Neutron capture mechanisms in the threshold region

Barry John Allen
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NEUTRON CAPTURE MECHANISMS IN THE THRESHOLD REGION

Thesis submitted for the degree of
Doctor of Philosophy
by
Barry John Allen M.Sc. (Hons)

School of Physics,
University of Wollongong
1978
In memoriam

LOUISA ARENDINA ALLEN, M.B., B.S.
ABSTRACT

Neutron capture mechanisms in the threshold region are investigated through the measurement of high resolution γ-ray spectra and resonance capture cross sections. Initial and final state width correlations are observed for many nuclides across the periodic table, and are indicative of dominant valence and doorway interactions.

A detailed investigation of resonant capture in the isotopes of iron provides strong evidence for these effects. Estimates of the magnitude of the doorway component are obtained for the 3s region when valence and statistical contributions are compared with the average s- and p-wave radiative widths.
SUMMARY

The manifest properties of the neutron capture mechanism have been investigated by high resolution measurements of γ-ray spectra and resonance capture cross sections for neutron energies up to 1000 keV.

The AAEC pulsed Van de Graaff accelerator and on-line computer were used in the γ-ray measurements and complementary capture cross sections were obtained at the Oak Ridge Electron Linear Accelerator (ORELA). Measurements have included a wide range of nuclides, but emphasis is placed on a detailed examination of resonance capture in the isotopes of iron. The properties of capture resonances are investigated and compared with calculations based on various models of the capture reaction.

Resonance capture theory is reviewed in Chapter 1, with emphasis on the definition and role of the statistical, valence and doorway mechanisms, and on their relationship to non-resonant capture and width correlations. Detailed valence calculations are made for nuclides across the periodic table, particularly in the regions of the magic neutron numbers.

The γ-ray spectroscopy system at Lucas Heights includes both NaI and Ge(Li) detectors, and is described in Chapter 2. The ORELA capture facility comprises total energy detectors and a $^6$Li glass neutron monitor. This system is reviewed in Chapter 3, and a detailed account of the resonance scattered neutron background is given. These experimental facilities have provided an abundance of data on the properties of capture reactions.

In Chapter 4 evidence is presented which shows that valence transitions in the threshold region are effectively exempt from relocation of El strength to the giant dipole resonance (GDR). This result is found to apply for s-, p- and d-wave capture.
Extensive evidence for width correlations across the periodic table is presented in Chapter 5. These results are interpreted in terms of statistical and valence effects and, together with evidence from γ-ray spectra, indicate the importance of doorway state interactions.

These general observations are confirmed in a detailed investigation of neutron capture in the isotopes of iron in Chapter 6. Results for $^{54}\text{Fe}$ show resonance-resonance interference, large correlations and a dominant valence component. However, this is not the case in $^{56}\text{Fe}$ where statistical and doorway contributions are strong.

The intrinsic El strength of neutron and γ-ray doorway states is reviewed in Chapter 7, with emphasis on particle-hole and particle-vibrator interactions. Evidence for the p-h interaction is found in a detailed investigation of the capture mechanism in $^{45}\text{Sc}$ and $^{139}\text{La}$.

General conclusions regarding the magnitude and variances of the capture mechanisms in the 3s region are presented in Chapter 8. A comparison of valence and statistical radiative widths, with average s- and p-wave radiative widths, provides strong evidence for doorway state mechanisms in several isotopes.

Finally, an overview of the mechanism of neutron capture is presented in Chapter 9, together with the general conclusions of this thesis.
ACKNOWLEDGEMENTS

The experimental work described in this thesis was undertaken at the Australian Atomic Energy Commission Research Establishment (Lucas Heights) and at Oak Ridge National Laboratory (Oak Ridge, Tennessee, USA), and was supported in part by the United States Energy Research and Development Administration, under contract with Union Carbide Corporation, and by the AAEC.

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Antediluvian Neutron Physics

The origins of neutron physics are found in the first man-made transmutations of the elements in 1919 when Rutherford bombarded nitrogen with α-particles from the radioactive decay of radium. Not long after these experiments, Rutherford was drawn to the conclusion that a neutral particle might exist which would bind the nucleus together, overcoming the repulsive Coulomb forces of the protons. His student, Chadwick, was receptive to this idea and initiated a search for the neutral particle, but without success.

In 1930, Bothe and Becker found that radiation was emitted after α-bombardment of light elements. Webster observed that the so-called Be-radiation (emitted after α-bombardment of Be) was most penetrating in the forward direction, but Chadwick failed to observe any tracks in the newly developed expansion chamber. However, when Curie and Joliot found that Be radiation ejected protons from paraffin, Chadwick was able to show that this radiation resulted from the interaction of neutral particles. By measuring the maximum recoil velocities of ejected H and N atoms, he demonstrated that the neutron had a mass comparable to that of a proton.

The Curie-Joliot team went on to discover artificial radioactivity in 1934 by observing positrons emitted after α-bombardment of Al. While this work attracted a great deal of interest, Enrico Fermi, realising the potential of the neutron as a nuclear probe, began a series of neutron activation experiments using Po-α-Be neutron sources. Fermi's group discovered many new isotopes and observed the slowing down of neutrons in paraffin and the consequent increase in reaction probability.

Neutron capture γ-rays from hydrogen were discovered by Lea in 1934 and Moon and Tillman were able to show that neutrons...
could reach equilibrium with the thermal vibrational energies of paraaffin molecules. Insights into resonance structure came from the observations of neutron capture absorption bands by Fermi and Amaldi (FA35) and Szilard (Sz35). The development of a mechanical velocity selector by Dunning et al. (Du+35) enabled the study of these bands by the neutron time of flight method. In 1936 Bohr (Bo36) and Breit and Wigner (BW36) were able to interpret these absorption bands as resonances or highly excited states of the compound nucleus. Strong absorption of the incident nucleon by the target resulted in an immediate coalescence of the neutron-target system into a many body 'compound state'. Bohr also pointed out (Bo38) that a direct transition from the entrance to the exit channel could also occur.

The Lorentzian form of the Breit-Wigner single level formula is a direct consequence of the assumption of a quasi-stationary state. An exponential decay of this state is obtained in the Fourier transform from energy to time, and the decay constant $\tau$ is related to the resonance width $\Gamma$ by the relation $\tau = h/\Gamma$. In the case of overlapping compound levels, general resonance theories were developed (KP38,We47) which yielded the Breit-Wigner formula in the single level approximation.

With the discovery of fission by Hahn and Strassman (HS39) and Meitner and Frisch (MF39) in 1939, the foundations of neutron physics had been established within seven years of Chadwick's discovery of the neutron. The first reactor achieved criticality in 1942 at Chicago under Fermi's leadership, and in the postwar years reactors became the major source of neutrons. In 1947 Zinn obtained a monochromatic neutron beam via Bragg reflection (Zi47), utilising the property that the wavelength of the thermal neutron (1.82 Å) is comparable to the crystalline lattice spacings.

Refinements in the design of velocity selectors led to the development of the fast choppers (e.g. Se54) which permitted the study of
resonances in the eV region. The development of the electron linear
accelerator as a photoneutron source at Harwell {Wi56} paved the way
for high resolution measurements of resonance reactions at higher
neutron energies. In 1960 Firk and Gibbons {FG60} used terminal
pulsing in a Van de Graaff accelerator to produce nanosecond bursts
of neutrons from the $^7\text{Li}(p,n)$ reaction for time of flight measurements
with keV neutrons. As a result of these developments, increasing
quantities of resonance data have been measured in capture, transmission
and fission experiments, providing data needed for reactor calculations
as well as for an improved understanding of the interactions of neutrons
with nuclei.

In parallel with these advances in neutron sources was the develop­
ment of γ-ray detectors. Thermal capture γ-ray measurements by Groshev
et al. at Dubna {Gr55,GDP60} used a magnetic pair spectrometer, while at
Chalk River Kinsey and Bartholomew employed a Compton spectrometer
{KBW51,BB57,BH58}. Electrons from pair or Compton photon interactions
were passed through a magnetic field and, on detection, provided
intensity and precise γ-ray energy information. These detectors suffered
from low efficiency and could not be used to study resonance capture.

The high efficiency of the NaI detector {Pr52,Dr59} (which converts
the γ-ray energy into light for detection by a photomultiplier tube),
permitted low resolution measurements of resonance capture γ-ray spectra.
These detectors were used for measurements at eV, keV and MeV neutron
energies. Fast capture γ-ray measurements were carried out at Oak Ridge
{BGG62a,b;Bi+65,Be+65,Bi+67} and at Studsvik {BS61a,b;BS62a,b,c} for a
wide range of nuclides and neutron energies.

In the late sixties neutron and γ-ray spectroscopy came of age with
the development of advanced neutron time of flight sources (e.g. BNL fast
chopper, Columbia synchrocyclotron, Oak Ridge Electron Linac Accelerator
(ORELA) and the high resolution Ge(Li) detector {ET64}). The coupling of
the high resolution and moderate efficiencies of these detectors with
the dramatic increase in neutron energy resolution and intensity, led
to a deluge of new data and a revival of interest in the capture
reaction.

The first high resolution $\gamma$-ray measurement after the capture of
neutrons in the keV range was made by the author \{Al68\} using a Ge(Li)
detector and pulsed Van de Graaff accelerator. The author also
participated in the commissioning of a high resolution capture cross
section facility at ORELA, and in the measurement and analyses of a
great many isotopic cross sections.

From this abundance of diluvian data, systematic effects have
emerged which have led to new insights into the nature of the neutron
capture mechanisms.
5.

CHAPTER 1
THEORY OF NEUTRON CAPTURE

Neutron reactions near threshold excite levels at \(5-10\) MeV in the compound nucleus. At these energies Bohr's compound nucleus model \{Bo36\} was initially considered to provide an adequate description of the neutron interaction which exhibited narrow resonances in the scattering and capture cross sections. Since the wave functions of these resonances are highly complicated, their properties are expected to be described by statistical theory.

However, the observation \{KBW51,Gr+58\} in thermal capture spectra of correlations between the reduced \(\gamma\)-ray intensities and the spectroscopic factors of final states, led to the proposal of a direct capture mechanism. Lane and Wilkinson \{LW55\} pointed out that the matrix elements of the \((n,\gamma)\) and \((d,p)\) reactions contain a parentage overlap factor which would yield final state correlations in the extreme case of a unique parent. For the above reactions this is the target ground state. Employing the fractional parentage concept, Bockelman \{Bo59\} calculated direct capture cross sections at thermal energies. In general, these cross sections were too low since that part of the nuclear wave function outside the nuclear radius was not included in the calculation. The important role of the external part of the nuclear wave function was first recognised by Thomas \{Th51\} using the formalism of Breit and Yost \{BY35\}. Calculations of \(3s+2p\) transition strengths in \(^{40}\text{Ca}\) \{MI60\} also demonstrated the importance of the external region.

Beginning with basic dispersion theory, Lane and Lynn \{LL60\} recognised three components of the resonance capture cross section: the compound nucleus (or resonance internal), channel (or resonance external), and direct capture (hard sphere potential and distant resonances). The last two components preferentially feed single particle final states. Furthermore, the resonance external part could give rise to enhanced
γ-ray transitions to single particle states if the resonance reduced neutron width is large.

The foregoing theoretical developments were applied to s-wave capture, mostly at thermal energies. With the advent of improved reactor choppers and pulsed Van de Graaff and linear accelerators, γ-ray spectra could be measured from both s- and p-wave resonances. These measurements also showed final state correlations in many cases {Ch69} and when sufficient resonances could be studied, initial state correlations between the reduced neutron widths and radiative widths were observed {Bl+71, BS75, Bi+76}.

For neutron energies in the giant dipole resonance region, the direct and compound nucleus models could not reproduce the measured capture cross sections, and a collective semi-direct capture mechanism {Br64, CLR65} was introduced to account for the data. This theory assumes that the incident nucleon is captured into a lower orbit, exciting the target nucleus into its giant dipole resonance. This intermediate state subsequently decays by enhanced γ-ray emission. Interference between the direct and semi-direct mechanism {LS58} and the introduction of a complex coupling function {Po76} are required for the description of high energy capture.

In the range of neutron capture reactions from thermal to MeV energies, three distinct processes are needed to account for the data. The statistical model describes resonances with small widths and corresponding long life times (τ ~ 10^{-15} s), while non-resonant thermal capture is accounted for by a short-lived direct process (τ ~ 10^{-22}). The life time for semi-direct capture lies somewhere between these two limits.

The resonance capture reaction can be pictured as a series of two-body interactions {FKL67} beginning with the entrance channel one particle-zero hole state (i.e. lp-0h with respect to the target nucleus),
Fig. 1.1
Schematic representation of non-resonant and resonant capture
exciting doorway states (2p-1h or collective modes) and through a succession of more complex p-h interactions, leading ultimately to the statistical interaction involving many nucleons. The resonance wave function therefore contains many configurations, any of which can undergo a radiative de-excitation when overlap with the final state wave function occurs. In the threshold region, precompound nucleon emission is energetically impossible (apart from potential neutron scattering), but radiative decay, which occurs in the first two stages of the interaction (Fig. 1.1a), is non-resonant in character as a consequence of the uncertainty principle.

The main emphasis of this thesis is on resonant processes which are divided into valence, doorway and statistical mechanisms (Fig. 1.1b). When the resonance is in the entrance channel state, the valence neutron can undergo a radiative transition without perturbing the core. Radiative decay can also occur from the doorway components of a resonance, either by particle-hole annihilation or by particle transition in the presence of an excited core. All other decay modes are grouped together under the heading of statistical interactions, which are reviewed in section 1.1.

A discussion of non-resonant neutron capture is given in section 1.2, where the role of direct capture at thermal energies is described, together with the effects of interference between resonant and non-resonant capture amplitudes. A major contribution to this field has come from threshold photoneutron experiments. Because of its single channel nature, the \((\gamma, n)\) reaction has provided information on resonant, non-resonant interference, as well as initial state correlations \(\rho^I(\Gamma_n, \Gamma_{\gamma})\), spin assignments and intermediate structure. The early Livermore work \((BBB71)\) provided a strong stimulus to neutron capture research, and extensive contributions have followed from Argonne \((Ja74b)\).

Attention is restricted to neutron energies below 1 MeV. At higher
energies the role of the giant dipole resonance (GDR) becomes more important, and Bergqvist \cite{Be76} has reviewed in detail the experimental data and theoretical developments in this energy region. The GDR dominates the $\gamma$-ray strength function at MeV neutron energies and contributes significantly at excitation energies as low as 5-7 MeV. However, additional structure has been observed in this range and Bartholomew and co-workers \cite{Ba+74} have made an intensive study of this effect, particularly in the 181 < A < 208 region.

The theoretical description of valence neutron transitions is presented in section 1.3, where it is shown that enhanced $E1$ transitions can occur in the regions of maxima of the $2p$, $3s$ and $3p$ neutron strength functions. The relationship between valence, doorway interactions and width correlations is formalised in section 1.4.

1.1 STATISTICAL THEORY

The basis of the statistical model is the Bohr condition \cite{Bo36} that the lifetime of the excited state is much longer than the time required for the neutron to traverse the nucleus. Thus the probability of forming the excited state $E_\lambda$ in channel $c$ is independent of the probability of decay in channel $c'$. Kapur and Peierls \cite{KP38} obtained the scattering matrix in the form

$$S_{cc'} = \frac{\gamma_{\lambda c}^* \gamma_{\lambda c'}}{E_\lambda - E - i/2 \Gamma_\lambda}, \quad \ldots \quad (1.1)$$

where $\gamma_{\lambda c}$ is the reduced width amplitude for resonance $\lambda$ in channel $c$. A similar form of the collision matrix can, with appropriate approximations, be obtained from later resonance theories such as the R-matrix theory \cite{LT58}.

The overlap of the resonance and channel wave functions is small, $<\lambda | c >^2 \approx 10^{-4}-10^{-7}$, and for neutron scattering is equal to the dimensionless reduced neutron width $\gamma_n^2/\gamma_{sp}^2$, where $\gamma_{sp}^2$ is the single particle width.

The fundamental assumptions of the statistical model are that
(a) $\gamma_{\lambda c}$ and $\gamma_{\lambda c'}$, have random phases such that $\sum_{\lambda} \gamma_{\lambda c} \gamma_{\lambda c'} = 0$ over many levels in the range $\Delta E$, and that the amplitude correlation coefficient $\rho(\gamma_{\lambda c}, \gamma_{\lambda c'}) = 0$;

(b) the strength function is independent of energy, i.e.

$$\frac{1}{\Delta E} \sum_{\lambda} \gamma_{\lambda c}^2 = \text{constant}.$$ 

Since $\gamma_{\lambda c}$ is the coupling amplitude of entrance channel $c$ with the compound state of energy $E_\lambda$, a zero correlation implies that the statistical properties of decay of the exit channel $c'$ contain no 'memory' of the entrance channel amplitude.

In the case of radiative neutron capture, the statistical description for the primary $\gamma$-ray spectrum is given by

$$f(E_\gamma, E_\lambda) dE_\gamma \propto E_\gamma^3 \frac{\rho(E_\lambda - E_\gamma)}{\rho(E_\lambda)} <\lambda|D|\mu> dE_\gamma,$$  \hspace{1cm} ...(1.2)

where $E_\gamma^3$ is the phase space factor for dipole transitions {BW52} from an initial state at excitation energy $E_\lambda$, $\rho$ is the level density of states at the specified energy, and $<\lambda|D|\mu>$ is the overlap integral for an El transition from resonance $\lambda$ to final state $\mu$. In the statistical model, the resonance wave functions $|\lambda>$ are assumed to be complex and have random overlaps with final states $|\mu>$ such that wide fluctuations in $\gamma$-ray intensities will occur.

Porter and Thomas {PT56} assumed that the partial radiative amplitudes $\gamma_{\lambda \mu}$ should have a normal distribution with zero mean, as is the case for the neutron amplitudes. Since the square of $\gamma_{\lambda \mu}$ is just the reduced partial radiative width $\Gamma_{\lambda \mu}/E_\gamma^3$, the distribution of $\Gamma_{\lambda \mu}/E_\gamma^3$ is described by a $\chi^2$ distribution with one degree of freedom (i.e. a Porter-Thomas distribution),

$$P(x) dx = \sqrt{\frac{2}{\pi}} x^{-1/2} e^{-x/2} dx,$$

$$x = \frac{\Gamma_{\lambda \mu}}{\Gamma_{\lambda \mu}}.$$

where
Many attempts have been made to measure these distributions for nuclides with $A > 100$ with level spacings less than 50 eV. Overall, the results are now consistent with $\nu \sim 1$ (Ch69,Bo70), but they are quite sensitive to the quality of the measurements. Some exceptions to $\nu = 1$ persist, for example, certain transitions in $^{238}\text{U}$ (Wa+71).

By averaging $\gamma$-ray intensities over many resonances, the Porter-Thomas fluctuations can be eliminated, and the relative standard deviation of the average value reduced to \( \sqrt{\frac{2}{N}} \) where $N$ is the number of resonances. In this way the energy dependence of $\gamma$-ray transitions can be investigated. Averaged $\gamma$-ray measurements in copper at keV neutron energies were found to show the predicted $E^3_\gamma$ energy dependence (Al68a), but averaged eV measurements in platinum (BT67) and other heavy isotopes (BT70) supported an $E^5_\gamma$ energy dependence. This latter result is consistent with the prediction by Axel (Ax62) of the influence of the tail of the GDR in the threshold region.

Bollinger (Bo73) has shown that the averaged $\gamma$-ray intensities are often independent of the structure of the low lying states, and that the partial radiative width can be written

\[
\Gamma_{\lambda \mu} = f(E_\gamma, \Delta \pi, \Delta J) R, \quad \ldots (1.4)
\]

where $R$ is a random variable which satisfies Porter-Thomas statistics, and $f$ is a function of the transition energy $E_\gamma$ and the change in parity ($\pi$) and spin ($J$). The averaging method was used to show that the ratios of average partial radiative widths for $E1$ and $M1$ transitions are independent of mass number and $\langle \Gamma_{\lambda \gamma} (E1) \rangle / \langle \Gamma_{\lambda \gamma} (M1) \rangle = 7 \pm 1$ for $A > 100$.

These results apply to the 'statistical' nuclides with level spacings <50 eV. However, for some of these nuclides, initial state correlations $\rho_I (\Gamma_{\lambda \pi}^0, \Gamma_{\lambda \gamma})$ have been observed, and for nuclides with $A < 100$ and those near closed shells, non-statistical spectra are the
norm. For these cases, the single particle structure of the low lying states results in excess strength for high energy γ-ray transitions. For $110 < A < 140$ and $181 < A < 208$, anomalous γ-ray spectra have also been observed which deviate from the γ-ray strength function $\langle \Gamma_{\lambda \mu}^\gamma \rangle \rho^\gamma_\lambda E^n_{\gamma}$ predicted by the statistical model ($n=3$) or extrapolated from the GDR ($n=5$) ($\rho_\lambda$ is the level density of states $\lambda$) {St64}. The presence of maxima in the s- and p-wave neutron strength functions, together with the above results, indicates that single particle motion persists within the long-lived resonances. These observations show that the statistical model does not provide a complete description of resonance neutron capture.

1.2 NON-RESONANT CAPTURE

As noted earlier, Lane and Lynn {LL60} separated collision matrix elements into internal and external parts (with respect to the nuclear radius), and the latter into resonant and non-resonant components, i.e.

$$U = U(\text{internal}) + U(\text{external})$$

$$U(\text{external}) = U(\text{resonant}) + U(\text{non-resonant})$$

These last terms were respectively called channel and hard sphere (pr potential) capture, and represent the non-statistical part of the capture process.

In the strong coupling model, the potential capture cross section $[\sigma_{\gamma\mu}(P) \text{ barn}]$ for an El transition to final state $\mu$ is given by:

$$\sigma_{\gamma\mu}(P) = \frac{0.062}{R\sqrt{E_n}} \cdot \left(\frac{Z}{A}\right)^2 \cdot \theta^2 \frac{y^2}{\nu^2} \left(\frac{y+3}{y+1}\right)^2$$

The partial radiative width $[\Gamma_{\lambda \mu}(\text{Ch}) \text{ eV}]$ resulting from channel capture is

$$\Gamma_{\lambda \mu}(\text{Ch}) = \frac{16\pi}{9} k^2 \gamma \left(\frac{2}{R}\right)^2 \theta^2 \frac{\theta_\lambda^2}{\nu} \left(\frac{eR}{k_\mu}\right)^2 <J_\lambda ||Y^{(1)}||J_\mu>^2 \frac{2J_\lambda + 1}{2J_\lambda + 1}$$

where $E_n$ is the neutron energy in eV, $R$ fm the nuclear radius,
13.

\[ y = k_\mu R, \]

\( k_\mu \) is the neutron wave number for the single particle state bound by energy \( E_Y \), and \( k_\mu^2 = \frac{2mE_Y}{\hbar^2} \).

\( k_\gamma \) is the photon wave number \( \left( \frac{E_\gamma}{\hbar c} \right) \),

\( \theta_\lambda^2 \theta_\mu^2 \) are the dimensionless reduced widths of the initial and final states,

\( e \) is the effective charge, \( \frac{Z e}{A} \), and

\[ \frac{\langle J_\lambda || Y^{(1)} || J_\mu \rangle^2}{2J_\lambda + 1} \]

is the angular part of the reduced matrix element.

The potential cross section is proportional to \( y^2 \) or \( E^1_y \). Since this cross section is proportional to \( \frac{\langle J_\lambda || Y || D || J_\mu \rangle^2}{k_\gamma^3} \), the dipole matrix element must be proportional to \( E^{-1}_y \).

Marisocotti et al. \( \{Ma+69\} \) identified direct capture for nuclides with \( N \sim 82 \), and Kopecky et al. \( \{KSL74,KS76\} \) have shown that direct capture occurs in the 3s region at thermal energies. The final state correlations between the \( (n,\gamma) \) and \( (d,p) \) cross sections are greatly improved by the use of an \( E^1_y \) energy dependence, rather than the classical \( E^3_\gamma \) factor, when the thermal capture cross section is comparable to that calculated from equation 1.5 (i.e. \( \lesssim 1 \) barn).

Large thermal capture cross sections (\( \gg 1 \) barn) are indicative of resonance capture, and in these cases the final state correlations tend to maximise for \( E^n_\gamma \) with \( n > 1 \). This result is expected since the energy dependence for resonant channel capture is \( E^2_\gamma \) (equation 1.6).

The mass dependence of the direct capture cross section in the 3s and 4s regions has been obtained from calculations using R-matrix theory with intermediate coupling \( \{LL60\} \) and, more recently, using optical and shell model formulations of the valence model \( \{CM75,Cu76\} \). In the 3s region at thermal energy, \( \sigma(P) \) peaks at 0.9 barn for \( A \sim 52 \), dropping
sharply to zero at $A = -58$ and recovering to a second maximum of 0.3 barn at $A = -67$. A similar pattern is found in the 4s region with peaks at $A = -150$ and 195, and a minimum at $A = -170$. These results were obtained with an imaginary potential $w_o = 3.30$ MeV and are sensitive to variations in this quantity {LL60}.

The potential capture cross section is a consequence of the scattering of the incident neutron by the nuclear potential into a bound, final, single particle state. However, if channel capture is significant and initial state correlations are observed, then a second contribution to the non-resonant collision matrix element arises from the tails of distant resonances, i.e.

$$U(NR) = U(P) + U(DR)$$

The distant resonance component $U(DR)$ is normally zero owing to the random nature of the sign and magnitude of the partial radiative width amplitudes, as predicted by the statistical model. But when channel capture is dominant, the amplitudes can add coherently. Lane {La71} derived an expression relating this background contribution to the initial state correlation coefficient,

$$o_{\gamma\mu}^{(DR)} = \rho I (\Gamma^2_{\lambda n}, \Gamma_{\lambda\mu}) \cdot \sigma_n \cdot \frac{<\Gamma_{\lambda\mu}>}{<D>} \cdot \frac{\pi}{2} \left( \frac{\text{Re} R_{\lambda\mu}}{\text{Im} R_{\lambda\mu}} \right)^2$$

where $\sigma_n = 2\pi^2k^2g<\Gamma_{\lambda n}>/<D>$ and $<D>$ is the average level spacing.

The real and imaginary parts of the R-matrix, $R_{\lambda\mu}$, are closely related and are comparable in the case of a common doorway (see section 1.4). The 'potential' and 'distant resonance' cross sections cannot be differentiated, and both add to give the total non-resonant capture cross section $\sigma(NR)$. This is not surprising since it is known that distant resonances change the hard sphere scattering phase shifts to optical model phase shifts.

The capture cross section can be written in terms of resonant and direct components, as well as resonance-direct and resonance-resonance...
interference terms \{LL60,Lo62,Au68,LFB74\}.

\[
\sigma_{\gamma \mu} = \pi \chi^2 q \sum_{\lambda} \frac{\Gamma_{\lambda n} \Gamma_{\lambda \mu}}{(E_\lambda - E)^2 + (\Gamma_{\lambda/2})^2} + \sum_{\lambda} \frac{\Gamma_{\lambda n} \Gamma_{\lambda \mu} U^{NR}_{\gamma \mu}(E_\lambda - E)}{(E_\lambda - E)^2 + (\Gamma_{\lambda/2})^2} + \left| U^{NR}_{\gamma \mu} \right|^2 + \sum_{\lambda \neq \lambda'} \frac{\Gamma_{\lambda n} \Gamma_{\lambda n'} \Gamma_{\lambda' n} \Gamma_{\lambda' \mu}}{(E_\lambda - E)^2 + (\Gamma_{\lambda/2})^2 (E_\lambda' - E)^2 + (\Gamma_{\lambda'/2})^2} \left\{ (E_\lambda - E)^2 + (\Gamma_{\lambda/2})^2 \right\} \left\{ (E_\lambda' - E)^2 + (\Gamma_{\lambda'/2})^2 \right\} \]

...(1.8)

where \(\Gamma_{\lambda n}\) is energy dependent and \(\Gamma_{\lambda n} = \Gamma_{\lambda n}(E_\lambda) \sqrt{\frac{E_\lambda}{E_\lambda}}\).

When the final state is a pure p-wave, single particle state, the \(\gamma\)-ray amplitude is correlated in sign with the neutron amplitude and the average cross section for s-wave resonance-resonance interference is

\[
\sigma_{\gamma \mu}(r-r) = \pi \chi^2 \frac{<\Gamma_{\lambda n}>}{<D>} \frac{<\Gamma_{\lambda \mu}>}{<D>} , \quad \text{...(1.9)}
\]

and is the same order as the potential capture cross section.

The sign of the interference between the resonant and direct components is of interest. The coefficient of the interference term, \(\frac{\Gamma_{\lambda n}}{\Gamma_{\lambda n} \Gamma_{\lambda n} U^{NR}_{\gamma \mu}}\), can be shown to have the same sign as \((E_\lambda - E)\) in the vicinity of the single particle resonance with energy \(E_{SP}\) \{Ly68\}. For nuclides below the s-wave size resonance, interference between resonant and non-resonant components is expected to be constructive below the resonance energy, i.e. opposite to that observed in neutron scattering. Nuclides above this size resonance would exhibit asymmetries similar to scattering.

Impurities in the final state \(\mu\) introduce a degree of randomness into the sign of the interference term, and eventually the sign of \(\frac{\Gamma_{\lambda n}}{\Gamma_{\lambda n} U^{NR}_{\gamma \mu}}\) becomes random with respect to \(\frac{\Gamma_{\lambda n} U^{NR}_{\gamma \mu}}{\Gamma_{\lambda n}}\).

Measurements of the non-resonant capture cross section can therefore be made by observing interference effects in partial capture \(\gamma\)-ray...
channels. Many measurements have been made by the Brookhaven group using high resolution Ge(Li) detectors to obtain γ-ray spectra at and between resonances in the eV energy range. Problems have arisen because of count rate effects, and the interpretation of the data must also take into account resonance-resonance interference \cite{Ch74}. A study of eV capture in the 4s region \cite{CC74} supports the presence of non-resonant capture cross sections at thermal energy which are comparable with predicted values. Note that only final states with large spectroscopic factors will have significant direct components, and when \( \rho \Gamma^{\lambda n,\lambda m}_{\lambda n,\lambda m} \) is small, an estimate of the potential cross section can be obtained.

An alternative method to the Ge(Li) measurements is provided by the threshold photonuclear reaction. Neutron energies are measured by the time of flight method and, in effect, the ground state partial capture cross section is obtained. The observations of asymmetries in s-wave resonances in \(^{28}\text{Si},^{52}\text{Cr},^{207}\text{Pb}\) have been interpreted in terms of interference effects \cite{Ja74b}. Again, experimental problems have arisen and results are complicated by the possible contribution of \( \ell > 0 \) resonances. Results for Cr and Ni, however, are found to be in agreement with theory since the 90 keV resonance in \(^{52}\text{Cr}\) is found to have a low energy tail \( \sigma^{\text{NR}} (90 \text{ keV}) = 200 \mu \text{b} \). On the other hand, little interference is observed in the 12 keV resonance in \(^{60}\text{Ni}\) \( \sigma^{\text{NR}} (12 \text{ keV}) = 12 \mu \text{b} \) which is above the 3s size resonance.

Measurements of interference in the 41 keV resonance in \(^{207}\text{Pb}\) have yielded conflicting results. Original \((\gamma,n)\) results \cite{BBB69} showed a strong asymmetry which was explained by postulating a large background cross section derived from the GDR. However, measurements of the \(^{207}\text{Pb}(n,\gamma)\) reaction \cite{AM70,71b} showed no indication of asymmetry. A later \((\gamma,n)\) measurement \cite{Ja73,Ja74a}, while not in agreement with the capture data, yielded a value of \( \sigma^{\text{NR}} (41 \text{ keV}) = 1.3 \text{ mb} \), which is at least
consistent with resonance-resonance interference and the expected potential cross section, without recourse to anomalous non-resonant processes. Direct measurements of the $^{207}\text{Pb}(n,\gamma)$ cross section with filtered neutron beams at 2 and 24 keV \cite{DR71,WC73} are also consistent with a small non-resonant cross section.

