \]

where \( P \) is the penetrability defined by Equation 2.49, Chapter VIII(8).

The level spacing for a one-dimensional square well of width \( b \) has been calculated by Blatt and Weisskopf(8) with the result

\[ W = \frac{m^2 \pi^2}{mb^2} \quad 2.46 \]

If equations 2.45 and 2.46 are substituted into equation 2.43 the escape width for the "resonance direct" emission of protons then becomes

\[ \Gamma_{esc} = \frac{2kPm^2}{Kmb} \quad 2.47 \]

If the Wilson-Sommerfeld rule is applied, this equation may be reduced to the form(109)

\[ \Gamma_{esc} = \frac{2kPm^2}{mb} \quad 2.43 \]

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FIGURE 2.1. Proton Transmission Coefficients.

The transmission coefficients for s, p, d, f and g protons are given for Ni, Nb and Sn for proton energies up to 16 MeV. The classical Coulomb barrier heights are indicated by the arrows.
The capture width leading to compound nucleus formation with subsequent decay by statistical emission is assumed by Wilkinson (109) to be 2D where D is the imaginary part of the optical model potential. The relative probability of a "resonance direct" emission is then (109)

\[ C = \frac{\Gamma_{\text{esc}}}{2D + \Gamma_{\text{esc}}} \]  

The cross section for the "resonance direct" emission of protons then becomes

\[ \sigma_d(\omega) = \sigma(\omega) \frac{\Gamma_{\text{esc}}}{2D + \Gamma_{\text{esc}}} \]  

where \( \sigma(\omega) \) is the absorption cross section for the I.P.M. transition.

This expression must be calculated for each of the shell-model transitions that are operative and then summed in order to compute the "resonance direct" emission of protons for a nucleus. Wilkinson (109) has shown that most of the absorption cross section is carried by shell-model transitions between states of high orbital angular momentum; however, the final states of high orbital angular momentum have a small escape width for the "resonance direct" emission of protons. Consequently, the shell-model transitions to low angular momentum states can make the greatest contribution to the "resonance direct" emission of protons although they are not those chiefly responsible for photon absorption.
2.3 Properties of the Reaction Products

In this section the angular distributions, energy distributions and yields of the photoprotons are discussed in terms of the two possible modes of de-excitation after photon absorption.

2.3.1 Angular distributions

2.3.1.1 Statistical de-excitation

In a compound nucleus process the energy of the incident particle or photon is spread over the many particles of the nucleus with the result that the incident particle or photon can no longer be distinguished. The decay products are then often said to be independent of the way in which the compound nucleus was formed. This loss of memory, as it is usually called, must be qualified to some degree, however, since the total angular momentum, \( j \), its component along a reference axis, \( m \), and the parity, \( p \), must be conserved in the statistical process.

When the angular distributions resulting from the decay of the compound nucleus are calculated with these conservation principles applied, it has been shown\(^{(47,111)}\) that they are isotropic if the statistical theory can be applied to both the compound and residual nuclei. Since a large amount of energy is available in photonuclear reactions this is expected to be the case for the emission of all low-energy photonucleons.

The emission of high-energy photonucleons may leave the residual nucleus in its ground state or one of its first few excited states so that statistical theory no longer applies to the residual nucleus. The angular distributions have then been shown\(^{(47)}\) to be anisotropic. For dipole
absorption the angular distribution is expected to be of the form
\[ a + b \sin^2 \theta \] (59).

Although the statistical angular distribution may be anisotropic, as discussed above, it is expected to be symmetric about 90 degrees. This follows from the fact that interference terms between outgoing waves of opposite parity are assumed to cancel out (9,111). An experimental forward-backward asymmetry consequently disproves the statistical nature of the reaction.

2.3.1.2 Independent particle model

The angular distribution of nucleons which are directly emitted following the absorption of El radiation between appropriate single-particle levels of the I.P.M. is given by Courant (23) as

\[ Y(\theta) = a + b \sin^2 \theta \]  \hspace{1cm} 2.51

where

\[ \frac{b}{a} = \frac{l + 2}{2l} \] \hspace{1cm} \text{for} \hspace{1cm} l \rightarrow l+1 \]  \hspace{1cm} 2.52

and

\[ \frac{b}{a} = \frac{l - 1}{2l + 2} \] \hspace{1cm} \text{for} \hspace{1cm} l \rightarrow l-1 \]  \hspace{1cm} 2.52

The anisotropies or b/a ratios for the important El transitions are given in Table 2.3 below.
From this table it can be seen that the $b/a$ ratio for $l \rightarrow l + 1$ transitions is always greater than 0.5 and decreases with increasing $l$ and for $l \rightarrow l - 1$ transitions is always less than 0.5 and increases with $l$. The above equations, however, do not include the effects of nucleon spin.

When nucleon spin is taken into account the angular distribution is still given by equation 2.51 but the $b/a$ ratios, derived by Heiss(50), depend on the total angular momentum, $j$, of the nucleons in the initial state. Due to interference terms between the different final $j$'s the simple predictions of Courant's approximations are lost. Since the prediction of the angular distribution using the equations of Heiss requires a knowledge of the proton wave functions and phase shifts, the equations of Courant are used in the discussion of the angular distributions in this work.

If the photon interaction involves any pure multipole the photoproton

<table>
<thead>
<tr>
<th>Transition</th>
<th>$b/a$ ratio</th>
</tr>
</thead>
<tbody>
<tr>
<td>$s \rightarrow p$</td>
<td>$\infty$</td>
</tr>
<tr>
<td>$p \rightarrow d$</td>
<td>1.5</td>
</tr>
<tr>
<td>$d \rightarrow f$</td>
<td>1.0</td>
</tr>
<tr>
<td>$f \rightarrow g$</td>
<td>0.8</td>
</tr>
<tr>
<td>$g \rightarrow h$</td>
<td>0.75</td>
</tr>
<tr>
<td>$g \rightarrow f$</td>
<td>0.3</td>
</tr>
<tr>
<td>$f \rightarrow d$</td>
<td>0.25</td>
</tr>
<tr>
<td>$d \rightarrow p$</td>
<td>0.17</td>
</tr>
<tr>
<td>$p \rightarrow s$</td>
<td>0.0</td>
</tr>
</tbody>
</table>
angular distribution exhibits no forward-backward asymmetry and, as already indicated, for pure dipole interaction is no more complex than \( \sin^2 \theta \). The inclusion of the next higher multipole, \( E2 \), involves radiation of another parity and consequently the inclusion of \( \cos \theta \) terms. The forward-backward asymmetry, observed in the experimental results on photoproton angular distributions\(^{(82,104)}\), is then the result of dipole-quadrupole interference. The \( E1-E2 \) interference causes the angular distribution to peak forward of 90 degrees. Large forward shifts in the angular distributions of photoprotons from nuclei irradiated with bremsstrahlen of maximum energy greater than 30 MeV\(^{(93,94,95)}\) have indicated that the quadrupole absorption increases significantly with energy. The \( E1-E2 \) interference is expected to be small for photoneutrons due to their small effective charge for \( E2 \) absorption\(^{(41)}\); however, small forward-backward asymmetries have even been observed in the angular distributions of photoneutrons resulting from irradiation with 22 MeV bremsstrahlen \(^{(1,2)}\). When this forward-backward asymmetry is present, the angular distributions are fitted to the expression

\[
Y(\theta) = a + b \sin^2 \theta (1 + p \cos \theta)^2.
\]

The constant \( p \) can be used to estimate the relative \( E2 \) to \( E1 \) absorption cross sections. For combined \( E1 \) and \( E2 \) absorption by an initial \( s \)-state, Levinger\(^{(59)}\) has shown that

\[
\frac{p^2}{C_q} = \frac{5C_q}{C_d}
\]

where \( C_q \) is the direct quadrupole absorption cross section and \( C_d \) is the direct dipole absorption cross section.
Levinger (59) has indicated that equation 2.54 overestimates the value of \( p \) for other cases so that this would represent the minimum relative quadrupole absorption.

From Table 2.3 it can be seen that the only pure transition which can lead to a \( b/a \) ratio greater than 1.5 is the \( \ell=0 \) to \( \ell=1 \) transition. Empirical evidence on the angular distributions of photoprotons from nuclei between the closed proton shells at \( Z=28 \) and \( Z=50 \) (58, 63, 72, 82) have indicated anisotropies far in excess of this in spite of the fact that the excitation of \( \ell=0 \) protons in these nuclei is in contradiction to the predictions of the shell model. According to the calculations of Eichler and Weidenmüller (33), however, interference terms between the \( \ell \rightarrow \ell + 1 \) and \( \ell \rightarrow \ell - 1 \) transitions can lead to greater anisotropies than are found in the simple theory of Courant. In particular, they were able to show that if the \( p \rightarrow d \) and \( p \rightarrow s \) transitions occurred with approximately the same energy and intensity, the interference could lead to a pure \( \sin^2 \theta \) distribution. Although the calculations were for spinless particles and of the same order of approximation as Courant's equations, their results were not dependent on the type of coupling between nucleons and thus apply equally well to transitions from either an initial \( p_{1/2} \) or \( p_{3/2} \) state.

2.3.1.3 Previous results

The results of previous work on the angular distributions of photoprotons emitted by nuclei with \( Z \)-values between the magic numbers of 28 and 50 have been summarized in Table 2.4 according to proton number. The results have been presented only for experiments in which the maximum energy of the bremsstrahlen was less than 30 MeV. The angular distribution parameters, \( b/a \)
and \( p \), from equation 2.53 have been given for a number of photoproton energy ranges. The photoprotons were detected by either nuclear emulsions indicated by \( E \) or scintillation detectors in which case the phosphor has been indicated. With the exception of the results of Mitchell and McNeill\(^{(72)}\), the experiments were detailed studies of one or two elements only.

<table>
<thead>
<tr>
<th>Target Element</th>
<th>Maximum Bremsstrahlung Energy (MeV)</th>
<th>Proton Energy Range (MeV)</th>
<th>Method of Detection</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>(^{28})Ni</td>
<td>21.5</td>
<td>( \geq 3 ) 0 - E</td>
<td>Lejkin et al., 1956</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>( \geq 10 ) 1 - E</td>
<td></td>
<td></td>
</tr>
<tr>
<td>(^{28})Ni</td>
<td>25.5</td>
<td>( \geq 3 ) 0 - E</td>
<td>Lejkin et al., 1956</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>( \geq 10 ) 0 - E</td>
<td></td>
<td></td>
</tr>
<tr>
<td>(^{28})Ni</td>
<td>17.5</td>
<td>( \geq 4 ) 0 - E</td>
<td>Spicer et al., 1957</td>
<td></td>
</tr>
<tr>
<td>(^{29})Ni</td>
<td>30</td>
<td>( \geq 13 ) 0 - CsI</td>
<td>Loken, 1959</td>
<td></td>
</tr>
<tr>
<td>(^{29})Ni</td>
<td>22</td>
<td>( \geq 8 ) 0.0 # NE102</td>
<td>Mitchell et al., 1963</td>
<td></td>
</tr>
<tr>
<td>(^{29})Cu</td>
<td>24</td>
<td>3.6-10 0 - E</td>
<td>Byerly et al., 1951</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>( \geq 10 ) 3 - E</td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

\[ \# \text{p loses significance as } b \rightarrow 0 \] - not calculated

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<table>
<thead>
<tr>
<th>Target Element</th>
<th>Maximum Bremsstrahlung Energy (MeV)</th>
<th>Proton Energy Range (MeV)</th>
<th>b/a</th>
<th>p</th>
<th>Method of Detection</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{29}$Cu</td>
<td>23</td>
<td>$\geq 8$</td>
<td>2.1</td>
<td>0.5</td>
<td>ZnS</td>
<td>Mann et al., 1952</td>
</tr>
<tr>
<td>$^{29}$Cu</td>
<td>19</td>
<td>$\geq 3$</td>
<td>1</td>
<td>-</td>
<td>E</td>
<td>Lejkin et al., 1956</td>
</tr>
<tr>
<td>$^{29}$Cu</td>
<td>24</td>
<td>$\geq 3$</td>
<td>1</td>
<td>-</td>
<td>E</td>
<td>Lejkin et al., 1956</td>
</tr>
<tr>
<td>$^{29}$Cu</td>
<td>22</td>
<td>$\geq 7$</td>
<td>1.3</td>
<td>-</td>
<td>E</td>
<td>Odera et al., 1963</td>
</tr>
<tr>
<td>$^{29}$Cu</td>
<td>22</td>
<td>$\geq 8$</td>
<td>1.0</td>
<td>0.6</td>
<td>NE102</td>
<td>Mitchell et al., 1963</td>
</tr>
<tr>
<td>$^{29}$Cu$^{65}$</td>
<td>20</td>
<td>$\geq 4$</td>
<td>0.68</td>
<td>-</td>
<td>E</td>
<td>Lit'kova et al., 1960</td>
</tr>
<tr>
<td>$^{29}$Cu$^{65}$</td>
<td>24.5</td>
<td>$\geq 4$</td>
<td>0.76</td>
<td>0.40</td>
<td>E</td>
<td>Lit'kova et al., 1960</td>
</tr>
</tbody>
</table>

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Table 2.4 (continued)

<table>
<thead>
<tr>
<th>Target Element</th>
<th>Maximum Bremsstrahlung Energy (MeV)</th>
<th>Proton Energy Range (MeV)</th>
<th>Method of Detection</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{30}_{\text{Zn}}$</td>
<td>20.8</td>
<td>$\geq 3$ 0.4 0</td>
<td>E</td>
<td>Osokina et al., 1957</td>
</tr>
<tr>
<td>$^{30}_{\text{Zn}}$</td>
<td>23.3</td>
<td>$\geq 3$ 0</td>
<td>E</td>
<td>Osokina et al., 1957</td>
</tr>
<tr>
<td>$^{30}_{\text{Zn}}$</td>
<td>22</td>
<td>$\geq 8$ 1.1 0.5</td>
<td>NE102</td>
<td>Mitchell et al., 1963</td>
</tr>
<tr>
<td>$^{40}_{\text{Zr}}$</td>
<td>25</td>
<td>$\geq 4$ 1</td>
<td>E</td>
<td>Dushkov et al., 1964</td>
</tr>
<tr>
<td>$^{40}_{\text{Zr}}$</td>
<td>22</td>
<td>$\geq 8$ 3.1 0.2</td>
<td>NE102</td>
<td>Mitchell et al., 1963</td>
</tr>
<tr>
<td>$^{41}_{\text{Nb}}$</td>
<td>19.5</td>
<td>$\geq 3$ 0.75</td>
<td>E</td>
<td>Osokina, 1963</td>
</tr>
<tr>
<td>$^{41}_{\text{Nb}}$</td>
<td>23.5</td>
<td>$\geq 3$ 0.54</td>
<td>E</td>
<td>Osokina, 1963</td>
</tr>
<tr>
<td>$^{41}_{\text{Nb}}$</td>
<td>22</td>
<td>$\geq 8$ 3.0 0.1</td>
<td>NE102</td>
<td>Mitchell, 1963</td>
</tr>
</tbody>
</table>

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<table>
<thead>
<tr>
<th>Target Element</th>
<th>Maximum Bremsstrahlung Energy (MeV)</th>
<th>Proton Energy Range (MeV)</th>
<th>b/a</th>
<th>Method of Detection</th>
<th>Reference</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{42}$Mo</td>
<td>22</td>
<td>$\geq 8$</td>
<td>1.2</td>
<td>0.2</td>
<td>NE102</td>
</tr>
<tr>
<td>$^{92}$Mo</td>
<td>22.5</td>
<td>3-5</td>
<td>0</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td></td>
<td></td>
<td>5-7</td>
<td>0</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td></td>
<td></td>
<td>7-9</td>
<td>1</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td>$^{100}$Mo</td>
<td>22.5</td>
<td>$\geq 9$</td>
<td>1</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td>$^{47}$Ag</td>
<td>20.8</td>
<td>4-7</td>
<td>0</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td></td>
<td></td>
<td>7-10</td>
<td>0</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$\geq 10$</td>
<td>3</td>
<td>-</td>
<td>E</td>
</tr>
<tr>
<td>$^{47}$Ag</td>
<td>30</td>
<td>$\geq 13$</td>
<td>2</td>
<td>-</td>
<td>CaI</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
<td>Lokan, 1959</td>
</tr>
<tr>
<td>$^{47}$Ag</td>
<td>22</td>
<td>$\geq 8$</td>
<td>1.7</td>
<td>0.2</td>
<td>NE102</td>
</tr>
<tr>
<td>$^{48}$Cd</td>
<td>22</td>
<td>$\geq 8$</td>
<td>0.6</td>
<td>0.6</td>
<td>NE102</td>
</tr>
<tr>
<td>$^{49}$In</td>
<td>22</td>
<td>$\geq 8$</td>
<td>1.5</td>
<td>0.4</td>
<td>NE102</td>
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<tr>
<td>$^{50}$Sn</td>
<td>22</td>
<td>$\geq 8$</td>
<td>0.1</td>
<td>-</td>
<td>NE102</td>
</tr>
</tbody>
</table>

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2.3.2 Energy distributions

2.3.2.1 Statistical de-excitation

Assuming compound nucleus formation, Blatt and Weisskopf(8) employed the Bohr assumption and the reciprocity theorem to derive an expression for the relative energy distribution of particles b resulting from the absorption of monoenergetic photons of the form

\[ I(\varepsilon_b)d\varepsilon_b = \text{const.} \frac{2m_b}{\hbar^2} \varepsilon_b \sigma_c(\varepsilon_b) W_x(\varepsilon_y)d\varepsilon_b \]  \hspace{1cm} 2.55

where \( I(\varepsilon_b)d\varepsilon_b \) is the relative number of particles with energy between \( \varepsilon_b \) and \( \varepsilon_b + d\varepsilon_b \),

\( \varepsilon_y \) is the excitation energy of the residual nucleus which equals \( (E_y - B_b - \varepsilon_b) \),

\( E_y \) is the energy of the absorbed photon,

\( B_b \) is the binding energy of a particle b in the nucleus,

\( \sigma_c(\varepsilon_b) \) is the capture cross section for particles of energy \( \varepsilon_b \) on the residual nucleus of excitation energy \( E_y \),

\( W_x(\varepsilon_y) \) is the level density of the residual nucleus with excitation energy \( E_y \) and

\( m_b \) is the mass of the particle.