It would be desirable to obtain improved agreement between $(\gamma_o,n)$ and $(n,\gamma_o)$ methods, and so gain confidence in the validity of these experiments. Investigations into the scattered neutron sensitivity of the detectors used in the capture measurement (Chapter 3.3) show that a degree of asymmetry can result from this effect \cite{Al+77}. Since this asymmetry is opposite to that found for interference, the net result may be the appearance of a symmetric resonance. Detailed capture calculations are needed to resolve this problem.

1.3 **THE VALENCE MODEL**

First formulated by Lane and Lynn \cite{LL60}, and Lynn \cite{Ly68}, and applied by Mughabghab et al. \cite{Mu+71}, the valence model describes the change of state of the incident neutron in the entrance channel by the emission of dipole radiation in the field of a spectator target.

The partial radiative width for an El transition from resonance $\lambda$ to final state $\mu$ is given by \cite{Ly68}

$$\Gamma_{\lambda\mu} = \frac{16\pi k^3}{9} \frac{|\langle X_{\lambda}(J_\lambda) \rangle| |H^1_E| |X_{\mu}(J_\mu)||^2}{(2J_\lambda+1)} \quad \ldots(1.10)$$

where $k_\gamma$ is the photon wave number, and $H^1_E$ is the irreducible tensor operator specialised to electric dipole radiation, i.e. the dipole operator $D$.

The basis functions $\chi$ can be expanded in terms of radial wave functions for a set of single particle states defined in the spin-orbit coupling scheme. Neglecting core transitions \cite{La59}, the partial valence radiative width becomes
\[
\Gamma^V_{\lambda\mu} = \frac{16\pi k^3}{9} \theta^2_\lambda \theta^2_\mu |\bar{e}|^2 \int dr u^{*}_\lambda r u_\mu |^2 \cdot \frac{|<j' I \alpha J_1 \gamma | y^{(1)}_{\alpha} |j'' I I I >|^2}{2J_{\lambda} + 1} \]

...(1.11)

where \(\theta^2_\lambda, \theta^2_\mu\) are the dimensionless reduced widths of the resonance and final state, \(\theta^2_\lambda = \Gamma_{\lambda n} [2kRP_{\lambda} \Gamma_W]^{-1}\), where \(\Gamma_W\) is the Wigner single particle limit, \(R\) the nuclear radius and \(P_\lambda\) the penetrability; \(\theta^2_\mu\) is the \((d,p)\) spectroscopic factor; \(u_\lambda, u_\mu\) are the resonance and final state single particle radial wave functions, \(\bar{e}\) is an effective charge equal to \(Z/A\) times the electron charge.

For a zero spin target \((I_\alpha = 0)\), the angular part of the matrix element for a transition from the initial state \((J_\lambda, j' I'')\) to the final state \((J_\mu, j'', I'')\) becomes
\[
\frac{3}{4\pi} \frac{(2J_{\mu} + 1)}{2 \left[1 + (-1)^{j' + j'' + 1}\right]} . \quad \ldots (11)
\]

Tabulations of the overlap integral given by Lynn \{Ly68\} can be incorrect by a factor of three. This error arises because the wave functions were normalised over all space rather than the interior region only. To compensate, the initial state reduced width \(\theta^2_\lambda\) can be calculated for a square well leading to a larger value for \(\theta^2_\lambda\) than for a diffuse-edged Woods-Saxon potential. Although these two factors tend to compensate \{LM74\}, the method must be regarded as somewhat unsatisfactory since it cannot account for the variation of the radial overlap integral with neutron energy and mass number.

Lane and Mughabghab \{LM74\} have derived an optical model formulation of the valence process which has been further investigated by Barrett and Terasawa \{BT75\}. Since the valence width is given in terms of the ratio of readily calculated optical model quantities, the normalisation problem is avoided, i.e.
\[
\Gamma^V_{\lambda\mu} = \left[ \frac{\text{Im} u^{*}_{\mu} |D| U_{E}^{(\text{opt})}}{\text{Im} \tan \delta^{(\text{opt})}} \right]^2 \cdot \Gamma_\lambda \quad . \quad \ldots (1.13)
\]
where \( u_\mu \) is the final state wave function and \( U_\mu (\text{opt}) \) is the optical model initial state wave function at neutron energy \( E \). The dipole operator is denoted by \( D \) and \( \delta(\text{opt}) \) is the optical model phase shift. Removing the energy dependence from \( \Gamma_\lambda \) and explicitly showing the \( E^3 \) dependence, we obtain

\[
\Gamma^V_\lambda \mu = q_\lambda \mu (E_n) \cdot E^3_\lambda \mu \cdot \theta^2_\mu \cdot Z^2/\Lambda^2 \cdot \Gamma^l_\lambda , \quad \ldots (1.14)
\]

where \( q_\lambda \mu (E_n) = \Gamma_n (P_{\text{opt}})^{-1} \) eV and the reduced partial valence width \( q_\lambda \mu (E_n) \) (MeV)\(^{-3} \) is an energy dependent parameter calculated from the optical model and contains the radial integration and geometrical factors given in equation (1.11).

The total valence width is obtained by summing over all final states \( \mu \) containing the single particle state configuration:

\[
\Gamma^V_\lambda \gamma = \sum_\mu q_\lambda \mu \cdot E^3_\lambda \mu \cdot \theta^2_\mu \cdot Z^2/\Lambda^2 \cdot \Gamma^l_\lambda \to
\]

\[
= Q_\lambda \cdot \Gamma^l_\lambda , \quad \ldots (1.15)
\]

The average total valence width for resonances of spin \( J \) and angular momentum \( \ell \) is given by

\[
\langle \Gamma^V_\gamma \rangle_{J,\ell} = Q_{\ell J} \cdot S_\ell \cdot <D_{\ell J}> , \quad \ldots (1.16)
\]

where \( S_\ell \) is the \( l \)-wave neutron strength function and \( <D_{\ell J}> \) the average spacing of resonances with the same \( J^\pi \).

The valence process is therefore expected to be important when the resonances (or initial states) have large reduced widths, corresponding to even-even target nuclides with large level spacings in the regions of the neutron strength function maxima; and when \( El \) transitions can excite final states with large spectroscopic factors which occur near closed neutron shells.

Cugnon and Mahaux \{CM75\} have further investigated the valence model and the mutual relationships between the R-matrix approach of Lane and
Fig. 1.2
Single particle neutron binding energies

Fig. 1.3
Reduced valence widths ($10^b q$) for El transitions between s, p, d and f-wave single particle states
Lynn (LL60), the optical model formulation of Lane and Mughabghab (LM74) and the shell model approach (MW69, BM73). The role of external capture and of the dependence of the scattering length on mass number were shown to be of central importance in determining the validity of the valence model.

Using the preceding formalism, calculations have been made for allowed El transitions between s-, p-, d- and f-wave states across the periodic table, with emphasis on the closed neutron shells or subshells at N = 14, 20, 28, 50, 82 and 126.

The locations of single particle states have been calculated using an optical model code with the Moldauer parameters (Mo63) given in Table 1.1 and their mass and energy dependence are shown in Fig. 1.2. s-wave resonances are linked by El transitions to low lying p-wave final states, and d-wave resonances to p- and f-wave final states at A = 50, 140. At A = 30 and 90, p-wave resonances decay by El transitions to s- and d-wave final states (Bi+73, Bi+76). It is readily apparent that the shell model provides the underlying basis for the valence capture process.

The reduced partial valence widths (q) and El transitions from resonances with the appropriate (λ, J) are calculated at thermal energies for s-, p- and d-wave capture assuming zero spin targets. These data are given in graphical form in Fig. 1.3 as a function of mass number to expedite valence calculations. For non-zero spin targets, the reduced valence width is given by

\[
q(I_a) = (2J_a + 1)(2J_a + 1) \left\{ J_{\lambda}^\mu \frac{I_a}{J_{\lambda}^\mu} \right\}^2 \cdot q(I_a = 0) \ . \ 
\]  

Calculated values for this coefficient for some cases of interest are given in Table 1.2.

Since the valence process predominantly occurs outside the nuclear radius, the reduced valence widths are dependent on the variation with
TABLE 1.1  
NEUTRON OPTICAL MODEL PARAMETERS

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<tr>
<td>$r_2$</td>
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with \[ R = r_1 A^{\frac{2}{3}} + r_2 \]

\[ V(r) = (V_c + iW)f(r) + V_{so} \left( \frac{1}{r} \frac{df}{dr} \right) e^{-\frac{r}{a}} \]

with \[ f(r) = \left[ 1 + \exp \left( \frac{r-R}{a} \right) \right]^{-1} \]
TABLE 1.2

SPIN FACTORS FOR NON-ZERO SPIN TARGETS

\[ G = (2J_{\alpha} + 1)(2J_{\mu} + 1) \left\{ j', J_\lambda, I_\alpha \right\}^2 \left\{ J_\mu, j'' \right\} \]

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<th>( j'' )</th>
<th>( J_\lambda )</th>
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Fig. 1.4
Energy dependence of s-, p- and d-wave valence transitions relative to thermal values for the (a) 2p, 3s, 2d, (b) 3s, (c) 3p and (d) $^{208}\text{Pb}$ regions.
energy of the external part of the initial and final state wave functions. As the final state becomes more tightly bound, the external part decreases which results in a decrease in the dipole overlap integral. Also, as the energy of the initial state increases, a cancellation in the matrix element can occur.

The dependence on the final state single particle binding energy has been investigated by observing the variation of the reduced valence width with changes in the central potential. A dependence of $E^{-1}$ is observed, resulting in an overall $E^{2}$ dependence for valence transitions, in agreement with the prediction for channel capture (equation 1.6).

The experimental confirmation of this result is not easy, in spite of the large amount of spectral data available, because accurate spectra are required for resonances with dominant valence contributions, since the presence of other non-statistical components can result in interference between the radiative amplitudes. Further, the final state single particle strength should be fragmented over a large energy range if meaningful results are to be obtained. Since these requirements have not been found in practice, adequate confirmation of the $E^{-1}$ dependence of the reduced valence width has not yet been achieved.

The dependence of $q$ on the difference between the calculated binding energy $BE(OM)$ and the experimental centroid energy $BE(EXP)$ (determined from $(d,p)$ measurements), can be accounted for in the approximation

$$q = \left[ \frac{BE(OM)}{BE(EXP)} \right] q(OM), \quad \ldots(1.18)$$

to an accuracy of better than 20 per cent and typically better than 10 per cent.

To obtain the $q$ value at a higher neutron bombarding energy ($E_n$) in terms of the thermal value, the ratio $R$ can be used (Fig. 1.4) for the appropriate mass region, viz,

$$q_{\lambda\mu}(E_n) = R(E_n) q_{\lambda\mu}(th), \quad \ldots(1.19)$$
More accurate results can be obtained in the 3s region by inter­
polating $R$ between the curves for $^{36}$Ar and $^{54}$Fe. Otherwise the curves
accurately represent the energy dependence within the specified mass
regions to several hundred keV. The minima in the curves are sensitive
to the mass number and specific optical model calculations should be
made for energies close to those of the minima.

The experimental confirmation of the marked energy dependence of
$q$ predicted in the 3s and 4s regions is extremely difficult to achieve.
A large energy range is required for nuclides which also exhibit a
dominant valence effect. To date, an adequate check of the theory
has been possible only in $^{54}$Fe (Al+77a) where the capture data have
been analysed to 500 keV. These results are reported in section 6.5.

Calculations of reduced valence widths have also been made by
Cugnon (Cu76) using the optical model approach. This study was
restricted to $p$-$s$ and $s$-$p$ transitions in the 2p, 3s, 3p and 4s mass
regions, for neutron energies up to 500 keV. Our results are in good
agreement with these calculations at thermal energy. However, the
calculated energy dependence is much higher in the 3s and 4s regions,
and lower in the 3p region. The small energy dependence predicted by
Cugnon for $A = 50$-$60$ is not supported by the $^{54}$Fe results (Al+77a).

1.4 DOORWAY STATES AND CORRELATIONS

Several authors have studied the underlying basis of width
correlations in resonance neutron capture. Beer (Be69) used a two
group expansion of the resonance wave function to account for the
Porter-Thomas distribution of the reduced neutron widths, the narrow
distributions of partial radiative widths and the correlations between
reduced neutron and partial radiative widths, and between pairs of
partial radiative widths. In the case of $^{169}$Tm, one group of orthogonal
basis functions contained the single particle component of the neutron
resonances, while the other contained a small number of doorway or
collective states.

Beer went on to apply the projection operator formalism (Be7l), utilising the doorway state assumption. The transition matrix divided up into direct, semi-direct and compound nucleus terms, and the same doorways contributing to the semi-direct term were also involved in the resonance term. The partial radiative amplitude was expanded into single particle, doorway and statistical components and expressions were obtained for the correlation coefficients $\rho_I(\Gamma^0_\lambda \Gamma_\mu)$, $\rho(\Gamma_\lambda \Gamma_\mu')$, and $\rho_F(\theta^2_\mu' \Gamma_\lambda \Gamma_{\mu}/E^3)$; ($\rho_I$ and $\rho_F$ are the initial and final state width correlations.) The first two correlations are predicted to be unity if the resonance reaction proceeds through an isolated radiative doorway state.

Soloviev (So7l) expanded the wave function into few quasi-particle components and noted that if these were considered as doorways, then a correlation between channels occurs for common doorways. In strongly deformed nuclides, large correlations should be expected between K-allowed transitions to the ground and low lying quasi-particle rotational levels.

A detailed treatment of doorway states and correlations has been developed by Lane (La70,71), and it is this approach which is followed in this section.

The reduced width amplitudes $\gamma_{\lambda c}$ for resonance $\lambda$ in a channel $c$ can be expressed in terms of a set of channel states $\phi^*_c U_{n\lambda}$, where $\phi_c$ is the product of internal states in channel $c$ and $U_{n\lambda}$ is the single particle state of relative motion with principal quantum number $n$ and orbital angular momentum $\lambda$. Thus

$$\gamma_{\lambda c} = \sum_n <\lambda|\phi^*_c U_{n\lambda} \gamma_{n\lambda}> \ldots (1.20)$$

where $\gamma_{n\lambda}$ is the amplitude for the state $n$. 
For a given resonance, a hierarchy of states with increasing complexity may exist, the simplest being the channel states. These states form a complete set and diagonalising the nuclear Hamiltonian $H$ in this set yields the resonances $\lambda$ as eigenstates. In an analysis of doorway states and correlations, Lane [La71] divided this hierarchy of states into three levels of complexity: the simple channel states $\phi_{n}\lambda^{'}$, the states $|i>$, say, which we extend here to include two quasi-particle (particle-hole) and phonon states, and all the more complex states $|b>$. If the resonance contains only a small number of states $|i>$, then the role of doorway states [FKL67] may be of physical significance. The doorway states $|d>$ are obtained by diagonalising $H$ in the set $|\phi_{c}\lambda^{'}>$ and $|i>$, to obtain

$$|d> = a|\phi_{c}\lambda^{'}> + \sum_{i} \beta_{i}|i>.$$ ... (1.21)

Since $Y = 0$ and $|d> + |b>$ form a complete set, then

$$Y_{\lambda} = \sum_{d} \lambda_{d}Y_{dc},$$ ... (1.22)

where

$$Y_{dc} = \sum_{n} \delta_{d}\phi_{n}\lambda^{'}Y_{n}\lambda^{'}.$$ ... (1.23)

In the extreme situation of an isolated doorway, a modulation in the energy dependence of $Y_{\lambda}^{2}$ will occur. Lane has shown that if the doorway state is present in channels $c$ and $c'$, then a complete correlation is observed ($\rho = 1$) between the amplitudes $Y_{\lambda}$ and $Y_{\lambda'}$, i.e.

$$\frac{Y_{\lambda}}{Y_{dc}} = \lambda_{d} = \frac{Y_{\lambda}}{Y_{dc'}}.$$ ... (1.24)

More generally, several doorway states may contribute to $Y_{\lambda}$, $Y_{\lambda'}$ (equation 1.22). In this case, the amplitude correlations reduce by a factor $\frac{1}{\sqrt{n}}$ (the width correlations by $\frac{1}{n}$) where $n$ is the number of doorways with significant contributions to the sum [La70].
Note that for an isolated common doorway, $\gamma_\lambda \gamma_{\lambda c}' \gamma_{\lambda c}$, will have the same phase for all resonances $\lambda$ (i.e. that of $\gamma_{dc}', \gamma_{dc}$) instead of the random fluctuations as predicted by the statistical model. Consequently, a tendency for destructive interference between resonances should be observed.

We now examine the effects of doorway states in the partial radiative width amplitude for an El transition from resonance $\lambda$ to final state $\mu$.

For brevity we write

$$|\phi_o U_{n\lambda} \lambda\rangle = |0\rangle \quad \text{and} \quad |\phi_o U_{n\lambda'} \lambda'\rangle = |0'\rangle \quad , \quad ...(1.25)$$

where $\phi_o$ is the ground state for the target. Then the doorway can be written as

$$|d\rangle = \alpha |0\rangle + \sum_i \beta_i |i\rangle \quad , \quad ...(1.26)$$

and

$$\gamma_{\lambda n} = \sum_d \langle d|\langle d|0\rangle \gamma_o \quad ..(1.27)$$

For the El single particle operator $D$

$$\gamma_{\lambda \mu} = \langle \lambda | D | \mu \rangle \quad ..(1.28)$$

Expanding $|\lambda\rangle$ in terms of the doorway states $d$, we obtain the non-statistical amplitude

$$\gamma_{\lambda \mu}^{\text{NS}} = \frac{\gamma_{\lambda \mu}}{\gamma_o} \langle 0 | D | \mu \rangle + \sum_d \sum_i \langle d | D | i \rangle \langle d | i \rangle \langle i | D | \mu \rangle \quad ..(1.29)$$

We expand the final state $\mu$ into ground state $\phi_o$ and excited target states $\phi_j$ coupled to the valence neutron, i.e.

$$|\mu\rangle = \theta_\mu |0'\rangle + \sum_j a_j |j\rangle \quad , \quad ...(1.30)$$

where $\theta_\mu$ is the (d,p) spectroscopic amplitude.
Separating the dipole operator into particle and target components

\[ D = D_p + D_t \quad \ldots (1.31) \]

\[
\gamma^{NS}_{\lambda \mu} = \gamma_{\lambda \mu}^{\mu} <0|D_p|0'> + \\
+ \gamma_{\lambda \mu}^{\mu} \sum_j a_j <0|D_t|j'> + \\
+ \sum_i \sum_j \gamma_{\lambda \mu}^{\mu} <\lambda|d><d|i><i|D_p|0'> + \\
+ \sum_i \sum_j a_j <\lambda|d><d|i>[<i|D_p|j> + <i|D_t|j>] \quad \ldots (1.32)
\]

This equation is similar to that obtained by Beer {Be71}. The first term is just the valence term, e.g. \(<\phi_0|D_p|\phi_2p>\). The second corresponds to transitions from the entrance channel to \(\lambda^-\) states of the core. Since the unperturbed energy of \(\lambda^-\) two-quasiparticle states is comparable with the neutron binding energy, such configurations are not expected at low excitation energies. Collective \(\lambda^-\) excitations are possible, but calculations by Harvey and Khanna {HK74} indicate that collective fragments of the giant dipole resonance do not occur at energies much below the neutron threshold. Consequently, the second term is not expected to be significant in neutron capture in the threshold region. These results are supported by measurements of the \(\gamma\)-ray strength function \({Ba+74,Be76}\) which show that extrapolation of the GDR to low energies overestimates the observed \(\gamma\)-ray strength to a considerable degree.

The third term results from particle-hole excitation of the core with subsequent radiative decay by p-h annihilation to low lying, single particle states. The strength of these transitions therefore depends on the spectroscopic factors \(|\lambda_{\mu}^2|\) of the final states; for example \(<\phi_{i U'_n l'},D_p|\phi_{0 U n l'}>\) where \(\phi_i \equiv |\phi_0,2p,2s^{-1}>\).
The fourth term corresponds to all other excited target state
particle and core transitions which have non-zero matrix elements;
examples are $\langle \phi_{\mathbf{i} n \ell} | D_{\mathbf{p}} | \phi_{\mathbf{i} n' \ell'} \rangle$ and $\langle \phi_{\mathbf{i} n \ell} | D_{\mathbf{c}} | \phi_{\mathbf{i} n' \ell'} \rangle$.

In addition to these radiative amplitudes, there also exists a
statistical amplitude $\gamma_{\lambda \mu}^S$ which is independent of the doorway states.
This amplitude results from transitions between the complex states
$|b\rangle$ and $|b'\rangle$ in the resonance and final states, i.e.

$$
\gamma_{\lambda \mu}^S = \langle b | D | b' \rangle \quad \ldots (1.33)
$$

For the low lying states with dominant single particle or particle-
vibration configurations, $|b'\rangle$ is expected to be small and $\gamma_{\lambda \mu}^S$ insignificant. Only at higher excitation energies (>few MeV) are the complicated
states expected to be significant, but then the statistical width is
relatively small because of the low $\gamma$-ray energy of the transition
$(\Gamma_{\lambda \mu}^S \propto E_\gamma^3)$.

In the limiting case of a single two-quasiparticle doorway, if core
terms are ignored, equation 1.29 reduces to

$$
\gamma_{\lambda \mu}^{NS} = \theta_{\mu} \left[ \sum_i \langle \lambda | d | i \rangle \langle i | D_{\mathbf{p}} | 0' \rangle \right] 
= \theta_{\mu} \left[ \sum_i \langle \lambda | d | i \rangle \langle i | D_{\mathbf{p}} | 0' \rangle \right] \quad \ldots (1.34)
$$

The radiative amplitude is therefore completely correlated with the
single particle strengths of the resonance and final states. Consequently,
initial and final state correlation coefficients ($\rho_1$ and $\rho_\mathbf{p}$) should be
large and symmetric, even for a negligible valence contribution. For a
given doorway, the radiative amplitude may be systematically larger or
smaller than the valence amplitude. If there are many states $|i\rangle$ contrib-
uting to the doorway, then their contribution will tend to cancel out on
average, reducing the interfering effect of the doorway state.
Returning to the case of overlapping doorways and ignoring term 2 of equation 1.32, we can write

\[
\gamma_{\lambda \mu}^{NS} = \theta \left[ \frac{\gamma_{\lambda \mu}^V}{\gamma_0} \gamma_{\lambda \mu}^V + \sum_d \gamma_{\lambda \mu}^a \right] + \sum_d \gamma_{\lambda \mu}^r
\]...

(1.35)

where the superscripts a and r correspond to the annihilation and retention of excited target state configurations, respectively.

We see that \( \sum_d \gamma_{\lambda \mu}^a \) introduce an asymmetry between the initial and final state correlations because the doorway contributions reduce the initial state correlation. The \( \gamma_{\lambda \mu}^a \) terms interfere randomly with the valence component, but have constant phases over resonances \( \lambda \) within the doorways \( d \). Interference between the valence and doorway contributions occurs without affecting the final state correlations.

The sum of \( \gamma_{\lambda \mu}^r \) terms reduces both initial and final state correlations, since it is independent of the single particle strengths of the resonance and final states. However, for many such states, this contribution will average out to zero.

In summary, large and symmetric initial and final state correlations are expected in the case of a single doorway with two-quasiparticle character. If the doorway has a particle-vibrator character, the final state correlations will be reduced.

When several overlapping doorways are present, \( \rho_I \) reduces by \( 1/n \), where \( n \) is the number of doorways, but in the case of two-quasiparticle doorways, the final state correlations will remain large.
Nine decades of neutron capture

Fig. 2.1

DISTRIBUTION OF $\Gamma_{\gamma_1}$ $\ell$-WAVE STRENGTH FUNCTIONS

\[
\begin{array}{c}
\sigma_{\text{mb}} \\
10^{-1} \quad 10^0 \quad 10^1 \quad 10^2 \quad 10^3 \quad 10^4 \quad 10^5 \quad 10^6 \quad 10^7 \quad 10^8
\end{array}
\]

KEV REGION

0.025 eV --- BORON --- Sc --- 14 MeV

DIFFRACTION MONOCHROMATOR

CHOPPERS LINACS

VAN DE GRAAFF

THERMAL EPITHERMAL RESONANCE REGION (HEAVY A) RESONANCE REGION (LIGHT A, CLOSED SHELLS) GIANT RESONANCE REGION

DISTRIBUTION OF $\Gamma_{\gamma_1}$ $\ell$-WAVE STRENGTH FUNCTIONS

\[
\begin{array}{c}
\frac{1}{2} \text{DIRECT CAPTURE} \\
\text{STATISTICAL} \\
\text{SEMI-DIRECT}
\end{array}
\]

PROPERTIES OF RESONANCES --- AVERAGE RESONANCE BEHAVIOR

CHANNEL OR DIRECT (1p-0h) --- DOORWAY (2p-th) --- STATISTICAL (MANY p-h EXCITATIONS) (INTERMEDIATE STRUCTURE)
CHAPTER 2
MEASUREMENT OF γ-RAYS FROM keV NEUTRON CAPTURE

Neutron capture research extends over nine decades of energy, from cold and thermal meV fluxes to the d(t,n) reaction near 15 MeV. Most measurements have been at thermal energies where capture cross sections are large. High flux reactors allow gamma ray spectra to be obtained with minute quantities of separated isotopes. In these nine decades of energy (Figure 2.1) the capture cross section varies from well-separated s-wave resonances, through unresolved resonance regions where p,d and higher l-waves are important, to the giant resonance region and beyond, where direct and semi-direct capture mechanisms are required to explain the magnitude of the cross sections. But even in the eV and keV regions evidence is accumulating which requires the presence of direct effects, as well as resonant valence and doorway configurations which may dominate the complex nuclear states described by the statistical model.

The thermal capture studies are important because they provide a vast amount of nuclear data on energy levels and cascade schemes, spins and parities of low-lying states, as well as the gamma ray spectra which gave the first indication of a systematic variation in the capture mechanism across the periodic table [Gr+59]. However, only a limited investigation of the mechanism involved in thermal capture is possible, since:

(a) Thermal capture can be either a non-resonant reaction or one in which an indeterminate mixture of the tails of bound and unbound excited states contributes.

(b) The observation of transitions to low lying excited states in the compound nucleus is limited by multipole selection rules operating on the possible s-wave spin states, the wide variation of partial radiative widths and the possibility of interference terms between the various radiative amplitudes.
(c) The correlation of capture radiative widths and the reduced neutron widths of final states (from the (d,p) reaction) is therefore not rigorous and it is difficult to separate direct effects from channel or valence capture.

These limitations are largely eliminated in the study of resonance neutron capture. Additional bonuses are the isotopic identification of resonances by γ-ray spectra, the observation of a greater range of transitions through higher ℓ-wave capture, and the determination of spin assignments of capture resonances. Resonance spins can be obtained from:

(a) Angular distributions of primary γ-rays to final states with known spin and parity \( J^\pi \) \{CBW69\}.

(b) Population of low lying states by the γ-ray cascade \{Po66\}, and

(c) Resonance-resonance interference analysis.

It is also possible to investigate valence capture by the observation of correlations between the partial radiative reduced widths and the product of the reduced neutron widths of both resonance and final states. The observation of asymmetry in the resonance partial capture cross section can lead to estimates of the direct capture component, provided that corrections are made for resonance–resonance interference and resonance scattered neutrons.

Resonance γ-ray spectra are analysed with the aim of obtaining information on their departure from the expected statistical behaviour. The statistical properties of interest are that:

(a) The primary γ-ray decay spectrum should be proportional to \( E_\gamma^3 \rho(E_\lambda - E_\gamma) \), where \( \rho(E_\lambda - E_\gamma) \) is the density of final states. Gamma ray intensities to discrete final states, however, should be proportional to \( E_\gamma^3 \), assuming only dipole transitions contribute significantly.

(b) The partial radiative widths should vary widely according to a chi-squared distribution with one degree of freedom.
Partial radiative widths and the reduced neutron widths of resonance and final states should be uncorrelated.

Most resonance capture studies have been in the eV region on medium and heavy A nuclides, where the resonance spacing is in the range 10-100 eV. However, there are many nuclides with resonance spacings in the keV region (light A or near closed shells) which warrant study and which, by virtue of their simpler shell structure, deviate more seriously from statistical behaviour. These nuclides are the subject of investigation in this thesis.

In referring to the keV region, a neutron energy range of a few keV to several hundred keV is implied. The lower bound of the keV region is reached by beams from fast choppers and linear accelerators. Resonance spectra have been measured up to 23 keV at Brookhaven and γ-flash problems appear to set a similar limit for measurements at ORELA.

The Sc reactor filter, developed at Idaho Falls (Gr+69), provides a monoenergetic neutron beam at 1.95 keV with a FWHM of 700 keV. A $^{56}$Fe (25 keV) filter is in operation at Brookhaven (GC76) and Si shows promise of a 144 keV beam. The Van de Graaff accelerator is the major source of fast neutrons with energies from 5 keV to the MeV region. Intense kinematically collimated monoenergetic fluxes are produced at 30 and 65 keV when Li and tritium targets are bombarded by protons at the threshold energy. Normally, the accelerator is pulsed (2-10 ns at ~1 MHz) and the time of flight method is used to measure neutron energies.

The first measurements of keV capture spectra were made with a pulsed Van de Graaff accelerator by Bergqvist and Starfelt at Studsvik (BS61a,b) and Firk and Gibbons at Oak Ridge (FG61). Sodium iodide detectors were used to obtain the main features of the capture spectra. The Swedish effort (BS62a,b,c) was directed at investigating the spectrum shape for heavy nuclei, while γ-ray intensities were obtained at Oak Ridge for a number of nuclei with simple decay schemes (BGG62,a,b;Bi+65;Be+65).
The Lucas Heights study was unique in that high resolution Ge(Li) detectors were used for the first time to obtain γ-ray spectra averaged over ~60 keV neutron energy range {Al68a, AB68, AKS68}, or spectra from resolved resonances {BKA68, BAK69}. Emphasis was placed on studies in the 40 ≤ A ≤ 65 region {ABK69} where the 3s neutron strength function reaches a maximum value. Other studies extended into the 2s-1d shell {Ke+74} and the 3p region {Bi+73}. A systematic survey of NaI spectra for a wide range of nuclides was obtained at 430 keV {Ba+77} and Ge(Li) measurements on iron were made up to 1 MeV in energy {Al+74}. The γ-ray spectrometry system used in these studies is described in this chapter.

2.1 NanoSecond Source of keV Neutrons

The 3 MeV Van de Graaff accelerator at Lucas Heights is equipped with a terminal pulsing and bunching system. Pulsing is achieved by sweeping the ion source beam through an ellipse and over an aperture. The transit time across the aperture, and hence the pulse length, is 10 ns. A phased r.f. longitudinal accelerating voltage is applied to each 10 ns pulse to produce bunched pulses of ~2 ns duration {Mo+67}. However, an energy spread of 3 keV results from the retardation and acceleration of the beam. An output pulse width of 2 ns FWHM can then be achieved under ideal conditions at a 1 or 2 MC rate. Application of the countdown system and external triggering of the pulsing via light pipes can produce almost any lower frequency {RF73}.

Some 10 WA can be delivered onto the target in the experimental area. However, if optimum timing resolution is required 5 WA is more typical. The beam passes through an analysing magnet and capacitive and inductive nanosecond pick-off devices which can be used to deliver an accurately timed pulse.

A ⁷Li target is evaporated onto a Cu or Ta backing in a separate chamber and transferred under argon to the beam line. The target is
spray cooled to reduce heating effects and high vacuum techniques such as metal seals, pumping impedance, liquid nitrogen cold traps and turbo-vacuum pumps are used to maintain a standard vacuum of 0.1 mPa \( \text{(A168b)} \).

A recent development is an off-axis target can with 5 cm diameter. Lithium is evaporated onto the end face and every 24 hours the can is rotated slightly to expose a fresh Li target to the proton beam. In this way carbon buildup and contamination of the target is restricted and the neutron yield is maintained within 10% of the maximum value for the duration of the run.

A feature of the \(^7\text{Li}(p,n)\) reaction is the angular and energy distribution of emitted neutrons when bombarded just above the threshold at \( E_p = 1.881 \text{ MeV} \). At the threshold, the centre of mass motion of the compound system results in the emission of neutrons with 30 keV energy at 0° to the beam. As the proton energy increases, a cone of neutrons centred at 0° opens up and a range of neutron energies occurs with the most intense component at the periphery of the cone \( \text{(GN60)} \). For an incident proton energy 5 keV above threshold, the cone half angle is 30° and the neutron energy range varies from 10 to 60 keV. This is the typical operating condition since it is possible to site the \( \gamma \)-ray detector just outside of the beam with a minimum of shielding. The neutron yield also increases rapidly with energy, since the threshold occurs on the low side of an excited state in \(^8\text{Be} \) \( \text{(GN60)} \). The cone breaks at 120 keV above threshold and an isotropic angular distribution of neutrons is observed in the centre of mass system. A computer program based on the reaction kinematics \( \text{(St72)} \) is used to calculate the intensity, energy and angle of neutrons as a function of incident proton energy.

A burst of protons onto the Li target generates a burst of neutrons into the forward cone. If a target is placed within the cone at a distance from the source (20-50 cm), \( \gamma \)-rays are emitted when the neutrons are captured in the target. These \( \gamma \)-rays are detected by the
**Fig. 2.2**
NaI response functions for 9.94 MeV γ-ray

**Fig. 2.3**
Shield for NaI detector
\(\gamma\)-spectrometer and a timing pulse starts a 'clock' which is stopped by a delayed beam pick-off pulse. Consequently, the time of flight is a measure of the energy of the incident neutrons. The clock is a time to pulse height converter which operates by generating a voltage ramp with height proportional to the time difference between the start and stop pulses. The neutron flux is monitored by a long counter or a thin \(^6\text{Li}\) glass scintillator mounted on a 5 cm photomultiplier tube with ancillary timing electronics.

2.2 **SPECTRUM MEASUREMENTS WITH A NaI DETECTOR**

A heavily shielded 20 cm x 15 cm diameter NaI detector was used in a survey of capture \(\gamma\)-ray spectra across the periodic table \(\{\text{Bo+77}\}\). Most measurements were made at 430 keV neutron energy in order to obtain new data at an intermediate energy, but some extended to 1 MeV. A number of measurements had been made at MeV energies \(\{\text{Be+66}\}\) and more recently up to the giant dipole region \(\{\text{BDM71,Ri+71,Li+77}\}\). The poor resolution of the NaI detector (8% at 0.667 MeV) is offset by its high efficiency. When well collimated with a 10 cm aperture, the line shape is dominated by the full energy peak. A larger aperture of 20 cm increases the single escape peak relative to the full energy peak (Figure 2.2).

The detector is well shielded (Figure 2.3) by an inner liner of litharge (PbO), glycerine (C\(_3\)H\(_8\)O\(_3\)), boron carbide (B\(_4\)C) and leadshot in the ratio by weight of 5:1:4:10. This liner is surrounded by a 10 cm lead sleeve, a 5 cm sleeve of B\(_4\)C in paraffin (70% by weight) and a 20 cm blanket of boric acid in paraffin (50% by weight). The front face is shielded by 10 cm of B\(_4\)C in paraffin and a 15 cm Pb aperture. The multilayer detector shield was designed by Broomhall \(\{\text{Br71}\}\), using neutron transport calculations to minimise the \(\gamma\)-ray and neutron background.

Measurements were taken at a pulse rate of 1 MHz and a pulse width of typically 3 ns FWHM, and constant fraction timing was employed to achieve 5 ns overall timing resolution. The energy loss in the Li target at
Fig. 2.4
Background subtracted spectra

Fig. 2.5
Capture γ-ray intensity distribution
threshold was ~30 keV. Capture samples were located 15 cm from the neutron source and the detector was placed 1 m away at right angles to the beam.

The neutron energy distribution at different incident proton energies was determined in a separate experiment with a 5 cm diameter plastic scintillator placed 2 m away from the neutron source. The mean energy and FWHM of the neutron distribution were then determined to be 40±20 keV, 135±25 keV, 210±30 keV, 430±30 keV and 1000±50 keV.

Coincident energy and timing data were digitised and stored in an on-line computer (see next section). Spectra were recorded for both prompt capture events and the time-independent background resulting from scattered neutrons. This background is present under the \((n,\gamma)\) peak and is subtracted from the capture spectrum after noramlisation. The background spectra also served as an energy calibration since they contain the 2.226 MeV hydrogen capture peak and the 6.797 MeV peak for thermal capture in \(^{127}\text{I}\), a constituent of the detector. Additional background runs were made with 'no sample' and with a carbon scatterer.

The background subtracted spectra (Figure 2.4) were then reduced by line-shape analysis to intensity distributions. A set of line shapes were obtained from measurements of standard sources and some \(^{27}\text{Al}(p,\gamma)\) resonances and covered the energy range of interest. These line shapes were interpolated by a non-linear parameter optimisation program, and used to reduce the capture spectra into 50 energy bins of equal width. The program was constrained so that only positive intensity components were allowed. An initial set of \(\gamma\)-ray intensities was necessary, but the final result was independent of the initial values. The data are summed into groups of three channels to yield the histograms shown in Figure 2.5.

2.3 **HIGH RESOLUTION \(\gamma\)-RAY SPECTROSCOPY AT keV NEUTRON ENERGIES**

The use of high resolution Ge(Li) detectors in measurements of keV
neutron capture $\gamma$-ray spectra posed problems of efficiency, resolution and background. The early Ge(Li) detectors were planar with volumes of about 10 cm$^3$. A more recent detector (Quartz and Silice) is coaxial and has an effective volume of 90 cm$^3$ and 21% peak efficiency at 1.33 MeV relative to the 7.6 cm x 7.6 cm NaI detector. The lower efficiencies, coupled with the generally much lower capture cross sections at keV energies, required the use of massive targets and long runs (-5 days) in order to obtain adequate statistics. Even then it was difficult to observe weak transitions in the $\gamma$-ray spectra.

The resolution of the detector (~0.2% at 1.33 MeV) is reduced under high count rate conditions, long running times where stability becomes a problem, and broadening of the $\gamma$-ray lines by capture over the range of neutron energies (e.g. 10-60 keV) \cite{Al68b}.

The background comprises both time independent and time dependent components and both can be discriminated by the time of flight method. However, part of the random background still contributes to the capture $\gamma$-ray spectra and limits measurements on nuclei with small capture cross sections since the scattering to capture ratio becomes too large.

In general, it is possible to overcome these problems by adequate shielding and electronics and the use of an on-line computer for data taking and analysis.

(a) Gamma ray response of Ge(Li) detectors

Before proceeding with a discussion of the experimental arrangements, it is useful to consider some properties of Ge(Li) detectors. A detailed study \cite{A167,ABE67} was made of a 30 cm$^3$ open ended coaxial detector (trapezoidal cross section 7.5 cm, 4.6 cm length, capacity 4.2 p.f. and leakage current $10^{-10}$ amp at 900 V).

A number of ($p,\gamma$) and ($p,\alpha\gamma$) reactions were used to provide essentially monochromatic $\gamma$-rays and pulse height spectra were recorded with the axis of the detector parallel to the incident radiation.
Response functions are shown in Figure 2.6 for γ-ray energies from 0.66 to 11.13 MeV. The spectra have been corrected for background and cascade γ-rays and have been normalised to the same level for the flat part of the continuum response. The main processes responsible for the structure of these response functions are shown in Figures 2.7 and 2.8 for 0.48 and 11.13 MeV γ-rays.

Strong backscattering effects are observed in the 0.48 MeV spectrum as well as a sharp rise at low energies. The energy distribution of Compton scattered electrons as calculated from the Klein-Nishina formula, is shown with area equal to that of the estimated experimental distribution of Compton events. Photoelectric absorption of scattered photons results in 9% of these events falling within the full energy peak. The Compton distribution is further reduced by multiple scattering (curves c', c" and c''' for single, double and triple scattering of photons). In each case, some photoelectric absorption takes place and approximately 70% of the full energy peak is due to this process.

Wainio and Knoll (WK66) determined the fraction of total absorption events resulting from primary photoelectric interactions for small volume planar detectors by Monte Carlo calculations. The above estimate of 30% at 0.48 MeV compares with 47% and 28% for 0.9 and 3.0 cm³ diodes. The central inert volume of the 30 cm³ coaxial detector therefore appears to inhibit the contribution of multiple scattering events to the full energy peak.

At higher γ-ray energies the pair production process becomes important and at 11.1 MeV the ratio of Compton and pair cross sections is 0.7. A variety of multiple processes are possible and the observed shape is complex (Figure 2.8). A Compton edge is not observed at the expected 10.88 MeV, but electrons of this energy have a 1 cm range in germanium so that β escape events become important. (The average path length is $R(\text{mm}) = 0.97 E(\text{MeV}) - 0.2$.) Further, about 22% of the β energy
Fig. 2.10
Peak efficiencies for Ge(Li) detector.
will be emitted as bremsstrahlung radiation.

When the positron and electron pairs are stopped within the sensitive volume of the detector, the total energy deposited in the detector depends on the probability for interaction of the 0.51 MeV annihilation $\gamma$-rays. The probabilities for capture $P_c$ and Compton scattering $P_s$ determine the shape of the spectrum in Figure 2.7, and from this $P_c/P_s = 0.21$. If the three peaks are due entirely to the pair process, their areas will be given by $F = P_c^2$, $S = 2P_cP_s$, $D = P_s^2$, where $P_e = 1 - P_c - P_s$ is the probability of escape from the sensitive region of the detector. A comparison of relative peak areas as a function of $\gamma$-ray energy indicates the constant ratio $S/D = 0.26$ (Figure 2.9). About one third of the full energy peak ($F$) is estimated to result from capture of both annihilation $\gamma$-rays, with two thirds of the area resulting from the absorption of scattered photons.

The structure between the peaks can be fitted using the results for the 0.48 MeV $\gamma$-ray. The areas above a smooth background are given by $2P_cP_s$ and $2P_sP_c$.

Bremsstrahlung radiation by the pair betas is also expected and the dashed curve in this region was obtained by subtracting a calculated bremsstrahlung escape spectrum with area derived from the double escape peak ($D$). This result is consistent with the slowly varying background of Compton events observed at higher and lower energies. The narrow dip below the double escape peak is attributed to photoelectric absorption of this radiation. Similar dips should be observed below the single escape ($S$) and full energy peaks ($F$), but are too small to be observed.

Full energy and double escape peak efficiencies are given for detectors with a range of active volumes in Figure 2.10. A number of authors have attempted to relate the full energy peak efficiency to the active volume \{FJ66,PH69,Va+75\}. The most general result has been obtained by Hnatowicz in Monte Carlo calculations \{Hn77\}. 


Fig. 2.11
Dual parameter system for radiative capture at keV energies

Fig. 2.12
Time of flight spectrum with Ge(Li) detector
(b) Geometry and shielding

Two detector and shielding geometries were used initially \{Al68b\}. In the standard geometry, similar to that used for NaI measurements, the Ge(Li) detector is placed out of the beam (Figure 2.11). More massive shielding can then be used and standard crystal-dewar configurations are practicable. This method has the lower, solid angle efficiency, and is less suited for large targets since the detector faces the sample at an oblique angle and the solid angle efficiency for the far side of the target is reduced.

The cone geometry comprises a $B_4C$/paraffin cone in which the Ge(Li) detector, at the end of a long cold arm from a specially designed cryostat, is buried. An annular target is placed over the cone to achieve maximum solid angle efficiency for large targets. This geometry fell into disfavour as larger Ge(Li) detectors became commercially available in standard cryostat dewar configurations.

In both systems the Ge(Li) detector is enclosed by a lead pot which is further surrounded by borated paraffin (boron carbide mixed with paraffin). A specially positioned lead shield is used to reduce the intensity of $\gamma$-rays from the $^7\text{Li}(p,\gamma)$ and $^7\text{Li}(p,p'\gamma)^7\text{Li}$ reactions.

Thermalised neutrons scattered back from the floor are reduced by shielding the floor with ~8 cm thick slabs of borated paraffin.

(c) Electronics

The Ge(Li) detector is cooled to liquid nitrogen temperature to minimise trapping of charge carriers. Since the detector is effectively a solid state ionisation chamber, the charge is collected by a charge sensitive preamplifier. The signal at the output of the preamplifier is a tail pulse, and is split for timing and pulse height analysis. Originally, a double delay line amplifier and crossover pick-off and an RC amplifier with long time constant were used (Figure 2.11). The current approach is a fast filter amplifier (50 ns) and constant fraction
Fig. 2.13
Neutron energy shift of primary γ-rays in Si

Fig. 2.14
keV average γ-ray intensities in copper
discriminator to provide the best timing pulse by minimising walk through
the spread in rise times in the detector. Improved timing resolution for
large volume diodes has been achieved recently by Engel et al. using a pulse
shape selection method with amplitude rise time compensated discriminators
{ESS77}. An active filter amplifier (4 µs time constant) with base line
restoration is used to achieve the ultimate pulse height resolution. A
pulse height stabilisator was used to eliminate drift during long runs of
up to five days.

(d) On-line computer

In the keV capture γ-ray measurements, a neutron energy range of
60-80 keV and a γ-ray range of 3-5 MeV was typical. The latter was
limited originally by the capacity of the analogue to digital converters
and later by the storage capacity of the computer. It was therefore
necessary to compress both time and γ-ray data. This was achieved by
recording the time of flight spectrum and setting windows over the
appropriate regions of interest (Figure 2.12). A corresponding γ-ray
spectrum could then be recorded for each time of flight window {AUT68}.
Measurements on Fe, Ti {BKA68} showed for the first time the shift in
primary γ-ray energy which occurs relative to thermal energy in resonance
neutron capture. This effect is illustrated in $\text{Si}(n,\gamma)$ {Ke+77} in
Figure 2.13. The method therefore provided a means of obtaining partial
capture cross sections when absolute intensities are measured.

When the resonance spacing is too close to be resolved by our limited
time of flight resolution, broadening of capture γ-rays would occur equal
to the resonance energy range being excited. The digital window technique
was therefore used to obtain spectra for discrete regions of the $(n,\gamma)$
envelope (Figure 2.12). These regions were adjusted to correspond to a
spread in γ-ray energy of ~10 keV. The spectra were then shifted by the
γ-ray energy difference above thermal and added to obtain average
intensities over a 60 keV neutron energy range with typically 10 keV
(e) Normalisation and calibration

Absolute measurements of keV capture γ-rays have been accomplished using the ground state γ-ray of the 35 keV resonance in $^{27}$Al as a secondary standard. This resonance has a radiative width $\Gamma_\gamma = 1.9 \pm 0.1$ eV [Al+75] and since the branching ratio to the ground state is $0.66 \pm 0.07$ [Ke+74,Be+67], the ground state radiative width is $\Gamma_{\gamma_0} = 1.39 \pm 0.20$ eV. This resonance is well placed for measurements just above the neutron threshold, and the 7761 keV ground state is readily observed. This γ-ray also provides a useful energy calibration, in addition to the 8799 keV γ-ray from thermal capture in nickel, and the 7632 and 7646 keV doublet from thermal capture in Fe.

The normal mode of operation is to obtain thermal γ-ray calibrations at the beginning and end of an experiment. An Al sample, with identical area and shape to the sample under measurement, is cycled into the beam for one hour every twenty four hours to obtain the absolute normalisation. The intensity and energy distribution of the neutron flux is derived from the time spectrum of capture γ-rays observed in an elemental cadmium target and a knowledge of the cadmium cross section. This measurement is also made periodically through the course of the experiment.

The observed capture yield for a γ-ray of primary energy $E_\gamma$ to final state $\mu$ is:

$$Y(E_\gamma) = \eta(E_\gamma) \cdot \phi_n(E_\lambda) \cdot A_{\gamma \lambda} \cdot C_\gamma \cdot n,$$

where

$$A_{\gamma \lambda} = 2\pi^2 \lambda^2 g_{\lambda \mu} \Gamma_{\lambda \mu} / \Gamma_{\lambda} \text{ b}\text{ev},$$

$n$ is the number of atoms b$^{-1}$ of isotopic target,

$\phi_n(E_\lambda)$ is the incident neutron flux at the resonance energy $E_\lambda$,

$\eta(E_\gamma)$ is the absolute Ge(Li) efficiency for $E_\gamma$,

$C_\gamma$ include corrections for resonance self shielding and multiple scattering, γ-ray attenuation in the sample and the angular distribution correction factor.
Self shielding and multiple scattering corrections were calculated using a modified version of the ORNL/RPI Monte Carlo code [Su+69]. This is done for both Al (0.137 atom b\(^{-1}\)) and sample targets.

The energy dependence of the Ge(Li) detector efficiency was determined earlier by a comparison of thermal capture γ-rays from nitrogen with values quoted by Motz [Mo62]. For a 14% Ge(Li) detector, the efficiency variation in the sum of all three peaks (i.e. full, single and double escape peaks) is less than 10% between 5 and 9 MeV.

(f) Analysis

Peak areas are currently assessed by either of two techniques. The summed intensity of all three peaks associated with a particular γ-ray can be obtained after a linear background is subtracted from each. Alternatively, the computer can be used to shift the peaks by 511.1 keV for the single escape peak and 1022.2 keV for the double escape peak, thus combining three peaks into one from which a linear background is then subtracted. Normally, the two methods agree within statistical uncertainties. The data are analysed by the PKANAL program on the IBM 365.

Spectra in separate windows may be analysed separately if overlapping of resonance shifted γ-rays is likely. The difference between the measured (\(E_γ\)) and thermal (\(E_γ^{th}\)) γ-ray energies is then the resonance energy in the centre of mass frame. The observed γ-ray energy is \(E_γ^{th} + (A/A+1)E_λ\).

Transitions with similar resonance energy shifts therefore are primary transitions of the same resonance (if the resonance spacing is >10 keV). Further, isotopic assignments of resonances in elemental targets can be made from the nature of the γ-ray spectra and resonance spins and parities can be deduced when final state spins and parities are known (but only by the presence of a transition, not by its absence). It is generally assumed that El transitions are much
stronger than M1 transitions to the same final state. Evidence for E2 transitions in primary capture γ-rays is extremely limited, and are assumed to be much weaker again.
CHAPTER 3
HIGH RESOLUTION MEASUREMENTS
OF RESONANCE CAPTURE CROSS SECTIONS

The measurement of resonance capture cross sections can contribute to an improved understanding of the mechanisms of neutron capture, as well as providing data needed for reactor shielding and astrophysical applications. The variation of resonance total radiative widths with spin and parity provides a signature of the neutron capture mechanism, and its dependence on the quantum numbers of the capture state. Similarly, the dependence of average radiative widths on mass numbers can reveal the influence of shell structure and pairing effects on the capture reaction.

A knowledge of such properties can then be used to estimate unknown cross sections of fission products which are of importance in the calculation of neutron economy in reactors \{In+77\}. The emphasis of this thesis is in the keV energy region which is most applicable to fast breeder reactors and to the nucleosynthesis of isotopes in stellar interiors \{AGM71\}.

Radiative capture cross sections can be measured by the detection of prompt $\gamma$-rays emitted within $10^{-12}$ s after neutron capture. If the compound nucleus so formed is unstable, then the cross section can also be obtained by measuring the radioactive decay (betas and/or gammas) with a characteristic half-life. Other methods are the absorption (shell-transmission) method, modification of pile reactivity and mass spectrometry of the product nuclei. The prompt method, used in conjunction with time of flight analysis of neutron energy, has seen the widest application in the past two decades, and a variety of methods and $\gamma$-ray detectors have been developed for this purpose.

Large liquid scintillator tanks surrounding the capture sample were first applied \{DTH60,Gi+61\} and have been further developed.
The Moxon-Rae detector (MR63) was devised to yield an efficiency proportional to the energy of the incident γ-ray. For a capture γ-ray cascade, the efficiency per capture is proportional to the total energy, i.e. neutron binding energy plus centre of mass neutron energy, and is independent of the capture spectrum. The inherent low efficiency (~3% including solid angle) and approximate linearity have led to further developments (MGI63) including multiple layer detectors (WCB67). The properties of these detectors have been studied in detail (IAJ74, MM75).

The Moxon-Rae detector approach to spectral independence has been generalised in the total energy weighting technique (MG67) whereby most γ-ray detector systems can be adapted to give exactly the required response if the pulse height spectrum is recorded. By assigning an importance or weight to detector events, which depends on pulse height alone, the average response can be made proportional to the total γ-ray energy emitted in the capture measurement. Since this quantity is also proportional to the product of the number of neutron captures and the total energy of the capture reaction, the neutron capture cross section of isotopic samples of known binding energy can be measured. The technique is therefore restricted to samples or isolated resonances when capture is predominantly due to a single isotope. For time of
Fig. 3.1
Properties of total energy detectors

Fig. 3.2
Observed and weighted pulse height spectrum in $^{103}$Rh($n,\gamma$)
flight measurements, a two parameter experiment is required with analogue or (usually) on-line digital computer data processing [CL70]. The properties of total energy γ-ray detectors are summarised in Figure 3.1.

The total energy detector system in operation at the Oak Ridge Electron Linear Accelerator (ORELA) is described in section 3.1. A $^6$Li glass neutron detector is used to monitor the neutron flux, and this is discussed in section 3.2, together with the saturated resonance method for determining the absolute efficiency of the capture detectors. In section 3.3, the neutron sensitivity of the total energy detector system at the 40 m station of ORELA is investigated, and a Monte Carlo approach is presented which includes the estimation of the resonance scattered neutron background, together with multiple scattering and self shielding effects in the capture sample.

3.1 TOTAL ENERGY DETECTORS

Maier-Leibnitz has shown [ML63,Ra63] that for a wide class of radiation detectors an average response can be generated which is proportional to the energy of the radiation. Macklin and Gibbons [MG67] adopted this philosophy in the design of a fast, capture γ-ray detector for high resolution capture cross section measurements. For optimum timing resolution a liquid scintillator was adopted. A code was written to calculate the pulse height distribution resulting from Compton interactions for incident γ-rays up to 10 MeV. Multiple Compton interactions were followed by taking the average Compton scattered γ-ray energy (and angle) at each step with a typical path length, until the residual γ-ray energy reduced to 25 keV. As the carbon photoelectric cross section is comparable to the Compton cross section at this energy, the calculation was terminated with a photoelectric interaction. The small pair cross section was also included.

The computed pulse height spectra $P(I,E_γ)$ were used to obtain the weighting function $G(I)$ from the chain of equations
\[ \sum_{i} G(i)P(i,E_{\gamma}) = E_{\gamma} \]  

where \( I \) is the pulse height channel number. The chain was terminated at the Compton edge \( E_{c} \).

The weighting function can be represented by

\[ G(X) = aX + bX^2 + cX^3 \]

where \( X \) is the pulse height in units (or channel widths) of 0.15 m \( e^2 \).

For a non-hydrogenous liquid scintillator (\( C_6F_6 \)) in a 10 cm diameter x 4 cm thick detector, the coefficients are

\[ a = 1.50, b = 0.086, c = -0.0034 \]

An isotropic point source, 1.8 cm from the axially symmetric pair of liquid scintillator cells is assumed.

The capture \( \gamma \)-ray spectrum observed for capture of 85 keV neutrons in \(^{103}\)Rh is shown in Figure 3.2. The weighting function is also shown (note right hand ordinate) together with the weighted pulse height spectrum.

The use of a \( \gamma \)-ray detector with essentially no full energy peak results in the capture yields being rather insensitive to the weighting function. This was demonstrated by comparing weighted counts with the unweighted counts. In measurements of the average weighting factors for various isotopes with binding energies in the range 6-8 MeV, the average weighting factor was found to be relatively constant, i.e.

\[ \frac{\Sigma G N}{\Sigma N} \sim 16 \pm 2 \]

The reliability of the weighting function for the pulse height weighting technique has been investigated in detail by Yamamuro et al. [Ya+76], by means of the black resonance method. The results of measurements and Monte Carlo analyses on black resonances in \(^{197}\)Au (4.9 eV), \(^{109}\)Ag (5.2 eV) and \(^{181}\)Ta (4.3 eV) showed that the weighting function used was reliable to 2%.
Fig. 3.3

Flight path at ORELA

- Shadow Bar Centered in First Collimator
- Water Moderator
- Tantalum Target Plates
- Pulsed Electron Beam (140 MeV)
- Evacuated Target Cell
- Borated Concrete Casing
- Typical Intermediate Collimator
- Resonance Blackout Filters (Al, CF₂, S)
- Final Collimator (Double Tapered)
- Flux Monitors
- Sample Changers (4)
- Gamma Ray Scintillation Detectors (2)
- Indicators and Controls
- To Sample Changers
- To Beam Stop
- Hydraulic Pump
- Future Experiment

Profile of ORELA Flight Path 7
The neutron capture cross section facility at the 40 m time of flight station at ORELA (Figure 3.3) uses two fluorocarbon liquid scintillators (NE-226, supplied by Nuclear Enterprises) for the detection of capture γ-ray events. The elimination of hydrocarbons from the detector system reduces the thermalisation of the scattered neutrons through collisions with hydrogen atoms, and eliminates the subsequent production of hydrogen capture γ-rays. The liquid scintillator is encased in a 10 cm diameter x 4 cm thick Al cell with faces of borosilicate glass. The cells are mounted on 12 cm RCA4522 fast photomultiplier tubes with fast discriminators to achieve a detection timing resolution of 1.86 ns FWHM for a 64:1 dynamic pulse height range. For the two detectors, the overall resolution is 2.8 ns.

The pulse pile up coincidence compensation method of Macklin and Gibbons (MG67) is not used since the smaller solid angle from sample to detector reduces the coincidence rate to between 1 and 2%. The sum of two pulses (either between or within detectors) is assigned more weight than the pulses taken separately, as a result of the non-linearity of the weighting function. This excess weight shows a broad peak of 25% overweight and, when combined with coincidence rates of 1-2%, results in a mean overweight of only 1/3%. The effect is therefore ignored.

While the efficiency, including solid angle, for a typical neutron capture cascade is 15%, the weighting process reduces the statistical accuracy. Since most pulses are small and receive less than average weight, the net efficiency is comparable to an ideal detector with Poisson statistics of 12%. For cascades of low multiplicity or single γ-rays, this may drop to 5%. Contrary to the case of the large liquid scintillator, the cross section results are not biased by low multiplicity, and the poorer statistics are revealed by the propagation of the variance. With total typical efficiencies of ~60% for scintillator tanks, the ORELA TED system is five times less efficient. However, the
Fig. 3.4
Capture electronics

Fig. 3.5
On-line data acquisition system
much lower background and superior time resolution reduce this dis-
advantage considerably. A block diagram of the electronics network
is shown in Figure 3.4. The time cycle is initiated by a flash
detector at a flight path of 47 m and capture events in the neutron
monitor or either γ-ray detector can stop the timing cycle within a
preset range. In subsequent developments the digital clock is now
triggered by a pulse from the linear accelerator.

The time dependent background in the neutron beam has been measured
with a $^{10}$BF$_3$ chamber and black resonance filters [Mu70]. At energies up
to 35 keV the overall beam backgrounds are less than 0.5%. At 105 keV
using a 5 cm sulphur filter and a $^{10}$B(n,αγ) detector, the upper limit is
1%. Similar results were achieved using a $^6$Li monitor, described in the
next section. This excellent performance results largely from the series
of copper collimators along the beam line in a scatter and trap configura-
tion, and the fact that the beam pipe is buried for a large fraction of its
length.

The neutron energy resolution has several components. The FWHM varies
from $\frac{E}{600}$ at 3 keV to $\frac{E}{400}$ near 2000 keV. The best electron beam resolution
is 2.1 ns (0.05 ns m$^{-1}$), but typically 5 ns bursts are employed. The
spread in time introduced by the neutron moderator surrounding the linac
neutron source can be expressed as a spread of 3.2 cm FWHM in the flight
path. This term dominates the low energy range (~3 keV) and is still an
appreciable component at 2 MeV.

Only neutrons which have scattered in the cooling jacket of the
neutron source can travel down the beam tube. A copper-lead shadow bar
shields out neutrons produced in the photoneutron reaction in the centre
of the linac target.

The main features of the on-line data taking system are shown in
Figure 3.5. A table of 128 channels of weights (G) and variances (V)
is included in the fast computer memory. When an event occurs at time
channel t and pulse height (PH), +1 is added to the time spectrum, G(PH) is added to the weighted spectrum and V(PH) to the variance spectrum. These spectra are stored on fast access disk and, to minimise time delays, time spectra are retained for each of 4 groups of pulse height, such that the dynamic range of any group is 64. The weights (G) are therefore reduced by 1, 1, 2.5 and 8 and the variances (V) by 1, 17, 330 and 4000.

A sharp 150 keV (electron energy loss) is imposed by suppressing the first two pulse height channels. This bias therefore eliminates carbon and fluorine recoils for neutrons below 2.5 MeV. The bias may be raised to 630 keV γ-ray energy by suppressing the first pulse height group and to 1.65 MeV by suppressing the second. Use has been made of the fourth pulse height group for discriminating against γ-rays below 4.08 MeV, and a fifth pulse height group can be set up prior to an experiment if desired. The ratio of high bias to total yield obtained in this way has been used successfully to determine λ-values of resonances.

The pulse height channel width of 76.7 keV enables a γ-ray energy range of 10 MeV to be covered. A Pu-Be neutron source produces γ-rays of 4.43 MeV and is used to calibrate the linear pulse height analysis system.

3.2 $^{6}$Li GLASS MONITOR AND ABSOLUTE NORMALISATION

The $^{6}$Li(n,α)T reaction offers a number of advantages as a neutron flux standard below 200 keV. At neutron energies below a few keV its cross section is smooth and nearly inversely proportional to the neutron velocity. The $^{6}$Li(n,α)T cross section at one keV is known to 1%, both experimentally and theoretically, and to 2.3% at 10 keV {Ca75}. At higher energies the uncertainty increases to 15-20% between 150 and 300 keV. Recent measurements {FM72,CHU72,Po74,Fr+74,SK75,Ba75} have resulted in substantial changes to the cross section at the 250 keV resonance and at higher energies. These results have been parameterised
Fig. 3.6

$^6$Li glass monitor

Fig. 3.7

Efficiency correction for 0.05 cm $^6$Li glass scintillator
and an accuracy of 10-20% is expected for energies up to 1 MeV \{MHW75\}. The latest results of Gayther \{Ga77\} are in excellent agreement with an R-matrix calculation with best fit parameters for total, scattering and (n,a) reactions on $^6$Li and $^4$He(t,t)$^4$He \{Ha77b\}. The overall accuracy of the (n,a) cross section now appears to be better than 10% up to 1 MeV.

Fast timing is possible when a $^6$Li glass scintillator (NE912) is used in conjunction with a fast photomultiplier tube. A timing resolution of less than 2 ns has been achieved with constant fraction discriminators.