The level density \( W_x(\varepsilon_y) \) is a strongly increasing function of the excitation energy of the residual nucleus which can be written in terms of the entropy \( S(E_y) \) of the residual nucleus according to the relation(8)

\[ S(E_y) = \log W_x(\varepsilon_y). \]  \hspace{1cm} 2.56

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The entropy can be expanded in a Taylor series about the maximum energy of the residual nucleus, $E_\gamma - B_b$, with the result

$$ S(E_\gamma - B_b - \varepsilon_b) = S(E_\gamma - B_b) - \varepsilon_b \left( \frac{dS}{dE} \right)_{E = E_\gamma - B_b} $$

2.57

Since $S(E_\gamma)$ is the entropy of the residual nucleus we may write

$$ \left( \frac{dS}{dE} \right)_{E_\gamma - B_b} = \frac{1}{T(E_\gamma - B_b)} $$

2.58

where $T(E_\gamma - B_b)$ can be interpreted as the nuclear temperature of the residual nucleus at the excitation energy $E_\gamma - B_b$.

The level density of the residual nucleus can then be written in the form

$$ W_\nu(E_\gamma) = \text{const.} \exp(-\varepsilon_b/T) $$

2.59

Livesey(62) has pointed out that the experimental data on level densities can best be fitted particularly over the first 10 MeV by assuming that the nuclear temperature is both independent of the excitation energy of the residual nucleus and given by the relation

$$ T = \left( \frac{100}{A} \right)^{1/2} $$

2.60

where $A$ is the atomic mass number.

The constant in equation 2.59 is then chosen to fit the conditions at 10 MeV.
The relative energy distribution of particles evaporated from the compound nucleus is then

\[ I(\varepsilon_b) d\varepsilon_b = \text{const.} \frac{2m_b}{h^2} \varepsilon_b \left( \frac{\varepsilon_b}{T} \right) \exp\left( -\frac{\varepsilon_b}{T} \right) d\varepsilon_b. \]  

2.61

The cross section for the formation of the compound nucleus by the inverse process can be calculated on the basis of continuum theory. In the case of neutrons \( \sigma_c(\varepsilon_n) \) is found to be a slowly varying function of the neutron energy and may be considered constant to a first approximation(22). The energy distribution of the statistical neutrons then becomes

\[ I(\varepsilon_n) d\varepsilon_n = \text{const.} \frac{2m_n}{h^2} \varepsilon_n \exp\left( -\frac{\varepsilon_n}{T} \right) d\varepsilon_n. \]  

2.62

If equation 2.62 is differentiated with respect to the neutron energy it can be seen that the energy distribution of the neutrons peaks at the nuclear temperature.

The average value of the emission probability per unit time of a particle, \( b \), from a particular state excited by monoenergetic photons is proportional to the integral of \( I(\varepsilon_b) d\varepsilon_b \) over the energy spectrum of the particle and is given by

\[ \frac{\Gamma_b}{\hbar} = \frac{2m_b}{h^2} \varepsilon_b \max \int_0^{\varepsilon_b \max} \sigma_c(\varepsilon_b) W_x(\varepsilon_y) d\varepsilon_b. \]  

2.63

For the continuous spectrum of photons used in this experiment the integral must be weighted with the bremsstrahlung spectrum. The energy
distribution of protons per nucleus resulting from irradiation of a target with bremsstrahlung then becomes (29)

\[ F(\varepsilon_p) \, d\varepsilon_p = \frac{2m_p}{\hbar^2} \varepsilon_p \, \sigma_c(\varepsilon_p) \, d\varepsilon_p \int \frac{\gamma(E)N(E,E_{\text{max}})\Gamma(E)}{\sum_b \Gamma_b} \, dE \]

where \( F(\varepsilon_p) \, d\varepsilon_p \) is the number of protons with energy between \( \varepsilon_p \) and \( \varepsilon_p + d\varepsilon_p \),

\( E_{\text{max}} \) is the maximum bremsstrahlung energy,

\( \sigma_c(E) \) is the cross section for the formation of the compound nucleus by the absorption of photons of energy \( E \),

\( N(E,E_{\text{max}}) \) is the number of photons of energy \( E \) per unit energy interval per roentgen in the bremsstrahlung spectrum of maximum energy \( E_{\text{max}} \) and

\( \sum_b \) is the summation over all the modes of disintegration of the compound nucleus.

Then since

\[ \frac{\sigma_c(E)}{\sum_b \Gamma_b} = \frac{\sigma_{\gamma,p}(E)}{\Gamma_p} = \frac{\sigma_{\gamma,n}(E)}{\Gamma_n} \]

the first of these equivalents can be replaced with either of the others.

The (\( \gamma \), n) cross section is usually the best known from experiment for most targets so that the above equation may be rewritten

\[ F(\varepsilon_p) \, d\varepsilon_p = \frac{2m_p}{\hbar^2} \varepsilon_p \, \sigma_c(\varepsilon_p) \, d\varepsilon_p \int \frac{\sigma_{\gamma,n}(E)N(E,E_{\text{max}})\Gamma(E)}{\Gamma_n} \, dE \]
where $\sigma_{\gamma,n}$ is the cross section for the $(\gamma,n)$ reaction.

Since the function $\mathcal{C}_c(\varepsilon_n)$ can be neglected the expression for $\Gamma_n$ upon substituting for $W_r(E_\gamma)$ assumes the form

$$
\Gamma_n = \frac{2m^3}{\hbar^2} \int_0^{\varepsilon_n \text{max}} xe^{-x/T} \, dx
$$

which is a slowly varying function and can be considered constant. Thus, the equation for the energy distribution of protons evaporated by a nucleus excited with bremsstrahlen reduces to the form (26)

$$
F(E_p) dE_p = \text{const.} E_p \mathcal{C}_c(\varepsilon_p) e^{-\varepsilon_p/T} dE_p \int_0^{E_{\text{max}}} \sigma_{\gamma,n}(E) N(E,E_{\text{max}}) dE
$$

2.3.2.2 "Resonance direct" emission

In order to calculate the energy distribution of photoprotons emitted in a "resonance direct" interaction (109) the contributions of the various possible I.P.M. transitions and the energies of the emitted protons must be individually calculated and summed. Since the probability of a proton being directly emitted depends on its ability to penetrate the potential barrier of the nucleus, the emission strength of direct protons from any I.P.M. transition is proportional to the product of the absorption strength for the transition and the transmission coefficient for protons from the excited shell-model state. In the case of irradiation with bremsstrahlen the strength of the "resonance direct" emission must also be weighted by the photon flux at the transition energy, as well as the transition energy.
The transition energies can be readily calculated from the separations between the various energy levels of the I.P.M. Many authors\[^{18,71}\] employ the energy scale shown in Figure 2.2 which was calculated by Schröder\[^{89}\] with spin-orbit coupling. The separations between the unoccupied levels of the nucleus must be increased by a factor of two in order to correct for the effective mass of the nucleon within the nucleus\[^{109}\]. The energy of the protons emitted by any I.P.M. transition is then the energy of the photon absorbed in the transition less the binding energy of the proton in the nucleus. For transitions from the uppermost occupied level the binding energy of the proton is equal to the separation energy\[^{37}\]. In the case of transitions from a lower level the binding energy is the sum of the proton separation energy and the energy difference between this level and the uppermost occupied level. Since the various shells of the I.P.M. are considered to be spread out over several MeV\[^{109}\] to allow for the interaction effects between nucleons and the uncertainty in the positioning and ordering of the shells, the protons from a given I.P.M. transition are not expected to be monoenergetic but rather are smeared out over a number of MeV about the above energy.

2.3.3 Yield of photoprotons

2.3.3.1 Statistical de-excitation

The yield of photoprotons predicted by the statistical theory of de-excitation can be readily calculated using the Bohr assumption which for \((\gamma, p)\) reactions can be written in the form

\[
\sigma_{\gamma,p}(E) = \xi \sigma(\gamma E) G_p
\]

2.69

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FIGURE 2.2. Schematic Diagram of the Proton Energy Levels

The energy scale was calculated by Schröder(89) with spin-orbit coupling. In the I.P.M. of Wilkinson(109) the separations between the unoccupied levels of the nucleus must be increased by a factor of two in order to correct for the effective mass of the nucleon within the nucleus.
where the branching ratio, $G_p$, is given by

$$G_p = \frac{\Gamma_p}{\sum_b \Gamma_b} \quad . \quad 2.70$$

The expression for $\Gamma_b$ can be written from equation 2.63 as

$$\Gamma_b = \frac{2m_b}{x^2} \int_{E_b}^{E_b \text{ max}} \varepsilon_b \Gamma_c(E_b) \frac{\varepsilon_x(E_y)}{dE_b} \quad . \quad 2.71$$

The photoproton yield per mole per roentgen is then

$$Y_p = N_0 \int_{E_p}^{E \text{ max}} \varepsilon(E) N(E, E_p) dE \quad . \quad 2.72$$

$$= N_0 \int_{E_p}^{E \text{ max}} \varepsilon(E) \frac{\Gamma_p}{\sum_b \Gamma_b} N(E, E_{\text{max}}) dE$$

where $N_0$ is Avogadro's number.

This expression has been evaluated by Weinstock and Halpern(103) as a function of Z-value. In computing the branching ratios it was assumed that protons and neutrons were emitted only. The $\Gamma$ functions were calculated for the photon energy corresponding to the peak position of the giant resonance and consequently could be removed from the integral. The binding energies were calculated from the mass formula and the level densities were assumed to be exponential. The bremsstrahlung distribution was also removed from the integral by using the intensity at the resonance energy. The
integrated cross section was then evaluated using the dipole sum rule of Levinger and Bethe(60), equation 1.3, with an exchange term of 0.4. The results of these calculations are indicated by the solid curve in Figure 2.3. The experimental results shown are those of Mann and Halpern (64), Diven and Almy(29), Butler and Almy(15), Toms and Stephens(101,102,103,104) and Weinstock and Halpern(108). The agreement in absolute yield between the experimental results and the predictions of the statistical theory is excellent in the region of Z between 28 and 50. The experimental photoproton yields are, however, definitely in excess of the statistical predictions for nuclei of high Z-value.

2.3.3.2 "Resonance direct" emission

The yields of "resonance direct" photoprotons have been calculated for nuclei of high Z-value by Wilkinson(109). In these calculations the probability C, equation 2.49, that the particle emerges from the single-particle component of the wave function rather than loses its energy in collective motion was evaluated for each of the possible I.P.M. transitions. The capture width, 2D, leading to compound nucleus formation was set equal to 3 MeV and the penetrability, P, was calculated from the appropriate Coulomb wave functions. The relative probabilities of "resonance direct" emission were first weighted by the strength of the absorption in each shell and then, since each of the shells was considered to be spread out over a range of 3 MeV, were summed over the appropriate depths in the nucleus. The results of these calculations which are shown by the broken curve in Figure 2.3 are in good agreement with the experimental data for high Z-values.
FIGURE 2.3. Yields of Photoprotons

The solid curve indicates the photoproton yields predicted by the statistical theory\cite{108}. The yields of "resonance direct" photoprotons are shown by the broken curve\cite{109}. The experimental points are taken from references which are listed in the text.
2.4 Particle-Hole Interactions

From the introduction it can be recalled that the main difficulty with the single-particle model is the prediction of a giant resonance energy which is about one half of that experimentally observed. In order to explain this discrepancy Wilkinson(109) introduced a velocity-dependent potential into the single-particle model which gave rise to an effective nucleon mass inside the nucleus of about one half of its normal mass. The energy spacing between single-particle levels belonging to neighbouring shells would then be increased by a factor of two. The positions of the single-particle levels which have been found more recently through (d,p) experiments(20,21) are, however, in conflict with those required by Wilkinson. This difficulty has been resolved by the suggestion that particle-hole interactions can shift the energy of the dipole state quite considerably from the single-particle value(11).

In the particle-hole nomenclature a single-particle excitation is called a particle-hole state. Such a state has one more particle above and one more hole below the Fermi level than the ground state. The particle-hole calculations require the usual central potential of the I.P.M. plus a residual interaction between pairs of nucleons (or holes) which cannot be readily reduced to a central potential. The Hamiltonian for closed-shell nuclei can then be written as the sum of the single-particle Hamiltonian, $H_0$, and the potential, $V$, of the residual particle-hole interaction(96) in the form

$$H = H_0 + V \quad 2.73$$
where \( H_0 = \sum_{i=1}^{A} H_i = \sum_{i=1}^{A} \left( \frac{p_i^2}{2m} + V_i \right) \) and

\[
V = \sum_{i < j} V_{ij}.
\]

The particle-hole interaction leads to a mixture of pure shell-model states so that the wave function of the excited level can be represented by the form (3)

\[
\psi = \sum_{\tilde{d}, \tilde{d}'} \phi_{\tilde{d}\tilde{d}'} C_{\tilde{d}\tilde{d}'}
\]

where \( \tilde{d}, \tilde{d}' \) are the final particle and hole states respectively,

- \( \phi_{\tilde{d}\tilde{d}'} \) give the mixing of the single-particle wave functions and
- \( C_{\tilde{d}\tilde{d}'} \) are the pure single-particle wave functions of the central potential.

The mixing of the single-particle excitations is restricted to transitions within one oscillator spacing of the Fermi level. In order to determine the perturbed dipole state energies, \( E_\nu \), it is necessary to diagonalize the secular matrix which has the form (11)

\[
(E - \varepsilon_{c'c}) \delta_{cd} \delta_{c'd'} = \sum_{\tilde{d}, \tilde{d}'} V_{cc'} \tilde{d}\tilde{d}' C_{\tilde{d}\tilde{d}'}
\]

where \( \nu, \nu' \) are the initial particle and hole states respectively,

- \( \varepsilon_{c'c} \) are the unperturbed energies of the single-particle excitations i.e. the energies of the dipole excitations of Wilkinson and
- \( V_{cc'} \tilde{d}\tilde{d}' \) are the matrix elements of the particle-hole interaction.
In the particle-hole calculations the energies of the single-particle states, calculated on the basis of the shell-model potential, are commonly replaced by the experimental values of the respective levels of the adjacent nuclei, which are determined by pickup and stripping reactions. Since the polarizability of the core by the hole is assumed to be only slightly affected by the presence of the particle, the nucleon interactions in the \( A + 1 \) and \( A - 1 \) nuclei can be applied to the nucleus of \( A \) particles.

The schematic model originally proposed by Brown and Bolsterli\(^{(11)}\) assumes both that the participating EL transitions are almost degenerate in energy and that the particle-hole interaction matrix element, \( V_{cc'dd'} \), is separable and of the form

\[
V_{cc'dd'} = \lambda D_{cc'} D_{dd'}
\]  

where \( \lambda \) is a universal strength constant and

\[
D_{cc'}, D_{dd'}
\]

are the matrix elements of the dipole operator between the ground state and the configurations \( cc' \) and \( dd' \) respectively.

When the secular equation is solved for the above degenerate case it is found that one state, called the dipole state, is considerably shifted in energy and that this state collects the total dipole transition strength. Since the assumptions on which the schematic model is based are not strictly reproduced by the nuclei, its predictions do not show up in the results of actual computations with realistic models. This is particularly true in the case of light nuclei where the initial assumption of the Brown-Bolsterli scheme is not expected to hold.

Extensive calculations have been performed on the light nuclei.
by Gillet(44) who found that the energy shifts of the states were small compared to the spread of the zero-order energies so that no single state was pushed away from the remainder as is the case in the schematic model. The prediction of the schematic model regarding dipole transition strength was fulfilled, however, except that two states rather than one carried practically the complete dipole strength. The splitting of the dipole state was interpreted as the effect of spin-orbit coupling. In the calculations of Balashov et al. (3) the correct resonance energy in \( ^{40}\text{Ca} \) was found by relying upon the empirical data on the nuclear levels of neighbouring nuclei with \( A=39 \) and \( A=41 \). The diagonal part of the particle-hole interaction increased the excitation energy by two or three MeV only while the off-diagonal elements, responsible for the mixture of configurations, did not lead to an energy shift. Although the large shift in energy of one state was not observed, the prediction of the schematic model with regard to the concentration of the dipole strength was fulfilled. Similar results to these have been obtained in other calculations for \( ^{12}\text{C} \) (107), \( ^{16}\text{O} \) (12), \( ^{16}\text{O} \) (12). In the case of the light nuclei the calculations have given good agreement not only for the position of the giant dipole resonance and its integrated value but also for the characteristics of the nucleon spectra(3).

At the other end of the periodic table similar calculations performed on \( ^{208}\text{Pb} \) (4) indicated that most of the dipole strength was collected in the highest energy state. Whereas the off-diagonal elements of the particle-hole interaction caused almost no energy shift in the light nuclei, these calculations indicated that they were responsible for the large energy shift which was required for agreement with the experimental giant resonance energy. Since the energy shift of the dipole state was dependent on the type of...
nucleon-nucleon interaction assumed, a high density of single-particle levels was found not to be a sufficient condition for the appearance of the strongly correlated dipole state of Brown-Bolsterli.

In the medium-weight range of elements, of particular interest in this thesis, particle-hole calculations have been performed on Zr$^{90}$(96). Since the same nucleon-nucleon interaction shifted the energy of the dipole state into agreement with the experimental giant resonance energy in Zr$^{90}$ as in the calculations on Pb$^{208}$, it was concluded that the particle-hole interaction plays qualitatively the same role in the medium-weight nuclei as in the heavy nuclei. The particle-hole calculations in the medium to heavy nuclei have correctly predicted only the position of the giant resonance.

The above discussion is concerned with doubly closed-shell nuclei only. The particle-hole calculations for non-magic nuclei are quite difficult and are available for only a few special cases(32,77).
CHAPTER III

EXPERIMENTAL DETAILS

Chapter III describes the experimental arrangement and procedures employed in this work to determine the yields and the energy and angular distributions of photoprotons emitted by several medium-weight nuclei under bombardment by high-energy photons.

3.1 Introduction

Up to the present time several methods have been employed in the detection of photoprotons including nuclear emulsions (101), cloud chambers (6), scintillation counters (64) and more recently solid state detectors (75, 87, 105). Any such detection system must be able to discriminate between the photoprotons and the intense background of scattered photons, secondary electrons and neutrons produced by ($\gamma$,n) reactions in the shielding material. This problem is made more difficult by the pulsed nature of the betatron which, in this experiment, had a pulse repetition frequency of 180 per second and a beam length of 4 microseconds. The background per second was consequently confined to the 720 microseconds of beam time per second. The scattered radiation from the abundant low-energy photons of the continuous bremsstrahlung beam added further to the background within which the photoprotons had to be detected.

The solid state detector offers many advantages in the detection of photoprotons including excellent energy resolution and linearity of response, as well as a fast rise time, very compact size and a relatively low sensitivity to photons and neutrons. As a result of its small size the detector can be

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mounted directly inside the target chamber and further can be housed in a shielding material both to collimate the charged particles from the target and to prevent ionizing background radiations from impinging on the detector in the direction of the long dimensions of the sensitive volume in which case they would deposit more energy. Since the complete detection system is confined to the target chamber the overall shielding problem is reduced considerably.

The main prerequisite for a photoproton detection system is a relatively low sensitivity to electrons, photons and neutrons compared with protons. The response of a solid state detector depends only on the energy deposited by the particle in the depletion layer, is linear with the amount of energy deposited and is completely independent of the type of charged particle which deposits the energy; consequently, the background radiations with their low specific ionization deposit much less energy and give smaller pulses than the protons. In order to obtain the maximum advantage from this fact, the depletion layer of the solid state detector can be readily adjusted to the depth corresponding to the range of the maximum energy photoprotons to be detected by applying a suitable reverse-bias voltage across the detector.

To prevent a number of electron pulses from piling up in the detector, within its resolving time, the rise time of the detector and the associated electronics should be kept as short as possible. The fast rise time of the solid state detector makes it a very desirable detection system in this respect.

A difficult problem in photoproton work, in which a direct counter technique has been used, is the adequate energy calibration of the detection system. With solid state detectors, the linearity of response and the independence of the response on particle type completely resolve this difficulty, since an alpha source with two or three peaks can be used as an

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absolute energy calibration.

Since the solid state detector has so many features suitable for photoproton work it was chosen for use in this experiment. In the remainder of this chapter the complete detection system and experimental procedures are discussed in detail.

3.2 Experimental Arrangement

3.2.1 General

A schematic diagram of the experimental arrangement is shown in Figure 3.1. The bremsstrahlung beam from the Princess Margaret Hospital's betatron, with a peak energy of 22 MeV, was collimated with a lead insert in the front of the machine to an angular diameter of 0.015 radians, or a lateral diameter of 1.3 centimeters, at the photoproton target position. Thin metallic foils of natural isotopic abundance were mounted as targets in the centre of an evacuated cylindrical target chamber approximately 115 centimeters from the internal bremsstrahlung target of the betatron. The photoprotons, emitted by the target elements under bombardment by the high-energy photon beam, were detected in solid state surface barrier detectors mounted in a lead collimator and shield which was free to rotate, under vacuum, about the target position. The detection system was calibrated using the four alpha peaks from the decay of Ra$^{226}$ and its daughter products.

3.2.2 Target chamber

The essential features of the target chamber are shown in Figure 3.2. The chamber was of a cylindrical brass construction with a 15.5 inch outside
FIGURE 3.1. Schematic Diagram of the Experimental Arrangement

The following notation is used:

D - betatron donut
B - bremsstrahlung target
M - monitor
S - lead shielding well
T - target chamber
A - detector assembly
The target chamber was of a cylindrical brass construction with a 15.0 inch inside diameter, ID, and a 0.25 inch wall thickness, WT. The interior of the chamber, which was lined with 0.25 inches of lucite, L, to reduce the electron background, was 7.5 inches in height, H. The chamber was fitted with collinear entrance and exit beam ports, EP, of 3.0 inch outside diameter, each of which was O-ring sealed to a thin window of 0.005 inch mylar, M. An additional viewing port, VP, set in the wall of the chamber at 90 degrees to the beam ports, was used for vertical alignment. The chamber, which could be evacuated through the nozzle, N, was mounted on three supports, S, which were threaded for leveling and vertical adjustment.
diameter and a 0.25 inch wall thickness. The interior height of the target chamber was 7.5 inches. The large chamber volume minimized the scattering of charged particles from the walls of the chamber into the sensitive area of the detectors. In order to reduce this source of background further the interior of the chamber was lined with 0.25 inches of lucite, which has a low effective atomic number and therefore low photoelectric and pair production cross sections.

The chamber was fitted with collinear entrance and exit beam ports of 5.0 inch outside diameter, each of which was O-ring sealed to a thin window of 0.005 inch mylar. An additional viewing port, set in the wall of the chamber at 90 degrees to the beam ports and opposite the detector mount, was used for vertical alignment of the bremsstrahlung beam by means of a light system mounted on the far end of the defining collimator. The three supports, on which the target chamber was mounted, were threaded for leveling and vertical adjustment of the chamber.

The lid of the chamber was vacuum sealed to the chamber wall with a large diameter O-ring. With the lid in position the chamber could be evacuated, by means of a mechanical pump to a pressure of about 1.0 millimeters, as measured with a Vacustat vacuum gauge. The lid of the target chamber which supported the detector housing and collimation assembly, the detector angular-position indicator, the adjustable target support, the adjustable Ra$^{226}$ alpha-particle source support and housing assembly and the preamplifier, is shown in Figure 3.3. The two surface barrier detectors, used for photoproton detection, were mounted in a hard rubber support. This support was further mounted in a thin-walled brass container to which one side of each detector was electrically connected. This complete assembly was then inserted into the lead housing, shown in Figure 3.4, which collimated...
FIGURE 3.3

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FIGURE 3.3. The Assembly of the Target Chamber Lid

The lid of the target chamber supported the detector housing and collimation assembly, D, the detector angular-position indicator, I, the adjustable target support, T, the adjustable Ra\textsuperscript{226} alpha-particle source support, S, the Ra\textsuperscript{226} source housing assembly, H, and the preamplifier, P. The detector housing assembly, which was counterbalanced by the caster, C, could be rotated about the target through 150 degrees under vacuum.
FIGURE 3.4. Lead Housing Assembly for the Detectors

The detectors were inserted into the lead housing assembly which collimated the photoprotons originating in the target to a circle of diameter 0.31 inches at the detector position. The lead housing greatly reduced the number of electrons impinging on the sensitive volume of the detectors during the beam pulse of the betatron. All dimensions are indicated in inches.
the photoprotons originating in the target to a circle of diameter 0.31 inches at the detector position 3.8 inches away. The lead housing greatly reduced the probability of electrons traversing a diameter of the detector, in which case, the electrons could give pulses of magnitude comparable to those of the protons. The lead housing support could be rotated under vacuum about the target position by means of a brass tube through the centre of the target chamber lid. The useful angle of rotation of the detectors with respect to the photon beam was purposely limited by the lead housing assembly to 45 degrees in the forward direction and 145 degrees in the backward direction. The detector angular-position indicator was a long brass pointer rigidly fixed to the top end of the supporting brass tube. The pointer rotated across a circular scale drawn on the top of the lid to indicate the angle of the detectors with respect to the direction of propagation of the bremsstrahlung beam.

The target holder shaft passed concentrically through the brass tube, centrally located in the chamber's lid for supporting the lead housing assembly containing the detectors. A double O-ring vacuum seal allowed the target to be both rotated in the beam and withdrawn from the beam for "no target" background runs.

A retractable rod, mounted through the lid of the chamber directly in front of the viewing port, allowed the Ra$^{226}$ source to be lowered, under vacuum, onto a horizontal line through the target position to the detection system for energy calibration. When the source, which was approximately 10 inches from the detector, was not being used for this purpose it was withdrawn into its cylindrical brass tube to prevent alpha particles from impinging on the detector.

The preamplifier was mounted directly on the lid of the chamber in
order to keep the connecting wire from the detectors as short as possible. A fine insulated wire lead was connected from the detectors through a Kovar glass-to-metal seal to the preamplifier. This wire supplied the reverse-bias voltage for the detectors and carried the signal from both detectors to the preamplifier. The necessity of a second wire lead was eliminated by running the chassis of the preamplifier, the target chamber, the lead housing support and the brass container for the detectors all at ground potential.

3.2.3 Targets

The targets were in the form of metallic foils of natural isotopic abundance with dimensions of 1.0 centimeters by 5.0 centimeters. The thickness of the foils, which was a compromise between the minimum acceptable yield and the maximum acceptable energy loss by a proton in traversing one-half the thickness of the target foil (half-value thickness), varied from 13 mg/cm² to 45 mg/cm². The larger thicknesses were employed for targets of higher Z-value for which the specific ionization(67) is less so that the effective spread in half-value thickness for protons of 3 MeV varies in the range from 0.6 to 1.2 MeV with seven of the targets in the range from 0.6 to 0.9 MeV. Nine target elements were studied in all with Z-values from 23 to 50. Nuclei with adjacent Z-values were studied when this was possible.

The target was mounted vertically on the end of the target holder shaft and was rotated with the detector assembly. The target was normally set perpendicular to the central axis of the detectors except at the 90 degree position of the detectors with respect to the photon beam where the target was set at 4 degrees to the beam. In this manner the half-value thickness correction was minimized and was not a function of the angle of detection.
The attenuation of the photon beam was also kept to a negligible value.

A slight tension was placed on the target by hanging a small weight on its lower end. Since there were no supports down the sides of the target, the photon beam was made larger in diameter than the width of the target. The irradiated area of the target was consequently independent of its orientation in the photon beam. The target could be removed from the photon beam for background measurements by withdrawing the target-holder shaft to its secondary position. The lid of the target chamber could be removed for changing the target.

3.3 Surface Barrier Detectors

3.3.1 Interaction mechanism

When an energetic charged particle passes through a semiconductor, inelastic collisions with the electrons excite these electrons into the nominally unoccupied conduction bands thereby creating holes in the nominally full valence bands. Interactions among the electrons and holes cause the electrons to fall to the bottom of the conduction band and the holes to rise to the top of the valence band within about $10^{-12}$ seconds. During this energetic Auger de-excitation process many more electrons and holes are created. On the average one electron-hole pair is created for each 3.6 electron volts of energy lost by the incident particle in silicon\(^{(70)}\). This quantity, which is independent of the particle type\(^{(39)}\), is approximately three times the 1.1 electron volt band-gap energy in silicon. According to Shockley\(^{(97)}\) the extra energy is wasted through strong coupling of the electrons to lattice vibrations of the solid. The final step in the
de-excitation process is the recombination of the electrons and holes. It is
of paramount importance in semiconductor radiation detectors that the
recombination time, commonly called lifetime, is not too short.

3.3.2 Characteristics

In the surface barrier detectors a thin p-type surface layer is
allowed to form spontaneously by the oxidation of a chemically etched surface
of high-purity n-type body material. Contact is made to the required
sensitive area of this inversion layer with a thin evaporated gold film and
to the n-type silicon through an "ohmic" contact. The application of a
reverse bias across the n-p junction establishes a depletion layer or space-
charge region on both sides of the junction in which the electric field
falls linearly from a maximum value at the junction to approximately zero at
the outside edges of the depletion layer(27).

The thickness of the depletion layer is then

\[ x = x_n + x_p \]  \hspace{1cm} 3.1

where \( x_n \) is the thickness of the depletion layer in the n-type material
and

\( x_p \) is the thickness in the p-type material.

If Poisson's equation is applied to the space-charge regions, it can be
shown that the thickness of the depletion layer in the two regions is given
by

\[ x_{n(p)} = \left[ \frac{k(V_o + V)}{2\pi eN_d(a)} \right]^{1/2} \]  \hspace{1cm} 3.2
where $K/4\pi$ is the rationalized permittivity of silicon, $V_0$ is the potential barrier at zero bias, $V$ is the applied reverse-bias voltage and $N_{d(a)}$ is the concentration of donor (acceptor) impurity atoms on the n(p) sides of the junction respectively.

Since the space-charge region must contain an equal number of positive donors and negative acceptors, it is almost completely in the n-type region on account of its low relative density of impurity atoms.

The capacitance of the junction arises from the existence of an insulating layer of width $x_n$ and rationalized permittivity $K/4\pi$ between conducting regions so that the diode capacitance per unit area is

$$C_d = \frac{K}{4\pi x_n}.$$  

The resistivity of the silicon is related to the impurity concentration by the equation

$$\rho = \frac{1}{Nue}.$$  

where $\rho$ is the resistivity and $\mu$ is the mobility of the majority carriers.

The depth of the depletion layer and the capacitance in a reverse-biased diode thus depend on the density of uncompensated carriers and consequently on the resistivity. The depth of the depletion layer is directly proportional to and the capacitance is inversely proportional to the square root of the
product of resistivity and bias voltage.

Blankenship and Borkowski(7) have prepared a useful nomograph for the graphical solution of $x_n$ and $C_d$. This nomograph, which is shown in Figure 3.5, also indicates the maximum energies of charged particles that will just be stopped in a given thickness of depletion layer.

### 3.3.3 Collection of charge

The electrons and holes created in the depletion layer by the primary charged particle drift in the electric field established across the depletion layer by the reverse-bias voltage on the diode and give rise to an external current. The time integral of the current is used as a measure of the energy deposited in the detector by the primary particle.

The contribution of each carrier to the time integral of the current is derived as follows. The electric field gives energy, $E$, to a carrier at a rate

$$\frac{dE}{dt} = e\, \vec{v} \cdot \vec{E}$$

where $\vec{v}$ is the carrier velocity and $\vec{E}$ is the local electric field.

This energy is extracted from that stored in the charged condenser so that

$$\frac{dE}{dt} = \frac{d\left(\frac{Q^2}{2C}\right)}{dt} = \frac{Q}{C} \, i = VI$$

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FIGURE 3.5. Solid State Detector Nomograph.

The barrier depth, \( E_D \), in microns, and the capacitance per \( \text{cm}^2 \), \( C \), in picofarads, of a surface barrier detector of resistivity, \( R \), in ohm-cm, are determined by drawing a straight line from the resistivity of the N-type silicon in the detector to the applied reverse-bias voltage, \( E \), in volts. The barrier depth and capacitance can then be read at the appropriate junction. In order to determine the maximum energy of electrons, \( E \), protons, \( P \), or alphas, \( A \), which can just be stopped in the barrier depth, a line must be extended horizontally from the above junction to the appropriate range-energy curve, \( R-E \), in MeV.
where $V$ is the total reverse-bias voltage across the diode. Thus

$$i = \frac{eV \cdot \bar{E}}{V} \quad \text{(3.7)}$$

which may be rewritten as

$$i = \left(\frac{e}{V}\right) \frac{\Delta V}{dt} \quad \text{(3.8)}$$

The time integral of the current due to the motion of one carrier through the fraction $\Delta V/V$ of the total potential across the diode then becomes

$$\int i \, dt = \frac{e}{V} \Delta V \quad \text{(3.9)}$$

where $\Delta V/V$ is the effective charge reaching the electrode in units of the electronic charge.

Thus the time integral of the current for an electron-hole pair is just one electronic charge independent of the position in the depletion layer in which the pair was produced. Recalling that 3.6 electron volts are required to produce one electron-hole pair in silicon the number of pairs produced by a particle losing energy $E$ MeV in the depletion layer is seen to be

$$N = \frac{E}{3.6} \times 10^6 \quad \text{(3.10)}$$
Since the collection efficiency for charge produced in the depletion layer is 100 percent, the total effective charge, $Q_{eff}$, reaching the electrodes is

$$Q_{eff} = N e$$  \hspace{1cm} (3.11)

and the voltage pulse to the input of the preamplifier then becomes

$$V_I = \frac{Q_{eff}}{C_d + C_I}$$  \hspace{1cm} (3.12)

where $C_I$ is the input capacitance to the preamplifier.

### 3.3.4 Rise time

The transit time for the carriers across the depletion layer of the detector is a measure of the rise time of the voltage pulse to the preamplifier. Charge generated within the depletion layer of a surface barrier detector drifts under the influence of the electric field, $E$, and is collected in a time, $t$, equal to the distance travelled divided by the average velocity. If the field is high enough a major part of the carrier transit occurs while the velocity is near the maximum carrier velocity, in which case the carrier velocity is approximately the same as in the uniform field case. Thus

$$v = \mu_p E = \mu_p \frac{V}{x_n}$$  \hspace{1cm} (3.13)
where \( \mu_p \) is the mobility of the holes since the rise time is determined by the motion of the slower charge carrier and \( V \) is the total reverse-bias voltage across the diode.

The transit time for the carriers across the depletion layer then becomes:

\[
 t = \frac{x_n^2}{\mu_p V} \tag{3.14}
\]

Substitution of equation 3.14 for the resistivity into equation 3.2 for the depletion layer thickness leads to the result:

\[
x_n = \left[ \frac{KV\mu_n}{2\pi} \rho_n \right]^{1/2} \tag{3.15}
\]

Thus:

\[
t = \frac{K}{2\pi} \left( \frac{\mu_n}{\mu_p} \right) \rho_n ^{1/2}
\sim 5.0 \times 10^{-12} \rho_n \text{ seconds} \tag{3.16}
\]

The transit time for the carrier in 3,000 ohm-centimeter resistivity silicon is then approximately 15 nanoseconds.

### 3.3.5 Dead layer

It may be recalled from equation 3.2 that the depletion layer extends a distance \( x_p \) from the junction into the p-type surface layer. If the
depletion layer does not extend to the surface of the detector, as is usually the case, then a thin "dead layer" remains on the surface of the detector and particles entering the depletion region through this "dead layer" lose a fraction of their energy before reaching the depletion layer. In a surface barrier detector with a negligibly thick gold electrode the "dead layer" is negligible and the detector is effectively "windowless".