In order to minimise multiple scattering effects a thin (0.5 mm) glass is used. This scintillator is thin enough to suffer only small efficiency perturbations from multiple scattering effects, and the transmitted neutron flux is only slightly perturbed by resonances in the constituents of the glass. The charged particle ranges of the alpha and triton particles are short enough to allow about 90% of the captures to deposit full energy in the glass. Further, the high energy release (4.79 MeV) permits excellent separation of capture events from background events. A peak to valley ratio of 40:1 is achieved in the ORELA configuration.

A major requirement is to minimise saturation effects from the $\gamma$-flash which occurs at time $t = 0$. It was therefore necessary to ensure that the photomultipliers are outside the neutron beam. This is achieved by mounting a 5 cm diameter RCA8575 photomultiplier on 3 mm thick quartz light pipes at each end of the rectangular glass scintillator (Figure 3.6). The $^6$Li glass scintillator is mounted in a light Al frame perpendicular to the slightly smaller, collimated neutron beam. The scintillator surfaces are polished to achieve total internal reflection. The light collected at each photomultiplier varies slowly with the distance to its point of origin in the scintillator. Thus the summed light from both ends is independent of position and a symmetric pulse height peak is observed.
The combined dynode signal is amplified and a window set to include the neutron induced pulse height peak as shown in Figure 3.6.

Using $({\text{CF}_2})_n$, S, Al filters, most of the neutrons at energies of 5.9, 27, 35, 49, 89 and 100 keV can be removed from the beam. Under these conditions the background in the detector is less than 1% of the normal yield. Some of this undoubtedly results from neutrons escaping the collimation system. The 'machine off' background is twenty times lower.

As the glass is upstream from the capture sample (flight path 39.5 m) corrections for flux perturbation as well as multiple scattering are required. The energy dependent correction factor is shown in Figure 3.7, where the ratio of the efficiency per transmitted neutron ($\eta$) to the macroscopic cross section is plotted. Without corrections, this ratio would be unity. Ninety degree scattering by the glass constituents results in increased ($n,a$) yields as shown at the silicon, oxygen and $^6\text{Li}$ resonances.

The transmitted neutron flux $\phi_n(E)$ is determined from the observed $^6\text{Li}$ yield $Y_\mu(E)$ by

$$\phi_n(E) = Y_\mu(E)/\eta(E) \quad , \quad ...(3.3)$$

where $E$ is the kinetic energy of the neutron in the laboratory frame of reference. The capture cross section is obtained in the following way.

In the measurement of neutron capture cross sections, the capture yield is defined as the number of captures per incident neutron. In the thin sample case, the yield is

$$Y(E) = \frac{\sigma_c(E)}{\sigma_T(E)} \left[1 - T(E)\right] \quad , \quad ...(3.4)$$

where $\sigma_c(E)$, $\sigma_T(E)$ are the capture and total cross sections.

$$T(E) = \exp(-n\sigma_T(E)) \quad , \quad ...(3.4)$$

is the transmission, where $n$ is the atom b$^{-1}$ of sample.
If \( T(E) \approx \left( n_\sigma T(E) \right) \),

then \( Y(E) = n_\sigma C(E) \). \( \ldots(3.5) \)

However, measurements are made with samples of finite thickness and the above approximations are rarely valid on resonance. The capture yield therefore comprises a primary component due to neutrons which are captured on their first interaction and a multiple scattering component which results from capture events following one or more scatters. The estimation of these events is considered in detail in the next section.

Assuming the thin sample approximation (i.e. \( n_\sigma T \leq 0.2 \)), and neglecting self shielding, multiple scattering and background corrections, the capture cross section is given by

\[
\sigma_C(E) = N_c(E) \left[ \eta_\gamma \cdot \phi_n(E) \cdot n \cdot c_\gamma \right]^{-1}, \quad \ldots(3.6)
\]

where \( N_c(E) \) is the number of capture events at energy \( E \), \( c_\gamma \) is a correction for \( \gamma \)-ray attenuation within the sample, and \( \eta_\gamma \) is the efficiency of the total energy detectors. \( N_c \) is obtained from

\[
N_c(E) = G(E) \left[ E_B + E \cdot A / (A+1) \right]^{-1}. \quad \ldots(3.7)
\]

The sum of all pulse height weights \( G(E) \) observed in the energy channel \( E \) to \( E + \Delta E \), is divided by the total energy of the capture reaction (i.e. binding energy plus centre of mass neutron energy).

The capture efficiency of the total energy detector is determined by the saturated resonance method for the 4.9 eV resonance in gold. Since \( \Gamma_\gamma \gg \Gamma_n \) for this resonance, the observed saturated \( \gamma \)-ray yield for a 0.0030 atom b\(^{-1}\) sample is proportional to the incident neutron flux because virtually all neutrons are captured.

The weighted detector efficiency is thus

\[
\eta_\gamma = N_c(E) / \phi_n(E). \quad \ldots(3.8)
\]

...
A correction of 2.4% is calculated to account for the transmission on resonance and the backscattering from the gold foil.

Originally, a value of $\eta_\gamma = 0.98\pm0.05$ was determined. This result is dependent on the discriminator settings of the associated electronics and needs periodic evaluation. Problems arising with the optical light seal on the $^6$Li glass resulted in substantial variations in $\eta_\gamma$ since the observed neutron flux was in error. However, the current value of $\eta_\gamma = 1.14\pm0.02$ has remained constant over a period of 10 months.

Each capture run is monitored by a time gated fission detector which is used for normalisation prior to the installation of the $^6$Li glass or during periods of inoperation of this monitor.

The on-line computer system records the following information during the course of a measurement:

- 128 channels (double precision) of pulse height data integrated over all time channels, 76.7 keV per channel.
- 1 x 18432 channels of time of flight data
- 4 x 18432 channels of G-weight data
- 4 x 18432 channels of variances.
- 1 x 18432 channels of monitor data

The pulse height groups, Compton energies and equivalent $\gamma$-ray energies are given below. Since the line shape of the total energy detector is a broadened Compton distribution, the actual limits set by the pulse height groups are larger than the appropriate channel energies.

<table>
<thead>
<tr>
<th>Pulse Height Group</th>
<th>Channels</th>
<th>$E_C$ (keV)</th>
<th>$E_\gamma$ (keV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0</td>
<td>1- 2</td>
<td>150</td>
<td>285</td>
</tr>
<tr>
<td>1</td>
<td>3- 6</td>
<td>450</td>
<td>630</td>
</tr>
<tr>
<td>2</td>
<td>7- 19</td>
<td>1425</td>
<td>1650</td>
</tr>
<tr>
<td>3</td>
<td>20- 51</td>
<td>3825</td>
<td>4080</td>
</tr>
<tr>
<td>4</td>
<td>52-128</td>
<td>9820</td>
<td>10080</td>
</tr>
</tbody>
</table>
The time range covered by the time of flight spectra, using the EGG TDC100 digital clock, is 57 µs. At the completion of a run, the data are dumped onto magnetic tape. A data reduction program, operating on the IBM 365, reads the tape, applies the scale factor to the G-weight and variance data, and calculates dead time corrections to all time of flight data, based on the combined monitor and total energy detector count rates. This correction utilises the total number of bursts in a run. Time independent background corrections are made for both the total energy detector and monitor.

The time of flight spectra are converted into 12 linear energy groups from 2.5 keV to 2.5 MeV, using an interpolative procedure. The channel widths are adjusted so that the average resolution width corresponds to about six energy channels. The γ-ray yield is then normalised relative to the $^6\text{Li}(n,\alpha)$ cross section to determine an effective capture cross section (in millibarns) as a function of neutron energy. This cross section does not represent the thin sample cross section assumed in the analysis, since resonance self shielding, multiple scattering and background corrections have not been made. The contributions of resonance scattered neutrons and the delayed, scattered (time dependent) background are also present. The estimation of these components is considered in the following section.

3.3 NEUTRON SENSITIVITY OF CAPTURE Γ-RAY DETECTORS

Measurements of the radiative widths of resonances with large neutron widths can be critically dependent on the sensitivity of the γ-ray detector to resonance scattered neutrons. This neutron sensitivity is dependent on the variation with energy of the capture cross sections of the detector and environs. The problem has been found to be of particular importance to capture measurements with the total energy detectors at the 40 m station of ORELA, but it is applicable to all low γ-ray resolution detectors in close proximity to the capture target.
Under these conditions neutron shielding may be minimal or even absent, and the time of flight method cannot be used to resolve capture \( \gamma \)-ray and 'prompt' scattered neutron events. The use of pulse shape discrimination for suitable liquid scintillators will not resolve the problem, since the capture of prompt scattered neutrons causes the emission of \( \gamma \)-rays which will be unresolved in time (and \( \gamma \)-ray energy) from the sample capture events. This problem is resolved in section (a) by the experimental determination of the neutron response function of the detector and environs.

The analysis of resonance capture yields is greatly facilitated by the use of Monte Carlo methods. The observed capture yield is divided into primary and multiply scattered components of both the resonance capture \( \gamma \)-ray yield and the \( \gamma \)-ray yield from the interaction of scattered neutrons. This method is described in section (b).

Scattering resonances in the keV region are analysed in section (c) for a range of nuclides. These resonances provide good tests for the neutron sensitivity of any capture detector, and are used to obtain the neutron sensitivity of the total energy detectors used at the 40 m station of ORELA.

(a) Neutron sensitivity

Neutrons scattered by the capture target interact with the detector and environs, and may be captured promptly within the timing resolution of the scattering resonance. If the average efficiency for resonance capture \( \gamma \)-ray detection is \( \eta_\gamma \), and that for neutron detection by prompt capture \( \eta_n \), then we define the neutron sensitivity

\[
\kappa = \eta_n / \eta_\gamma
\]

...(3.9)

In the thin sample approximation, the measured radiative width \( (\Gamma')^\gamma \) is the sum of both capture \( \gamma \)-ray and prompt scattered components, i.e.

\[
(\Gamma')^\gamma = (\Gamma)^\gamma + \langle k \rangle \cdot \Gamma_n
\]

...(3.10)
where $\Gamma_\gamma$ and $\Gamma_n$ are the radiative and neutron widths of a resonance at energy $E_R$, and $<k>$ is the average sensitivity of the capture detector to resonance scattered neutrons.

Direct measurements of the detector neutron sensitivity were made by Hockenbury et al. (Ho+69) for a liquid scintillator tank. A carbon scatterer was used with a $^{24}$Mg target to determine the neutron sensitivity at the 83.5 keV resonance in $^{24}$Mg. Since the ratio of capture to scattering is small (WMA76) ($\Gamma_\gamma / \Gamma_n = 6.6 \times 10^{-4}$), the $\gamma$-ray yield can be a sensitive indicator of the prompt background effect. Measurements were made with and without the carbon scatterer, and the difference in resonance areas gave the neutron sensitivity directly. Since the macroscopic carbon cross section ($n\sigma_p$) was comparable to the macroscopic scattering cross section ($n\sigma_o$) on resonance, the prompt background component was large at all energies (not just on resonance) and provided a new background upon which the resonance capture $\gamma$-ray yield was measured.

Other estimates of the neutron sensitivity of total energy detectors have been reported (MA71, AM71a, AL+75, Bo+75b, AL+77d) for the 40 m station of ORELA and (RAT75) for the Cadarache facility.

The energy dependence of the neutron sensitivity can be obtained from measurements on targets with few capture resonances, e.g. $^{12}$C, $^{208}$Pb (AL+73). Neutrons scattered by the potential cross section of $^{208}$Pb provide the energy dependent background, shown in Figure 3.8, from which Pb resonances can be removed by interpolation of a background fitting curve. The structure observed here illustrates the role of the constituents of the detector and environs. Scattered neutrons are seen to capture in the 27 keV resonance of F, a constituent of the liquid scintillator (CgFg) and in resonances in the Al beam pipe and structural assembly, at 5.9, 35 and 88 keV. Capture at higher energies occurs in many other resonances in F and Al, and inelastic scattering occurs in F.
Fig. 3.8
Observed γ-ray yield ($\gamma/E$) per eV for $^{208}$Pb. Unfiltered and filtered beam yields are shown, together with the time dependent background.

Fig. 3.9
Energy dependence of TED neutron sensitivity function for 8 MeV binding energy
above 110 keV.

While it is clear that the structure seen in Figure 3.8 results from the prompt detection of scattered neutrons, calculations show that a considerable degree of multiple scattering also occurs in the detector and assembly. The effect of this is to provide a step in the yield above the 6 keV peak, and to smear out the F and Al resonance responses above 30 keV. While this response could be obtained from Monte Carlo calculations on an idealised detector assembly, it is sufficient to use the empirical results obtained for $^{208}$Pb. The only target mass dependence observed for a range of nuclides is the structure at 6 and 27 keV, and this results from $-90^0$ scattering from targets with mass number $A$, given by $\frac{2}{(A+1)}E_n^2$. However, this mass dependence can be obtained directly from the Monte Carlo analysis described in section (b).

It remains to determine the time dependent component in the observed background, i.e. those events which occur out of time with the accelerator pulsing, and which add randomly to the background. Measurements with black Al and S filters on $^{208}$Pb scattering samples show dips in the detector yield which permit estimates to be made of the time dependent background component. By definition, the background level at an energy corresponding to zero incident neutron flux is the time dependent background which remains after subtraction of the normalised 'sample out' background. It is obtained in the following way. The γ-ray yield from the $^{208}$Pb sample (curve A in Figure 3.8) is observed for an unfiltered neutron beam. An assumed time dependent background ($\alpha E_n^{-x}$) is subtracted and the residue multiplied by the calculated transmissions through the S or Al filters. This result is compared with the observed γ-ray yield using the S or Al filters (curve B is with the Al filter). The background is then varied in magnitude and energy dependence until good fits to the filtered beam results are obtained. The background is then subtracted from the unfiltered beam γ-ray yield to obtain the prompt back-
ground yield.

The effective prompt background cross section for the total energy detectors is derived in the following way. The observed $\gamma$-ray yield is the sum of sample capture ($G_\gamma$), prompt ($G_{PBG}$) and time dependent ($G_{TDBG}$) backgrounds. For a given time channel,

$$G = G_\gamma + G_{PBG} + G_{TDBG}, \quad \ldots (3.11)$$

where

$$G_\gamma = \frac{TE(A) \cdot \sigma_\gamma \cdot \eta_\gamma \cdot n \cdot \phi}{\sigma_s \cdot \eta_n \cdot n \cdot \phi}, \quad \ldots (3.12)$$

where $\sigma_\gamma$ and $\sigma_s$ are the sample (mass number $A$, $n$ atom $b^{-1}$) capture and scattering cross sections, $\phi$ is the incident neutron flux and $TE(A)$ and $TE(PBG)$ the total energies for capture in the sample and in the detector and environs (e.g. Al and F).

For $^{208}$Pb, $G_\gamma$ is subtracted and $G_{TDBG}$ is treated in the code as sample capture.

Defining the effective prompt background capture cross section as

$$\sigma_{PBG} = \frac{TE(PBG) \cdot \sigma_\gamma \cdot \eta_\gamma \cdot n \cdot \phi}{TE(208) \cdot \eta_\gamma \cdot n \cdot \phi \cdot \sigma_s}, \quad \ldots (3.14)$$

and substituting for $G_{PBG}$

$$\sigma_{PBG} = \frac{TE(PBG) \cdot k \cdot \sigma_p}{TE(208)} \cdot \sigma_s \quad \ldots (3.15)$$

Therefore the modified neutron sensitivity function for a $^{208}$Pb sample is

$$k(208) = \frac{\sigma_{PBG}}{\sigma_p} = k \cdot \frac{TE(PBG)}{TE(208)}, \quad \ldots (3.16)$$

where $\sigma_p(A) = 4\pi\lambda^2\sin^2(R/\lambda)$ for nuclear radius $R$ and neutron wave length $\lambda$. For target $A$ with total energy $TE(A)$

$$k(A) = k(208) \cdot \frac{TE(208)}{TE(A)}, \quad \ldots (3.17)$$
The variation and magnitude of the neutron sensitivity of the ORELA γ-ray detector as a function of energy is shown in Figure 3.9 for a typical sample with 8 MeV binding energy.

(b) Monte Carlo analysis

The analysis of resonance capture cross sections normally includes Monte Carlo techniques to determine the multiple scattering contribution to the capture yield, using a modified version of the ORNL/RP1 Monte Carlo code {Su+69}. The purpose of the code is to calculate the capture yield as a function of neutron energy for specified target geometry parameters, temperature and resonance parameters. Single level Breit-Wigner formulae for the capture and total cross sections are used with Doppler broadening to include the effects of thermal motion of the target nuclei.

For the case of an isolated resonance $\lambda$ with angular momentum $l$, spin $J$ and energy $E_\lambda$; total width $\Gamma_\lambda = \Gamma_\lambda n(E) + \Gamma_\lambda Y$; where $\Gamma_\lambda n(E)$ is the energy dependent neutron width and $\Gamma_\lambda Y$ the total radiative width; the capture and scattering cross sections are:

$$\sigma_c = \frac{\sigma_0 \Gamma_{\lambda Y}}{\Gamma_\lambda} \frac{1}{1+x^2}$$  \hspace{1cm} \ldots (3.18)

$$\sigma_s = \sigma_p + \frac{\sigma_0}{1+x^2} \left[1+x\tanh 2\delta_1\right]$$  \hspace{1cm} \ldots (3.19)

where

$$\sigma_p = \pi/k^2 \delta_l (2l+1) \sin^2 \delta_l = 4\pi/k^2 \sin^2 kR$$  \hspace{1cm} \ldots (3.20)

$$\delta_0 = -kR, \quad \delta_1 = -(kR)^3/3$$  \hspace{1cm} when $kR < 1$, and $k$ is the neutron wave number.

$$\sigma_0 = 4\pi/k^2 g \frac{\Gamma_n}{\Gamma_\lambda} \cos 2\delta_l = \frac{2608}{E(\text{kev})} \cdot \left(\frac{A+1}{A}\right)^2 \frac{g\Gamma_{\lambda n}}{\Gamma_\lambda}$$  \hspace{1cm} \ldots (3.21)

$$x = \frac{E-E_\lambda}{\Gamma_\lambda/2}$$

The capture area of the resonance is obtained by integrating over the resonance, and is

$$A_{\lambda Y} = 2\pi^2/k^2 g \frac{\Gamma_n}{\Gamma_{\lambda Y}} \frac{E_\lambda}{\Gamma_\lambda} \frac{\Gamma_{\lambda n}}{\Gamma_\lambda}$$  \hspace{1cm} \ldots (3.22)
Random number techniques are used to simulate a neutron history in the sample. Neutrons are assumed to undergo isotropic scattering in the centre of mass system, and a new direction is chosen randomly, so determining the new energy of the scattered neutron. The neutron history is followed, summing all weighted capture contributions until its weight drops to a preset minimum or until its energy falls below a given threshold.

In order to treat the prompt neutron scattering problem explicitly, the Monte Carlo code has been modified to include these events.

For each neutron considered, the capture yield at the first collision is given by:

\[
Y_1(E_1) = \frac{\sigma_c(E_1)}{\sigma_T(E_1)} \left[ 1-T(E_1) \right], \quad \ldots (3.23)
\]

where \( \sigma_T = \sigma_c + \sigma_s \) and is the total cross section. The first collision probability is \( 1-T(E_1) \), where \( T(E_1) \) is the transmission calculated for neutrons of energy \( E_1 \) normally incident on the target.

At each subsequent collision in the target, a contribution to the multiple scattering yield is accumulated. For example, after one scattering, the capture yield is:

\[
Y_2(E_1) = \frac{\sigma_s(E_1)}{\sigma_T(E_1)} \left[ 1-T(E_1) \right] \left[ 1-T(E_2) \right] \frac{\sigma_c(E_2)}{\sigma_T(E_2)} \quad \ldots (3.24)
\]

where \( E_2 \) is the neutron recoil energy after scattering and \( (1-T(E_2)) \) is the probability that the second collision occurs within the target. Also, from each scattering event in the target, a weighted contribution to the prompt neutron correction is estimated for those neutrons which escape the target in the direction of the capture detector. After one scattering event, we obtain

\[
Y_p(E) = \frac{\sigma_s(E_1)}{\sigma_T(E_1)} \left[ 1-T(E_1) \right] T(E_2) \cdot k(E_2) \quad \ldots (3.25)
\]
Fig. 3.10
Mass dependent structure in the prompt background

Fig. 3.11
55 keV s-wave resonance in $^{28}$Si

Fig. 3.12
53 keV p-wave resonance in $^{23}$Na
The factor $k(E_2)$ is the probability that the escaping neutron with energy $E_2$ is promptly captured in the detector. Additional terms also arise from multiple scattering events in the sample, leading to the escape of neutrons with reduced energy. Since the neutron sensitivity function shown in Figure 3.9 is effectively that for zero energy loss (actually -1% energy loss for $^{208}$Pb), $k(E_1)$ is obtained directly from this function at each stage of the calculation for scattered energy $E_1$.

The scattered neutrons take a finite time to reach the $\gamma$-ray detector and/or the beam tube walls. Consequently, prompt background events occur at slightly later times (i.e. lower energies) than the sample capture events. This effect has been simulated in an approximate manner by defining a minimum interaction distance and allowing an empirically determined average distance to have a Gaussian distribution about the mean. These constants are then used in a random sampling method to determine the modified line shape of the prompt background components. The average interaction distance and variance are dependent on the incident neutron energy since they are affected by the presence of resonance structure in the capture cross sections of the detector and environs.

(c) Analysis of capture yields

Once a neutron sensitivity function has been established, it can be used in the Monte Carlo analyses to reproduce the mass dependent structure observed in the energy ranges 6-10 and 27-65 keV. Results for $^{23}$Na, $^{28}$Si, $^{56}$Fe, $^{88}$SrCO$_3$ and $^{206}$Pb are shown in Figure 3.10 for the range 2.5 to 60 keV. The broken line corresponds to the calculated prompt background which is superimposed on a slowly varying ($E^{-h}$) background. The lower histogram corresponds to multiple scattering events resulting from scattering by nearby resonances.

The mass dependence of the peaks at 6 and 27 keV is largely reproduced by the calculations. These peaks are broadened out in oxide
Fig. 3.13
Time delayed background in $^{140}\text{Ce}$

Fig. 3.14
27 keV resonance in $^{140}\text{Ce}$
and carbonate targets, and the Monte Carlo method satisfactorily accounts for this effect. It is therefore possible to confirm the mass and energy dependence of the neutron sensitivity function.

In the analysis of scattering resonances, the multiple scattering capture yield is subtracted from the observed capture yield. The calculated primary yield, together with the prompt background component, is then fitted to the residual capture yield.

When the prompt background component is sufficiently large with respect to the capture yield, a strong asymmetry is introduced for s-wave resonances and the calculated yield becomes proportional to the product of the scattering cross section and the neutron sensitivity when the ratio $\Gamma_n / \Gamma_\gamma = 0$. Asymmetry resulting from s-wave scattering in $^{28}$Si at 55 keV is shown in Figure 3.11. This effect is not observed for the 53 keV p-wave resonance in $^{23}$Na which has a symmetric scattering cross section (Figure 3.12). The scattering shape is modified, however, by the time-straggling effect discussed earlier. This is most apparent in Figure 3.13 for the 12.47 keV resonance in $^{140}$Ce. Because the neutron width of this resonance is small, the time-straggling effect is most pronounced and is responsible for the low energy asymmetry.

The effect of neutron sensitivity depends not on the absolute magnitude of the capture and scattering widths, but on their ratio. This is evident in Figures 3.14 and 3.15, where capture resonances in $^{140}$Ce and $^{56}$Fe at similar energies are shown. The neutron widths of the 27.7 keV $^{56}$Fe resonance is a factor of twenty larger than that of the 28.1 keV $^{140}$Ce resonance. However, the ratio $\Gamma_\gamma / \Gamma_n$ is smaller for $^{140}$Ce ($<6 \times 10^{-4}$) than for $^{56}$Fe ($9 \times 10^{-4}$), and the scattered neutron yield dominates the Ce resonance. Both resonances are seen to be strongly asymmetric.

The predicted resonance asymmetry resulting from the prompt background can be used to determine the neutron sensitivity ($k$) of the
Fig. 3.15
27.7 keV resonance in $^{56}\text{Fe}$

Fig. 3.16
Dip in γ-ray yield at 500 keV in $^{208}\text{Pb}$
detector system. Table 3.1 lists a number of resonances with large neutron widths and their measured radiative widths. Ideally, a fit to both area and shape should yield the recommended radiative and neutron widths as well as a value of k which is consistent with the expected energy dependence.

The magnitude of the correction for resonances in Table 3.1 can be quite large. In the case of the 27 keV resonance in Al, the radiative width reduces from the observed value of -7 eV to ~3.7 eV. However, for the p-wave resonance at 5.9 keV, the correction is negligible. This is because neutrons scattered by this resonance fall below the resonance energy and are captured with much lower probability in the detector environs.

Our radiative width for the 35 keV resonance in $^{27}$Al is still much higher than other measurements. The problem, however, is unlikely to lie in the magnitude of the neutron sensitivity function, since a larger value would exceed upper limits set by other resonances in $^{27}$Al, as well as in $^{28}$Si, $^{140}$Ce and $^{208}$Pb. It is also unlikely that the Monte Carlo analysis is in serious error since the measured neutron width is within 10% of the reported total cross section value. It is possible though that assumptions in the Monte Carlo calculation are inadequate when sample scattering and background capture resonances are the same. Note that our radiative width is consistent with the thermal capture cross section.

By fitting the observed γ-ray yields and/or published radiative widths for selected resonances, the accuracy of the prompt background is confirmed on average to about 20%. (The 35 keV Al resonance is the only exception.) However, if the prompt background constitutes the larger proportion of the capture yields, unacceptably large uncertainties may result in the radiative widths.

Note that some instances occur where the reported neutron widths are entirely inconsistent with our shape analyses. The 55 and 180 keV
### TABLE 3.1

PARAMETERS OF SCATTERING RESONANCES

<table>
<thead>
<tr>
<th>A</th>
<th>El</th>
<th>$E_n$ (keV)</th>
<th>$\Gamma_n$ (eV)</th>
<th>$\Gamma_\gamma$ (eV)</th>
<th>Ref.</th>
</tr>
</thead>
<tbody>
<tr>
<td>23</td>
<td>Na</td>
<td>53.39</td>
<td>0.72±0.05</td>
<td>RAT76</td>
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<tr>
<td></td>
<td></td>
<td>52.2 (0.70)</td>
<td>1.58±0.24</td>
<td>Ho+69</td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>53.15</td>
<td>1.08</td>
<td>Ra+73</td>
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<td></td>
<td></td>
<td>53.1±0.1</td>
<td>1.15±0.30</td>
<td>this work, MAM77</td>
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<tr>
<td>27</td>
<td>Al</td>
<td>34.7</td>
<td>2.0 ±0.3</td>
<td>MG73</td>
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<tr>
<td></td>
<td></td>
<td>34.8</td>
<td>2.0 ±0.5 b)</td>
<td>!+75</td>
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<td></td>
<td>34.8</td>
<td>5.0 ±1.0</td>
<td>this work</td>
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<td>28</td>
<td>Si</td>
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<td>0.45</td>
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<td></td>
<td></td>
<td>1.5 ±0.3</td>
<td>0.11±0.03</td>
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<td></td>
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<td>(1.5)</td>
<td>0.06±0.06</td>
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<td>(56)</td>
<td>1.0 ±2.0</td>
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<td>33.2</td>
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a) Supersedes results given in refs. Ke+76, Ba+75

b) Corrected value
resonances in $^{28}_{\text{Si}}$ require neutron widths about half the values given in the literature {MG73}. Since in these cases the prompt background constitutes most of the $\gamma$-ray yield (Figure 3.11), the capture measurements are, in effect, neutron scattering measurements. Our values for these neutron widths for $^{28}_{\text{Si}}$ have been confirmed by recent ORELA measurements of the total cross section {Ha77a}.

In $^{20}_{\text{Pb}}$ the 515 keV resonance is observed as a minimum in the total cross section. Since the scattered neutron flux is also a minimum at resonance, a dip in the prompt background is observed. A fit to this dip (Figure 3.16), and to adjacent resonances, confirms the magnitude of the neutron sensitivity at high energies, but requires $E_R = 500$ keV and $\Gamma_n \sim 60$ keV. This result is in disagreement with that given in {MG73}, but confirms the Duke result {Fa+65}.
EXEMPTION OF VALENCE TRANSITIONS FROM
THE GIANT DIPOLE RESONANCE

4.1 S- AND p-WAVE VALENCE CAPTURE

Many measurements of s- and p-wave radiative widths have been made
in the AAEC-ORNL collaborative program, mostly in the region of the 3s,
3p and 4s strength function maxima. These data, together with results
from other laboratories, are listed in Table 4.1. An insight into the
importance of the valence model can be gained through a comparison of
these average radiative widths with those predicted by the valence model.
Following the procedure of section 1.3, average total valence radiative
widths can be calculated from equation 1.16 using the Q values, average
strength functions and l-wave spacings listed in Table 4.1.

Many of these Q values are obtained on the assumptions that the
calculated binding energies are in good agreement with the experimental
centroid energies, that the complete spectroscopic strengths are present
in the low lying states, and that for the non-zero spin targets,
\( \langle Q(I_\alpha) \rangle = Q(I_\alpha = 0) \). When detailed valence calculations have been
reported in the literature, these widths have been used in preference
to the average values obtained above.