3.3.6 Sensitivity to background radiations

Although the surface barrier detectors have a relatively low sensitivity to neutrons, photons and electrons these ionizing radiations can still interact with the detectors to produce background pulses. The neutrons can produce pulses in the detectors through \((n,p)\) and \((n,\alpha)\) reactions in both the silicon and the gold surface layer. The latter is negligible since the layer is so thin and the high Z-value inhibits the emission of charged particles. Finally background pulses can be produced by recoil protons from adsorbed hydrogen and water vapour in the surface of the detector. Since the gold has a very low adsorption it is particularly suitable in this respect. Measurements\(^{(23)}\) at A.E.R.E. showed only a very low counting rate when surface barrier detectors were exposed to a flux of approximately \(10^5\) neutrons per square centimeter per second of \(2\) MeV, an energy which corresponds to the peak of the statistical photon neutron energy distribution.

As the thickness of the depletion layer increases the response of the surface barrier detectors to electrons and photons also increases. It can be seen from the solid state detector nomograph, Figure 3.5, that the range of a 0.6 MeV electron in silicon is 1,000 microns.
In order to minimize the background from neutrons, photons and electrons, sufficient bias is applied to the detector to give a depletion depth which is only slightly in excess of the range of the charged particles to be detected.

3.3.7 Description of the detection system

Since the rise time of the surface barrier detectors is dependent on the transit time of the majority carriers across the depletion layer, it was possible to obtain a faster rise time and still achieve the required thickness of depletion layer by using two detectors mounted back to back. The photoprotons which were sufficiently energetic passed through the front detector which was a totally depleted or transmission detector with a very thin front and back "dead layer" and deposited their remaining energy in the second detector. The lower energy protons deposited all of their energy in the front detector. The two Ortec surface barrier detectors used in this experiment were connected in parallel electrically with the same reverse-bias voltage supply. The output signals were added directly at the input to the preamplifier. The detector area, defined by the collimator, subtended an angle of 4.6 degrees and a solid angle of $5.5 \times 10^{-3}$ steradians at the target position which was 3.8 inches away.

The transmission detector was an Ortec model number TMEJ100-500 with an active area of 100 square millimeters and a depletion depth of 500 microns. The front "dead layer" of the detector was 9.1 keV for 5.5 MeV alpha particles and the depletion depth extended through to the back at 240 volts. A check on the back "dead layer" made under the experimental conditions of reverse-
bias voltage with a Po\textsuperscript{210} alpha source revealed no detectable change in position on the pulse height analyser for alpha particles impinging on the front or back surface of the detector. It was therefore assumed that the detector was fully depleted throughout the experiment. According to the solid state detector nomograph, Figure 3.5, a 500 micron depletion depth gives a linear response for protons with energy up to 8.5 MeV. Protons with more than 8.5 MeV penetrated through the transmission detector and impinged on the back detector.

The back detector, Ortec model number SBCJ150-500, had a sensitive area of 150 square millimeters and a depletion depth of 500 microns. The larger area of the back detector prevented the protons passing through the transmission detector at an angle from missing the sensitive area of the back detector. The front "dead layer" of this detector was less than 20 keV for a 5.5 MeV alpha particle which corresponds to about 5 keV for a proton of the same energy. The total depletion depth of the two detectors was 1000 microns which has a linear response up to 12.5 MeV of proton energy. For protons of energy greater than 12.5 MeV the output pulse from the detectors drops off as shown in Figure 3.6 because the region of maximum specific ionization occurs outside of the depletion depth.

Although the gold-silicon surface barrier detectors are slightly photoconductive the relative darkness of the interior of the target chamber was sufficient to prevent this from becoming a problem.

3.4 Electronics

3.4.1 General

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FIGURE 3.6

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The total depletion depth of the two detectors was 1000 microns which gives a linear response up to 12.5 MeV of proton energy. For protons of greater energy the output pulse from the detectors falls off as shown. Since 4.2 microns of silicon is equivalent to an areal density of $1\text{ mg/cm}^2$, the total depletion depth is equal to $238\text{ mg/cm}^2$.
The extremely high photon flux impinging on the target during the four microsecond beam pulse of the betatron caused the emission of large numbers of electrons as well as the desired photoprotons. Since the detection system was sensitive to all charged particles some means had to be found to discriminate against these electrons. The rise time of the detection system for a single charged particle was found to be 50 nanoseconds. The pile-up pulses of the electrons, due to a number of electrons impinging on the detectors within their resolving time, rose more slowly to their peak value with a rise time between 300 and 400 nanoseconds. In order to prevent this build-up of the electron pulses to amplitudes equivalent to the protons' amplitudes the pulses were clipped with a shorted delay line. In this way, fewer electron pulses could add their amplitudes in a pile-up pulse and these pulses could be kept below the equivalent of a four MeV proton pulse for targets of low Z-value. On account of an increased number of electrons for the target elements with the higher Z-values the discrimination level had to be increased to 5.0 MeV.

A block diagram of the electronics is shown in Figure 3.7. The output pulses from the surface barrier detectors were added linearly and fed into a single-stage cathode follower preamplifier mounted on the lid of the target chamber. The signal from the preamplifier was then amplified in a bank of four commercially available distributed amplifiers (three Hewlett-Packard type 404A and one Hewlett-Packard type 404B). The type 404A amplifiers had a continuously variable gain from 5 to 10 and did not invert the input pulse. The type 404B amplifier had a voltage gain of 6 and did invert the input pulse. Both types of amplifiers had a rise time of 2.7 nanoseconds. The pulses were clipped after the second distributed amplifier to a length of 100 nanoseconds with 15 feet of 100 ohm cable type RG62A/U.
The reverse bias, $V$, was applied to the detectors, $D$. The output pulses from the detectors were added linearly and fed into the cathode follower preamplifier, CFP. The pulses from the preamplifier were then amplified in a bank of five Hewlett-Packard distributed amplifiers, HPA. The pulses were clipped to a length of 100 nanoseconds after the second distributed amplifier. The output pulses from the bank of distributed amplifiers was split to feed both the fast discriminator, FD, and the fast linear gate, FLG. The discriminator pulses opened the gate whose output pulses were fed into the pulse lengthener, PL. After further amplification in a conventional amplifier, CA, the experimental data were recorded on the pulse height analyser, PHA.
The clipped pulses were then passed through a long cable which carried the pulses to the control room of the betatron and into another type 404A distributed amplifier. The signal from this amplifier was split to feed both a fast discriminator and a fast linear gate. The output pulse from the discriminator was fed through a variable delay and opened the gate. The output pulse from the fast linear gate was fed into a pulse lengthener and low noise preamplifier. After further amplification with a conventional 0.5 microsecond pulse amplifier the pulses were recorded on a multichannel pulse height analyser. All connections carrying fast pulses were made with low capacity 200 ohm cable.

3.4.2 Preamplifier

The task of the preamplifier was to transmit the signal from the detectors to the bank of distributed amplifiers with impedance matching and without affecting the rise time of the pulse. Since a voltage sensitive preamplifier has both a faster rise time and a larger output voltage than a charge sensitive preamplifier, this type was chosen.

The single-stage cathode follower preamplifier used in this experiment, is shown in Figure 3.8. The R.C.A. high-μ Nuistor triode, 6CW6, was chosen for its high frequency response which did not effect the leading edge of the pulse from the detectors and for its compact size which enabled the preamplifier to be conveniently mounted on the lid of the target chamber. The detectors could then be connected to the preamplifier with a minimum input capacitance through a Kovar glass-to-metal seal. The trailing edge of the pulse from the preamplifier was shaped to a length of 10 microseconds by the time constant of the 10 megohm leak resistor and the stray capacitance...
FIGURE 3.8. Cathode Follower Preamplifier.

The single-stage cathode follower preamplifier employed a Nuvistor triode type 6CW6, V1. The following components were used:

- V - 300 volt dry cell
- R1 - 5 Mohm
- R2 - 5 Mohm
- R3 - 5 Mohm
- R4 - 10 Mohm
- R5 - 6.8 Kohm
- R6 - 47 ohm
- R7 - 3.3 Kohm
- C1 - 0.01 μfarad
- C2 - 0.001 μfarad
- C3 - 0.047 μfarad
- C4 - 0.01 μfarad
The fast pulse-amplitude discriminator supplied a standard output pulse of six volts, to trigger the gate circuit, for each input pulse from the distributed amplifiers whose amplitude exceeded the threshold bias. The design of the fast discriminator was basically that of Farley (36) and consisted of three stages (i) the threshold amplifier, (ii) the Schmidt trigger circuit and (iii) the amplifier and cathode follower. The threshold bias which was adjustable from 1 to 20 volts by means of a helipot was set to discriminate against the electron pile-up pulses. In this way the extremely high counting rates in the previous sections of the electronics were prevented from passing the gate circuit. Since less than a 0.1 volt difference in the amplitude of the clipped 100 nanosecond pulses from the distributed amplifiers at the discrimination threshold was sufficient to drop the standard discriminator output pulse from 90 to 10 percent of its maximum amplitude, the trigger action of the discriminator was considered sharp. Two input pulses of 100 nanoseconds length whose peaks were separated by 0.2 microseconds were handled completely independently.

The output pulses from the discriminator were used to trigger the fast linear gate circuit.

When the trigger pulses from the fast discriminator were applied to the fast linear gate circuit, it transmitted the output signals from the
distributed amplifiers without modifying their pulse height spectrum. The basic element of the gating circuit which was based on the design of Valckx and Rymanus (106) consisted of two 1:1 pulse transformers in series and separated by two fast crystal diodes as shown in Figure 3.9. The diodes were biased off by a positive voltage when the gate was not triggered. The trigger pulses from the discriminator after suitable shaping were applied to the centre tap of the secondary winding of the first transformer driving both diodes into a highly conducting state so that the gate was opened. Since the circuit was well balanced the contributions to the fluxes in the second transformer, caused by the currents in the loops due to the gate pulse, were equal and opposite. Thus the pedestal arising from the gate pulse was negligible. The reverse-bias voltage on the diodes kept the feed through of the input pulses when the gate was closed to a negligible value. By choosing the appropriate operating point in the conducting state the variations in the impedances of the two diodes were made almost equal and opposite so that the linearity was maintained over a large range of input pulse heights. It was determined that the gate was linear up to pulse heights corresponding to 13 MeV which was greater than the maximum proton energy loss in the depletion layer. This was done by increasing the gain of the distributed amplifiers until the high-energy alpha peak from Ra$^{226}$ was located in the position on the pulse height analyser that normally corresponded to 13 MeV during the experimental runs.

The pulse transformers in the gating element were made with ferrite potcores (Philips Ferroxcube 3B). Each coil of the transformer consisted of 6 windings of high frequency Litze wire (10 x 0.07 mm). A copper static shield was placed between the coils to reduce capacitive feed-through. The transformers were wound with great care in order to achieve the highest
FIGURE 3.9. Basic Diagram of the Gate Circuit.

The basic element of the gating circuit consisted of two 1:1 pulse transformers, T1 and T2, in series and separated by two fast crystal diodes, D. The diodes were biased off by the positive voltage, V, when the gate was not triggered. The trigger pulses, TP, from the discriminator after suitable shaping were applied to the centre tap of the secondary winding of the first transformer driving both diodes into a highly conducting state so that the gate was opened.
possible symmetry. The diodes used in the gating element were Transitron S555G.

In order to analyse the energy spectrum of the output from the gate, it was necessary to stretch the pulses to about 2 microseconds before feeding them into the pulse height analyser. A passive R-C stretcher was used to maintain good linearity. Since the pulses suffered the inevitable loss in pulse height in the passive network, they were then amplified with a low noise cascode preamplifier with a gain of 12 and a conventional microsecond pulse amplifier before being fed into the pulse height analyser.

3.5 Bremsstrahlung Control

3.5.1 Energy control and calibration

Since the \((\gamma,p)\) absorption cross sections for medium-weight nuclei are high in the energy range about 22 MeV(92), the peak bremsstrahlung energy at which this experiment was run, the yield of photoprotons is a very sensitive function of this energy. Consequently, adequate control of this energy had to be maintained during the experimental runs to ensure that the errors arising from the instability of the maximum bremsstrahlung energy were negligible compared to the statistical errors. It was, in fact, relatively simple to maintain the energy to within ±100 keV at 22 MeV. If a 20 percent increase in the photoproton yield results from a 1 MeV increase in the peak energy of the bremsstrahlung(46) then the 100 keV variation in peak energy corresponds to a ±2 percent variation in the photoproton yield, which is well within the best statistics of the experiment.

The maximum energy of the bremsstrahlung spectrum was controlled by
adjusting the amplitude of the magnet current by means of a motor driven variable transformer. The magnet current was proportional to the peak energy and induced a 180 cycle per second voltage in a leg turn of the magnet, which was filtered and rectified to a preselected reference level. By maintaining the rectified voltage within ±0.5 volts the energy could be kept constant to within ±0.5 percent from 10 MeV to 22 MeV.

Since a linear relationship exists between the rectified voltage and the peak energy of the bremsstrahlung spectrum (69), the energy scale of the betatron was determined by two calibration points, the (γ,n) thresholds in Cu^{63} at 10.8 MeV and in C^{12} at 18.7 MeV. The thresholds were determined by counting, with a thin-walled Geiger tube, the positron activity induced in uniform hollow cylinders of copper and polystyrene irradiated for a series of energies near their respective thresholds. By extrapolating the results to the low background counting rate the thresholds could be determined to within 200 keV. The results of these measurements were reproducible within this limit after both adjustments of the donut and replacement of the donut. Thus the peak bremsstrahlung energy was 22±0.4 MeV and was constant to within ±0.1 MeV.

3.5.2 Dosimetry

The photon flux from the betatron was monitored during the experimental runs by a built-in ionization chamber through which the bremsstrahlung beam passed before collimation. Both the instantaneous beam intensity and the relative integrated dose were indicated by monitors on the control panel of the betatron. The integrated dose monitor was calibrated daily against the response of a standard Victoreen 100-roentgen thimble chamber embedded in
five centimeters of lucite and placed at the target position in the chamber. The excellent reproducibility of this calibration indicated the constancy of the beam quality.

3.6 Experimental Procedure

3.6.1 Alignment of the bremsstrahlung beam

The lead shielding wall and target chamber were supported on a platform of concrete blocks set up on a truck which ran on rails set in the floor of the betatron room. The irradiation facility was used for medical therapy during the days so that it was necessary to wheel the apparatus up to the betatron before commencing the alignment procedures and away from the betatron after completing the evening's run. Since the rails ran perpendicular to the face of the machine, the lead collimator in the shielding wall, which had been previously centred both vertically and horizontally on the entrance and exit ports of the stationary target chamber, using x-ray photographs, was never shifted out of horizontal alignment with the collimator in the betatron, when the truck was moved up to its location in front of the machine. Further alignment of the bremsstrahlung beam was accomplished by manoeuvring the betatron which could be adjusted both in vertical height above the floor and angle of rotation about a horizontal axis perpendicular to the beam.

The machine was first rotated approximately into the position of a horizontal beam and then was raised or lowered until the collimated beam of light from a flashlight bulb, mounted on the end of the lead collimator inside the betatron, passed through the lead shielding wall and target chamber.
casting its outline on a piece of paper set in the exit beam port window. When the target was lowered into position, the location of the shadow on the exit beam port window was noted. The lid of the target chamber along with the target and detectors was then removed so that the final alignment of the bremsstrahlung beam could be effected with the aid of a Victoreen 100-roentgen thimble chamber embedded in five centimeters of lucite.

By viewing the thimble, located at the target position, through a window set in the target chamber at right angles to the beam ports its sensitive volume was positioned at the vertical height of the sensitive area of the detectors. The thimble was located horizontally so that its shadow cast on the exit beam port window was superimposed on the location of the target's shadow as previously noted. The thimble was then exposed to the bremsstrahlung beam for five monitor counts* in this position as well as in positions 3/16 inches above and below this position. When the maximum reading was obtained in the central position and the two extreme positions registered the same reading the beam was vertically aligned. As a check on the horizontal alignment the thimble was exposed to the bremsstrahlen in the target position as well as positions about 3/16 inches to the left and right of the central axis of the beam. Once again the above conditions satisfied the requirement for horizontal alignment. These conditions were generally satisfied without touching the target chamber; however, it was occasionally necessary to shift the chamber slightly to satisfy them. The target chamber was never shifted by more than about 1 degree. Checks on the alignment

* 1 monitor count was a measure of the integrated dose through the built-in ionization chamber mounted inside the betatron corresponding to about 12 roentgens at the target.
after completion of the experimental runs revealed that the beam shift
during the course of the runs was negligible. Reproducible counting rates
from the same target on different days indicated that the beam alignment
was reproducible.

3.6.2 Reverse-bias voltage

The reverse-bias voltage for the surface barrier detectors was
supplied by an Eveready "Mini-Max" 300 volt dry cell type number 493. The
bias voltage was applied gradually to the detectors through a potentimeter
mounted in the preamplifier. A ten megohm resistor was inserted between the
potentiometer and the detectors to act as a voltage drop in case of a breakdown
in the detectors. Under normal operating conditions this resistor dropped
the voltage by approximately ten volts.

As the reverse-bias voltage was being raised the noise from the
detectors was monitored, with the clipping line removed, on a Tektronix
oscilloscope type number 545A. If the noise from the detectors was greater
than the usual 200 keV energy equivalent noise level the reverse-bias voltage
was dropped until the excess noise disappeared and was then raised slowly
over a period of several minutes. The detectors were normally run at 240
volts. At this voltage the transmission detector was fully depleted and
the total depletion depth of the detection system was 1000 microns.

3.6.3 Energy calibration of the electronics

The electronics were calibrated with the target chamber under vacuum
using Ra\(^{226}\) and its daughter products with alpha energies of 4.78 MeV,
5.50 MeV, 6.00 MeV and 7.68 MeV. With the target-holder shaft in its upper position so that the target did not block the sensitive area of the detectors, the Ra-226 source was lowered by means of the source-supporting rod into the horizontal plane of the detectors which were normally biased with 240 volts. The source was rotated until its active face was looking at the sensitive area of the transmission detector. Since a previous comparison of the alpha-particle spectra taken with the source incident on the face of each of the detectors under the experimental conditions had indicated no dependence on the detectors, the transmission detector was always used for the energy calibration. The alpha particles were stopped in the transmission detector and the signals were passed through the entire electronic system including the clipping line. The energy spectra were then displayed on the 512-channel pulse height analyser.

In order to obtain a consistent calibration of the electronics the gain of the distributed amplifiers was adjusted until the 7.68 MeV alpha peak was located in the range of channels 87±4 of the pulse height analyser. The system was then considered to be ready for an experimental run. Due to the large number of noise pulses in the electronics when the betatron was on, there was a base line shift of 7±2 channels so that it was necessary to recalibrate the energy scale. If the above condition for the high-energy alpha peak had been satisfied this shift generally lead to a desirable energy calibration of the high-energy peak within the range of channels 80±5. In this way the electronic system was calibrated to approximately 10 channels per MeV.

Since the gain drift in the wide band amplifiers was excessive it was necessary to calibrate the electronics before and after each experimental run. If the drift during the course of the run was less than 1±4 channels
3.6.4 Detection of photoprotons

Both the calibrations of the electronics and the experimental runs for the detection of photoprotons were made with the target chamber evacuated to a pressure of 1 millimeter of mercury. The absorption of the alpha energy from the source to the detector and of the proton energy from the target to the detector was consequently negligible. When the system had been properly calibrated the alpha source was withdrawn into its cylindrical brass housing and the target foil was lowered into position in the path of the bremsstrahlung beam.

Since the greatest pile-up of electron pulses occurred at the most forward angle, a new target element was first checked by moving the detection system to this angle which was limited to 45 degrees with respect to the photon beam. This angle was a compromise between an acceptable level of the electron pile-up pulses which increased rapidly as the detectors were moved forward and a reasonable amount of shielding in the lead housing support which could be kept out of the photon beam at the most forward angle while preventing electrons from passing through the long sensitive dimensions of the detectors. If the counting rate in the low-energy channels was not excessively high, the discriminator threshold was assumed to be satisfactory. A high counting rate in the low-energy range was eliminated by raising the bias level of the discriminator. The discriminator threshold varied between 3.5 MeV for the lower-Z target elements and 5.0 MeV for the higher-Z target elements.
When the pile-up of the electron pulses had been reduced to a reasonable level in this way a thorough calibration of the electronic system was made in order to check both the energy linearity and the energy resolution. The experimental runs to detect the photoprotons were then commenced.

Due to the low yield of photoprotons it was necessary to take several runs at each angle for each target in order to achieve the desired statistical precision of 10% standard deviation for protons of energy greater than or equal to 8 MeV, which is the generally accepted dividing line between the "resonance direct" and statistical protons. The integrated dose in roentgens at the target position required to accumulate 100 photoprotons in this energy range for each angle was thus determined in order to achieve this precision for the Poisson counting statistics. To eliminate systematic errors the experimental runs were both taken in such a way that approximately the same number of counts were accumulated at each angle on each day and cycled through the angles in a different order each day.

The electronics system was calibrated before and after each run to maintain the control on the drift which occurred during the run. The target foil received an integrated dose of approximately 500 roentgens in lucite at a dose rate which was maintained between 55 and 60 roentgens per minute throughout the course of the experimental runs. During the run care was taken to maintain the energy constant to within 0.5%. No significant differences were observed either between runs with the same target on different days or between the two sets of results taken three months apart for the copper target. Thus it was shown that the results of the experimental runs, which were taken at 45°, 60°, 90°, 120° and 145°, were reproducible.

Since the yields of photodeuterons(112) and photoalphas(49) are expected to be only a small percentage of the photoproton yield in the range.
of atomic numbers employed in this experiment, no effort was made to discriminate against these photoreaction products.
CHAPTER IV

EXPERIMENTAL RESULTS

In chapter IV the results of the experimental measurements of the angular and energy distributions as well as the yields of photoprotons from nine medium-weight nuclei, with atomic number between 28 and 50, are presented in detail. The first section of this chapter deals with the alpha-particle spectra which were used for calibration of the energy scale. The remaining sections deal directly with the results and analysis of the measurements pertaining to photoprotons alone.

4.1 Alpha-Particle Spectra

The alpha-particle spectra were taken principally for the purpose of calibrating the energy scale and checking the drift in gain during the course of the experiment; however, these spectra were used also to check both the energy linearity and resolution of the detection system. The alpha-particle source, as previously mentioned, was Ra\(^{226}\) with peak energies of 4.78 MeV, 5.50 MeV, 6.00 MeV and 7.68 MeV\(^{53}\). Since the "dead layers" of the two detectors were negligibly thick in comparison with the energy resolution of the experimental set-up, no corrections were applied for energy loss in these regions. A typical energy spectrum of the Ra\(^{226}\) alpha peaks along with the corresponding energy calibration curve is shown in Figure 4.1. The most striking feature is the excellent linearity of the electronic system for charged particles up to about 3 MeV. In order to check the linearity for pulse heights of greater amplitude the gain of the distributed amplifiers...
FIGURE 4.1
FIGURE 4.1. Typical Alpha-Particle Spectrum of Ra$^{226}$

The alpha-particle spectrum of Ra$^{226}$ with peak energies of 4.78 MeV, 5.50 MeV, 6.00 MeV and 7.68 MeV is shown. Since the "dead layers" of the two detectors were negligible thick, no corrections were applied for energy loss in these regions. The corresponding energy calibration curve is seen to be linear.
was increased until the 7.68 MeV alpha peak corresponded in position on the pulse height analyser to that normally occupied by a charged particle of 13 MeV. This energy was greater than the maximum proton energy which could be deposited in the detection system. Since the linearity was found to be excellent for pulse heights of this magnitude it was assumed that the energy calibration curve could be linearly extrapolated to the maximum energy of the protons detected. The agreement between the theoretical range-energy relationship for protons in silicon and the maximum proton energy detected during the experimental runs was further evidence of the linearity of the detection system at the higher energies.

Since the noise in the electronics, in particular in the distributed amplifiers, was high it was not possible to completely resolve the low energy alpha peaks at 4.78, 5.50 and 6.00 MeV, in order to check the energy resolution. The high-energy alpha peak at 7.68 MeV was well separated from the other peaks, however, and its energy resolution was found to be 5 percent independent of its relative position on the pulse height analyser, between positions normally corresponding to 5 and 13 MeV. On account of the relatively thick targets in terms of the proton energy lost in traversing one half of the target's thickness such energy resolution was quite satisfactory for this experimental work.

Alpha-particle spectra were taken before and after each experimental run to check on the gain drift during the course of the run and to determine a mean energy calibration scale if some drift had occurred.

4.2 Corrections to the Experimental Data

In the analysis of the angular and energy distributions of photoprotons...
it is generally necessary to correct the data both for the attenuation of the photon beam in the target element and for the absorption of photoproton energy in the target. Since the photon beam was both uniform in intensity across its diameter and larger in diameter than the target's width the target could be rotated with the detectors without affecting the geometrical efficiency of the experimental set-up.

Under these circumstances the correction for the attenuation of the photon beam was dependent on angle and was a maximum for the 90 degree position of the detectors with respect to the photon beam, in which case, the beam traversed the one centimeter width of the target. By rotating the target by four degrees from its normal position perpendicular to the central axis of the detectors, when the detectors were set at the 90 degrees position, the attenuation of the photon beam was reduced to a negligible value for the thin targets, which were used in this experiment, without affecting the amount of the target element seen by the detectors in comparison with the other angles of observation. Thus it was not necessary to make corrections for the attenuation of the photon beam.

Since the target element was rotated with the detectors the correction to the proton energy spectrum for the energy lost by the proton in the target was independent of the angle of observation even at the 90 degrees position where the small four degree rotation had a negligible effect. The approach, taken in this experiment, was to calculate the energy lost by a proton of specific energy, produced at the centre of the target, in leaving the target and to apply this correction to each proton of that energy independent of its point of origin in the target. The specific energy losses for protons in the various elements were taken from Marion(67). If these data were not available for a specific element, they were determined by interpolation.
As the target nuclei were all of medium weight the corrections for the motion of the centre of mass are small and none have been made. All angles are expressed in the laboratory system.

4.3 Normalization of Data

On account of the drift during the course of an experimental run as indicated by the energy calibrations before and after the run and the difficulty in adjusting the gain so that the calibration peaks always appeared in the same channels of the pulse height analyser, the criterion for an appropriate gain setting was reduced to adjusting the highest energy alpha calibration peak, 7.68 MeV, to fall within the range of 80±5 channels under the experimental conditions. Since the experimental data for each target consisted of a number of runs taken at each of the five angles of observation it was necessary to standardize the energy calibration of the data before carrying out any mathematical operations. This energy calibration correction was carried out on the IBM 7094 at the University of Toronto.

Before the data could be analysed by the program it was necessary to draw the energy calibration curve for each run in order to obtain both the zero-energy intercept and the number of channels from 0 to 12 MeV. The calibration correction program then corrected the data of the experimental runs to a standard energy calibration curve which passed through the origin and had a slope of 100 keV per channel. The curves were adjusted to pass through the origin by adding to or subtracting from the experimental photoproton spectrum the appropriate number of channels corresponding to the zero-energy intercept, which was usually less than 5 channels. This led to a negligible zero error when the slopes were corrected by the program.
The program was capable of handling all of the experimental data for a single target which consisted of (i) the number of experimental runs taken at each angle for that target, (ii) the unnormalized photoproton spectra, corrected for the zero-energy intercept and arranged according to angle, (iii) the number of channels from 0 to 12 MeV on the energy calibration curve for each run, (iv) the normalized "no-target" background for each angle which was less than 10 percent of the yield even at the lowest photoproton energies and (v) the target half-value thickness corrections as a function of the photoproton energy. The program expanded or contracted the photoproton spectra until each contained 120 channels in the 0 to 12 MeV energy range, added the spectra at each angle together channel by channel and subtracted the appropriate "no-target" background.

When all of the data for a given target had been so analysed the target half-value thickness corrections were applied in the following manner, so that the data for the energy and angular distributions could be extracted from the normalization program directly. Since the photoproton spectra had a slope of 100 keV per channel after normalization, a proton with X MeV at the detector position was then located in channel 10X. In order to determine the channel, Yx, in which X MeV protons, emitted at the centre of the target, would have been recorded, the half-value thickness expressed as a number of channels, using the above slope, was subtracted from 10X according to the relation

\[ Y_x = 10X - T_x - S_x/10 \]
where $T_x$ is the integral number of 100 keV units in the target half-value thickness and

$S_x$ is the integral number of 10 keV units in the target half-value thickness.

Thus $Y_x$ was calculated to 0.1 channel accuracy in 0.5-MeV intervals from 4 MeV to 13.5 MeV. The counts in the observed spectra between $Y_x$ and $Y_x+0.5$ were then summed for each of the observed angles in 0.5-MeV intervals in order to generate the corrected energy spectra at each of the observed angles. As $Y_x$ usually corresponded to some intermediate point in a channel of the observed spectra, the counts in that channel were linearly proportioned between the two 0.5-MeV intervals involved. The energy spectrum at 90 degrees to the photon beam was then printed out for each target as it was analysed.

4.4 Angular Distributions and Integrated Yields of Photoprotons

The energy normalization program also calculated the total yield of photoprotons at each of the five observed angles for a number of different energy intervals. These yields were then normalized both to the number of experimental runs taken at the 90 degree position of the detectors with respect to the photon beam and to 100 events at this position. These data were also printed out as part of the energy normalization program. In order to determine if pile-up pulses were present in the energy spectra at the most forward angle of observation, the yields of photoprotons measured using one-half beam intensity were compared with those measured using full beam intensity for each of the targets. If the yields per roentgen differed by more than the 10 percent statistical error, it was assumed that pile-up was present and that angle was eliminated from the angular distribution.
For targets which showed significant pile-up at 45 degrees, a similar test was made at the 60 degree angle of observation. This angle was similarly eliminated if pile-up was observed. From observation of the normalized data it was possible, in some cases, to detect the possible occurrence of pile-up, which had been missed by the previous tests. Any angles which fell into this category were similarly eliminated from the angular distributions in order to ensure that the results were not weighted by the presence of pile-up.

The normalized yields of photoprotons for the various energy intervals were then fitted by the method of least squares to the expression

\[ Y(\theta) = a + b \sin^2 \theta (1 + p \cos \theta)^2. \]  

The angular distributions were fitted to the five analysed angles unless pile-up was present at the most forward angle of observation, in which case, the fit was made to four angles only. The angular distributions were not calculated when pile-up was present at the 60 degree position as was the case for the low-energy intervals with the nuclei of higher Z-value. The details of the calculations, which were made on the IBM 7094 computer, are given in Appendix A.

The total yield of photoprotons within a specified energy interval is obtained by integrating its angular distribution over a sphere. The resulting expression for the integrated yield, \( Y \), then becomes

\[ Y = \frac{1}{2\pi T} \left( 2a + \frac{b}{3} + \frac{4}{15} b p^2 \right) \]
where $N$ is the multiplication factor used to normalize the proton yield for the specified energy interval to 100 events at the 90 degree position,

$M$ is the number of moles of the target element in the photon beam which can be seen by the detectors,

$R$ is the dose in roentgens given to the target at the 90 degree position,

$a, b, p$ are the angular distribution coefficients derived from the least squares fit for the specified energy interval and

$G$ is the geometrical efficiency of the detection system.

This efficiency is defined by the equation

$$G = \frac{A}{4\pi r^2}$$

where $A$ is the sensitive area of the detectors defined by the collimation of the lead housing support and

$r$ is the distance from the target to the inside edge of the collimator.

The geometrical efficiency for the experimental set-up was $4.4 \times 10^{-4}$.

**4.5 Statistical Energy Distributions**

The energy distributions predicted by the statistical theory for seven of the target nuclei, irradiated with 22 MeV bremsstrahlen, were calculated using equation 2.68. In these calculations the cross sections for the formation of the compound nucleus by the inverse process, $\sigma_c(\varepsilon_p)$, were taken from the calculated results of Shapiro(90), based on the continuum theory of
Feshbach and Weisskopf(58). Since the inverse process adds one proton to the nucleus the atomic number of the nucleus for which the cross sections were determined was one less than that of the nucleus for which the energy distribution was to be determined. The cross sections for nuclei which were not given by Shapiro were found by interpolation. The level densities were assumed to be exponential with nuclear temperatures as given by Livesey(62), equation 2.60. The proton separation energies were taken from Nuclear Data Sheets(78) and Feather(37). In order to evaluate the remaining integral the (γ,n) cross sections of Katz et al.(54,55,73) were employed along with the Schiff bremsstrahlung spectrum, $N(E,E_{\text{max}})$, tabulated by Katz and Cameron(55).

The statistical energy distributions are shown in Figures 4.2 and 4.4; they were approximately normalized to the low-energy part of the observed spectrum. Since the (γ,n) cross sections were not available for cadmium and tin the statistical energy spectra were not calculated for these nuclei.

4.6 Experimental Results

The results of the experimental measurements made on the nine target elements, listed in Table 4.1, are summarized in Tables 4.2 to 4.5 and Figures 4.2 to 4.7. The photoproton energy distributions measured at 90 degrees to the bremsstrahlung beam are presented in 0.5-MeV. intervals in Table 4.2. Both these energy distributions and the calculated statistical energy distributions are shown in Figure 4.2. Table 4.3 contains the angular distribution coefficients, fitted by the method of least squares to equation 2.53, and the integrated yields, calculated from equation 4.3, for a number of 2-MeV proton energy intervals. These results are plotted as a function
of the mid-range energy of the photoprotons in Figures 4.3 and 4.4. The indicated errors are standard deviations determined from counting statistics. The statistical energy distributions are again presented in Figure 4.4. The photoprotons are divided into high- and low-energy groups in Table 4.4 and the angular distribution coefficients and integrated yields are given for each group. Since only about 10 percent of the photoprotons in the statistical energy distribution have energies greater than 8 MeV, this energy was considered to be a satisfactory boundary between the groups. In Figure 4.5 the anisotropy ratios are given for both groups of photoprotons as a function of Z. The integrated yields of the high- and low-energy photoprotons are plotted against Z-value in Figure 4.6. The total integrated yields of photoprotons, measured in this experiment, are shown in Table 4.5 and are plotted along with the predictions of statistical emission and "resonance direct" emission in Figure 4.7.
<table>
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<tr>
<th>Target Element</th>
<th>Z-value</th>
<th>Thickness (mg/cm²)</th>
<th>Half-value Thickness at 8 MeV (MeV)</th>
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<td>30</td>
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<td>Niobium</td>
<td>41</td>
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<td>Molybdenum</td>
<td>42</td>
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<td>Tin</td>
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<td>35</td>
<td>0.9</td>
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Reproduced with permission of the copyright owner. Further reproduction prohibited without permission.
<table>
<thead>
<tr>
<th>Energy Range (MeV)</th>
<th>Nickel</th>
<th>Copper</th>
<th>Zinc</th>
<th>Zirconium</th>
<th>Niobium</th>
<th>Molybdenum</th>
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</tr>
<tr>
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<td>57.0</td>
<td>40.0</td>
<td>34.2</td>
<td>31.1</td>
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<td>76.2</td>
<td>39.9</td>
<td>30.0</td>
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<td>105.6</td>
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<td>77.6</td>
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Reproduced with permission of the copyright owner. Further reproduction prohibited without permission.
<table>
<thead>
<tr>
<th>Energy Range (MeV)</th>
<th>Energy Distributions of Photoprotons at 90 Degrees</th>
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<td>24.3</td>
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<tr>
<td>5.0-5.5</td>
<td>15.2</td>
</tr>
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</table>
FIGURE 4.2. Energy Distributions of the Photoprotons.