The average, total radiative width can be expressed as the sum of
statistical \( \langle \Gamma^S \rangle \) and residual components \( \langle \Gamma^R \rangle = \langle \Gamma \rangle - \langle \Gamma^S \rangle \), the latter
including single particle and collective effects which are non-statistical.
For most nuclides in the 3p region, s-wave resonances do not exhibit any
significant initial state correlations \( \{Bo+76a\} \), and can only decay via
M1 transitions to the low lying states with strong single particle
class. s-wave capture resonances will therefore be largely statistical
in nature. However, Johnson \( \{Jo77\} \) has noted that the statistical component
of the p-wave radiative widths will be larger than that for the s-waves
since E1 transitions to the low lying states are allowed. When such
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<td>3.1</td>
<td>0.07</td>
<td>0.7</td>
<td>86</td>
<td>(40)</td>
<td>46</td>
</tr>
<tr>
<td>144</td>
<td>Nd</td>
<td>( s_{1/2} )</td>
<td>3.3</td>
<td>3.9</td>
<td>0.45</td>
<td>5.8</td>
<td>47</td>
<td>(40)</td>
<td>7</td>
</tr>
<tr>
<td>145</td>
<td>Nd</td>
<td>( s_{1/2} )</td>
<td>3.2</td>
<td>5.2</td>
<td>0.04</td>
<td>0.7</td>
<td>87</td>
<td>(40)</td>
<td>47</td>
</tr>
<tr>
<td>146</td>
<td>Nd</td>
<td>( s_{1/2} )</td>
<td>3.1</td>
<td>3.7</td>
<td>0.27</td>
<td>3.0</td>
<td>51</td>
<td>(40)</td>
<td>11</td>
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<tr>
<td>148</td>
<td>Nd</td>
<td>( s_{1/2} )</td>
<td>3.1</td>
<td>2.7</td>
<td>0.17</td>
<td>1.4</td>
<td>46</td>
<td>(40)</td>
<td>6</td>
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<tr>
<td>154</td>
<td>Sm</td>
<td>( s_{1/2} )</td>
<td>4.0</td>
<td>1.8</td>
<td>0.12</td>
<td>0.8</td>
<td>65</td>
<td>(40)</td>
<td>15</td>
</tr>
<tr>
<td>169</td>
<td>Tm</td>
<td>( s_{1/2} )</td>
<td>8.5</td>
<td>1.4</td>
<td>0.009</td>
<td>0.11</td>
<td>87</td>
<td>(40)</td>
<td>47</td>
</tr>
<tr>
<td>181</td>
<td>Ta</td>
<td>( s_{1/2} )</td>
<td>12.8</td>
<td>1.8</td>
<td>0.009</td>
<td>0.2</td>
<td>53</td>
<td>(40)</td>
<td>13</td>
</tr>
<tr>
<td>184</td>
<td>W</td>
<td>( s_{1/2} )</td>
<td>5.4</td>
<td>2.6</td>
<td>0.10</td>
<td>1.3</td>
<td>72</td>
<td>(40)</td>
<td>32</td>
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<tr>
<td>197</td>
<td>Au</td>
<td>( s_{1/2}, 2 )</td>
<td>6.4</td>
<td>2.1</td>
<td>0.03</td>
<td>0.4</td>
<td>132</td>
<td>(40)</td>
<td>92</td>
</tr>
<tr>
<td>198</td>
<td>Hg</td>
<td>( s_{1/2} )</td>
<td>6.2</td>
<td>1.2</td>
<td>0.10</td>
<td>0.8</td>
<td>147</td>
<td>(40)</td>
<td>107</td>
</tr>
<tr>
<td>203</td>
<td>Tl</td>
<td>( s_{1/2} )</td>
<td>710</td>
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<td></td>
<td></td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>A</td>
<td>El.</td>
<td>$(l,J)_\lambda$</td>
<td>$10^5Q$</td>
<td>$10^4\delta$</td>
<td>$D_{l,J}$ (keV)</td>
<td>$\Gamma_Y^V$ (meV)</td>
<td>$\Gamma_Y$ (meV)</td>
<td>$\Gamma_Y^V$ (meV)</td>
<td>$\Gamma_Y^P$ (meV)</td>
</tr>
<tr>
<td>-----</td>
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<td>----------------</td>
<td>---------</td>
<td>-------------</td>
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<td>----------------</td>
<td>------------------</td>
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</tr>
<tr>
<td>204</td>
<td>Pb</td>
<td>$s_{1/2}$</td>
<td>5.8</td>
<td>0.65</td>
<td>4.2</td>
<td>16</td>
<td>680</td>
<td></td>
<td></td>
</tr>
<tr>
<td>206</td>
<td>Pb</td>
<td>$s_{1/2}$</td>
<td>0.22</td>
<td>(0.1)</td>
<td>30</td>
<td>0.7</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>207</td>
<td>Pb</td>
<td>$s_{1/2}$, 1</td>
<td>9.8</td>
<td>(3.3)</td>
<td>(13)</td>
<td>420</td>
<td>3200</td>
<td>2000</td>
<td>1200</td>
</tr>
<tr>
<td>208</td>
<td>Pb</td>
<td>$s_{1/2}$</td>
<td>0.35</td>
<td>1.0</td>
<td>11.6</td>
<td>4.1</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$P_{1/2}$</td>
<td>0.13</td>
<td>(1)</td>
<td>500</td>
<td>64</td>
<td></td>
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<tr>
<td></td>
<td></td>
<td>$P_{3/2}$</td>
<td>0.35</td>
<td>(1)</td>
<td>175</td>
<td>62</td>
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<tr>
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<td>$d_{3/2}$</td>
<td>0.12</td>
<td>(1)</td>
<td>65</td>
<td>8</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
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</tr>
<tr>
<td>209</td>
<td>Bi</td>
<td>$s_{1/2}$, 5</td>
<td>0.4</td>
<td>0.65</td>
<td>14</td>
<td>3.6</td>
<td>164</td>
<td>34</td>
<td>13</td>
</tr>
<tr>
<td></td>
<td></td>
<td>$P_{3/2}$, 5</td>
<td>4</td>
<td>0.25</td>
<td>14</td>
<td>14</td>
<td></td>
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<td></td>
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<tr>
<td></td>
<td></td>
<td>$d_{3/2}$, 5</td>
<td>0.2</td>
<td>1.0</td>
<td>14</td>
<td>20</td>
<td></td>
<td></td>
<td></td>
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<tr>
<td>232</td>
<td>Th</td>
<td>$s_{1/2}$</td>
<td>4.5</td>
<td>0.84</td>
<td>0.007</td>
<td>0.03</td>
<td>21</td>
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<td></td>
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<tr>
<td></td>
<td></td>
<td>$P_{3/2}$</td>
<td>1.5</td>
<td>1.6</td>
<td>0.004</td>
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<td></td>
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<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$P_{1/2}$</td>
<td>0.5</td>
<td>1.6</td>
<td>0.007</td>
<td>0.006</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>238</td>
<td>U</td>
<td>$s_{1/2}$</td>
<td>0.5</td>
<td>1.1</td>
<td>0.018</td>
<td>0.01</td>
<td>22.5</td>
<td></td>
<td></td>
</tr>
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<td></td>
<td>$P_{3/2}$</td>
<td>2.0</td>
<td>1.6</td>
<td>0.009</td>
<td>0.03</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>239</td>
<td>Pu</td>
<td>$s_{1/2}$, 1</td>
<td>1.2</td>
<td>1.3</td>
<td>0.003</td>
<td>0.005</td>
<td>40</td>
<td></td>
<td></td>
</tr>
<tr>
<td></td>
<td></td>
<td>$P_{3/2}$, 1</td>
<td>2.1</td>
<td>2.3</td>
<td>0.003</td>
<td>0.01</td>
<td></td>
<td></td>
<td></td>
</tr>
</tbody>
</table>

a) Valence width calculated for the valence neutron and resonance spins $(l,J)_{\lambda}$.  
b) At thermal energy - accurate only to a factor of two if unadjusted optical model parameters are used.  
c) Strength function and level spacing data obtained from Ref. MG73 or from the cited references.  
d) Average value obtained from detailed calculations reported in the literature (Ref.) $D_{l,J}$ is the average level spacing per spin state.  
e) $\Gamma_Y^R = \Gamma_Y^V (s) - \Gamma_Y^V (p)$ in the 3s region  
$\Gamma_Y^V (p) - \Gamma_Y^V (s)$ in the 3p region
transitions are to strong single particle states, they are excluded from the statistical component by definition. Nevertheless, detailed spectral calculations would be preferred, but these require a knowledge of the \( \gamma \)-ray strength function. Since this is not well known in the 3p region, such calculations have not been made. It is therefore assumed that \( \langle \Gamma^S_\gamma \rangle \approx \langle \Gamma^s_\gamma \rangle \) in the 3p region, and similarly for the 3s region \( \langle \Gamma^S_\gamma \rangle \approx \langle \Gamma^s_\gamma \rangle \), but it is noted that the statistical widths in each region may be larger by up to a factor of two. A constant value of 40 meV is assumed for \( \langle \Gamma^S_\gamma \rangle \) in the 4s region, since few p-wave radiative widths have been measured.

It is the non-statistical component which should be compared with the valence calculations, and the ratio \( \frac{\langle \Gamma^V_\gamma \rangle}{\langle \Gamma^R_\gamma \rangle} \) is shown in Figure 5.1c. For even-even targets with 3s and 3p regions, the valence process accounts for a large fraction of the residual, non-statistical widths. Should the statistical widths be up to a factor of two larger, then \( \frac{\langle \Gamma^V_\gamma \rangle}{\langle \Gamma^R_\gamma \rangle} > 1 \) in a number of cases, and evidence for moderate GDR depletion could be advanced. However, there would be many more nuclides with \( \langle \Gamma^V_\gamma \rangle \approx \langle \Gamma^R_\gamma \rangle \).

Save for possibly Ce, Ba and some Nd isotopes, the valence process can only play a very minor role in the 4s region.

A case for depletion of the valence El strength by the GDR would be found if the calculated and valence total radiative widths were very much larger than the observed non-statistical component. This is clearly not the case. When, as in the 4s region, the valence component is estimated to be very small, it is not possible to comment either way. However, we find that many cases exist in the 2p, 3s and 3p regions where the calculated valence strength is a large fraction of the non-statistical radiative width. In the 3s region there are no unequivocal cases where the calculated valence width for individual resonances exceeds the measured value (see Figure 5.3). In the 3p region a number of cases are found, but recent measurements of neutron widths in \(^{90}\text{Zr}\) have eliminated
In addition, there are instances \{Ch+76,Bh+75,Mu+76\} where the partial valence width ($\Gamma_{V}^{\lambda\mu}$) has been found to be larger than the measured partial width ($\Gamma_{\lambda\mu}$) but these may well result from interference rather than depletion effects.

In general, the calculated total valence widths are comparable with or less than the non-statistical components. The data are therefore consistent with the decoupling of valence strength to the threshold region. This result was initially unexpected because the Brown and Bolsterli schematic model \{BB59\} predicts the elevation of El strength into the GDR.

Early attempts \{La71\} to explain this effect in terms of the diagonalisation of El states in the 3s and 4s regions showed that decoupling of the $2p^{-1}$ 3s and $3p^{-1}$ 4s states from the GDR could occur for zero-range two-body forces. However, results for more realistic forces did not support this conclusion. Another cause for the decoupling of low $\ell$-orbits arises when the particle-hole states are combined with a dense set of 'complicated' states to generate the observed fine structure resonances. In general, these resonances have boundary conditions at the nuclear surface that are different from those of the p-h states, resulting in a boundary condition mixing \{GLZ74\} which can modify the coupling to the GDR. However, the effect is more to distort the normal shape of the single particle state rather than to strengthen it, and to decrease the El strength retained in the threshold region \{La76\}.

Lane has shown in a simple two-state model \{La74b\} that it is possible to have the El strength from the internal part of the wave function elevated to the GDR when the external component is largely exempt. The interaction between an El excitation and the GDR is treated explicitly in S-matrix theory to yield the El particle strength which contains both coupled and uncoupled components. For low $\ell$-values for which the external part dominates, appreciable El strength can be
Fig. 4.1
Capture γ-ray spectra for $^{40}\text{Ca}(n,\gamma)$ at 40, 135, 210 and 430 keV
decoupled from the GDR.

4.2 **VALENCE TRANSITIONS FROM d-WAVE CAPTURE IN $^{40}$Ca**

The existence of non-statistical effects in neutron capture $\gamma$-ray spectroscopy has been attributed to the decoupling of El strength from the GDR. However, such decoupling is only expected to occur for low angular moments \{GLZ74, La74b\} and it is important to establish whether it occurs beyond $l = 0,1$ where prominent single particle $\gamma$-ray transitions and correlations between different reaction channels have often been observed.

Experiments with intermediate energy neutrons \{BAK69\} have provided evidence for the occurrence of $\gamma$-ray transitions following d-wave neutron capture, particularly in the region of the d-wave strength function maximum at $A = 40-70$, and led to the prediction of d-wave valence effects at relatively low neutron energies \{BT75\}.

The $^{41}$Ca nucleus is particularly suitable for the study of d-wave effects since its ground state is an almost pure single particle $f_{7/2}$ state \{BSB65\}. Of the possible reaction modes for the $^{40}$Ca(n,$\gamma$)$^{41}$Ca ground state reaction, only the capture of a $d_{5/2}$ neutron followed by El decay is likely. Capture of neutrons with other angular momenta requires a $\gamma$-decay mode of order E2 or higher. As the first excited state of $^{41}$Ca is at 1.943 MeV, a NaI detector can readily resolve ground state transitions and separate d-wave capture from s-wave.

The AAEC 3 MeV pulsed Van de Graaff accelerator was used to produce a neutron beam with an energy spread of ~50 keV at energies of 40, 135, 210 and 430 keV. A 20 cm x 15 cm NaI crystal was employed as a photodetector and standard fast timing techniques were applied (section 2.2).

In Figure 4.1 capture $\gamma$-ray spectra are displayed after background subtraction and unfolding by least squares analysis with the response functions of the NaI detector. Spectra are normalised to equal area above 2.9 MeV. Weak transitions to the $f_{7/2}$ ground state are observed.
for 40 keV neutron capture with an increase in strength for higher neutron energies. Other prominent transitions in the spectra at energies of 6.1 MeV and 4.5 MeV correspond to transitions to the 1.943 MeV $p_{3/2}$ state and the 3.945 MeV $p_{1/2}$ state in $^{41}$Ca. Both of these states are strongly single particle in nature. The latter two transitions are explained by s- or d-wave neutron capture followed by E1 radiative decay, but the transitions to the $f_{7/2}$ ground state can only be explained by d-wave capture unless a multipolarity of order E2 (or higher) occurs.

The spectra confirm the presence of ground state transitions as observed in early Ge(Li) measurements below 100 keV [AKB69]. However, the observed intensity at 40 keV (0.8%) is below the sensitivity of recent Ge(Li) measurements at this energy [Mu+76].

The statistical model strength to the $7/2^-$ ground state in $^{41}$Ca is expected to be negligible since this is almost a pure single particle state (spectroscopic factor ~1.0). Consequently, E1 transitions can only occur from appropriate valence or (2p-1h) configurations in the initial state.

Because the statistical model is inadequate to predict the γ-ray strength to the single particle final states, optical model valence neutron calculations were made. This theory has been successful in accounting for that component of the radiative widths which is correlated with the entrance channel width in a number of closed shell nuclei. In the optical model, the partial valence capture cross section $\sigma$ to final state $\mu$ can be calculated [LM74,BT75] from the equation

$$\sigma_{\gamma \mu} (\ell, J) = 2\pi^2 \chi^2 (2J+1) \left[ \frac{F_{\gamma \mu}^2}{1 + \sum_{\mu} F_{\gamma \mu}^2} \right] \text{Im}(\tan \delta_{\ell J})$$

where $\delta_{\ell J}$ is the optical model phase shift for the partial wave $(\ell, J)$ and

$$F_{\gamma \mu} (\ell, J) = \text{Im}\langle u_{\ell J} | H_E^{(1)} | \mu \rangle \cos \delta_{\ell J} \text{Im}(\tan \delta_{\ell J})$$

... (4.1)
Here \( u_{LJ} \) is the optical model wave function, \( \gamma \) the final state wave function and \( H^{(1)}_E \) is the electric dipole operator with the usual factors included. Optical model calculations of the cross sections \( \sigma_{\gamma_0} (d_{5/2}) \) and \( \sigma_{\gamma_1} (s_{1/2}) \) were made using these formulae and the results are contained in Table 4.2.

To analyse the experimental data for comparison with the optical model predictions, the following procedure was used at each energy. The spectral fractions \( (R_0 \text{ and } R_1) \) respectively were first obtained from the capture spectrum assuming the sum of \( \gamma\)-ray intensities above 2.9 MeV was proportional to the total capture cross section. The total capture cross section, \( \sigma_{\gamma_0} \), for \(^{40}\text{Ca}\) was then calculated using the usual formula for average neutron capture cross sections and the average resonance parameters for \(^{40}\text{Ca}\) reported by Musgrove et al. (Mu+76):

\[
(\bar{\Gamma}_Y (s) = 1.5 \text{ eV, } \bar{\Gamma}_Y (p) = 0.36 \text{ eV, } \bar{\Gamma}_Y (d) = 0.74 \text{ eV, } D_0 = 37 \text{ keV})
\]

The product of the ground state spectral fraction, \( R_0 \), with the total neutron capture cross section, \( \sigma_{\gamma_0} \), gave directly the partial d-wave capture cross section to the ground state, \( \sigma_{\gamma_0} (d_{5/2}) \). These quantities are also given in Table 4.2.

The partial s-wave cross section, \( \sigma_{\gamma_1} (s_{1/2}) \) to the \( p_{3/2} \) first excited state could not be directly obtained in this way. It was first necessary to estimate the d-wave contribution to the observed experimental peak. Optical model calculations using the parameters of Moldauer (Mo63) indicate that the ratio \( \sigma_{\gamma_1} (d_{5/2})/\sigma_{\gamma_0} (d_{5/2}) \) should be 2.5±0.1 over the entire experimental range and that the d_3/2 contribution to the decay to the first excited state should be small. This ratio is largely independent of quantities such as the d-wave strength function in which considerable experimental uncertainty exists. The cross section, \( \sigma_{\gamma_1} (s_{1/2}) \) is then given by
TABLE 4.2
SPECTRAL FRACTIONS ($R_\mu$) AND PARTIAL CAPTURE CROSS SECTIONS ($\sigma_{\gamma u}$) IN $^{40}$Ca(n,γ)

<table>
<thead>
<tr>
<th>E(n) (keV)</th>
<th>$\sigma_{n\gamma}$(total) (mb)</th>
<th>$R_0$</th>
<th>$R_1$</th>
<th>$\sigma_{\gamma 0}(d_{5/2})$ mb</th>
<th>$\sigma_{\gamma 1}(s_{1/2})$ mb</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td></td>
<td>(±20%)</td>
<td>(±30%)</td>
<td>O.M.</td>
<td>Exp.</td>
</tr>
<tr>
<td>40</td>
<td>10.0</td>
<td>0.008</td>
<td>0.50</td>
<td>0.58</td>
<td>0.08±0.04</td>
</tr>
<tr>
<td>135</td>
<td>7.2</td>
<td>0.034</td>
<td>0.34</td>
<td>0.41</td>
<td>0.24±0.10</td>
</tr>
<tr>
<td>210</td>
<td>5.1</td>
<td>0.065</td>
<td>0.34</td>
<td>0.30</td>
<td>0.33±0.15</td>
</tr>
<tr>
<td>430</td>
<td>3.0</td>
<td>0.068</td>
<td>0.21</td>
<td>0.2</td>
<td>0.2±0.1</td>
</tr>
</tbody>
</table>
The results obtained are listed in Table 4.2.

From the Brown-Bolsterli \{BB59\} schematic model of the GDR, the dipole strength in the threshold region is expected to be gathered up by the residual interaction into one coherent state at 10-20 MeV. Gyarmati et al. \{GLZ74\} have investigated ways whereby states of low orbital angular momentum (i.e. \(l = 0,1\)) may retain their strength in the threshold region despite the effect of the residual interaction, thus explaining the applicability of the valence neutron model in this region. The present experiment investigates the extent of depletion of d-wave strength over valence model predictions near threshold.

In Table 4.2 it can be seen that the optical valence model severely overestimates the d-wave cross section at 40 keV, to a lesser extent at 135 keV, but predicts the experimental results accurately at 210 keV and 430 keV.

The substantial disagreement at 40 keV may point to depletion of the d-wave dipole strength by the GDR at low energies. However, fluctuations in the d-wave resonance spacings and neutron widths probably account for this result and must necessarily qualify the depletion interpretation. On the other hand, it is apparent that above 200 keV the valence neutron model adequately accounts for the d-wave dipole strength to the ground state. In the case of s-waves, although the experimental error is large, the valence model adequately explains the data at all energies. It is apparent, therefore, that significant depletion of the d-wave E1 strength by the giant resonance does not occur in the \(^{40}\text{Ca}(n,\gamma)^{41}\text{Ca}(\text{g.s.})\) reaction at neutron energies above 200 keV.

\[ \sigma (s_{1/2}) = (s_1 - 2.5 R_0)^\gamma \sigma \]
INITIAL AND FINAL STATE CORRELATIONS

5.1 MEASUREMENTS

The valence partial radiative width is proportional to the reduced neutron widths of both the initial and final states. Consequently, for transitions with large valence components, comparable correlations are expected for the initial states \[ \rho_I (\Gamma_{\lambda n}, \Gamma_{\lambda \mu}) \] and final states \[ \rho_F ((2J +1)\theta, \Gamma_{\lambda \mu}/E^2) \]. The valence correlation \( \rho_V (\Gamma_{\lambda \mu}, \Gamma_{\lambda \mu}) \) combines both effects and takes into account the detailed dependence of valence transitions on spin factors and the variation of the radial integrals with neutron energy. The linear correlation coefficient is given by

\[
\rho(a,b) = \frac{\sum_i [(a_i - \langle a \rangle)(b_i - \langle b \rangle)]}{\sqrt{\sum_i (a_i - \langle a \rangle)^2 \sum_i (b_i - \langle b \rangle)^2}} \quad \ldots(5.1)
\]

where \( a, b \) are pairs of parameters such as those given above.

Final state correlations can be obtained from the measurement of relative γ-ray intensities, but absolute values are needed for the valence correlation. High resolution Ge(Li) γ-ray detectors are required in order to ensure adequate resolution of γ-ray transitions. Similarly, high quality (d,p) measurements are needed to ensure definite \( \ell_n \) assignments and the observation of the total spectroscopic strength, since correlations should be restricted to a specific \( \ell_n \) population. Problems can arise if dominant transitions are not observed, since the inherent large errors on weak (n,γ) or (d,p) intensities can result in spurious correlations. A case in point is the final state correlation for thermal capture in \(^{142}\text{Nd}\), where \( \rho_I \) varies from 0.05 to 0.53 depending on the choice of data. A further drawback, particularly for Van de Graaff measurements in the keV energy region, is that often only the strongest transitions are observed above the relatively high backgrounds. While these may be well correlated with the stripping strengths
(as in the 3s region), large errors can result from the limited sample size.

While similar problems also occur with the determination of initial state correlations, an additional source of error must be noted. Most capture cross section measurements employ poor γ-ray resolution, high efficiency detectors, which may be sensitive to scattered neutrons. Resonances with the largest neutron widths scatter the most neutrons which will cause secondary γ-ray events in the detector. If these events are unresolved in time from the genuine capture events, an enhanced radiative width will be observed, resulting in a large initial state correlation. Consequently, the determination of the neutron sensitivity of the capture detector is of utmost importance in correcting for the prompt detection of resonance scattered neutrons. This effect has been studied in detail in section 3.3 for fluorocarbon scintillators in close proximity to the capture detector. It was found that when $\Gamma_{\text{ln}} > 1000 \Gamma_{\text{ly}'}$, a significant correction is required. Further when the scattered neutron energy coincides with a resonance energy of the isotopes which constitute the detector and environs, the prompt background can enhance the observed γ-ray yield by a factor of three. Monte Carlo analyses are required to account for single and multiple scattering events which result in capture γ-rays from both the sample and environs.

In some cases, published data have been corrected for prompt background effects, with a subsequent reduction in the initial state correlation. A notable case is $^{138}$Ba where the average s-wave width was reported to be a factor of three greater than other isotopes in the 4s mass region [Mu+75]. Subsequent re-analysis of this data using the Monte Carlo method showed that the prompt neutron background was severely underestimated because the target thickness (1.37 cm) had not been adequately taken into account. Since for the largest resonances, ≥80% of the observed capture yield can be ascribed to prompt background effects, no worthwhile estimate
of the magnitude of the s-wave radiative width nor of the correlation coefficient can be made.

These comments apply only to low resolution γ-ray detectors. When Ge(Li) detectors are used to obtain partial initial state correlations, the prompt background contribution is normally separated out by γ-ray analysis. Van de Graaff measurements near the Li(p,n) threshold are also exempt since scattered neutrons can often be separated out in time from capture events in the sample.

An additional source of information on final state correlations is available in the averaging experiments of Bollinger and Thomas (BT70). Averaged γ-ray intensities are measured for resonance capture in an energy range determined by the absorption of neutrons in a boron filter and the 1/E reactor spectrum. This combination limits the energies of neutrons captured by the sample to a FWHM ~700 eV, with a maximum intensity near 90 eV. This range is sufficient to ensure good averaging of γ-ray intensities over many resonances in high level density nuclides, and so dramatically reduce the expected Porter-Thomas fluctuations (variance = 2/n). An important result of these measurements is that most nuclides with high level densities do exhibit a uniform banding of reduced γ-ray intensities to final states with the same spin and parity. This effect, as noted in section 1.1, supports a statistical capture mechanism since the reduced γ-ray intensities are independent of the nuclear structure of the final states. While final state correlations have not been directly obtained, the above result implies that such correlations do not exist. We therefore assign $\rho_f = 0$ for the relevant nuclides, but refrain from estimating the error in this result.

The initial and final state correlations calculated from the available data are listed in Table 5.1. The correlation data for $^{19}$F to $^{243}$Am are plotted in Figure 5.1a,b. In most cases for $\rho_i$, total radiative widths are used rather than partial widths. However, this limita-
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<td>0.21</td>
<td>MG73</td>
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</table>

a) Number of resonance total radiative widths - n_λ x n_μ corresponds to n_λ resonances with n_μ partial widths

b) Large thermal cross section (barn) signifies dominant resonant component.

- \( \bar{P} = \) unweighted average of \( P \) over n_λ resonances.
- \( \bar{\rho_f} = \) correlation of average intensities over n_λ resonances.
- \( \rho_v = \rho(\Gamma_\lambda, \Gamma_\mu, \lambda, \mu) \) is the valence correlation.

(\( \langle \gamma, \gamma \rangle \)) - reaction type; note compound nucleus is mass A, not (A+1) as for capture.

- \( \bar{\rho} = \) averaged intensities in eV range imply zero correlation [BT70].

- \( N \times n_\mu \) is the average over N unresolved resonances for n_μ final states.

d) \( E^3 \gamma \) energy dependence assumed.

- Results invalid according to ref. Mu72.
Fig. 5.1
Variation of (a) initial and (b) final state correlations and (c) of the relative valence radiative width $\frac{\langle \Gamma^V \rangle}{\langle \Gamma^R \rangle}$ with mass number.
tion appears to be of little consequence since the statistical component of the radiative width is expected to have a narrow, Gaussian distribution, which will have little effect on the correlation coefficient. Further, statistical amplitudes \( (\Gamma^S_{\lambda \mu})^\frac{1}{2} \) to low lying final states with large, single particle components are very small and are unlikely to interfere significantly with the corresponding valence amplitude \( (\Gamma^V_{\lambda \mu})^\frac{1}{2} \). Decomposing the partial widths into valence and statistical amplitudes, the expected initial state correlation coefficient is then

\[
\rho_i (\Gamma^S_{\lambda n}, \Gamma^V_{\lambda Y}) = \sqrt{1 + \frac{1}{n} \frac{\langle \Gamma^S_{\lambda Y} \rangle^2}{\langle \Gamma^V_{\lambda Y} \rangle^2}}, \quad \ldots (5.2)
\]

where \( n \) is the number of effective channels in the statistical width \( \Gamma^S_{\lambda Y} \). For \( n \sim 10 \) and \( \langle \Gamma^V_{\lambda Y} \rangle \sim 0.1 \langle \Gamma^S_{\lambda Y} \rangle \), \( \rho_i \sim 0.3 \). The use of total radiative width correlations therefore provides a valuable tool in the study of resonance capture mechanisms, even when valence effects are expected to be small.

The final state correlations are of the type \( \rho_F (\Gamma^S_{\lambda n}, \Gamma^V_{\lambda Y}) \) where the partial widths are averaged over \( n \) resonances, or \( \bar{\rho}_F \) which is the unweighted average of single resonance, final state correlation coefficients. In order to extend the mass range of final state correlations, extensive use has been made of thermal data where a large thermal cross section indicates a dominant resonance contribution. An earlier review of final state correlations was presented by Mughabghab \{Mu74\}.

It has been customary to quote the probability that an observed correlation was significantly different from zero. This was achieved by Monte Carlo sampling from a Porter-Thomas distribution of reduced neutron widths and of uncorrelated partial radiative widths for a given sample size. However, Figure 5.1a,b shows that correlations are widespread and interest now lies in testing the correlation against the value predicted by theory. We adopt, therefore, the convention
of quoting a standard deviation derived from the calculated distribution of \( \rho \) for a given sample size \( n \), taken from an uncorrelated parent population. The standard deviation of this distribution is approximated by \( (n-1)^{-1} \).

The uncertainties resulting from errors in measured values of the partial or total radiative widths \( \{a_i\} \) have not been estimated. These errors can be obtained from the relation \( \{CL71\} \)

\[
(\Delta \rho)^2 = \frac{\sum_i (b_i - \overline{b})^2}{\sum_i (a_i - \overline{a})^2 \sum_i (b_i - \overline{b})^2}, \quad \text{...(5.3)}
\]

where \( \Delta \rho \) is the standard deviation of \( \rho \) due to the experimental standard deviation \( \epsilon_i \) in \( a_i \).

The significance level of a measured correlation coefficient (i.e. probability that \( \rho \) is consistent with zero) can be obtained using Fisher's transformation \( \{KS69\} \)

\[
z = \frac{1}{2} \ln \left[ \frac{(1+\rho)}{(1-\rho)} \right], \quad \text{...(5.4)}
\]

with variance \( \text{var}(z) = (n-3)^{-1} \),

which reduces the \( \rho \) distribution to near Gaussian form. Detailed tests of this transformation have been made \( \{Ch+76\} \) which show its equivalence with Monte Carlo calculations down to a sample size of \( n = 5 \). Additional properties and tests of correlation distributions are given by Baudinnet-Robinet \( \{Ba74\} \).

5.2 SYSTEMATICS

Large values of \( \rho \) for \( p \)-wave resonances in the 2p and 3p regions and for \( s \)-wave resonances in the 3s and 4s regions are found (Figure 5.1). These results are mirrored by the final state correlations below \( A < 100 \), but there is a serious shortage of final state data in the 4s region and above, and \( (d,p) \) measurements are needed to fill this gap.

In the region \( 60 < A < 88 \) there are very little data on \( p \)-wave radiative widths. A combination of smaller level spacings away from
the closed shells, the filling of final p-wave states, and the passing of the peak in the s-wave neutron strength function, leads to the expectation that statistical radiative theory should apply to these nuclides.

A further insufficiently charted region is $A \sim 100-130$. There is evidence for non-statistical capture in the isotopes of tin at the $Z = 50$ magic number \cite{Bh+68} and anomalous capture $\gamma$-ray spectra have also been observed in Ag, Sn, I and Cs \cite{BS62b}. However, these results have led to few estimates of initial or final state correlations in this region.

Returning to the regions of the 3s and 3p strength function size resonances, it is apparent that $\langle \rho_I \rangle - \langle \rho_F \rangle \sim 0.5$ for most nuclides. However, if valence transitions were the only non-statistical effect, much larger correlations would be expected from equation 5.2 since $\langle \Gamma_Y^V \rangle - \langle \Gamma_Y^S \rangle$. To assist in the interpretation of these data we therefore assume that there exists a non-statistical capture mechanism which is uncorrelated with the neutron widths. We further assume that the statistical component of the decay to final states with large spectroscopic factors is negligible. If the uncorrelated width is denoted by $\Gamma_{\lambda \mu}^U$, we write

$$
\Gamma_{\lambda \gamma} = \sum_{\mu} \left[ (\Gamma_{\lambda \mu}^V)^2 + (\Gamma_{\lambda \mu}^U)^2 \right] + \Gamma_{\lambda \gamma}^S \quad \text{(5.5)}
$$

where

$$
\Gamma_{\lambda \gamma}^S = \sum_{\gamma} \Gamma_{\lambda \gamma}^S,
$$

and $\mu$ designates those final states $\gamma$ with large spectroscopic factors. The uncorrelated component interferes with the valence widths and reduces the observed correlation which can be expressed \cite{Mu+76d} as

$$
\rho_I = \frac{\langle \Gamma_Y^V \rangle}{\langle \Gamma_Y^V \rangle + \langle \Gamma_Y^U \rangle} \cdot \frac{\sqrt{\sigma_{V}^2}}{\sqrt{(\sigma_{S}^2 + \sigma_{U}^2 + \sigma_{V}^2)}} \quad \text{(5.6)}
$$

where $\sigma_{\cdot}^2$ are the variance of the $S$, $U$ and $V$ components. Since the
Fig. 5.2
Mass dependence of radiative width components in (a) 3s, (b) 3p, (c) 4s regions
valence widths are distributed as chi-square with one degree of freedom, we note that \( \sigma_{V}^{2} = 2 \langle \Gamma_{V} \rangle^{2} \) whereas \( \sigma_{S}^{2} = \frac{2}{n} \langle \Gamma_{S} \rangle^{2} \).