The histograms give the numbers of photoprotons per 0.5-MeV energy interval emitted at 90 degrees to the photon beam. The smooth curves represent the energy distributions based on the statistical theory with level densities proportional to \( \exp(-\frac{E_p}{T}) \) where \( T \) is the nuclear temperature. The statistical curves have been approximately normalized to the low-energy photoproton yields.
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Number of Protons vs. Proton Energy (MeV) for Zirconium, Niobium, and Molybdenum.
### TABLE 4.3

<table>
<thead>
<tr>
<th>Target Element</th>
<th>Proton Energy Range (MeV)</th>
<th>Angular Distribution Coefficients</th>
<th>Photopronot Yield $10^5$ protons per mole.R</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>a</td>
<td>b</td>
<td>p</td>
</tr>
<tr>
<td>Nickel</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>$&gt;$10</td>
<td>123±25</td>
<td>1±30</td>
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<tr>
<td>9 to 11</td>
<td>76±19</td>
<td>31±27</td>
<td>*</td>
</tr>
<tr>
<td>8 to 10</td>
<td>61±12</td>
<td>37±18</td>
<td>0.22±0.20</td>
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<tr>
<td>7 to 9</td>
<td>75±10</td>
<td>23±15</td>
<td>0.22±0.28</td>
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<td>2±3</td>
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<tr>
<td>5 to 7</td>
<td>97±5</td>
<td>4±4</td>
<td>*</td>
</tr>
<tr>
<td>4 to 6*</td>
<td>90±4</td>
<td>5±3</td>
<td>*</td>
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<td>$&gt;$9</td>
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<td>67±19</td>
<td>0.34±0.15</td>
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<tr>
<td>8 to 10</td>
<td>17±11</td>
<td>71±19</td>
<td>0.29±0.13</td>
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<tr>
<td>7 to 9</td>
<td>50±10</td>
<td>46±16</td>
<td>0.27±0.17</td>
</tr>
<tr>
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<td>71±11</td>
<td>33±16</td>
<td>0.12±0.21</td>
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<tr>
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<td>58±9</td>
<td>40±15</td>
<td>0.00±0.12</td>
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<tr>
<td>4 to 6</td>
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<td>57±6</td>
<td>37±10</td>
<td>0.33±0.14</td>
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<td>4 to 6</td>
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<td></td>
<td>97±3</td>
<td>3±4</td>
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* p loses significance as b → 0 + 4 angles analysed only

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<table>
<thead>
<tr>
<th>Target Energy Range (MeV)</th>
<th>Proton Energy</th>
<th>Angular Distribution Coefficients</th>
<th>Photoproton Yield $10^5$ protons per mole. R</th>
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<tbody>
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<td>43±16</td>
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<td>60±22</td>
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<td>Zirconium</td>
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<td>25±14</td>
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<td>Molybdenum &gt;10</td>
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<td>0.47±0.11</td>
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### TABLE 4.3 (Cont'd)

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<tr>
<th>Target Element</th>
<th>Proton Energy Range (MeV)</th>
<th>Angular Distribution Coefficients</th>
<th>Photoproton Yield $10^5$ protons per mole.R</th>
</tr>
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<tbody>
<tr>
<td></td>
<td>a</td>
<td>b</td>
<td>p</td>
</tr>
<tr>
<td><strong>Silver</strong></td>
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<td></td>
<td></td>
</tr>
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<td>7 to 9</td>
<td>30±8</td>
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<td>0.01±0.07</td>
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The procedure for obtaining the error estimates on the angular distribution parameters is described in Appendix A.
FIGURE 4.3. b/a Ratios of the Photoproton Angular Distributions.

The b/a ratios, which have been determined by the method of least squares for a number of 2-MeV energy intervals, are plotted as a function of the mid-range energy of the photoprotons. The indicated errors are standard deviations determined from counting statistics.
The diagrams represent the b/a ratio as a function of proton energy for Zirconium, Niobium, and Molybdenum. Each graph shows data points and error bars for different proton energy values (in MeV). The graphs illustrate the variation of the b/a ratio with proton energy for each element.
SILVER

PROTON ENERGY (MeV)

b/a RATIO

PROTON ENERGY (MeV)

CAIUMIUM

PROTON ENERGY (MeV)

b/a RATIO

PROTON ENERGY (MeV)

TIN

PROTON ENERGY (MeV)

b/a RATIO

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FIGURE 4.4. Integrated Photoproton Yields.

The total yields of photoprotons for 2-MeV energy intervals are plotted as a function of the mid-range energy. The indicated errors are standard deviations determined from counting statistics. The smooth curves represent the energy distributions based on the statistical theory and are approximately normalized to the low-energy photoproton yields.
<table>
<thead>
<tr>
<th>Target Element</th>
<th>Proton Energy Range (MeV)</th>
<th>Angular Distribution Coefficients</th>
<th>Photoproton Yield $10^5$ protons per mole. R</th>
</tr>
</thead>
<tbody>
<tr>
<td>Nickel</td>
<td>$\geq 8$</td>
<td>$70\pm12$  $32\pm17$  $0.21\pm0.23$  $0.46\pm0.25$  $3.82\pm0.54$</td>
<td></td>
</tr>
<tr>
<td>Copper</td>
<td>$\geq 8$</td>
<td>$44\pm9$   $50\pm14$  $0.31\pm0.14$  $1.13\pm0.39$  $1.71\pm0.23$</td>
<td></td>
</tr>
<tr>
<td>Zinc</td>
<td>$\geq 8$</td>
<td>$63\pm8$   $36\pm14$  $0.81\pm0.35$  $0.57\pm0.23$  $1.76\pm0.19$</td>
<td></td>
</tr>
<tr>
<td>Zirconium</td>
<td>$\geq 8$</td>
<td>$18\pm8$   $76\pm14$  $0.17\pm0.08$  $4.24\pm2.04$  $1.45\pm0.20$</td>
<td></td>
</tr>
<tr>
<td>Niobium</td>
<td>$\geq 8$</td>
<td>$65\pm8$   $31\pm13$  $0.71\pm0.33$  $0.48\pm0.21$  $2.06\pm0.22$</td>
<td></td>
</tr>
<tr>
<td>Molybdenum</td>
<td>$\geq 8$</td>
<td>$50\pm9$   $52\pm15$  $0.53\pm0.19$  $1.04\pm0.35$  $1.67\pm0.21$</td>
<td></td>
</tr>
<tr>
<td>Silver</td>
<td>$\geq 8$</td>
<td>$26\pm8$   $75\pm13$  $0.26\pm0.09$  $3.68\pm1.64$  $1.14\pm0.15$</td>
<td></td>
</tr>
<tr>
<td>Cadmium</td>
<td>$\geq 8$</td>
<td>$38\pm10$  $58\pm16$  $0.37\pm0.16$  $1.52\pm0.58$  $0.84\pm0.12$</td>
<td></td>
</tr>
<tr>
<td>Tin</td>
<td>$\geq 8$</td>
<td>$74\pm16$  $26\pm24$  $0.25\pm0.41$  $0.35\pm0.33$  $0.88\pm0.17$</td>
<td></td>
</tr>
</tbody>
</table>

| Nickel         | 4 to 8$^+$               | $93\pm3$   $4\pm2$   $*\neq$  $0.04\pm0.02$  $18.05\pm0.46$ |
| Copper         | 4 to 8                   | $64\pm6$   $33\pm9$  $0.00\pm0.09$  $0.51\pm0.15$  $6.15\pm0.43$ |
| Zinc           | 4 to 8                   | $87\pm5$   $11\pm7$  $*\neq$  $0.13\pm0.08$  $10.53\pm0.60$ |
| Zirconium      | 4 to 8                   | $45\pm6$   $59\pm9$  $0.07\pm0.07$  $1.28\pm0.26$  $4.91\pm0.39$ |
| Niobium        | 5 to 8$^+$               | $74\pm12$  $27\pm16$  $0.12\pm0.30$  $0.37\pm0.22$  $3.50\pm0.42$ |
| Molybdenum     | 4 to 8$^+$               | $87\pm5$   $11\pm7$  $*\neq$  $0.13\pm0.08$  $7.45\pm0.42$ |
| Silver         | 6$^+$                    | $36\pm6$   $57\pm10$  $0.30\pm0.09$  $1.58\pm0.38$  $2.29\pm0.21$ |
| Cadmium        | 6$^+$                    | $64\pm10$  $34\pm15$  $0.41\pm0.31$  $0.54\pm0.25$  $1.63\pm0.13$ |
| Tin            | 6$^+$                    | $92\pm8$   $64\pm11$  $*\neq$  $0.07\pm0.12$  $1.64\pm0.19$ |

$*\neq p$ loses significance as $b \rightarrow 0$

$+ 4$ angles analysed only.

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FIGURE 4.5. The Dependence on $Z$ of the $b/a$ Ratios.

The $b/a$ ratios are shown as a function of the atomic number, $Z$, for the two energy groups (i) $\geq 8$ MeV and (ii) $< 8$ MeV. The indicated errors are standard deviations determined from counting statistics. The anisotropies of Mitchell and McNeill(72) are indicated by the closed circles for protons of energy $\geq 8$ MeV.
FIGURE 4.6
FIGURE 4.6. The Dependence on Z of the Integrated Protoproton Yields

The integrated photoproton yields for the two energy ranges, (i) \( \geq 8 \text{ MeV} \) and (ii) \( < 8 \text{ MeV} \), are plotted as a function of the atomic number, Z. The indicated errors are standard deviations from counting statistics.
<table>
<thead>
<tr>
<th>Target Element</th>
<th>Total Photoproton Yield $10^5$ protons per mole R</th>
</tr>
</thead>
<tbody>
<tr>
<td>Nickel</td>
<td>$21.85 \pm 0.71$</td>
</tr>
<tr>
<td>Copper</td>
<td>$7.86 \pm 0.53$</td>
</tr>
<tr>
<td>Zinc</td>
<td>$12.23 \pm 0.63$</td>
</tr>
<tr>
<td>Zirconium</td>
<td>$6.36 \pm 0.44$</td>
</tr>
<tr>
<td>Niobium</td>
<td>$5.05 \pm 0.47$</td>
</tr>
<tr>
<td>Molybdenum</td>
<td>$9.12 \pm 0.47$</td>
</tr>
<tr>
<td>Silver</td>
<td>$2.29 \pm 0.21$</td>
</tr>
<tr>
<td>Cadmium</td>
<td>$1.63 \pm 0.13$</td>
</tr>
<tr>
<td>Tin</td>
<td>$1.64 \pm 0.19$</td>
</tr>
</tbody>
</table>

TABLE 4.5

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FIGURE 4.7
FIGURE 4.7. The Total Integrated Yields of Photoprotons.

The total yields of photoprotons from the nine target elements are shown as a function of the atomic number, Z. The protons were integrated in the two groups (i) $\geq 8$ MeV and (ii) $< 8$ MeV. The theoretical predictions of statistical emission (smooth curve) and "resonance direct" emission (broken curve) are shown for comparison.
5.1 Introduction

Although the collective model has been successful in describing the properties of the giant dipole resonance, only the I.P.M. has been successful in describing the properties of the products of photonuclear reactions; consequently, the yields and energy and angular distributions of photoprotons from the medium-Z nuclei, measured in this experiment, will be discussed in terms of the latter model. It may be recalled from chapter II that the excited nucleon may be directly emitted in a "resonance direct" interaction or may interact with the rest of the nucleus and be emitted through a statistical process following compound nucleus formation. The range of Z-values studied in this experiment is of particular interest in this respect since the yield of statistical photoprotons decreases rapidly and the yield of "resonance direct" protons becomes significant between the closed proton shells at Z=28 and Z=50.

In the I.P.M. the nucleons are initially in nuclear shells of angular momentum, $l$, and make transitions to levels of angular momentum $l+1$ and $l-1$ through the absorption of $E1$ radiation. The dipole strength of an I.P.M. transition is then proportional to the product of the enhancement factor (Table 2.1) and the square of the radial overlap integral (Table 2.2). Both of these factors increase with the angular momentum, $l$, and favour the $l \rightarrow l+1$ transitions over the $l \rightarrow l-1$ transitions. While the enhancement factor only slightly favours the $l \rightarrow l+1$ transitions, however, the square of the
radial overlap integral favours these transitions by an order of magnitude. The dipole transition strength is consequently an increasing function of the angular momentum of the excited nucleon.

From equation 2.43 it can be seen that the escape probability for a proton in a "resonance direct" interaction is proportional to the transmission coefficient through the potential barrier of the nucleus. The transmission coefficients for protons, however, decrease with increasing $l$ so that the $l \rightarrow l-1$ transitions are favoured for "resonance direct" emission. The portion of the dipole strength which does not lead to "resonance direct" emission leads to compound nucleus formation from which the photoprotons are emitted in competition with the photoneutrons.

This competition can be strongly affected by nucleon pair correlations which cause an energy gap in even-even nuclei (74) equal to the sum of the proton and neutron pairing energies. The pair correlations are taken into account by calculating the nuclear excitation energy from a level lying above the ground state by the amount of the energy gap (81). The level density is then a function of the $Z$ and $N$ of the final nucleus. In particular, when odd-even nuclei are excited by photons the emission of a proton results in an even-even nucleus with the above energy gap while the emission of a neutron results in an odd-odd nucleus with no energy gap. The existence of the energy gap in the even-even nucleus leads to increased competition of the $(\gamma, n)$ reaction and therefore to a reduced yield of statistical photoprotons. When pair correlations were taken into account in the odd-even nucleus, Nb$^{93}$, the predicted statistical photoproton to photoneutron yield ratio was reduced by an order of magnitude (81). These calculations also indicated that the shape of the statistical energy distribution was independent of nucleon pair correlations.
In even-odd nuclei pair correlations lead to a considerably enhanced relative yield of statistical photoprotons since the emission of a neutron now leads to an even-even nucleus. For even-even initial nuclei the emission of either a proton or a neutron leads to a final nucleus of odd A-value so that neither photoproton nor photoneutron emission is enhanced by the inclusion of nucleon pair correlations.

In the following section the results of the measurements made in this experiment of the angular distributions, energy spectra and integrated yields of photoprotons are discussed with reference to the above facts. In order to supplement this discussion the strengths of the various E1 transitions are calculated in Appendix B for the copper, zirconium and cadmium nuclei. The relative yields of photoprotons from the various I.P.M. transitions are also calculated. Reference should be made to these calculations since they are helpful in the interpretation of the results.

5.2 Discussion

5.2.1 Angular distributions

In accordance with the shell model(68), the nickel nucleus with a total of 28 protons is a closed shell nucleus for protons with its outermost $f_{7/2}$ shell containing 8 protons. The photoprotons emitted by nickel in the low-energy range below 8 MeV were found to be isotropic in agreement with the statistical theory. The high-energy photoprotons developed a slight anisotropy in which the $b/a$ ratio reached a maximum.
value of 0.6 at 9 MeV. This anisotropy is probably the result of "resonance direct" transitions from the closed $f_{7/2}$ shell. Although El photon absorption is predominantly through the $l \rightarrow l + 1$ transition, $f \rightarrow g$, with a b/a ratio of 0.8, emission through these transitions is strongly inhibited by the centrifugal barrier so that the contribution of the $l \rightarrow l - 1$ transition, $f \rightarrow d$, with a b/a ratio of 0.25, to the angular distribution of the high-energy photoprotons is probably quite considerable. The calculations in Table B.1 on copper would seem to indicate that I.P.M. transitions from the lower proton energy levels are of negligible intensity for both nickel and copper.

A pronounced shell model effect is indicated by the radically increased anisotropy of the angular distributions of photoprotons from copper in which one $2p_{3/2}$ proton is added to the closed $f_{7/2}$ shell of nickel. Even the low-energy photoprotons from copper were anisotropic with the b/a ratio gradually increasing from 0.4 at 5 MeV to 0.9 at 9 MeV. Above this energy the b/a ratio suddenly increased to a value of 4. The anisotropy of the low-energy photoprotons appears to be the result of the combined contributions of statistical emission and direct emission by way of the $f \rightarrow g$ transition. From Table B.1 the gradual increase in anisotropy with photoproton energy then seems to be indicative of the increased relative strength of the direct emission due to the contributions from the $p \rightarrow d$ and $f \rightarrow d$ transitions.

Although the large b/a ratios for the high-energy photoprotons may not be statistically significant, the calculations in Table B.1 indicate that this high anisotropy might be expected for photoprotons of this energy, from copper, since the conditions of Eichler and Weidenmüller(33)
are fulfilled. This high anisotropy is then attributed to interference between the $p \rightarrow d$ and $p \rightarrow s$ transitions.

These results are in good agreement with the work of Lejkin et al.\textsuperscript{(58)} at 19 and 24 MeV. These authors also found that the angular distributions became more isotropic when the bremsstrahlung energy was increased to 30.5 MeV. The increasing isotropy was attributed to the increased contribution of proton transitions from the $f_{7/2}$ level.

The addition of a second $2p_{3/2}$ proton results in a marked decrease in the anisotropy of the photoprotons from the zinc nucleus. The low-energy photoprotons returned to the isotropic angular distribution which was observed for the nickel nucleus. Above 6 MeV a definite anisotropy developed which reached a maximum $b/a$ ratio of 1.1 at 8 MeV and then declined monotonically to 0.4 at 11 MeV.

As was the case with nickel the low-energy photoprotons appear to be the result of a statistical de-excitation process. The increased statistical yield over that observed in copper could be attributed to the odd-even effects of Osokina.\textsuperscript{(81)} It is also reasonable to expect that the strong $p \rightarrow d$ transition seen in copper, would be shifted down in energy by one or two MeV due to the increased separation energy of the two $2p_{3/2}$ protons from the zinc nucleus. The results of this work give no evidence of the pure $\sin^2 \theta$ distribution which was found by Osokina et al.\textsuperscript{(82)} for the high-energy photoprotons using a 20.8 MeV bremsstrahlung spectrum; however, they are in good agreement with their results at 23.3 MeV. It then appears that the high anisotropies are lost very quickly as the increasing bremsstrahlung energy reaches down strongly.
into the $f_{7/2}$ level. It can be recalled from the discussion of the results on copper that the contribution from the $f_{7/2}$ level did not become too significant until the energy of the bremsstrahlung increased to above 25 MeV. Since this effect is indicated with a 22 MeV bremsstrahlung energy in the case of the zinc nucleus, it appears that the $f_{7/2}$ level in zinc is several MeV higher than in the copper.