In the 3s, 3p and 4s regions a definite mass dependence occurs for \( \langle \Gamma_{s} \rangle, \langle \Gamma_{p} \rangle, \langle \Gamma_{V} \rangle, \langle \Gamma_{U} \rangle \) (Figure 5.2a,b,c). The regional averages of these components of the total radiative width are given in Table 5.2.

5.3 3s REGION

Low lying states in the 40 < A < 70 mass region are all of positive parity and many have large \( p_{1/2} \) and \( p_{3/2} \) single particle amplitudes. Since M1 transitions to these states will be hindered with respect to E1 transitions, the \( E_{γ}^{3} \) energy dependence will result in significant differences between the s- and p-wave total radiative widths. On average, we find that \( \langle \Gamma_{s} \rangle \approx 3 \langle \Gamma_{p} \rangle \) in the 3s region.

The statistical component of the s-wave radiative width is expected to be somewhat larger than \( \langle \Gamma_{p} \rangle \) since statistical transitions to the low lying states will be E1 rather than M1. (While the ratio of E1/M1 strength is found to be \( \approx 7 \) for A >100 (Bo73), a reliable value has not yet been established for the 3s region.) On the other hand, final state correlations have been observed for p-wave resonances in \( ^{56}\text{Fe} \) (Al+76b) which are indicative of non-statistical effects. This result suggests that for p-wave resonances the statistical component may be substantially less than \( \langle \Gamma_{p} \rangle \). We therefore assume that \( \langle \Gamma_{s}^{S} \rangle \approx \langle \Gamma_{p} \rangle \). If extra statistical strength is present, it will occur to the low lying states and will therefore interfere with the valence component in the same way as does the uncorrelated component. From equation 5.6 and Table 5.2, the expected initial state correlation is estimated to be \( \rho_{I} \approx 0.5 \), which compares favourably with the regional average \( \rho_{I}(3s) = 0.4 \).

These results therefore support the presence of a second, non-statistical mechanism which is uncorrelated with the reduced neutron widths. Since final state correlations are observed \( \{\text{Ko73,Bi+73,Ba+77}\} \), this mechanism may also favour transitions to final states with large,
TABLE 5.2

AVERAGE RADIATIVE WIDTH COMPONENTS IN THE 3s, 3p AND 4s REGIONS

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<th>3p</th>
<th>4s</th>
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<td>86</td>
<td>63</td>
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<td>(40)</td>
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<td>0.6</td>
<td>~0.3</td>
<td>0.5</td>
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</table>

$\bar{\Gamma}_S$ is the statistical component

$\bar{\Gamma}_R = \bar{\Gamma}_\gamma (\ell) - \bar{\Gamma}_S$

$\bar{\Gamma}_U = \bar{\Gamma}_R - \bar{\Gamma}_V$
single particle amplitudes. The observed initial and final state correlations can therefore be accounted for by the assumption that the valence process interferes with a second single-particle capture mechanism. This process appears to exhibit a size resonance at $A \sim 54$, as is evident from Figure 5.2.

5.4 3p REGION

A similar situation prevails in the $88 \leq A \leq 98$ mass range, except that the average magnitude of the valence width is twice that of the uncorrelated width. The assumption that $\langle \Gamma^S \rangle - \langle \Gamma^s \rangle$ appears valid because s-wave resonances exhibit initial state correlations which are consistent with zero. Evidence for final state correlations in s-wave resonances is not found in $^{93}$Nb [HT75a,Ri+69], but could be present in $^{92}$Mo [WC70].

The 3p region affords the opportunity for a detailed comparison of results for total and partial radiative widths. The best way to study valence effects is to measure absolute partial radiative widths for a large number of resonances and final states. While relative measurements can give correlation data, absolute results are needed for the comparison of valence widths with the measured values. The initial success of the valence model [Mu+71] in reproducing the partial widths of $^{92}$Mo and $^{98}$Mo could not be maintained when absolute measurements were made which extended to higher neutron energies [WS73,Ch+76]. These later results, together with the $(\gamma,n)$ data on $^{91}$Zr [TJ74], and the extensive total capture data, have indicated the presence of additional non-statistical capture mechanisms in the 3p region.

(i) $^{91}$Zr Ground State

Measurements of the $(d,p)$ spectroscopic factor of the ground state of $^{91}$Zr [BH70] has shown that it is a single particle $d_{5/2}$ state with $\Theta^2 = 1.0$. Consequently, transitions to this state can only occur from single particle components of the neutron resonances. Toohey and
Jackson (TJ74) measured ground state partial radiative widths in the $^{91}\text{Zr}(\gamma,n)$ reaction for 36 $p_{3/2}$ resonances. After some ($\ell,J$) assignments are updated (MCH77), the average ground state width is $\langle \Gamma \rangle = 143 \text{ meV}$, of which the valence process accounts for 95 meV. Toohey and Jackson assumed the residual component (48 meV) resulted from statistical effects, but since $\theta^2 = 1.0$, we can assume $\langle \Gamma^S \rangle \approx 0$ and that a further non-statistical component, with width about half the valence width, is present. If this component is uncorrelated with the reduced neutron widths, then equation 5.6 gives $\rho = 0.6$. This result is in excellent agreement with the measured correlation $\rho = 0.59 \pm 0.17$.

Measurements of the total radiative widths in the $^{90}\text{Zr}(n,\gamma)$ reaction (Bo+75a) also revealed initial state correlations with $\rho = 0.58$ for 37 $p_{3/2}$ resonances. Thus the excellent agreement between the partial and total radiative width measurements confirmed for the first time the value of total width correlations.

By analogy with the valence width equation, the reduced width ($\alpha$) for the uncorrelated component is defined as

$$\Gamma^U_{\lambda\mu} = a\theta^2 <D> E^3 Z^2 / A^2 \quad (\text{meV}, D \text{ eV}, E \text{ MeV}) \quad \cdots (5.7)$$

For the $^{91}\text{Zr}$ ground state $\alpha = 1.5 \times 10^{-4}$. This result is comparable to that obtained from the same state in $^{93}\text{Mo}$ ($\theta^2 = 0.64$) and $^{99}\text{Mo}$ ($\theta^2 = 0.21$), i.e. $\alpha = 2.1 \times 10^{-4}$ and $1.4 \times 10^{-8}$ respectively.

However, for the second $d_{5/2}$ excited state in $^{93}\text{Mo}$ ($\theta^2 = 0.12$) we find the average partial width would require $\alpha = 15.5 \times 10^{-4}$, an order of magnitude larger than the above values. Since this decay probably feeds the excited target state component of the second $d_{5/2}$ state (i.e. $(1-\theta^2)$), excited target state components in the low lying states may be of importance in this region.

(ii) $^{92}\text{Mo}$

A comparison between the partial and total radiative width correlations in $^{92}\text{Mo}$ reveals a complex situation. Referring to Table 5.1,
Fig. 5.3
Correlation of valence and observed total radiative widths in the (a) 3s and (b) 3p regions
\( \rho_I(\Gamma^I_{\lambda n}, \Gamma_{\lambda Y}) \) and \( \rho_I(\Gamma^I_{\lambda n}, \Gamma_{\lambda \mu}) \) are both large and consistent, but only one transition to an \( s_{1/2} \) state (albeit the strongest) is significantly correlated with the reduced neutron widths.

If this transition is responsible for the total width initial state correlation \( \rho_I(\Gamma^I_{\lambda n}, \Gamma_{\lambda Y}) \), then a significant asymmetry in the initial and final correlations in the other partial channels appears to exist.

(iii) \( ^{98}\text{Mo} \)

Chrien et al. (Ch76) have found that transitions to \( d_{5/2} \) final states are significantly less correlated with the reduced neutron widths (\( \rho_I \approx 0.06 \)) than are transitions to the \( s_{1/2} \) and \( d_{3/2} \) final states (\( \rho_I > 0.8 \)). However, in contrast to the results for \( ^{92}\text{Mo} \), initial and final state correlations appear to be symmetric (see Table 5.1).

It is apparent that the \( p \rightarrow s_{1/2}, d_{3/2} \) transitions are responsible for the observed total width initial state correlations. Further, the large \( \rho_I \) values are determined by the \( 1/2^- \) resonances with the largest reduced neutron widths and these, in turn, exhibit large final state correlations. Consequently, the valence process is more applicable to \( ^{98}\text{Mo} \), with the exception of transitions to \( d_{5/2} \) states, than to \( ^{92}\text{Mo} \).

The partial and total width data therefore indicate the presence of additional, non-statistical mechanisms which appear to depend on the single particle strength of those final states with either ground or excited target state components.

The limitations of the valence model in the 3s and 3p regions are apparent in Figure 5.3, where the correlations of the total valence and measured radiative widths are shown. Many values lie outside the bounds set by the solid lines representing \( (\Gamma^S_{\gamma} + \Gamma^V_{\lambda Y}) \) and \( (\Gamma^S_{\gamma} + \Gamma^U_{\lambda Y} + \Gamma^V_{\lambda Y}) \), and reflect the variance of the data resulting from the mass dependence of the total width components.

Note that the valence widths do not often exceed the measured values and in many cases agreement with errors if found. However, in
in both regions there exists a large number of resonances in which the
valence contribution is exceedingly small (i.e. <10%). Thus, while
the valence model accounts for most of the radiative strength of
resonances with large reduced neutron widths, there clearly exists a
body of resonances which have large radiative widths and negligible
valence components.

5.5 4s REGION

The correlation data differ from the 3s and 3p regions in that a
number of negative values are found, most cf which are not significantly
different from zero. On the other hand, many large and positive values
are also found, although the calculated total valence widths are generally
only a few per cent of the measured s-wave widths. Exceptions are 138Ba
and 140Ce with N = 82, where the valence width is about one third of the
estimated s-wave radiative width. Accurate measurements of s-wave radiative
widths are needed for these nuclides to confirm this result and the large
correlations which would be expected. Since few measurements of p-wave
widths have been made, and these only at N ~82 and 126, we assume the
average statistical width given in Table 5.2, and find that zero initial
state correlations are expected. This result is inconsistent with the
average measured value $\bar{\rho}_I(4s) = 0.3$ obtained by including both positive
and negative correlations for isotopes below Pb. A more definite study
can be made for the Nd isotopes in the 142 $\leq A \leq 148$ region, where the
observed $\bar{\rho}_I$ is more than twice the expected value (Table 5.2). It there­
fore appears that valence processes are not strong enough to account for
the observed initial state correlations in this mass region. An additional
capture mechanism, correlated with the s-wave reduced neutron widths,
appears to be required which could also account for the non-statistical
spectra observed in 139La {Al+76a} and for 181 $< A < 205$ {Bs62a,LS65,Ea+75,
Ba+77}. However, in a measurement designed to test this possibility, no
initial state correlation was observed for 57 s-wave resonances in
$^{139}$La (MAM77). In particular, the anomalous $\gamma$-ray transitions to the $l_n = 3$ final states in $^{140}$La were also found to be uncorrelated with the reduced neutron widths, suggesting that different capture mechanisms occur for the even and odd-$Z$ nuclides near $N \approx 82$. New results for $^{141}$Pr (Ta+77) also show a zero correlation, confirming this conclusion.

5.6 ACTINIDES

While low initial state correlations are found in many of the actinides, the statistical errors are such that a zero correlation cannot be ruled out. Exceptions are $^{235}$U and $^{238}$U, the latter being the more significant result. The initial state correlation for 71 resonances (MG73) in $^{238}$U is 0.40±0.12, a value which is fully consistent with the correlation observed over different energy ranges, but which conflicts with the quoted 'zero' correlation observed by Wasson et al. (Wa+71) for 15 partial radiative widths from 23 resonances, and the preliminary results of Cornelis (Co+77) who observed a zero correlation in $^{238}$U.

The total width correlation is also inconsistent with other statistical properties of $^{238}$U as manifested in the observed distribution of partial radiative widths with $1.2^{+0.20}_{-0.15}$ degrees of freedom (Wa+71), the average capture $\gamma$-ray spectra at 40, 128 and 300 keV (Be62), and the narrow distribution of total radiative widths with $\sim$70 degrees of freedom (Ra+72).

Nevertheless, the average $\gamma$-ray spectra over 28 resolved resonances (Wa+71), and somewhat more in the average resonance capture measurement of Bollinger and Thomas (BT72), exhibit fluctuations of a factor of two between averaged $\gamma$-ray intensities to like parity states, when a variation of only $\sim$25% would be expected. The strongest transitions are also found to be distributed with two degrees of freedom. One of these is the 4060 keV $\gamma$-ray transition which populates a strong (d,p) final state, and requires the presence of a direct capture cross section to explain the constructive interference observed below the 6.7 eV resonance (Pr+68).
Other strong transitions are the 3991 and 3982 keV γ-rays which Wasson et al. find to be strongly correlated ($\rho_{\mu\nu} = 0.81$). The final states involved here are not observed in the (d,p) reaction (Sh+66). These states have been tentatively assigned as the $\frac{1}{2}^-$ and $\frac{3}{2}^-$ members of an octopole vibrational band built on the $\frac{1}{2}^+$ (631) Nilsson orbital. It is just this type of example that has been considered by Soloviev (So71) as being a likely candidate for correlations between pairs of partial γ-ray widths as well as initial state correlations.

Since the valence contribution to capture in the actinides is negligible, it is apparent that additional, non-statistical processes occur. The nature of these processes, and their relationship with initial and final state correlations, is considered in Chapter 7.

The question of correlations in $^{238}$U is not solely of academic importance. While they undoubtedly exist in the BNL data, the possibility remains that their origin may lie in the neutron sensitivity of the capture detectors. If so, the radiative widths will be overestimated and the capture cross section must be reduced, a result which could have a significant impact on calculations for the doubling time of fast breeder reactors (GHW68,Hu71).
CHAPTER 6

NEUTRON CAPTURE MECHANISMS IN THE ISOTOPES OF IRON

The neutron capture mechanism in the region of the 3s strength function peak is dominated by non-statistical processes. Correlations between the reduced neutron widths and radiative widths of s-wave resonances have been reported \( \text{[SHB71, BS75]} \), and direct capture effects at thermal neutron energies are now well established \( \text{[Mu74]} \).

Of particular significance are the \( \gamma \)-ray spectra for keV capture in this mass region \( \text{[ABK69, Bi+73]} \) which exhibit intense transitions to the low lying p-wave single particle states. Since these states account for most of the 2\( p_{1/2} \), 2\( p_{3/2} \) single particle strength and the s-wave strength function maximises at Fe, valence neutron capture would be expected to dominate the reaction mechanism. To test this hypothesis, a detailed analysis of capture \( \gamma \)-ray spectra and cross section results over a wide energy range was undertaken.

In the inverse \( ^{57}\text{Fe}(\gamma, n) \) reaction, Jackson and Strait \( \text{[JS71]} \) have reported intermediate structure for p-wave resonances at 230 and 606 keV which account for much of the M1 single particle strength. These results were interpreted as narrow doorway states and it is of interest also to examine these states in the total capture cross section.

6.1 GAMMA RAY SPECTRA IN \( ^{56}\text{Fe} \)

Both NaI and Ge(Li) detectors had been used for a detailed study of \( ^{56}\text{Fe} \) capture \( \gamma \)-ray spectra for energies up to 70 keV neutron energy \( \text{[ABK69, Ke71]} \). Data are also available for the 1.167 keV p-wave resonance which has a spectrum similar to thermal s-wave capture. Spectra at 7.4 MeV and 14 MeV \( \text{[Be+66, Co+65]} \), which also exhibit strong components to low lying states, were interpreted in terms of semi-direct theory because of the proximity of the E1 giant resonance.

To obtain more data at intermediate energies, measurements have been made at 40, 135, 210, 430 and 1000 keV with a 20 cm x 15 cm NaI detector.
Fig. 6.1
NaI γ-ray intensity distributions in $^{57}$Fe up to 1 MeV, normalised to equal area. Dashed lines are upper and lower bounds of low energy data.

Fig. 6.2
Ge(Li) spectra for $^{56}$Fe up to 430 keV. Low lying states of $^{57}$Fe are shown in the inset.

Fig. 6.3
Statistical model calculation of primary γ-ray spectra for $s_{1/2}$ and $p_{1/2}$ resonances.
and up to 430 keV with a 17.6% Ge(Li) detector. Neutron energy spread was 50 keV and the sample masses for the NaI and Ge(Li) experiments were 1.5 kg and 2.9 kg of natural iron, respectively.

Intensity distributions were obtained from the NaI spectra, using a library of NaI γ-ray line shapes. The results are shown in Figure 6.1 and are seen to be quite similar in shape up to 430 keV. The high energy component broadens considerably at 1000 keV, without any real reduction in its strength. As the s-wave level spacing in $^{56}$Fe is 25 keV, only 2 to 3 s-wave resonances would contribute to each spectrum and substantial fluctuations in the capture spectra might be expected. The data are, therefore, conspicuous by their overall similarity over a wide energy range. The increasing p- and d-wave cross sections relative to the s-wave component do not have a dramatic effect on the spectral shape.

To clearly identify the final states involved in the high energy component, Ge(Li) measurements were made at 135, 185, 230 and 430 keV (Figure 6.2). The energies of the primary γ-rays are seen to shift with increasing neutron energy. Transitions are observed to all low lying states up to 706 keV and to the 1265 and 1627 keV states. The ground state doublet is always prominent, while transitions to the other states become more significant at higher neutron energies.

Estimates of intensity ratios of γ-ray transitions to the four lowest energy states, relative to the summed intensities for $E_\gamma > 3$ MeV, are given in Table 6.1. The results are remarkably consistent from thermal to 1 MeV, but estimates could not be made at 7.4 and 14 MeV.

The NaI spectrum at 430 keV has been analysed for both statistical and valence model components. The theory of Troubetzkoy (Tr61) was used to calculate the statistical component. When normalised to the observed intensity at 3 MeV, insufficient strength was predicted at high energies. Valence calculations were then made for transitions from the 27.7 keV s-wave resonance ($I^+_\gamma = 1.45$ eV, $I^0_n = 9.6$ eV) and for
<table>
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<tr>
<th>Energy</th>
<th>%</th>
<th>Detector</th>
<th>$I_{1}%$</th>
<th>$I_{\bar{1}}%$</th>
<th>B/A</th>
<th>C/A</th>
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<td>Thermal</td>
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<td>Ge(Li)</td>
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<td>7</td>
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<td>0.39</td>
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<td>0.07</td>
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<td>$s_{1/2}$</td>
<td>$E_{\gamma}^3$</td>
<td></td>
<td></td>
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<td>0.47</td>
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<tr>
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<td>$E_{\gamma}^3$</td>
<td></td>
<td></td>
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<td>1.49</td>
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<td>$E_{\gamma}^3$</td>
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<td>1.03</td>
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<tr>
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<td>$E_{\gamma}^5$</td>
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<td>0.17</td>
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<td>$E_{\gamma}^5$</td>
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Results of the calculations are shown in Table 6.1 where the valence strengths $l_1^V$ (per cent) are compared with the observed intensity $I_y$ (per cent) to the $l_n = 1$ states at 0.0, 0.014 and 0.366 MeV. The valence model cannot account for the observed high energy strength. This conclusion was also suggested by Bhat et al. (Bh+71) who calculated the ground state valence component for s-wave resonances in $^{56}$Fe up to 220 keV. In most cases, the valence component falls well below the ground state radiative widths ($\Gamma^0_{\gamma}$) measured in the inverse $^{57}$Fe(\gamma,n) reaction (JS71). In this experiment, intermediate structure in p-wave resonances excited by M1 transitions was observed at 230 keV ($J^\pi_C = 3/2^-$) and 606 keV ($J^\pi_C = 1/2^-$), and this was interpreted in terms of narrow doorway states with a $(f_5/2)(f_7/2)^{-1}$ particle-hole coupled to the $^{57}$Fe ground state. Only our results at 230 keV could be affected by the 230 keV doorway state, but the observed capture spectrum is similar to those at other energies.

Statistical calculations of the primary $\gamma$-ray spectra from $s^{1/2}$ resonances show that strong transitions are expected to three groups (A,B,C) of low lying states (Figure 6.3). An $E3\gamma$ energy dependence and an El/M1 ratio of 7 was assumed. The dotted histogram represents the spectrum calculation for a $p^{1/2}$ resonance, normalised to equal area. The energies and spin and parities of discrete states were provided up to 5.36 MeV and a continuum calculation used level densities for $J_\mu = J_\lambda - 1$, $J_\lambda$, $J_\lambda + 1$ spin states which were consistent with the low lying level density and that observed at threshold (see section 8.1). Normalising to group A transitions, the relative intensities of groups B and C are given in Table 6.1 for valence and statistical calculations and for the Ge(Li) and NaI data. The NaI values are accurate to only ~20%. All measurements have been normalised to the thermal capture $\gamma$-ray strength above 3.0 MeV $\gamma$-ray energy.

Note that the ratio B/A is independent of parity but varies with
TABLE 6.2

OPTIMISED CORRELATION COEFFICIENTS IN $^{56}\text{Fe}$

$$\rho \left[ (2J + 1) \theta_n^2 \mu, I^- \mu \gamma \right]$$

<table>
<thead>
<tr>
<th>Energy (keV)</th>
<th>$\ell$</th>
<th>$n$</th>
<th>$\rho$</th>
</tr>
</thead>
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</tr>
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<td>26</td>
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<td>0-1</td>
<td>0.17</td>
</tr>
<tr>
<td>36</td>
<td>1</td>
<td>0-1</td>
<td>0.59</td>
</tr>
<tr>
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<td>1</td>
<td>0-1</td>
<td>0.26</td>
</tr>
<tr>
<td>72</td>
<td>1</td>
<td>0-1</td>
<td>0.52</td>
</tr>
<tr>
<td>Valence</td>
<td>0</td>
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<td>1.0</td>
</tr>
<tr>
<td>calculation</td>
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</tbody>
</table>
resonance spin (as a result of transitions to the $^{5/2}_-^-$ state at 136 keV). The ratio $C/A$ is dependent on parity, and only if $p$-wave capture is comparable to $s$-wave can the data be explained at 30 keV.

The statistical values should be compared with results averaged over a number of resonances as fluctuations in $\gamma$-ray intensities are otherwise expected. The ratio $B/A$, being independent of parity, is the best gauge available to compare statistical and valence models. Generally the values for the former are too large and those for the latter too small. Only for an $E^5_\gamma$ dependence and $J_\lambda = 3/2$ does the statistical model come close to the experimental data. Since the $E^3_\gamma$ energy dependence is favoured in Cu {A168} and Ni {AKS68}, the $E^5_\gamma$ calculation is discounted.

The correlation coefficient between final state spectroscopic strengths and reduced $\gamma$-ray intensities $(\rho(2J^+1)S_\mu, I_{\gamma}/E^{-n}_\gamma)$ can be optimised with respect to the exponent $n$. Results are given in Table 6.2, where $n \sim 0.1$.

The low value for the $\gamma$-ray energy dependence found for thermal capture could be regarded as evidence of direct capture {KSL74}. However, the thermal capture cross section can be accounted for by positive and negative energy resonances and since the exponent is comparable to that observed in the resonance data, a direct capture explanation cannot suffice alone.

Consequently, the $\gamma$-ray data from thermal up to 1 MeV neutron energy are not accounted or by direct, statistical or valence capture models.  

6.2 CAPTURE CROSS SECTION MEASUREMENT IN $^{56}$Fe

The capture cross section measurements were made at the 40 metre station of the Oak Ridge Electron Linear Accelerator. The target was an enriched (99.7%) sample of $^{56}$Fe metal with dimensions 2.6 x 5.2 x 0.5 cm, and thickness 0.0410 atom b$^{-1}$. ORELA operating conditions were 6 ns pulses at 800 pulses s$^{-1}$, and running time was 71 hours. A second, thin sample measurement was made to obtain more accurate results for the lower energy resonances. In this case, the target thickness was 0.0082 atom b$^{-1}$.
Fig. 6.4

$^{56}$Fe capture cross section

Fig. 6.5

Staircase plot for $l > 0$ resonance in $^{56}$Fe
and enrichment was 98.8%.

The capture data were analysed using a modified version of the ORNL/RPI Monte Carlo code. Breit-Wigner single level theory was used to generate capture and total cross sections, and the observed capture areas were fitted by an iterative process after subtraction of a calculated multiple scattering component. A prompt background correction was made to account for the detection of scattered neutrons.

Resonance parameters for observed resonances below the inelastic scattering threshold at ~870 keV are given in {Al+76}, together with energies and $g\Gamma_n$ values from Pandey et al. {Pa+75}, capture results from Hockenbury et al. {Ho+69} and ($\gamma,n$) data from Jackson and Strait.

In general, agreement with the total cross section energies and neutron widths is good. However, 58 new resonances are observed in capture up to 500 keV. A systematic energy difference of 0.1% at 100 keV, increasing to 0.2% at 300 keV, is observed.

(a) Average resonance parameters

s-wave Resonances

The s-wave population has been well defined by Pandey et al. in the multilevel analysis of total cross section data up to 470 keV. However, for neutron energies above 300 keV, s-wave resonances in the capture data are masked by nearby $l > 0$ resonance. Energies and neutron widths from Pandey et al. were used and s-wave radiative widths, listed in Table 6.3, were obtained by fitting the valleys between the $l > 0$ resonances. Representative portions of the data are shown in Figure 6.4. A further difficulty arises because background corrections also become rather uncertain at higher energies. A constant background cross section of 2 mb was assumed above 350 keV, which was consistent with the minima in the higher energy data.

Pandey et al. note that the strength function varies from $(1.88 \pm 0.94) \times 10^{-4}$ up to 200 keV to $(2.6 \pm 0.86) \times 10^{-4}$ up to 500 keV.
<table>
<thead>
<tr>
<th>( E_n ) (keV)</th>
<th>( \Gamma_0^n ) (eV)</th>
<th>( Q(P_{1/2}) )</th>
<th>( Q(P_{3/2}) )</th>
<th>( \Gamma_\gamma ) (eV)</th>
<th>( \Gamma_\gamma ) (eV)</th>
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<tbody>
<tr>
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<tr>
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<td>0.012</td>
<td>0.032</td>
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<td>0.005</td>
<td>0.015</td>
<td>0.09</td>
<td>1.0</td>
</tr>
</tbody>
</table>
Variation of $q^\Gamma \Gamma /\Gamma$ with neutron energy

$d$-wave resonances in $^{56}$Fe
This effect arises solely from increasing $\Gamma_n^0$ values since the average level spacing shows no energy dependent structure.

\( \ell > 0 \) Resonances

The observed level density for $\ell > 0$ resonances (Figure 6.5) increases up to 300 keV, is approximately constant up to 600 keV before decreasing at higher energies. The initial increase results from the observation of increasing numbers of d-wave resonances above the detectability limit. Below 80 keV $\Gamma_n^0 \ll \Gamma_\gamma$ for most d-wave resonances, although some of these are observed below 20 keV where statistics and resolution most favour their detection. Assuming that the larger resonances are p-wave with $\Gamma_n^0 \gg \Gamma_\gamma$, we find $\langle \Gamma_\gamma(p) \rangle \approx 0.3$ eV. Above 300 keV, $g_\gamma \Gamma_n^0/\Gamma_\gamma$ values begin to increase (Figure 6.6), coinciding with the reduction in level density (Figure 6.5) in the same energy range. We therefore conclude that if the p-wave radiative width is energy independent, p-wave resonances become increasingly missed because of their low value for $g_\gamma \Gamma_n^0$.

Below 20 keV, resonances are observed with small values of $g_\gamma \Gamma_n^0/\Gamma_\gamma$ which are assumed to equal $g_\gamma \Gamma_n^0$. The 1.15 keV resonance is known to be $p_{1/2}$, and a Bayes' theorem analysis suggests that the others are d-wave resonances with a level spacing of -5 keV.

Additional information on the d-wave level spacing can be obtained from the $\gamma$-ray yield just above the inelastic threshold. As the spin of the 847 keV first excited state of $^{56}$Fe is $2^+$, only d-wave resonances allow an s-wave inelastic channel which will be uninhibited by penetrability. Assigning the strongest resonances in the 860 to 900 keV region as d-wave, the level spacing is found to be $D(d) \approx 10$ keV.

An estimate of the average d-wave radiative width $\langle \Gamma_\gamma(d) \rangle$ is obtained from $g_\gamma \Gamma_n^0/\Gamma_\gamma$ values in the region 400 to 600 keV. In this range, many p-wave resonances are undetected as their areas fall below the sensitivity of the experiment. The anomalously large values of $g_\gamma \Gamma_n^0$ observed at higher energies are also excluded. Assuming all $\ell > 0$
resonances are d-wave and $g = 2.6$, we find $\langle \Gamma (d) \rangle \sim 1.2$ eV after a small correction for unresolved p-wave resonances. The d-wave level spacing is $\geq 4.6$ keV, a result comparable to the spacing estimates below 20 keV and above the inelastic threshold.

At intermediate energies, it is not possible to distinguish directly between individual p- and d-wave resonances. Angular distribution measurements by Jackson and Strait in the $^{57}$Fe($\gamma$,n) reaction lead to $1/2, 3/2$ assignments for that small fraction of resonances (~20%) seen in the inverse reaction. These authors assumed that d-wave resonances would not be observed below 300 keV. However, on the basis of our estimated values of $S_2$ and $\langle \Gamma (d) \rangle$, more than half the d-wave population will have $\Gamma > \langle \Gamma (d) \rangle$ above 150 keV. The $p_{3/2}$ assignments by Jackson and Strait above this energy are therefore open to doubt as is their postulated $p_{3/2}$ doorway state at 230 keV. As the radiative widths of these resonances are comparable to $\langle \Gamma (d) \rangle$, a d-wave assignment is suggested.

The capture $\gamma$-ray data below 80 keV suggest that most $J = 1/2, 3/2$ resonances will have ground state transitions. Since only 20% of the capture resonances are observed in the ($\gamma$,n) measurement, we conclude that the sensitivity of that measurement was adequate to observe only about one third of the $J = 1/2, 3/2$ resonances. A more unsatisfactory aspect of the ($\gamma$,n) data is the paucity of $J = 3/2$ resonances as only three were identified up to 300 keV. Jackson and Strait note that the 220 keV s-wave resonance was used to normalise the angular distribution data. However, this normalisation could have included a contribution from a probable $p_{3/2}$ resonance at 221.2 keV which would be unresolved in the ($\gamma$,n) measurement. It is possible, therefore, that a number of $J = 1/2$ assignments are incorrect, although Jackson (private communication) considers that this is unlikely. If this is the case, then the anomalous $p_{1/2}$ resonance at 629 keV could well be $d_{3/2}$; $\Gamma_0$ is thereby reduced by a
factor of two, and the postulated $p_{1/2}$ doorway becomes a less favoured interpretation of the data.

Above 400 keV, the $g_{Y}^{T}n_{Y}^{n}/T$ values (Figure 6.6) increase with increasing neutron energy, and near 640 and 780 keV anomalously large values are observed ($g_{Y}^{T} - 15$ eV). If these are $d_{5/2}$ resonances with $g = 3$, then $\Gamma_{Y}^{n} \sim 5$ eV. An $f_{7/2}$ assignment is unlikely as the neutron widths are far too large for $f$-wave resonances. It is unlikely that the resonances at 640 keV are unresolved doublets as the resonance shapes are good fits to Breit-Wigner curves (see Figure 6.7), and the large capture areas are observed as intermediate structure in the average cross section.

(b) Average Capture Cross Sections

Average capture cross sections are shown in Figure 6.8 up to the inelastic scattering threshold. The increasing error with neutron energy results from a decreasing confidence in the magnitude of the background correction. The dotted curves in Figure 6.8 are $s$-, $p$- and $d$-wave cross sections calculated from the average resonance parameters, and the sum of all partial cross sections (solid curve) is compared with the data.