For the high-energy photoprotons from zinc the value of $p$ indicates that the E2 absorption may be as high as 10 to 15 percent. The values of $p$ determined by the least squares fit for nickel and copper indicate only about 2 percent E2 absorption.

The next group of target elements which were studied have proton occupation numbers 40, 41 and 42. Zirconium with 40 protons is the first of these elements and has its outermost shell filled with two $2p_{1/2}$ protons. The angular distributions of the photoprotons from this nucleus are extremely anisotropic, in fact, for protons greater than 8 MeV, the $b/a$ ratios are all greater than 1.5. Once again the outermost protons belong to a p-shell so that these large anisotropies may be caused by the dipole-dipole interaction mechanism of Eichler and Weidenmüller(55). At low photoproton energies between 4 and 7 MeV the $b/a$ ratio varies between 1.0 and 1.5 indicating that for this nucleus the contribution of the statistical emission is small even at these energies. This is in sharp contrast to the previously analysed even-Z nuclei whose low-energy photoprotons appeared to be the result of a statistical process. The anisotropies of these protons indicate that they arise from a $p \rightarrow d$ transition with a $b/a$ ratio of 1.5.
In order to determine if the large anisotropies observed for the high-energy photoprotons from zirconium were indeed reasonable on the basis of dipole-dipole interference, the relative emission strengths of the various I.P.M. transitions for protons were calculated in Table B.2. It can be seen from these calculations that the high-energy photoprotons are the result of transitions from the \(2p_{1/2}\) and \(2p_{3/2}\) levels. Since the transitions from both of these levels to the \(s\)- and \(d\)-levels have approximately the same transition energy and relative emission strengths, the protons resulting from these transitions would be expected to show an almost pure \(\sin^2\theta\) distribution according to Eichler and Weidemüller(33). The high anisotropies observed for the high-energy photoprotons is then expected to be the result of interference between the \(p \rightarrow d\) and \(p \rightarrow s\) transitions from these levels.

The E2 absorption is negligible for this nucleus except at the highest photoproton energies where it may contribute a few percent.

For the niobium nucleus one proton is added to the closed-shell configuration of zirconium. This proton which goes into the \(g_{9/2}\) shell causes a marked change in anisotropy as in the case of the copper nucleus when one proton was added to a closed-shell configuration. This time, however, the b/a ratios are considerably reduced to values between 0.4 and 0.8 for most of the proton energy range. It appears then that the high anisotropies of zirconium are significantly decreased by contributions from \(f \rightarrow l\) transitions. The \(g \rightarrow f\) transition is obviously suspected; however, its strength is expected to be limited as the \(g_{9/2}\) shell is only one-tenth filled. Since the binding energy of niobium is about three
MeV less than in cadmium, the photoprotons emitted from a given shell in niobium would be expected to be about three MeV more energetic than those emitted from the same shell in cadmium. It can then be seen from Table B.3 that the $f \rightarrow d$ transitions from niobium would be expected to contribute strongly due to the increased transmission coefficients of the higher-energy photoprotons. This fact is substantiated by the results of Osokina (81) who found that photoprotons in the energy range from 6 to 10 MeV had their $b/a$ ratios reduce from 1.23 for a maximum bremsstrahlung energy of 19.5 MeV to 0.84 for a 23.5 MeV bremsstrahlung spectrum. Osokina associated this reduction in anisotropy with the onset of $f \rightarrow l - 1$ transitions. The weak anisotropy found for the low-energy photoprotons is also probably the result of $f \rightarrow l - 1$ transitions from the $f$-levels. The forward-backward asymmetry of the angular distributions of the high-energy photoprotons indicates between 5 and 10 percent $E_2$ absorption.

In the molybdenum nucleus there are two $g_{9/2}$ protons outside of the filled $2p_{1/2}$ shell for a total of 42 protons. For this nucleus the low-energy photoprotons are practically isotropic and probably result from a statistical process. Since the statistical yield of photoprotons from the previous even-Z nucleus, zirconium, was negligible, it is interesting to see the return of a strong statistical component from molybdenum. This effect is discussed in more detail in the following section on photoproton energy distributions.

Above 8 MeV the $b/a$ ratio increases steadily to a value of 1.1 at 10 MeV. This steady rise in anisotropy is probably the result of the
increasing contribution of direct photoprotons from the $p \rightarrow d$ transitions relative to the statistical contribution. Due to the increased binding energy in the molybdenum nucleus, the $l \rightarrow l - 1$ transitions from the $f$-levels are not expected to contribute significantly as they did in niobium. The $g \rightarrow f$ transitions are also expected to be weak since the shell is only 20 percent filled. Since the very high anisotropy observed for photoprotons greater than 10 MeV cannot be considered significant, it is not associated with dipole-dipole interference.

Although the $E2$ absorption was small for most of the photoprotons, the values of $p$ in the proton energy range from 7 to 10 MeV indicates about 10 percent $E2$ absorption.

The last group of elements studied in this experiment were in the vicinity of $Z=50$. The first of these elements is silver with a total of 47 protons, eight of which are in the $g_{9/2}$ shell. The ground state configuration of this nucleus according to Klinkenberg(56) is a $2p_{1/2}$ state. The angular distributions of the photoprotons from this nucleus are again extremely anisotropic, in fact, the $b/a$ ratios for all protons above 3 MeV are greater than 1.5. It can be recalled that the angular distributions of photoprotons from copper and zirconium were both highly anisotropic and that this anisotropy was associated with transitions from filled or partially filled p-shells. This large anisotropy is then attributed to interference between the $p \rightarrow d$ and $p \rightarrow s$ transitions by the single $2p_{1/2}$ proton and the protons from the filled $2p_{3/2}$ level.

The $b/a$ ratio of the 6 to 3 MeV photoprotons is 0.7 indicating that even the lower-energy photoprotons are the result of direct transitions, in particular $l \rightarrow l + 1$ transitions for which the $b/a$ ratio is always greater than
0.5. Since the direct emission of \( g \rightarrow h \) photoprotons is strongly inhibited by the centrifugal barrier, the lower-energy photoprotons are likely the result of \( p \rightarrow d \) transitions from the \( 2p_{3/2} \) level. These transitions give a \( b/a \) ratio of 1.5; however, \( \ell \rightarrow \ell - 1 \) transitions from the \( f \)-levels could be expected to decrease this value significantly. In the case of silver E2 absorption is negligible except for the highest-energy photoprotons where it may be as high as 10 percent.

The addition of another proton refills the \( 2p_{1/2} \) level and results in the cadmium nucleus. The low-energy photoprotons from this nucleus have an isotropic distribution and would appear to result from a statistical process. Since the energy distribution of the photoprotons from cadmium does not resemble the predictions of the statistical theory, it seems more likely that the \( p \rightarrow s \) transitions, with an isotropic distribution, from the \( 2p_{1/2} \) and \( 2p_{3/2} \) levels are causing the isotropy. Although the calculations of Table B.3 indicate that the \( p \rightarrow d \) transitions emit photoprotons of slightly lower energy than the \( p \rightarrow s \) transitions, this result would indicate that the converse is true. It can also be seen from these calculations that the \( f_{5/2} \rightarrow 2d_{5/2} \) transitions, with a \( b/a \) ratio of 0.25, are contributing to the relative isotropy of the low-energy photoprotons.

At 8 MeV the \( b/a \) ratio increased to a value in excess of that expected from a pure \( p \rightarrow d \) transition. It is interesting to see that the calculations, presented in Table B.3, indicate that dipole-dipole interference between the \( p \rightarrow s \) and \( p \rightarrow d \) transitions could reasonably be expected to produce a large anisotropy for photoprotons of this energy, particularly, if the photoprotons from the \( p \rightarrow d \) transitions were slightly more energetic than those from the \( p \rightarrow s \) transitions. Above 8 MeV the \( b/a \) ratio decreases to about 1.0. It appears then that the \( p \rightarrow s \) transition strength is disappearing and that
the high-energy photoprotons result mainly from $p \rightarrow d$ transitions. The anisotropy of these photoprotons is probably further reduced by the contribution of the $g \rightarrow f$ transitions, with a $b/a$ ratio of 0.3. The forward-backward asymmetry of the photoproton angular distributions indicates that the E2 absorption is about 5 percent.

The tin nucleus, with $Z=50$, has a closed $g_{9/2}$ shell. The photoprotons in the 6 to 8 MeV energy range from tin have an isotropic angular distribution as in the case of cadmium. It again seems most probable that the $p \rightarrow s$ transitions are causing the isotropy. The $b/a$ ratios of the high-energy photoprotons varies between 0.4 and 0.6. With the closing of the $g_{9/2}$ shell, transitions from this shell are expected to become more probable so that the angular distributions of these photoprotons are probably strongly influenced by the $g \rightarrow f$ transitions. The influence of the $p \rightarrow d$ transitions which has been quite noticeable up to now for the nuclei in this Z-range seems to have decreased significantly. E2 absorption by this nucleus was restricted to proton energies above 9 MeV where it may be as high as 15 percent.

5.2.2 Energy distributions

The energy distributions of the photoprotons emitted at 90 degrees to the photon beam are shown in Figure 4.2 along with the predicted statistical spectra, approximately normalized to the low-energy part of the photoproton energy spectrum. The fit of the statistical theory to the experimental results is reasonably good in the region of Z about 28; however, in the Z-range around 40 the fit is becoming rather poor. The only nucleus in the range of atomic numbers near 50 for which the statistical spectrum was calculated was silver for which the fit can be seen to be poor. The energy distributions for both
Cadmium and tin also bear little resemblance to a typical statistical spectrum. The increasingly poor fit of the statistical theory of proton evaporation from a heated nucleus with increasing atomic number is a strong indication of the increasing importance of "resonance direct" emission.

In the case of all nine target elements the experimental energy distributions tend to increase above the predictions of the statistical theory as the proton energy increases above about seven MeV. This increased yield of high-energy photoprotons which cannot be accounted for on the basis of the statistical theory becomes more marked as the atomic number increases for these medium-weight nuclei.

In addition to this effect which is a function of increasing Z-value, it can be seen that the statistical energy spectra give a poorer fit to the experimental results for the odd-Z nuclei, copper and niobium, than for the neighbouring even-Z nuclei. This can be attributed to the decreased yield of statistical photoprotons from odd-even nuclei according to Osokina(81).

It is interesting to note that some of the energy spectra taken at 90 degrees to the photon beam have indications of peaks. The statistical significance of these peaks is generally questionable, since they are not seen in the integrated yield data and the yields at the single angle are low. This is particularly true in the case of the nuclei near Z=50 where the targets were thicker and the yields lower. It is interesting, however, to correlate these peaks with the angular distribution data. For the first group of elements near Z=28, the only indication of a peak appears between 8 and 10 MeV in the copper spectrum. From Table B.1 it is seen that the relatively strong p → d and f → d transitions could be expected at this energy. The measured angular distribution is in accord with this result.

In the second group of target elements, the 90 degree energy spectra
from zirconium and niobium appear to have some peaks. In particular, zirconium has a peak centred at about 8 MeV. It is interesting to note that the b/a ratio also peaks in this photoproton energy range. It can be recalled that this large anisotropy was attributed to interference between the strong \( p \rightarrow d \) and \( p \rightarrow s \) transitions from the p-levels in zirconium.

The energy spectrum of the photoprotons from niobium has a number of interesting features. The small yield of photoprotons below 5 MeV is in agreement with other results (81,98). This low yield is likely the result of a small contribution of statistical photoprotons due to the odd-even effects described by Osokina (81). The energy spectrum also appears to contain two peaks, the one centred at 7 MeV and the other centred at 9 MeV. The lower of these two peaks corresponds to the maximum b/a ratio in the neighbourhood of its energy while the upper peak corresponds to the minimum b/a ratio in its neighbourhood; consequently, these energies may correspond to the maximum contributions from the \( \ell \rightarrow \ell + 1 \), \( p \rightarrow d \), and \( \ell \rightarrow \ell - 1 \), \( f \rightarrow d \), transitions respectively.

The energy spectrum of photoprotons from molybdenum gives the best fit to the statistical theory of the elements in this group. This is rather interesting because the abundances of the even Z-odd N isotopes total 25 percent in natural molybdenum. It can be recalled from section 5.1 that the relative statistical photoproton to photoneutron yield ratio is increased in isotopes of this type (81). The isotropic angular distributions, measured in this experiment, for all photoproton energies below 8 MeV, add further support to the contention that these photoprotons result from a statistical process.

In the last group of elements near \( Z=50 \), a peak can be seen in the energy spectrum from cadmium at 8 MeV. This peak corresponds to the sharp
anisotropic peak in the b/a ratios which was attributed to interference between the p→d and p→s transitions.

There is also indication of a high-energy peak present in the energy spectrum of photoprotons from tin. This peak, which lies between 8 and 10 MeV, has a b/a ratio of 0.33. This is in excellent agreement with the b/a ratio of the g→f transition which is expected to make a large contribution with the closing of the g_{9/2} shell.

The small yield of low-energy photoprotons from these nuclei is attributed to the decreasing yield of statistical photoprotons due to the increasing Coulomb barrier height.

The integrated yields of photoprotons for a number of overlapping two-MeV intervals are shown in Figure 4.4. The statistical energy distributions, approximately normalized to fit the yields of low-energy photoprotons, are also shown. As in the case of the energy distributions at 90 degrees, the experimental integrated yields of photoprotons tend to increase above the statistical curves for all of the target nuclei as the photoproton energy increases. The anomalous yield of high-energy photoprotons becomes more and more apparent as the atomic number increases from 28 to 50. These photoprotons, which are in excess of the statistical predictions, must arise from a direct interaction mechanism. The reduction of the statistical photoproton yield from the odd-even nuclei, copper and niobium, according to Osokina (81) again results in a poorer fit to the statistical theory.

5.2.3 Integrated yields

In Figure 4.6 the integrated yields of photoprotons for the two energy ranges (i) > 8 MeV and (ii) < 8 MeV are presented. The high photoproton

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yield from nickel relative to copper has been previously observed for both the high-energy photoprotons(72) and the total photoprotons(58). The increased yield of high-energy photoprotons from nickel indicates enhanced absorption by the closed $f_{7/2}$ shell. The yield of low-energy photoprotons from nickel is then favoured over copper both by the enhancement from the closed shell and the odd-even effects which make the emission of statistical photoprotons highly unfavoured relative to statistical photoneutrons in copper(81). It can be seen from Figure 4.6 that the low-energy photoproton yield from zinc is significantly higher than from copper. This fact is also attributed to the odd-even effects of Osakina(81).

In agreement with the results of Mitchell and McNeill(72), no enhancement is found in the yield of high-energy photoprotons from the closed $2p_{1/2}$ shell in zirconium. In this case, however, the enhancement factors are expected to be smaller due to the small $J$ value. It can also be seen from Table 2.1 that the enhancement factors are smaller for the $J-1/2$ levels than for the $J+1/2$ levels. The small yield of photoprotons below 8 MeV from niobium relative to its even-Z neighbours is again attributed to the odd-even effect.

It can be recalled from section 5.1 that the statistical yield of photoprotons from even-odd isotopes is expected to be significantly increased due to the odd-even effects. It is then interesting to observe that the 25 percent abundance of even-odd molybdenum isotopes leads to a significantly increased statistical photoproton yield. It is, however, rather anomalous that the 11 percent abundance of even-odd zirconium isotopes does not lead to a significant statistical yield of photoprotons. This is probably associated with the relatively large transmission coefficients of the photoprotons from the strongest transitions in zirconium resulting in strong "resonance direct"
In the range of Z near 50 it can be seen from Figure 4.6 that the yield of low-energy photoprotons is significantly reduced. This decrease is associated with a low yield of statistical photoprotons due to the increased height of the Coulomb barrier. The odd-even effects, previously observed, are not noticeable now since they affect only the statistical yields. With the closing of the $g_{9/2}$ shell, no enhancement is observed in the integrated yield of photoprotons from tin. It can be recalled, however, that enhanced absorption by the closed $g_{9/2}$ shell was observed in the angular distribution data where the effect of transitions from the p-levels appeared to be considerably reduced. The lack of an enhanced yield from tin is then associated with the low barrier penetrability of the photoprotons resulting from the transitions of the $g_{9/2}$ level.

The total integrated yields of photoprotons are shown along with the predictions of the statistical and "resonance direct" theories in Figure 4.7. The errors indicated are due to counting statistics only. The yields are seen to be in reasonably good agreement with the predictions of the statistical theory for the nuclei with $Z \leq 42$. In the calculation of the theoretical statistical curve, however, odd-even effects were not taken into account so that the strong fluctuations, observed in the experimental yields, are smoothed out in this curve. Although the theoretical curves indicate a strong predominance of statistical over "resonance direct" yields in this Z-range, there is strong evidence presented in previous sections of this thesis to suggest that the contribution of the "resonance direct" interactions is significant, particularly, from zirconium and the odd-even nuclei, copper and niobium. In the Z-range near 50, the theoretical curves of Figure 4.7 suggest that the "resonance direct" interactions are expected to predominate.
The poor fit of the experimental energy distributions to the predictions of the statistical theory and the lack of odd-even effects, present for the lower Z-values, certainly suggest that the photoprotons from these elements arise mainly from "resonance direct" interactions.
CHAPTER VI

CONCLUSIONS

Previous to the work described in this thesis, extensive measurements of the photoproton angular distributions for medium-weight nuclei with $Z$ between 28 and 50 were available(72). In the present experiment, however, the simultaneous measurement of the energy spectra and the angular distributions provides a clearer understanding of these nuclear reactions.