Above 500 keV the calculated cross section begins to fall well below the average capture cross section. The inclusion of $f$-wave resonances in the calculation can account, in part, for the cross section above 500 keV. However, despite the $f$-wave contribution, the statistical calculation using energy independent resonance parameters cannot reproduce the intermediate structure observed in the average cross section.

Average resonance parameters for $s$-, $p$- and $d$-wave resonances are summarised in Table 6.4.
### TABLE 6.4

AVERAGE RESONANCE PARAMETERS IN $^{56}$Fe

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<tr>
<th>$\ell$</th>
<th>$D_\ell$ (keV)</th>
<th>$\Gamma_\ell$ (eV)</th>
<th>$10^4 \cdot S_\ell$</th>
</tr>
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<td>1.50</td>
<td>2.60</td>
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<tr>
<td>&lt;80 keV</td>
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<td>±0.2</td>
</tr>
<tr>
<td>2</td>
<td>≥5</td>
<td>1.2</td>
<td>1.0</td>
</tr>
<tr>
<td></td>
<td></td>
<td>±0.4</td>
<td>±0.5</td>
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</table>

### TABLE 6.5

OPTICAL MODEL PARAMETERS

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<tr>
<th>MeV</th>
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<th>$P_{1/2}$</th>
<th>$f_{5/2}$</th>
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<tr>
<td>Calculated energy</td>
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<table>
<thead>
<tr>
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<th>$V_{SO}$ (MeV)</th>
<th>$R_1f$</th>
<th>$A_1f$</th>
<th>$W$ (MeV)</th>
<th>$R_2f$</th>
<th>$A_2f$</th>
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<tbody>
<tr>
<td>-45</td>
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<td>1.24</td>
<td>0.58</td>
<td>-7.11</td>
<td>1.29</td>
<td>0.78</td>
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</table>
6.3 **VALENCE CAPTURE IN $^{56}$Fe**

In $^{56}$Fe, $\alpha$ valence neutron transitions can occur between s-wave resonances with large reduced neutron widths ($\Gamma_n^0$) and low lying $2p_{3/2}$, $2p_{1/2}$ states with large spectroscopic factors ($\theta_\mu^2$).

An optical model code was used to estimate the strength of valence transitions. This code was used with the fringe absorption optical model parameters of Cox and Cox (CC72) for valence calculations in $^{56}$Fe. However, these parameters fail to reproduce the experimentally observed centroid energies of the $2p_{3/2}$, $2p_{1/2}$, $1f_{5/2}$ shell model states (Th74). Since the valence radiative widths are sensitive to the binding energies of these shells, the central potential parameters were adjusted to obtain the agreement shown in Table 6.5. The valence widths were then relatively insensitive to small changes in the optical model parameters. The reduced valence width $\Gamma$ rapidly decreases at higher energies with a minimum near 1 MeV, as a result of a cancellation in the dipole matrix element overlap integrals for the initial (s) and final (2p) states. Valence effects for s-wave resonances with the largest reduced neutron widths above 300 keV are therefore substantially reduced.

The s-wave resonance parameters, $Q$ values for summed transitions to the $p_{1/2}$ and $p_{3/2}$ final states, and $\Gamma^V_\gamma$ estimates are given in Table 6.5. For those resonances with the largest $\Gamma_n^0$ values, the valence model accounts for up to half the total radiative width. However, the valence component contributes only 19% of the average s-wave radiative widths. The correlation coefficient $\rho(\Gamma_\gamma^{(s)}, \Gamma_n^0) = 0.5 \pm 0.4$ reflects the limited role of valence neutron capture for the s-wave resonances, and the decreasing valence effect at higher neutron energies. An estimate of the magnitude of the uncorrelated component ($\langle \Gamma_U^\gamma \rangle$) which reduces the initial state correlation, can be obtained from equation 5.6. For $\rho = 0.5$, $\langle \Gamma_U^\gamma \rangle \sim <\Gamma^V_\gamma > \sim 0.3$ ev.

Capture $\gamma$-ray spectra data (section 6.1) also show the limitations
of the valence model in accounting for anomalous high energy transitions to the low lying $p$-states ($^{1/2}_{-}$ ground state, 0.014 and 0.365 MeV $^{3/2}_{-}$ excited states). The calculated valence width to these three states is $0.73 \Gamma_Y^V$. The spectra data show that at most neutron energies these three states account for about half of all observed transitions, but the ratio of the calculated valence widths to the measured radiative width is much less than this value. Only at 27 keV does the valence model predict a significant fraction ($^{2/3}$) of the observed high energy strength (Table 6.1).

At higher neutron energies, the $\gamma$-ray results are averages over a 40-60 keV neutron energy spread. For these cases, the cross sections will be dominated by $p$- and $d$-wave resonances which also must have enhanced transitions to the low lying $p$-wave states. This is certainly the case for the 1.15 keV $p_{1/2}$ resonance which has 66% of the $\gamma$-ray strength going to these states {CBW70}. An estimate of the valence contribution to these $M1$ transitions can be obtained following the method of Lynn {Ly68}. The $M1$ valence width for the 1.15 keV resonance is given by

$$\Gamma_{\lambda \gamma}^{V}(M1) = 4.9 \times 10^{-7} \cdot a \cdot E^3 \cdot \langle D_{\lambda} \rangle \cdot 0^2 \cdot \Gamma^{1}/\Gamma^{1} = 60 \text{ meV} \quad \ldots(6.1)$$

where $\Gamma_{n}^{1} = 1.28 \text{ ev}$, $\langle \Gamma_{n}^{1} \rangle \sim 0.8 \text{ ev}$, $\langle D_{1/2}^{-} \rangle \sim \langle D_{1/2}^{+} \rangle \sim 2.5 \times 10^{5} \text{ ev}$, and $a$ = nuclear radius (fm). The valence width is only 20% of the average $p$-wave value and therefore cannot account for the spectral shape.

Unfortunately, $\gamma$-ray spectra have not been measured for resolved $d$-wave resonances. However, the average valence width can be calculated from deduced resonance parameters using the optical model. Taking $<\Gamma_{n}^{2}> \sim 1.15 \text{ ev}$ and $Q(d) \sim 0.05$ ($Q(d)$ varies slowly with energy; see Figure 1.4):

$$<\Gamma_{\lambda \gamma}^{V}(d)> \sim Q(d) <\Gamma_{n}^{2}> \sim 0.1 \text{ ev} \quad \ldots(6.2)$$

This value is -10% of the deduced average $d$-wave radiative width of 1.2 ev.
In summary, the valence model is found to account for only a small fraction of the radiative widths of s-, p- and d-wave resonances. Yet s-wave capture at thermal, p-wave capture at 1.15 keV and below 80 keV, and predominantly d-wave average capture spectra up to 1 MeV show anomalous intensities to the low lying 2p states.

The valence model also fails to reproduce the relative γ-ray intensities and the correlation with the spectroscopic factors of the final states (see section 6.1).

The capture cross section and γ-ray data point to a capture mechanism which correlates to a large degree with the final state spectroscopic factors, but which, in s-waves, is independent of the reduced neutron widths.

6.4 DOORWAY MECHANISM IN $^{56}$Fe

The concept of doorway states was introduced by Block and Feshbach {BF63} to explain the observation of intermediate structure in reaction cross sections. Ikegami and Emery {IE64} employed this concept to account for an observed anticorrelation between the thermal capture reduced γ-ray intensities $\langle I/\gamma E^3\rangle$ and stripping strengths $\langle(2J+1)\delta^2\rangle$ in the Fe isotopes. These authors suggested that (2p-1h) states are excited in capture and decay strongly to the levels at ~1.7 MeV which have small stripping strengths, but may have substantial (2p-1h) components.

Additional evidence for doorway states has been found in the observation of intermediate structure by Monahan and Elwyn {ME68} at $\varepsilon_d = 360$ and 700 keV in differential elastic scattering and polarisation data. The width of the 360 keV state was found to be $\Gamma\approx 140$ keV, with escape width $\Gamma^\gamma \approx 46$ keV and damping width $\Gamma^\gamma \approx 95$ keV. The resonances comprising the doorway in the range $|E_\lambda - \varepsilon_d| < \Gamma$ satisfied the requirement that $\Gamma^\gamma = \text{constant}$.

An important test of the doorway state hypothesis is to observe
intermediate structure in two reaction channels. For example, in the MeV range, Tomito \{To73\} has observed similar structures in the elastic and inelastic scattering channels. However, below 860 keV only the \(\gamma\)-ray channel is available and Baglan et al. \{BBB71\} in the \(^{57}\text{Fe}(\gamma,n)\) reaction, found some evidence for an envelope of \(\Gamma_{\gamma_0}\) strength centred at 250 keV. The rank correlation coefficient \(p(\Gamma_{\gamma_0}^{\gamma},\Gamma_{\gamma_0}^{n}) = 0.2\) observed for the s-wave resonances is consistent with that observed in capture. Jackson and Strait in the same reaction also found evidence of s-wave intermediate structure, but centred at 200 keV. The postulated s-wave doorway state at 360 keV is therefore not evident in the ground state \(\gamma\)-ray channel.

Jackson and Strait also found evidence for doorway states for \(p_{3/2}^3\) resonances at 235 keV, and \(p_{1/2}^1\) resonances at 629 keV, which would account for a large part of the M1 single particle strength. These data were explained in terms of \(f_{5/2}^1f_{7/2}^{-1}\) particle-hole pairs coupled to the \(^{57}\text{Fe}\) ground state which would decay by enhanced M1 spin flip transitions.

It was shown in section 6.2 that the \(p_{3/2}^3\) assignments may be suspect and that \(d_{3/2}^{-}\) assignments are possible. If this is the case, the small localisation of El strength at 230 keV can be readily attributed to statistical fluctuations. Gamma ray measurements with a Ge(Li) detector and <5 keV neutron energy resolution are needed to resolve this question. The observation of transitions to both \(1/2^-\) and \(5/2^-\) final states will indicate a \(3/2\) spin assignment.

In the region of 640 keV, \(g_{\gamma_0}^\gamma\) and \(g_{\Gamma_{\gamma_0}}\) values remain anomalously large. A \(d_{3/2}^{-}\) assignment at 620 keV and \(d_{5/2}^{-}\) assignments at 644 and 655 keV are indicated as the last two resonances are not seen in the \((\gamma,n)\) measurement. These resonances may be part of d-wave doorway states with small damping widths.

We will utilise the doorway concept to express the radiative
Fig. 6.8

$^{56}\text{Fe}$ average capture cross section

Fig. 6.9

p-h excitations in $^{56}\text{Fe}$
capture amplitude of a resonance in terms of its components in the manner of Beer [Be71]. For resonance $\lambda$ and final state $\mu$ we have

$$\Gamma_{\lambda \mu}^h = (\Gamma_{\lambda \mu}^V)^h + \sum_{Da} (\Gamma_{\lambda \mu}^{Da})^h + \sum_{Dr} (\Gamma_{\lambda \mu}^{Dr})^h + (\Gamma_{\lambda \mu}^S)^h$$

(6.3)

where $\Gamma_{\lambda \mu}^V$ is the valence width in the entrance channel; $\Gamma_{\lambda \mu}^{Da}$ is the radiative width of a doorway component $Da$ which decays by p-h annihilation; $\Gamma_{\lambda \mu}^{Dr}$ corresponds to a doorway state which retains the p-h configuration on decay; $\Gamma_{\lambda \mu}^S$ is the radiative width of all other mp-nh configurations which represent the statistical interaction.

The valence width can be calculated quantitatively using the optical model. However, at this time even qualitative estimates of the $\gamma$-ray strength of neutron doorway states are not possible.

Nevertheless, it is informative to consider possible doorway configurations which could contribute to the radiative widths of s-, p- and d-wave resonances. Referring to Figure 6.9, possible (2p-1h) configurations which can decay by El annihilation are, for example, 

(2p$_{1/2}$ or 2p$_{3/2}$; 2s$_{1/2}$)$^{-1}$, (2p$_{1/2}$ or 2p$_{3/2}$; 2d$_{3/2}$)$^{-1}$ proton or neutron pairs coupled to the 2p$_{1/2}$, 2p$_{3/2}$ or 1f$_{5/2}$ neutron orbit. The energies of the (p-h) components are 8-10 MeV, and reduce by ~2 MeV when two neutrons are paired. These energies are comparable to the observed $\gamma$-ray energies, but exact agreement is not necessary because of the spreading of configuration energies by component particles. However, when the annihilation energy is comparable to the observed $\gamma$-ray energies (i.e. for paired neutron 2p-1h configurations) the corresponding doorway states are expected to be strongest.

For M1 transitions after p-wave capture, the configuration 

$$(1f_{5/2}^{-1}, 1f_{7/2}^{-1})2p_{1/2}, 2p_{3/2}$$

will lead to enhanced transitions to the strong single particle states. The (p-h) energy is 6.6 MeV and is close to the observed $\gamma$-ray energies.
In addition, there are radiative contributions from doorways (Dr) which retain the (p-h) configuration and decay to low lying states with similar (p-h) configurations (e.g. the 1.7 MeV states referred to by Ikegami and Emergy {IE64}).

The neutron doorway state at 360 keV may carry no intrinsic El strength since Kirouac {Ki75} has proposed a particle vibrator coupling model. On the other hand, since non-statistical \( \gamma \)-ray spectra extend to 1 MeV and intermediate structure is observed in the average capture cross section, we postulate that broad El/Ml doorways may exist which are only weakly coupled to the entrance channel.

The capture mechanism in \( {}^{56}\text{Fe} \) is summarised as follows:

(a) The \( \gamma \)-ray spectra for s-, p- and d-wave resonances indicate non-statistical processes with a strong preference for \( p_{1/2} \) final states with strong single particle character.

(b) As the valence model accounts for only \( \sim 20\% \) of the \( \langle \Gamma^\gamma \rangle(s) \) and \( \rho(\Gamma_n^0,\Gamma^\gamma(s)) \sim 0.5 \), an additional process, uncorrelated with resonance neutron widths, is required which contributes, on average, \( \sim 40\% \) of the average s-wave radiative width.

(c) Possible 2p-1h doorway states are available across the \( 2p_{1/2} \), \( 2p_{3/2} \), \( 1f_{5/2} \) and \( 2s_{1/2} \) shells with intrinsic El strengths which may account for the excess \( \gamma \)-ray strength. For Ml transitions, \( 1f_{5/2} \), \( 1f_{7/2}^{-1} \) configurations may contribute.

(d) A second class of doorways which retain the p-h configuration may select final states with small spectroscopic factors near 1.7 MeV and reduce the final state correlations. Interference between valence and annihilation doorway components will also reduce these correlations.

6.5 VALENCE CAPTURE IN \( {}^{56}\text{Fe} \)

Thin and thick sample capture cross section measurements were made at ORELA to minimise self-shielding and multiple scattering corrections.
Fig. 6.10
$^{54}\text{Fe}$ resonance capture cross section

Fig. 6.11
$^{54}\text{Fe}$ average capture cross section

Fig. 6.12
Energy dependence of $\frac{I^V}{I^\gamma}$ in $^{54}\text{Fe}$. Hatched histograms calculated for zero energy dependence in valence widths
at low energies, and to obtain accurate statistics for energies up to 500 keV. Target thicknesses and $^{56}$Fe isotopic abundances were 0.0020 atom b$^{-1}$, 98.80% and 0.0197 atom b$^{-1}$, 97.59%, respectively.

In the $^{56}$Fe data, statistical errors are normally <5% and for resonances with large neutron widths, uncertainties in the linear background estimate and corrections for overlapping resonances dominate. Representative samples of the capture data are given in Figure 6.10.

The prompt background correction accounts for the sensitivity of the capture detectors to resonance scattered neutrons, and is shown as the broken line. This correction is seen to be quite large and for this measurement is the limiting factor in the accuracy of the data.

Asymmetric resonances were observed at 130, 148 and 174 keV after subtraction of the multiple scattering component (38%, 17%, 7% respectively, of the primary yields). For the first two resonances this asymmetry is unlikely to result from the prompt, resonance scattered background, and could be indicative of interference effects. However, if the neutron widths are in error, then oversubtraction of the multiple scattering yield can also result in asymmetry.

Excellent agreement in energy (to better than 0.1%) is found between the capture and total cross section data [Pa+75], but many more resonances are detected in capture. Neutron widths measured by Pandey et al. [Pa+75] have generally been used in the capture analysis.

Below 100 keV, neutron widths and spins were adjusted to yield comparable $\frac{g^* \Gamma_n}{\Gamma}$ values for both thin and thick sample measurements. The resonance parameter set is therefore consistent, but not necessarily unique. Detailed results and average cross sections are given in [Al+77a], and average resonance parameters are summarised in Table 6.6. These parameters satisfactorily fit the measured average cross section, as shown in Figure 6.11.

The valence radiative widths for El transitions from s-wave resonances
TABLE 6.6
AVERAGE RESONANCE PARAMETERS IN $^{54}$Fe

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<tr>
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<td>N</td>
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<td>22</td>
<td>133</td>
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$^a$ Corrected for missed resonances $\{\text{Pa+75}\}$

$^b$ Pa+75
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<thead>
<tr>
<th>E (keV)</th>
<th>$\Gamma_\lambda^0$ (eV)</th>
<th>$Q_\lambda$</th>
<th>$\Gamma_\lambda^V$ (eV)</th>
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<td>147.8</td>
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<td>0.122</td>
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<td>2.3 ±0.5</td>
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<td>174.0</td>
<td>8.82</td>
<td>0.119</td>
<td>1.1</td>
<td>3.5 ±1.1</td>
</tr>
<tr>
<td>192.2</td>
<td>91.2</td>
<td>0.111</td>
<td>10.1</td>
<td>10.0 ±4.0</td>
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<tr>
<td>223.5</td>
<td>1.54</td>
<td>0.103</td>
<td>0.16</td>
<td>1.5 ±0.3</td>
</tr>
<tr>
<td>246.8</td>
<td>39.7</td>
<td>0.095</td>
<td>3.8</td>
<td>5.7 ±1.4</td>
</tr>
<tr>
<td>258.0</td>
<td>not observed</td>
<td></td>
<td></td>
<td></td>
</tr>
<tr>
<td>291.1</td>
<td>2.04</td>
<td>0.084</td>
<td>0.17</td>
<td>1.1 ±0.2</td>
</tr>
<tr>
<td>308.1</td>
<td>9.73</td>
<td>0.080</td>
<td>0.78</td>
<td>2.7 ±0.8</td>
</tr>
<tr>
<td>326.3</td>
<td>35.0</td>
<td>0.075</td>
<td>2.6</td>
<td>6.0 ±3.0</td>
</tr>
<tr>
<td>332.4</td>
<td>41.6</td>
<td>0.073</td>
<td>3.0</td>
<td>3.6 ±1.8</td>
</tr>
<tr>
<td>371.0</td>
<td>15.4</td>
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<td>1.0</td>
<td>2.1 ±0.3</td>
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<tr>
<td>414.1</td>
<td>41.0</td>
<td>0.052</td>
<td>2.1</td>
<td>3.7 ±1.5</td>
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<tr>
<td>426.0</td>
<td>12.9</td>
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<td>0.70</td>
<td>2.6 ±1.0</td>
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<td>3.6 ±1.4</td>
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<tr>
<td>487.5</td>
<td>22.1</td>
<td>0.043</td>
<td>0.95</td>
<td>3.2 ±1.3</td>
</tr>
</tbody>
</table>

$N=11$ $\rho_\Gamma(<300) = 0.97$ $<\Gamma_\gamma > = 3.2±2.5$

$N=15$ $\rho_\Gamma(<400) = 0.94$ $<\Gamma_\gamma > = 3.3±2.3$

$N=20$ $\rho_\Gamma(<500) = 0.90$ $<\Gamma_\gamma > = 3.2±2.0$
to the $\frac{1}{2}^+$, $\frac{3}{2}^+$ final states are given in Table 6.7, together with the relevant s-wave data. The centroid energies and the spectroscopic factors $s^\lambda$, are taken from [KH72]. Since $q^\lambda$ is strongly dependent on the energies of the s-wave resonances, the valence width was evaluated at each resonance.

The calculated valence widths are compared with the measured radiative widths in Figure 6.12 as a function of energy, and account for large fractions of the radiative widths of most resonances with large neutron widths. The data are also consistent with the predicted energy dependence of the valence component. The energy dependence of $\Gamma^v/\Gamma_\gamma$ is shown with and without (hatched histogram) the dependence predicted by the optical model calculations (section 1.3). This is the first experimental confirmation of this effect. Without this dependence, the valence widths of resonances at 414 and 433 keV would be 8-10 eV, and several times larger than the measured values. The correlation between the calculated and observed widths is $\rho_v(\Gamma^v_\gamma, \Gamma_\gamma) = 0.93^{+0.07}_{-0.23}$, where the error is the standard deviation of a zero correlation distribution calculated for the appropriate sample size. This result is similar to the initial state correlation of $\rho^o_\gamma(\Gamma^o_\gamma, \Gamma_\gamma) = 0.94$. The large correlation is a far more significant result than the calculated valence magnitudes which are subject to uncertainties of at least 30%. It shows that the variation of the total radiative width is completely consistent with that predicted by the valence model. Note that the value of $\rho$ is not heavily dependent on the 192 keV resonance. Exclusion of this resonance reduces the correlation to $\rho = 0.81$.

Further confirmation of the dominant role of the valence process is found in the thermal capture $\gamma$-ray spectrum. The thermal capture cross section of 2.3 barn is much larger than the estimated direct capture cross section, and is attributed to the 7.6 keV resonance [Mu74].
### TABLE 6.8

**FINAL STATE $l_n = 1$ SPECTROSCOPIC DATA IN $^{54}$Fe**

<table>
<thead>
<tr>
<th>$E_\mu$ (MeV)</th>
<th>$J^{\pi a)}_{\mu}$</th>
<th>$10^4 \times q_\mu$</th>
<th>$\theta_{2a)}^{\pi}_{\mu}$</th>
<th>$\Gamma_{\gamma\mu}^{V}$ (meV)</th>
<th>$\Gamma_{\gamma\mu}^{b)}$ (meV)</th>
</tr>
</thead>
<tbody>
<tr>
<td>0.00</td>
<td>$^{3/2-}$</td>
<td>7.82</td>
<td>0.73</td>
<td>1260</td>
<td>1190</td>
</tr>
<tr>
<td>0.41</td>
<td>$^{1/2-}$</td>
<td>4.18</td>
<td>0.59</td>
<td>476</td>
<td>220</td>
</tr>
<tr>
<td>1.93</td>
<td>$^{1/2-}$</td>
<td>4.18</td>
<td>0.07</td>
<td>33</td>
<td>2</td>
</tr>
<tr>
<td>2.06</td>
<td>$^{3/2-}$</td>
<td>7.82</td>
<td>0.08</td>
<td>65</td>
<td>36</td>
</tr>
<tr>
<td>2.48</td>
<td>$^{3/2-}$</td>
<td>7.82</td>
<td>0.15</td>
<td>101</td>
<td>34</td>
</tr>
<tr>
<td>3.04</td>
<td>$^{3/2-}$</td>
<td>7.82</td>
<td>0.03</td>
<td>17</td>
<td>57</td>
</tr>
<tr>
<td>3.56</td>
<td>$^{3/2-}$</td>
<td>7.82</td>
<td>0.11</td>
<td>44</td>
<td>41</td>
</tr>
<tr>
<td>3.80</td>
<td>$^{1/2-}$</td>
<td>4.18</td>
<td>0.50</td>
<td>90</td>
<td>43</td>
</tr>
<tr>
<td>5.78</td>
<td>$^{1/2-}$</td>
<td>4.18</td>
<td>0.04</td>
<td>2</td>
<td>18</td>
</tr>
</tbody>
</table>

*a) Taken from KH72*

*b) Thermal partial widths derived from Ar67 assuming $\Gamma_\gamma = 1.8$ eV for the 7.7 keV resonance*
Our value for the radiative width of this resonance is fully consistent with this interpretation, and calculated partial valence widths are in excellent agreement with the measured values (Table 6.8).

The $l_n = 1$ final state correlation is $\rho_F (\Gamma^V_{\gamma\mu}, \Gamma^\mu_{\gamma\mu}) = 0.94^{+0.06}_{-0.39}$, reflecting the dominance of valence transitions to the $3/2^-$ ground state and $1/2^-$ first excited state. As the total valence width equals the measured radiative width within the errors, statistical and other mechanisms must be quite small for this resonance.

Average $\gamma$-ray spectra {Bi+73} are available at 30 keV, but are not representative of s-wave capture. Nevertheless, these spectra also show strong transitions to the single particle final states, including the $5/2^-$ second excited state, confirming that $l > 0$ capture occurs.

The valence model is seen to account for the final state correlation at 7.7 keV, the initial state correlation for 20 s-wave resonances, and over half of the average value of the total radiative widths. This result is outstanding for the $A = 40-70$ mass region where initial state correlations of less than 0.5 have been frequently observed, together with valence magnitudes which are often much less than the observed widths.

In $^{56}$Fe evidence was found for a substantial uncorrelated component in the radiative widths. It is probable that this component is still present in $^{54}$Fe, but its effect is diminished because of the valence component which is six times larger in $^{54}$Fe than $^{56}$Fe. The increased valence strength in $^{54}$Fe arises from the larger binding energies of the low lying p-wave single particle states, and the greater s-wave strength function.

6.6 RESONANT AND BACKGROUND INTERFERENCE IN $^{54}$Fe

The initial state correlation coefficient observed in $^{54}$Fe is related to the distant resonance component of the non-resonant capture cross section by equation 1.7. If the s-wave resonances are dominated by a common doorway, then the real and imaginary parts of the R-matrix
may be comparable, and the distant resonant background capture cross section can be written as:

\[
\sigma_{\gamma\mu}^{\text{BG}}(r) \approx \frac{\pi^2}{4} \cdot \frac{2.608}{\sqrt{E_n}} \cdot 10^6 \cdot g \cdot S_0 \cdot S_{\gamma\mu} \cdot \rho_{I\mu} (\Gamma_0^{\lambda_n, \Gamma_\mu}) \cdot \ldots (6.4)
\]

where \( S_0 \), \( S_{\gamma} \) are the neutron and radiative strength functions, \( E_n \) is in eV and \( \mu \) designates the final state. For the ground state (\( \mu = 0 \)), \( S_{\gamma 0} = 0.62 \times 10^{-4} \), \( S_0 = 8.6 \times 10^{-4} \) and \( \rho_{I0} \sim \rho_I = 0.94 \), since 75% of the valence strength is found in the ground state transition. Then \( \sigma_{\gamma 0}^{\text{BG}} = \frac{0.32}{\sqrt{E_n}} = 2 \text{ barn} \) at thermal, a value comparable to the measured thermal capture cross section of 2.25 barn. Since 1.74 barn is accounted for by the tail of the 7.7 keV resonance (\( \Gamma = 1.9 \text{ eV} \)), this background estimate is at best a factor of four too large.

In the strong coupling model, the non-resonant cross section arises from the scattering of the incident neutron by the nuclear potential into a bound, single particle, final state. This potential capture cross section is calculated to be 0.3 barn at thermal energy for \( A = 54 \) \{LL60\}.

The mass dependence of the direct capture cross section in the 3s region has been obtained from calculations using R-matrix theory with intermediate coupling \{LL60\} and, more recently, using optical model and shell model formulations of the valence model \{CM75\}. These calculations are particularly sensitive to the potentials used, and the calculated thermal background capture cross section for \(^{54}\text{Fe}\) varies from 0.1 to 1.0 barn. Accordingly, a significant background cross section is expected in \(^{54}\text{Fe}\) which will interfere with the partial resonance capture cross sections.

Interference between non-resonant and resonant capture can be observed in the total capture cross section only if one or two \( \gamma \)-ray channels are involved or if valence capture dominates. In the latter case, the signs of the \( \gamma \)-ray transition amplitudes to single particle,
Fig. 6.13
Asymmetric resonances in $^{54}\text{Fe}$

Fig. 6.14
7.7 keV resonance in $^{54}\text{Fe}$. Dashed curves are expected line shapes for thermal background of 0.3 barn
final states are correlated with that of the neutron amplitude, and the valence partial capture cross sections to these states therefore add coherently. Both conditions are satisfied in $^{54}$Fe since the ground state transition takes 75% of the valence strength, which is itself a large fraction of the observed radiative widths. Interference effects are therefore expected in the total capture cross section.

Asymmetric resonances are in fact observed at 130, 148 and 174 keV in Figure 6.10, and these are re-analysed in detail in this section. Careful attention has been given to asymmetries resulting from prompt resonance scattered neutrons, multiple scattering and resonance-resonance interference. The prompt background is treated in a Monte Carlo simulation of the neutron history in the target, detector and environs. It is derived from the product of the multilevel scattering cross section at the incident neutron energy and the neutron sensitivity at the scattered energy, this being determined from filtered beam measurements on $^{208}$Pb. The capture yield from multiply scattered events in the target is calculated and subtracted from the observed $\gamma$-ray yield, which is then fitted with the sum of the primary capture and prompt background yields.

Since valence capture is expected to allow the observation of resonance-direct interference in the total capture cross section, it will also result in interference between s-wave resonances. It is therefore necessary to calculate the valence $\gamma$-ray cross section using a multilevel formalism. Results for resonances at 7.72, 129.6, 146.8, 174.4 and -184 keV are shown in Figures 6.13 and 6.14. The calculated curves were obtained for $\sigma_y (BG) = 0$. The multiple scattered component has been subtracted (lower histogram) and the prompt background component (broken line) is added to the primary $\gamma$-ray yield. All resonance cross sections are well reproduced, although the energy of the 184 keV resonance is -8 keV below the total cross section value.

The asymmetries at 129.6, 146.8 and 174.4 keV are accounted for
by interference with the broad \((\Gamma_n = 40\ \text{keV})\) resonance at 184 keV. This resonance exhibits an asymmetry opposite to that of the lower energy resonances, a result which is inconsistent with a resonance-direct interference hypothesis. The estimated valence width of 10 eV accounts for ~90% of the total radiative width of this resonance, and supports the validity of the single capture channel, multilevel calculation.

After subtraction of a 25% multiple scattering contribution, the 7.7 keV resonance is found to be quite symmetric (Figure 6.14). (The peak at 6 keV results from the capture of scattered neutrons by the 5.9 keV resonance in the aluminium beam tube and detector assembly.) The expected line shapes for a thermal background cross section of 0.3 barn (for positive and negative amplitudes) are shown for comparison. Since this resonance is estimated to be completely valence in character, an upper limit of \(<0.15\) barn can be set for the thermal background cross section.

The large initial state correlation is therefore not associated with the expected background capture cross section. The upper limit is a factor of four lower than that reported for the ground state transition in the \(^{53}\text{Cr}(\gamma,n)\) reaction \((\text{Ja74})\), whereas a factor of one to ~1.5 would be expected from theory. Optical model calculations, using the Moldauer potential \((\text{CM75})\), are in good agreement with this result, but others which predict a peak in the background cross section near A = 54 are not so favoured.

The Moldauer potential calculation is also consistent with the reported background cross section in \(^{63}\text{Cu}\) \((\text{Au68})\). However, the near zero cross section at A = 60 is insensitive to the choice of potential and is inconsistent with results for \(^{59}\text{Co}(n,\gamma)\) \((\text{Au68,Wa+66})\).

As a result of the strong valence effect in \(^{54}\text{Fe}\) it has been possible for the first time to observe resonance-resonance interference
in the total capture cross section, and to set a much lower limit on the magnitude of the background capture cross section.
CHAPTER 7

INTRINSIC EL STRENGTH OF NEUTRON AND γ-RAY DOORWAY STATES

Strong correlations are associated with the existence of isolated doorway states which must be common to both the neutron and γ-ray channels, and these may be observable through the occurrence of intermediate structure in the reaction cross section.