The comparison of these experimental data with the theoretical models of nuclear photodisintegration, in particular, the "resonance direct" model of Wilkinson(109), indicates that no single model is presently capable of giving a complete description of photoproton reactions. The "resonance direct" model can account satisfactorily for many of the features observed in these data including the energy and angular distributions of the low-energy statistical component of the photoprotons from the even-$Z$ nuclei with $Z \leq 42$, the anomalous yield and the majority of the angular distributions of the high-energy photoprotons, the predominance of "resonance direct" interactions in the nuclei with $Z > 46$ and the enhanced absorption by nucleons in filled shells. There are, however, a number of results which cannot be explained within the framework of this model.

The smooth variation with $Z$ of the yield of statistical photoprotons, which the "resonance direct" model predicts, is not observed in the results of these measurements, which reveal strong pairing force interactions. These interactions lead to the striking odd-even effects observed in the odd $Z$-even $N$ nuclei, copper and niobium. These odd-even effects are the fact that the photoprotons from these nuclei result almost completely from
"resonance direct" interactions has not previously been observed.

The large anisotropies, which have been previously reported for nuclei with Z between 28 and 50(55,63,72,82) and are observed in this experiment for the nuclei, copper, zirconium and silver, cannot be reconciled on the basis of the "resonance direct" model. These anisotropies must be explained in terms of the dipole-dipole interference of Eichler and Weidenmüller(33). It is interesting to observe that the uppermost level in each of these three nuclei is a filled or partially filled p-shell for which the anisotropies are expected to be the largest on the basis of dipole-dipole interference.

In emphasizing the enhanced absorption by nucleons in a closed shell, the "resonance direct" model ignores the effect of the valence nucleons which have been shown, in this experiment, to have a striking effect on both the energy and angular distributions. The copper nucleus is a good case in point since the angular distributions of all the photoprotons from copper are dominated by the transitions of the single $2p_{3/2}$ proton.

In order to obtain agreement between the predicted and experimental giant resonance energies it may be recalled that the "resonance direct" model resorted to an effective nucleon mass within the nucleus; however, the results of particle–induced direct reactions indicate the contrary. This suggests that perhaps only one level or two, in the case of spin-orbit coupling, may be shifted up in energy as in the case of particle-hole calculations(11,12,44). Due to the difficulties involved in these calculations they have been restricted to doubly closed-shell nuclei; however, it is interesting to postulate on the basis of these experimental results, which levels could be expected to form the dipole state. In the closed shell nuclei,
nickel and tin, the angular distributions found in this experiment seem to indicate that the dipole state would involve the \( f_{7/2} \rightarrow 2d_{5/2} \) and the \( g_{9/2} \rightarrow 2f_{7/2} \) transitions respectively. In the case of the closed shell nucleus, zirconium, the dipole state could be split into \( p \rightarrow d \) transitions from the \( 2p_{1/2} \) and \( 2p_{3/2} \) levels. For the copper nucleus it would appear that the \( 2p_{3/2} \rightarrow 2d_{5/2} \) transitions would be shifted up in energy and form the dipole state. The results of this experiment suggest, however, that there is strong configurational mixing for these medium-weight nuclei.

Although a number of new and interesting features have been observed in this experiment, which have strongly indicated the limitations of the I.P.M., the interpretation of the results is still rather tentative. The basic limitation in the experiment was the relatively thick targets, which had to be used in order to obtain a reasonable statistical accuracy. With the advent of high-intensity machines, however, future experiments with the use of thin targets should enable a more critical test of the current theories.
ANGULAR DISTRIBUTION COEFFICIENTS AND ERRORS BY THE METHOD OF LEAST SQUARES

The photoproton angular distributions were fitted to equation 2.53, which can be written in the form

\[ Y = a + bs^2(1 + pc)^2 \]  \hspace{1cm} A.1

if the functions \( \sin \theta \) and \( \cos \theta \) are replaced by the letters \( s \) and \( c \) respectively. The angular distribution coefficients \( a \), \( b \) and \( p \) were determined by the method of least squares, in which the sum over the angles of observation of the squares of the differences between the experimental yields and the calculated curve, \( N \), is minimized. Since the coefficients are to be calculated to make \( N \) a minimum, the derivatives with respect to the three coefficients must vanish, that is

\[ \frac{\partial N}{\partial a} = \frac{\partial N}{\partial b} = \frac{\partial N}{\partial p} = 0 \]  \hspace{1cm} A.2

These three equations, known as the normal equations, are readily calculated to be

\[
\begin{align*}
\sum a + b \sum s^2(1 + pc)^2 & = \sum y & = 0 \\
\sum a s^2(1 + pc)^2 + b \sum s^4(1 + pc)^4 & = \sum y s^2(1 + pc)^2 & = 0 \\
ab s^2 c(1 + pc) + b^2 \sum s^4 c(1 + pc)^3 & = \sum y s^2 c(1 + pc) & = 0
\end{align*}
\]  \hspace{1cm} A.3

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where $n$ is the number of angles of observation and

$y$ is the experimental yield at a particular angle of observation.

The summations are to be taken over all the angles for which experimental measurements were made. This set of equations is linear in $a$ and $b$ if the third equation is divided through by $b$ (the case for $b=0$ is excluded).

Solutions then exist for the coefficients $a$ and $b$ only if the coefficient $p$ satisfies the following equation

\[
\begin{vmatrix}
    n & \sum s^2(l + pc)^2 & \sum y \\
    \sum s^2(l + pc)^2 & \sum s^4(l + pc)^4 & \sum ys^2(l + pc)^2 \\
    \sum s^2c(l + pc) & \sum s^4c(l + pc)^3 & \sum ys^2c(l + pc)
\end{vmatrix} = 0 \quad A.4
\]

If the third row of equation A.4 is multiplied by $p$ and subtracted from the second row then this equation reduces to the form

\[
\begin{vmatrix}
    n & \sum s^2(l + pc)^2 & \sum y \\
    \sum s^2(l + pc) & \sum s^4(l + pc)^3 & \sum ys^2(l + pc) \\
    \sum s^2c(l + pc) & \sum s^4c(l + pc)^3 & \sum ys^2c(l + pc)
\end{vmatrix} = 0 \quad A.5
\]

The coefficient $p$ is then one of the zeros of a fourth degree polynomial.

After solving the polynomial, the physically unacceptable values for $p$ were discarded. The values of $a$ and $b$ were then determined by solving the first
two equations of A.3.

The standard deviations of the angular distribution coefficients were calculated using the standard deviations due to the counting statistics, $\sigma_y$. Since the coefficient, $\sigma_p$, is a function of the yields at the various angles of observation, its standard deviation, $\sigma_p$, was determined from the approximate formula (71)

$$
(\sigma_p)^2 = \sum \left( \frac{\partial p}{\partial y} \right)^2 (\sigma_y)^2
$$

A.6

The values of $\frac{\partial p}{\partial y}$ were evaluated from the polynomial A.5.

The standard deviations of the coefficients $a$ and $b$ were calculated using the formula

$$
(\sigma_a)^2 = \left( \frac{\partial a}{\partial p} \right)^2 (\sigma_p)^2 + \sum \left( \frac{\partial a}{\partial y} \right)^2 (\sigma_y)^2
$$

A.7

with a similar equation for $\sigma_b$.

The coefficient, $a$, is inherently positive and consequently asymmetric errors were quoted for the cases where $\sigma_a > a$ (e.g., $a + \sigma_a$).

The errors in $b/a$ were then calculated using the equation

$$
(\sigma_{b/a})^2 = \left( \frac{\partial b/a}{\partial a} \right)^2 (\sigma_a)^2 + \left( \frac{\partial b/a}{\partial b} \right)^2 (\sigma_b)^2
$$

A.8

In the cases where asymmetric errors were quoted on $\sigma_a$, these were employed in equation A.8 to obtain the asymmetric errors quoted on $b/a$. The effect of the statistical correlation between $\sigma_b$ and $\sigma_a$ has been omitted in these calculations. As this omission leads to an overestimate of the errors on
b/a, these errors can be considered as upper limits on $\Sigma_{b/a}$.

Since it was necessary to determine the angular distribution coefficients for a large number of least squares fits, all of the calculations were carried out on the IBM 7094. For those targets in which the yields at the most forward angle were not reliable, due to the pile-up of electron pulses, a separate program had to be written, which fitted the least squares curve to four instead of the regular five experimental points.
APPENDIX B

CALCULATION OF THE PHOTOPROTON YIELDS

In this appendix the relative probability of photoproton emission from the various I.P.M. transitions is calculated in order to lend support to the qualitative discussion of the experimental results given in section 5.2. The calculations are carried out for copper, zirconium and cadmium, since their energy and angular distributions are particularly interesting and one of these elements is contained in each of the three groups of elements which were studied.

The E1 transition strengths, $S_{J^D_n}^{\nu',\nu}$, are calculated using Tables 2.1 and 2.2. When the proton shells are not filled, as in the case of both copper and cadmium, the appropriate correction\(^{(14)}\) is made to the transition strength. In order to determine the energy of the I.P.M. transition, $E_\gamma$, it is necessary to have an estimate of the proton energy levels. The level ordering for a modified square well\(^{(68)}\) is shown in Figure 2.2. The energy scale used in this diagram was calculated by Schröder\(^{(89)}\) with reasonable assumptions about spin-orbit coupling. In order to correct for the effective nucleon mass within the nucleus\(^{(109)}\), the energy spacing between unoccupied levels of the nucleus is doubled. The ground state configurations are taken from Klinkenberg\(^{(56)}\).

The separation energies of the protons are required to calculate the energy of the emitted photoprotons, $E_p$. This information was obtained from Nuclear Data Sheets\(^{(78)}\) and Feather\(^{(37)}\). The transmission coefficients, $T_p(E_p)$ for the protons through the potential barrier are taken from Figure 2.1. The relative intensity of the bremsstrahlung spectrum at the various

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transition energies, \( P(E_{\gamma}) \), is taken from the 22 MeV spectrum calculated by Katz and Cameron (55). The relative emission strengths for the various proton transitions are then calculated using the following equation

\[
Y_{j \rightarrow j'} = S J N_{\gamma} E_{T}\gamma P(E_{\gamma})
\]

The results of these calculations are presented in the Tables B.1, B.2 and B.3.
<table>
<thead>
<tr>
<th>Transition</th>
<th>Absorption Strength</th>
<th>$E_\gamma$ (MeV)</th>
<th>$E_p$ (MeV)</th>
<th>$T(E_p)$</th>
<th>$P(E_\gamma)$</th>
<th>Relative Emission Strength</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2p_{3/2} \rightarrow 2d_{5/2}$</td>
<td>0.084</td>
<td>15.2</td>
<td>9.2</td>
<td>0.51</td>
<td>0.69</td>
<td>0.452</td>
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<tr>
<td>$\rightarrow 2d_{3/2}$</td>
<td>0.009</td>
<td>18.8</td>
<td>12.8</td>
<td>0.73</td>
<td>0.43</td>
<td>0.054</td>
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<tr>
<td>$f_{7/2} \rightarrow g_{9/2}$</td>
<td>1.09</td>
<td>7.6</td>
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<tr>
<td>$\rightarrow g_{7/2}$</td>
<td>0.031</td>
<td>17.6</td>
<td>7.6</td>
<td>0.02</td>
<td>0.53</td>
<td>0.006</td>
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<tr>
<td>$d_{5/2} \rightarrow f_{5/2}$</td>
<td>0.53</td>
<td>8.1</td>
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<tr>
<td>$\rightarrow 2p_{3/2}$</td>
<td>0.009</td>
<td>8.1</td>
<td>-</td>
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<tr>
<td>$2s \rightarrow 2p_{3/2}$</td>
<td>0.153</td>
<td>10.4</td>
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<tr>
<td>$\rightarrow 2p_{1/2}$</td>
<td>0.077</td>
<td>12.5</td>
<td>-</td>
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TABLE B.2

Zirconium Z = 40

<table>
<thead>
<tr>
<th>A</th>
<th>Abundance</th>
<th>(γ,p)</th>
<th>(γ,n)</th>
</tr>
</thead>
<tbody>
<tr>
<td>90</td>
<td>51.5</td>
<td>8.4</td>
<td>12.0</td>
</tr>
<tr>
<td>91</td>
<td>11.2</td>
<td>8.7</td>
<td>7.2</td>
</tr>
<tr>
<td>92</td>
<td>17.1</td>
<td>9.4</td>
<td>8.6</td>
</tr>
<tr>
<td>94</td>
<td>17.4</td>
<td>10.3</td>
<td>8.2</td>
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<tr>
<td>96</td>
<td>2.8</td>
<td>12.1</td>
<td>7.8</td>
</tr>
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</table>

EL Transitions for Protons

<table>
<thead>
<tr>
<th>Transition</th>
<th>Absorption Strength $S_{2J}^{2J', n', l'}$</th>
<th>$E_\gamma$ (MeV)</th>
<th>$E_p$ (MeV)</th>
<th>$T_{J', l'}$($E_p$)</th>
<th>$F(E_\gamma)$</th>
<th>Relative Emission Strength</th>
</tr>
</thead>
<tbody>
<tr>
<td>$2p_{1/2} \rightarrow 2d_{3/2}$</td>
<td>0.187</td>
<td>17.0</td>
<td>8.5</td>
<td>0.32</td>
<td>0.58</td>
<td>0.59</td>
</tr>
<tr>
<td>$\rightarrow 3s$</td>
<td>0.04</td>
<td>17.5</td>
<td>9.0</td>
<td>0.55</td>
<td>0.55</td>
<td>0.21</td>
</tr>
<tr>
<td>$2p_{3/2} \rightarrow 2d_{5/2}$</td>
<td>0.336</td>
<td>14.3</td>
<td>4.8</td>
<td>0.02</td>
<td>0.78</td>
<td>0.075</td>
</tr>
<tr>
<td>$\rightarrow 2d_{3/2}$</td>
<td>0.037</td>
<td>17.9</td>
<td>8.4</td>
<td>0.31</td>
<td>0.51</td>
<td>0.105</td>
</tr>
<tr>
<td>$\rightarrow 3s$</td>
<td>0.08</td>
<td>18.3</td>
<td>8.8</td>
<td>0.54</td>
<td>0.47</td>
<td>0.37</td>
</tr>
<tr>
<td>$f_{5/2} \rightarrow e_{7/2}$</td>
<td>0.84</td>
<td>12.7</td>
<td>3.2</td>
<td>-</td>
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<tr>
<td>$\rightarrow 2d_{5/2}$</td>
<td>0.004</td>
<td>14.2</td>
<td>4.7</td>
<td>-</td>
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<tr>
<td>$\rightarrow 2d_{3/2}$</td>
<td>0.06</td>
<td>17.9</td>
<td>8.4</td>
<td>0.31</td>
<td>0.51</td>
<td>0.17</td>
</tr>
<tr>
<td>$f_{7/2} \rightarrow e_{9/2}$</td>
<td>1.09</td>
<td>6.6</td>
<td>-</td>
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<td>-</td>
<td>-</td>
</tr>
<tr>
<td>$\rightarrow e_{7/2}$</td>
<td>0.031</td>
<td>16.7</td>
<td>3.2</td>
<td>-</td>
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<td>-</td>
</tr>
<tr>
<td>$\rightarrow 2d_{5/2}$</td>
<td>0.086</td>
<td>18.2</td>
<td>4.7</td>
<td>-</td>
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</table>
# TABLE B.3

**Cadmium** \( Z = 48 \)

**El Transitions for Protons**

<table>
<thead>
<tr>
<th>Transitions</th>
<th>Absorption Strength ( S_{J^D} )</th>
<th>( E_\gamma ) (MeV)</th>
<th>( E_p ) (MeV)</th>
<th>( T_{p'}(E_p) )</th>
<th>( P(E_\gamma) )</th>
<th>Relative Emission Strength</th>
</tr>
</thead>
<tbody>
<tr>
<td>( g_9/2 \rightarrow h_{11/2} )</td>
<td>1.16</td>
<td>13.0</td>
<td>4.0</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow h_{9/2} )</td>
<td>0.021</td>
<td>21.0</td>
<td>12.0</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow 2f_{7/2} )</td>
<td>0.069</td>
<td>&gt;22.0</td>
<td>&gt;13.0</td>
<td>0.46</td>
<td>0.09</td>
<td>0.061</td>
</tr>
<tr>
<td>( 2p_{1/2} \rightarrow 2d_{3/2} )</td>
<td>0.187</td>
<td>16.3</td>
<td>6.6</td>
<td>0.07</td>
<td>0.62</td>
<td>0.130</td>
</tr>
<tr>
<td>( \rightarrow 3s )</td>
<td>0.04</td>
<td>16.7</td>
<td>7.0</td>
<td>0.17</td>
<td>0.58</td>
<td>0.065</td>
</tr>
<tr>
<td>( 2p_{3/2} \rightarrow 2d_{5/2} )</td>
<td>0.336</td>
<td>13.6</td>
<td>2.8</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow 2d_{3/2} )</td>
<td>0.037</td>
<td>17.2</td>
<td>6.4</td>
<td>0.06</td>
<td>0.55</td>
<td>0.021</td>
</tr>
<tr>
<td>( \rightarrow 3s )</td>
<td>0.08</td>
<td>17.6</td>
<td>6.8</td>
<td>0.16</td>
<td>0.52</td>
<td>0.113</td>
</tr>
<tr>
<td>( f_{5/2} \rightarrow g_{7/2} )</td>
<td>0.84</td>
<td>11.9</td>
<td>-</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow 2d_{5/2} )</td>
<td>0.004</td>
<td>13.4</td>
<td>2.6</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow 2d_{3/2} )</td>
<td>0.06</td>
<td>17.2</td>
<td>6.4</td>
<td>0.06</td>
<td>0.55</td>
<td>0.034</td>
</tr>
<tr>
<td>( f_{7/2} \rightarrow g_{7/2} )</td>
<td>0.031</td>
<td>16.0</td>
<td>1.2</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>( \rightarrow 2d_{5/2} )</td>
<td>0.086</td>
<td>17.5</td>
<td>2.7</td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

* Due to the large number of cadmium isotopes, they are not shown.

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