There are reported instances of intermediate structure in total cross section data which have been interpreted as doorway states in the neutron channel (although few authors have tested their data for significance by the Wald-Wolfowitz method (Ja71)). Since the widths of these isolated neutron doorways are expected to range from 50-200 keV, it is of interest to seek evidence for them in the capture channel. Ideally, a partial capture channel is most appropriate in order to maximise the effect (i.e. through Ge(Li) spectral measurements or the (γ,n) reaction). Total radiative widths can also be studied, although there remains the possibility of obscuring the effect if intermediate structure occurs in different partial capture channels at different energies.

The El strength of a neutron doorway state arises from valence effects as well as intrinsic components which could result from:

(a) particle transitions coupled to collective core excitations;
(b) collective transitions of the core;
(c) annihilation of 2p-1h configurations.

Before reviewing the experimental and theoretical evidence for these components, it is useful to consider a schematic model for the role of excited target states in the neutron capture reaction. Resonance and final states can be described as single neutron states coupled to the target ground state (φ₀) as well as quadrupole and octopole excitations (φ₁). These collective states are just coherent superpositions of 2p-1h quasiparticle states. If certain p-h configurations are dominant (φₚₘ),
Fig. 7.1
Schematic representation of excited target states in the 3s region

Fig. 7.2
Schematic representation of 2p-1h excitations in the s and p strength function regions
they may occur with unperturbed energies much greater than those of the collective states. These effects are shown in Figure 7.1, for the case of s-wave capture by an even-even target, with γ-ray decay of the \((A+1)\) nucleus to the low lying \(p_{3/2}^+\) final states, such as is observed in the 3s region. The hatched areas indicate the fragmentation of these states by the residual interactions in the \((A+1)\) system.

Strong, single particle El transitions can only occur between states with the same core configuration. Since \((d,p)\) experiments measure the strength of the valence neutron coupled to the target ground state, large final state correlations between the γ-ray intensities and spectroscopic factors are expected only when excited target state amplitudes in the capture resonance are small. Lane \(\{La74a\}\) has shown that near threshold the contribution of excited target state components is greatly reduced in the region of the strength function size resonance.

Nevertheless, there are certain states \(|\psi_{ij}\rangle\) with energies quite close to the threshold energy. In the example shown in Figure 7.1 for the 3s region and using the nomenclature of equation 7.1, the state \(|12^+,d_{5/2}^+;1/2^+\rangle\) occurs at an energy similar to the entrance channel state \(|00^+,s_{1/2}^+;1/2^+\rangle\), and would reduce the final state correlation as a result of transitions to the \(|12^+,p_{3/2}^+;1/2^-,3/2^-\rangle\) components of the low lying states.

7.1 PARTICLE-VIBRATOR DOORWAYS

The capture state wave function can be expanded in terms of the basis functions of the unified model \(\{KR71\}\), i.e. coupled particle-vibrator (p-v) states. For s-wave capture by an even-even target

\[
\psi_{JM} = \sum_{N,R,j} A_{NR}^j |NR,j;JM\rangle,
\]

where \(N\) denotes the number of phonons, \(R\) is the spin of the vibrational state; \(j\) the spin of the valence nucleon; \(J,M\) are the resultant moment of the level and its projection; \(A_{NR}^j\) are the amplitude coefficients
for the entrance channel $|00,s_{1/2};1/2^+\rangle$; the single phonon doorways $|12,d_{5/2};1/2^+\rangle$ and $|12,d_{3/2};1/2^+\rangle$; and the two phonon doorways $|20,s_{1/2};1/2^+\rangle$, $|22,d_{5/2};1/2^+\rangle$, $|22,d_{3/2};1/2^+\rangle$.

The matrix element for the radiative decay of the resonance is therefore a linear combination of the expansion coefficients:

$$M_{\lambda \rightarrow \mu} (\text{El}) = \alpha_{\lambda} A_{\lambda 1} + \beta_{\lambda} A_{\lambda 2} + \gamma_{\lambda} A_{\lambda 2}^*,$$

where $\alpha_{\lambda}$, $\beta_{\lambda}$, $\gamma_{\lambda}$ are expressed in terms of the calculated coefficients $A_{j}^{NR}$.

Particle-vibrator states (single particle coupled to a quadrupole excitation for s-wave and an octopole core excitation for p-wave capture) appear to account for the observed neutron doorway states in the 2p

$\{\text{Ha}+76,77\}$, 3s $\{\text{Ca}+75,\text{Ki}76\}$, 3p $\{\text{Ma}+75\}$ and 4s $\{\text{Fa}+65,\text{BD70, LB76}\}$ regions, and can cause interference with the valence component through the neutron amplitudes. When this occurs, Halderson $\{\text{Ha}+76\}$ has shown for $^{28}\text{Si}$ that simple valence calculations are in error and detailed initial and final state wave functions should be used. Additional calculations for $^{28}\text{Si}$ $\{\text{HCD76}\}$ have been made employing the Feshbach and Boridy-Mahaux formalisms $\{\text{BM75,CM75}\}$. Excellent agreement was achieved for the integrated strength of the ground state $\gamma$-ray over the proposed $3/2^-$ p-v doorways. Predicted neutron widths were also in agreement with experiment.

Neutron doorway states in $^{56}\text{Fe}$ have also been reported at 360 kev $\{\text{ME68}\}$ and higher energies $\{\text{To73}\}$. Kirouac $\{\text{Ki76}\}$ has interpreted these using a p-v model where a $d_{5/2}$ neutron is coupled to the first excited, single phonon state of the target ($2^+, 0.847$ meV). This model adequately describes the observed doorway states in the neutron channel.

Noting that the $^{57}\text{Fe}$ ground state has the probable configuration $a|00^+,p_{1/2};1/2^\sim\rangle + b|12^+,p_{3/2};1/2^\sim\rangle$, Kirouac calculated the ground state radiative widths ($\Gamma_{\gamma 0}$) assuming interference between the valence and doorway state amplitudes. While improved agreement over valence estimates was obtained for both $\Gamma_{\gamma 0}$ and $\Gamma_{\gamma}$ in the region of the 360 kev doorway,
consistent agreement over a 400 keV energy range could not be achieved. Kirouac concluded that the p-v doorway does not play a major role in the radiative decay of s-wave resonances.

Lev and Beres (LB76) have made calculations in $^{206}$Pb of the intrinsic radiative strength of p-v doorways. The most important of these is the postulated $|4^+,2g_9/2;1/2^+>$ doorway at ~0.5 MeV, but the calculated value of the capture cross section for this state was found to be only ~1/1000 of the observed cross section.

Additional evidence for phonon-particle doorways can be gained from an analysis of thermal capture γ-ray spectra. Calculations by Knat'ko and Rudak (KR71) for $^{54}$Fe, $^{138}$Ba and $^{140}$Ce indicate that thermal decay proceeds primarily through the initial state $|00,s_{1/2};1/2^+>$. This is not surprising since direct capture is expected for the last two nuclides. For $^{50}$Ti, single phonon doorways make the predominant contribution, while two phonon doorways contribute to both $^{52}$Cr and $^{142}$Nd.

L. V. Rudak (Ru76) has investigated the decay mechanism of the 1.167 keV $P_{1/2}$ resonance in $^{56}$Fe(n,γ). The experimental intensities of the M1 transitions are accounted for by valence neutron $|00^+,p_{1/2};1/2^->$ and single phonon doorways $|12^+,p_{3/2};1/2^->$, $|12^+,f_{5/2};1/2^->$, if comparable valence and single phonon configurations are present. Thus some γ-ray spectra are found to show the influence of p-v components, but if the magnitude of these effects are comparable with the valence widths, they are too small to account for the measured partial γ-ray widths.

An important property of transitions in the presence of excited phonon states of the core is that (d,p) and p-v γ-ray strengths will tend to be anticorrelated. For example, in $^{56}$Fe, the decay of the p-v configuration is given by

$$|12^+,d_{5/2};1/2^+> \rightarrow |12^+,p_{3/2};1/2^->$$

while for valence capture
As the final state may be given by

\[ a | 12^+, s_{1/2}^{1/2} > + b | 00^+, p_{1/2}^{1/2} > \]

the valence and (d,p) strengths will be correlated through the \( p_{1/2} \) ground state component, but the p-v strength will be correlated with the \( p_{3/2} \) excited state component.

The p-v component is therefore a 'retention doorway' as described in section 6.4, and in view of the extensive final state correlations found in the 3s and 3p regions, probably plays a minor role. However, in the deformed 4s region, evidence is accumulating \{KR71,CF76,R076\} which indicates a more dominant role for p-v interactions, as predicted by Soloviev \{So71\}.

7.2 COLLECTIVE CORE TRANSITIONS

Martsynkevich and Rudak \{MR76\} have investigated the role of El collective \( \gamma \)-ray transitions at the N=28 and N=82 closed neutron shells.

The capture state wave functions were expanded into the basis states of the unified model. Valence El transitions in the core field and collective El transitions of the core, with no change in the single particle state of the valence neutron, were then calculated.

A comparison with the thermal capture \( \gamma \)-ray data revealed that single particle \( \gamma \)-ray transitions make the major contribution to the capture reaction. In a number of cases, inclusion of collective \( \gamma \)-ray transitions improved the agreement with experiment, but the contribution was only \( -5-10\% \). Presumably such collective El transitions will also be largely depleted by the gathering of strength in the El giant dipole resonance.

7.3 2p-1h STATES

Evidence for p-h core excitations in neutron capture \( \gamma \)-ray spectra has been found in the 3s \{Al+76,a,b\}, 3p \{Ma+75,Ri+69,Ra+76\} and
A > 170 {Ba+67} regions, and 2p-lh doorway state calculations in $^{48}$Ca, $^{88}$Sr and $^{90}$Zr have been found to be consistent with the experimental scattering data on these nuclides {DBN72}. A schematic model is presented here for s-, p- and d-wave interactions of this type. The binding energies of single particle states appear as bands of like-parity states as a function of mass number (Figure 1.2). Since the energy gap between alternate parity bands (~5-7 Mev) often corresponds to the energies of intense El γ-ray transitions observed in capture spectra, the annihilation of p-h pairs across the opposite parity shells, as presented schematically in Figure 7.2, could provide the underlying basis for the 2p-lh state model. It is of interest to determine the energy range over which these doorway states are effective, and their radiative strengths.

In the 3s region, $^{56}$Fe is the most studied nuclide {Al+76b, Bi+73, Ke71, Al+74, Ba+77} and is considered typical of adjacent even-Z nuclides. s-, p- and d-wave resonances have similar spectra which change only slowly in overall shape over a neutron energy range from thermal to 1 MeV. Intense γ-rays are observed to low lying $p_{1/2}$, $p_{3/2}$ states which are similar in energy to the unperturbed energies of, for example, $2p(2s)^{-1}1f(1d)^{-1}1f_{5/2}(1f_7)^{-1}$ particle-hole pairs (see section 6.4).

With the exception of $^{54}$Fe, where the valence component reaches its maximum value, the even-Z nuclides all exhibit moderate correlations and low valence fractions of the radiative widths, suggesting that 2p-lh strengths vary slowly with mass number and hence energy. A similar conclusion is obtained from the γ-ray spectra in $^{56}$Fe over 1 MeV {Al+74} (section 6.1). The overall consistency of these data suggests that the El doorway states are broad (>1 MeV) with average resonance El strengths comparable with the average valence component (see Table 5.2).

Soloviev and Voronov {SV74,76} have expanded the capture state wave functions in the Zr, Mo and Sn isotopes into one and three quasiparticle
states (i.e. valence and 2p-1h states) as well as quasiparticle plphonon components. The energies of those states were calculated for which El and Ml transitions can proceed to the single particle components of the low lying states. While the effect of these p-h states depends on their spreading widths which were not calculated, it is apparent that the γ-ray spectra can be explained by the presence, or absence, of these states.

An important property of the p-h interaction is that the (d,p), p-h and valence strengths will be correlated because the single particle component in the final states is common to all (section 1.4). For example, the El decay of p-h configurations in the 3s and 4s regions may be given by

\[ |p_{3/2}^{-1},s_{1/2}^{-1},p_{1/2}^{-1/2}^{+} \rangle \rightarrow |p_{1/2}^{1/2}^{-} \rangle \]

which compares with the valence transition.

\[ |s_{1/2}^{1/2}^{+} \rangle \rightarrow |p_{1/2}^{1/2}^{-} \rangle \]

In a real nucleus, the 2p-1h states would of course be mixed with other states with the same J^π near the same energy. However, it is required to show that significant El strength can be retained in the threshold region and not taken up into the GDR.

Lane {La71} has reviewed the results of calculations in 208Pb by Pal, Soper and Stamp {PSS64}, and in Ni by Soper {So70}. Using zero range forces, these authors diagonalised the Hamiltonian using a basis of 35 l^- p-h states which carry all the El strength for 208Pb. It was found that certain of the p-h states decoupled from the giant resonance motion - specifically the 3p^-14s, 3p^-13d, and the 2f^-13d states.

These states interacted to form a 'pygmy' dipole resonance near 5.5 MeV which carried nearly all the dipole strength in the threshold region (~7% of the total dipole strength). Significantly, the 'pygmy' state was also found to carry considerable neutron strength and was one
of only two states in this energy region to carry significant common
strength for the two channels. No obvious feature of the three states
decoupled in this calculation suggested the mechanism or underlying
principle behind this decoupling, so the extension of these ideas to
lower neutron magic numbers was difficult, particularly since the lp-lh
states there are quite different.

The pygmy resonance also appeared to give a natural explanation for
the observation of anomalous capture γ-ray spectra above Ta. However, a
number of doubts concerned with this calculation were pointed out by Lane.
The calculation predicted much more El strength in the threshold region
than was experimentally detected, and hence could even then be taken as
illustrative only of how strong common doorway effects could arise in
practice.

However, further calculations with more realistic forces did not
support this conclusion. A study of p-h calculations in $^{208}$Pb with finite
range forces by Khanna and Harvey {KH73} showed that the $s_{1/2}p_{1/2}^{-1}$
neutron state is not uncoupled from other p-h configurations as a general
rule. These calculations were made with radial integrals for a Gaussian
interaction with range 1.6 fm, harmonic oscillator functions, realistic
p-h energies and Gillet exchange mixtures.

They indicated the presence of a state at 4.5 MeV and a group of
states between 6.5 and 8 MeV with a few per cent of the dipole strength.
Overall agreement with the $(\gamma,\gamma)$ measurements of Knowles and Khan {KK73}
was qualitative only. The calculations showed that states with large
$s_{1/2}p_{1/2}^{-1}$ amplitudes carry little El strength, and that $s_{1/2}p_{1/2}^{-1}$ and
d$_{3/2}p_{1/2}^{-1}$ fragments were not well localised, but spread over several
MeV and centred at 7.5 MeV. Since the widths of these states were
calculated to be 0.5 to 3 MeV, a number of overlapping doorway states
could be available in the threshold region.

Further calculations {HK74} were made for $^{132}$Sn, $^{90}$Zr and $^{56}$Ni.
In $^{132}\text{Sn}$, about 1% of dipole strength was found at ~6 MeV, and ~7% near 10 MeV. In $^{90}\text{Zr}$ no significant strength was observed below 12 MeV, although 5% of the dipole strength was found at this energy. The minimum energy for dipole strength increased to 11-12 MeV in $^{56}\text{Ni}$, where ~8% is found. Harvey and Khanna concluded that the dipole strength is brought low in the spectrum by the fluctuations in the radial integrals and that regrouping into two regions results from the difference in the mean energies of the neutron and proton p-h states. This dependence on mean energies is reflected in the wave functions as the states that contain dipole strength can be characterised as mainly neutron or proton p-h states. When the mean neutron and proton p-h energies are equal, as in $^{56}\text{Ni}$, no low lying structure is found.

Since the states in $^{208}\text{Pb}$ with significant dipole strength do not contain any p-h components with intensity greater than 15%, they are therefore collective in nature.

If a difference in the mean energies of the neutron and proton p-h pairs is required to retain E1 strength in the threshold region, then the 2p, 3s and, to a lesser extent, the 3p region should show little evidence of doorway state effects.

### 7.4 VALENCE DOORWAYS

If a neutron doorway carries no intrinsic E1 (or M1) strength, the neutron and radiative widths may still be highly correlated as a direct consequence of the valence model.

The reported doorway state in $^{88}\text{Sr}$ \{Bo+76b\} may be such a state, since the calculated valence widths account for most of the non-statistical component, and the initial state correlation is close to unity. This interpretation is consistent with a p-v description for the doorway state. The low lying states of $^{89}\text{Sr}$ are expected to contain insignificant excited target state components (the ground state of $^{89}\text{Sr}$ has $\theta^2 \approx 1.0$), and the $E^3_\gamma$ factor weighs heavily against high energy excited states which
might have a large overlap with the p-v configuration. Consequently, the contribution of excited core components to the total resonance radiative width is expected to be small.

The 800 keV doorway in $^{28}$Si has also been described as a p-v doorway [Ha+76,77a]. In this case it has been shown that destructive interference occurs in the neutron channel, resulting in incorrect estimates of partial valence widths. The inclusion of single particle transitions in the presence of excited target states, specifically the 1.78 MeV $2^+$ state in $^{28}$Si (i.e. inelastic valence capture), is found to increase significantly the overall $\gamma$-ray strength [HCD76]. Consequently, El strength originates from both valence and intrinsic components of this neutron doorway.

A p-v description has been applied [Ki76] to describe the 360 keV doorway in $^{56}$Fe. Since the $^{57}$Fe ground state contains a significant excited target state component, the intrinsic El strength of the doorway for the ground state transitions should be quite large. Structure is observed in the $(\gamma_0,n)$ channel, but this does not coincide with that observed in scattering. The total radiative widths do not exhibit intermediate structure and, as such, the data do not support a significant intrinsic El strength for this doorway. The total and partial width initial state correlations $\rho_{\gamma}^{p}(\Gamma_{\gamma},\Gamma_{\gamma})$ and $\rho_{\gamma}^{n}(\Gamma_{\gamma},\Gamma_{\gamma})$ are relatively small and the non-statistical total radiative width exceeds the calculated valence width. In this instance, neither valence nor p-v doorways are observed since they are presumably overshadowed by several overlapping El doorway states, with dominant particle-hole configurations.

7.5 CAPTURE MECHANISMS IN $^{45}$Sc AND $^{139}$La

The capture mechanism in two odd-Z nuclei, $^{45}$Sc and $^{139}$La, is investigated in this section. These nuclides are both located at or near closed neutron shells and have the common property of low-lying states with $l_n = 3$ neutron angular momenta in the residual nucleus.
Fig. 7.3
Reduced NaI γ-ray spectra for capture of 40 to 430 keV neutrons in Sc and La. Smooth curves are statistical calculations, normalised to equal intensity at 3 MeV.
NaI spectral measurements were made with natural samples \[ \{Al+76a\} \]. and, after background subtraction, the spectra were reduced by least squares analysis with NaI line shapes (see section 2.2). The reduced spectra (Figure 7.3) are compared with statistical calculations using the method of Troubetskoy \[ \{Tr61\} \] after normalisation to equal intensity between 2.7 and 3.4 MeV. Histogram errors are typically \(-20\%\). In each case, the spectra are seen to be anomalous with respect to the statistical model for the high energy transitions at 8.5 MeV in Sc and 5 MeV in La.

The GDR model predicts an \( E^5 \gamma \) energy dependence and the inclusion of this term gave an improved fit to the high energy bump. However, the minima in the Sc spectra at 7.8 MeV are overestimated, indicating that the enhancement is localised to certain final states.

As the calculations are sensitive to the level density parameter, the s-wave level density was calculated at the binding energy, and the level density parameter adjusted to give the observed s-wave spacing in the neutron resonance region.

For \(^{45}\text{Sc} \ (I^\pi = 7/2^-)\), the energy of the high energy bump corresponds to the centroid of \( l_n = 3\), positive parity states in \(^{46}\text{Sc} \). The intensity of the high energy transitions is found to be independent of neutron energy in the range 40 to 430 keV, varying between 10 and 12\% of the measured spectral area.

The anomalous bump is more apparent in \(^{139}\text{La} \ (I^\pi = 7/2^+)\) at a neutron energy of 210 keV and corresponds to the centroid of \( l_n = 3\), negative parity states below 0.6 MeV. While \( l_n = 1\) states occur in the range 0.6 to 1.0 MeV, they receive little \( \gamma \)-ray strength as is evident from the dip in the \( \gamma \)-ray yield at 4.3 MeV.

The statistical model being inadequate, we consider the applicability of the valence model of neutron capture. Its essential feature is that valence neutrons change state in the presence of a spectator core. For Sc and La the change in orbital angular momentum quantum
number is greater than one, and $E1$ valence neutron transitions to $l_n = 3$ states following $s$-wave capture are forbidden by the triangle rule for the addition of angular momenta. However, valence $E1$ transitions after $d$-wave capture can occur to $l_n = 3$ states.

Fortunately in Sc it is possible to directly estimate the significance of $d$-wave capture from averaged Ge(Li) measurements (Bi+73) in the 10 to 60 keV energy range.

These results show that transitions are not observed to the $6^+$ (0.051 Mev) and $7^+$ (0.975 MeV) excited states in $^{46}$Sc (CF73). The only $E1$ transitions which could reach these states must come from $5^-$ or $6^-$ resonances excited after $d$-wave capture. On a $(2J+1)$ basis, these resonances represent about half the total number of $d$-wave resonances and the absence of transitions to the $6^+$, $7^+$ states from more than 50 resonances strongly suggests a negligible $d$-wave component at 40 keV.

At 410 keV, the $d$-wave cross section is calculated to be 50% of the total, yet the relative intensity of the high energy bump remains comparable to that at 40 keV and the expected $d$-wave penetrability energy dependence is not observed.

In $^{139}$La both $s$- and $d$-wave valence transitions can occur to the $l_n = 1$ states. The $\gamma$-ray intensity at the 4.3 MeV minimum therefore provides an upper limit for the possible $d$-wave valence component to the $l_n = 3$ states and only a small fraction of $l_n = 3$ strength can derive from this source. Furthermore, Wasson et al. (WCG69) find strong transitions to the $l_n = 3$ states from the bound and 74.4 eV $s$-wave resonances, indicating the $s$-wave nature of the capture mechanism at low energies. Recent measurements at Lucas Heights with a Ge(Li) detector in the 10-60 keV energy range also show that $s$-wave capture accounts for most of the strength of the anomalous transitions (KA77) since transitions from $d$-wave resonances to the $0^-$ and $7^-$ states at 581 and 284 keV are absent.
The γ-ray spectra cannot be explained by the statistical model alone, and we have shown that s-wave valence capture is not applicable in $^{45}$Sc and $^{139}$La. While d-wave valence capture can occur, it appears to play only a minor role as the strength of the high energy bump in Sc remains approximately constant from 40 to 430 keV.

In the $^{45}$Sc target, the ground state is strongly single particle with a proton in the $f_{7/2}$ shell [YS64]. It is thus reasonable to expect that the low-lying states with $l_n = 3$ in $^{46}$Sc have the character $[\pi(f_{7/2}^{1})\nu(f_{7/2}^{1})]^J$ where $\pi,\nu$ represent protons and neutrons respectively, and $J$ is the spin of the $^{46}$Sc state. Because of the one-body nature of the El operator, strongly single particle final states of this type can be populated by El decay only from initial states of the form $[\pi(f_{7/2}^{1})\nu(f_{7/2}^{1}),\nu], [\pi(f_{7/2}^{1})\nu(f_{7/2}^{5}),\nu-\nu^{-1}]$ and $[\pi(f_{7/2}^{1})\nu(f_{7/2}^{5}),\pi-\pi^{-1}]$ where $\nu-\nu^{-1}, \pi-\pi^{-1}$ represent neutron and proton particle-hole pairs. The first of these three brackets corresponds to valence neutron capture which has already been excluded for s-wave neutrons; the last two correspond to a mechanism involving intermediate 2p-1h states in the capture process. This mechanism has been proposed to explain some aspects of neutron capture in $^{93}$Nb [Ri+69] and $A > 181$ [Ba+67].

Consider possible 2p-1h excitations in Sc and La. These configurations must fulfill requirements for energy, spin and parity conservation and can be derived from the energies of single particle states given in Table 7.1.

In $^{45}$Sc, the $s_{1/2}$ neutron can excite neutron p-h pairs such as $\pi(f_{7/2}^{1})\nu(f_{7/2}^{5},2p_{1/2}^{1},d_{3/2}^{\pi-1})$, with subsequent El annihilation of the $(2p_{1/2}^{1},d_{3/2}^{\pi-1})$ pair to form the
### TABLE 7.1
SINGLE PARTICLE ENERGIES FOR Sc AND La

<table>
<thead>
<tr>
<th>n(\lambda J)</th>
<th>Energy (MeV)</th>
<th>n(\lambda J)</th>
<th>Energy (MeV)</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>p</td>
<td>n</td>
<td></td>
</tr>
<tr>
<td>1f(_{5/2})</td>
<td>0</td>
<td>2.3</td>
<td>3p(_{1/2})</td>
</tr>
<tr>
<td>2p(_{1/2})</td>
<td>1.5</td>
<td>4.2</td>
<td>2f(_{5/2})</td>
</tr>
<tr>
<td>2p(_{3/2})</td>
<td>3.0</td>
<td>5.9</td>
<td>3p(_{3/2})</td>
</tr>
<tr>
<td>1f(_{7/2})</td>
<td>7.8</td>
<td>7.8</td>
<td>2f(_{7/2})</td>
</tr>
<tr>
<td>1d(_{3/2})</td>
<td>11.2</td>
<td>13.2</td>
<td>2d(_{5/2})</td>
</tr>
<tr>
<td>2s(_{1/2})</td>
<td>11.9</td>
<td>14.0</td>
<td>3s(_{1/2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>1h(_{1/2})</td>
</tr>
<tr>
<td></td>
<td></td>
<td></td>
<td>1g(_{7/2})</td>
</tr>
</tbody>
</table>
final state $\pi(f_{7/2}^1)\nu(f_{7/2}^5)$, with $\lambda_n = 3$ and positive parity. Other neutron and proton p-h pairs are $(f_{5/2}^1, l_d3^1)$ and $2p_{1/2}, 2p_{3/2}$ particles coupled to $1d_{3/2}^1, 2s_{1/2}$ holes. Of all these, the energy of the $(p_{3/2}^1, s_{1/2}^1)$ pair (-8.1 MeV) is closest to the observed $\gamma$-ray energy (8.5 MeV).

Transitions to the $3^-$ state at 0.59 MeV are also observed in keV capture and dominate the thermal spectrum. This state has most of the $\pi(f_{7/2}^2, d_{3/2}^1)$ strength {YS64}, and since this component is not present in the $^{45}$Sc ground state (opposite parity), the single neutron component has not been identified in the (d,p) reaction.

Assuming that a $\nu(f_{7/2}^5)$ component is present, a possible mechanism for enhanced transitions to the $3^-$ state after s-wave capture is

$$
\pi(f_{7/2}^1)\nu(f_{7/2}^4, s_{1/2}^1) \rightarrow \pi(f_{7/2}^2, d_{3/2}^1)\nu(f_{7/2}^4, f_{5/2}^1)
$$

$$
M1 \rightarrow \pi(f_{7/2}^2, d_{3/2}^1)\nu(f_{7/2}^5)
$$

The $f_{5/2}$ shell is near threshold and as the pairing energy released in the proton configuration is ~4-5 MeV {MH73}, the centroid energy for M1 transitions will be close to that observed (8.2 MeV). As the resonance and $3^-$ state have common parents, the M1 spin flip transition is enhanced.

For $^{139}$La, possible neutron p-h pairs are $(2f_{5/2}^1, 2d_{3/2}^1)$ and $3p_{1/2}, 3p_{3/2}$, particles coupled to $2d_{3/2}^1, 3s_{1/2}$ holes, e.g.

$$
\pi(lg_{7/2}^1)\nu(2f_{7/2}^1, 3p_{1/2}^1, 2d_{3/2}^1)
$$

The energies of many of these (p-h) pairs lie close to the observed $\gamma$-ray energy (5 MeV). However, exact agreement in energies is not necessary because of the spreading of configuration energies by the interaction of component particles.

The model predicts that only s-wave resonances in Sc and La will exhibit the anomalous spectra. Results for the 0.74 eV p-wave and 72.4 eV s-wave resonances in La {WCG69} are consistent with this interpretation.
The low lying $l_n = 3$ states in Sc and La are known to correspond to single-neutron states coupled to the odd proton in each case \{RSB66,In70\}. It is therefore possible to exclude a collective description for the proposed doorway states (sections 7.1, 7.2).
CHAPTER 8

CAPTURE MECHANISMS IN THE 3s REGION

In preceding chapters, evidence for valence and doorway components in resonance capture has been presented. In this chapter a careful analysis of the capture mechanisms in the 3s region is undertaken, using the measured values for the average s- and p-wave radiative widths and detailed valence and statistical model calculations.

8.1 STATISTICAL CALCULATIONS

The average partial radiative width between resonance $\lambda$ and final state $\mu$ is related to the $\gamma$-ray strength function $S(E_\gamma)$ for dipole transitions by

$$<\Gamma_{\lambda\mu}> = S(E_\gamma)E_\gamma^3/\rho(J^\pi)$$ \hspace{1cm} \ldots (8.1)

In the statistical model, the strength function is independent of $\gamma$-ray energy, but this is not so if the extension to the giant dipole resonance model is assumed. In that case, $S(E_\gamma)$ follows the energy dependence of the (assumed) Lorentzian (i.e. the Brink hypothesis \{Br55,Ax62\}).

The total radiative width is the sum of partial widths which can be divided into allowed dipole transitions to the discrete and continuum states (above excitation energy $E_c$) \{Jo77\}, i.e.

$$<\Gamma_{\lambda}> = \Gamma_{dJ^\pi} + \Gamma_{cJ}$$ \hspace{1cm} \ldots (8.2)

where

$$\Gamma_{dJ^\pi} = \sum_{E_{1},M_{1}} S_{E_{1}}(E_\gamma)E_\gamma^3 + \sum_{E_{1},M_{1}} S_{M_{1}}(E_\gamma)E_\gamma^3/\rho_{dJ^\pi}(E_\lambda)$$ \hspace{1cm} \ldots (8.3)

and is independent of parity. The sum is over only those final states $I$ which can be reached by dipole transitions.
The average total radiative width for a statistical capture mechanism reduces to:

\[ <\Gamma_{Y}^{S}> = \sum_{\mu} \Gamma_{\lambda \mu} = N \cdot D_{J}^{Y}(E_{\lambda}) \left[ \sum_{E} C_{E}^{Y}(E_{1}) + \sum_{E} C_{E}^{Y}(M_{1}) + \sum_{E} E_{Y}^{p}(E_{\lambda} - E_{Y}) dE_{Y} \right] \]

\[ \ldots \text{(8.5)} \]

where \( N \) is a normalisation constant, the energy dependence of the dipole transition is given by the exponent \( n \), and \( R \) is the ratio of average \( E_{1} \) and \( M_{1} \) partial radiative widths.

The level density formula \( \{\text{GC65}\} \) is adjusted to fit the spacing \( (D_{J}^{Y}) \) of resonances above threshold. All allowed spin final states are included in the statistical calculation since transition strengths averaged over many resonances are assumed to be independent of the final state configuration. The exception to this rule are those final states with large single particle amplitudes \( (\text{i.e. } E_{\mu}^2 \sim 1) \). Transitions to these states can only occur from the valence or appropriate 2p-1h component of the resonance wave function. If the giant dipole resonance dominates in the threshold region, then \( n \sim 5 \) is expected \( \{\text{Ax62}\} \).

Energy dependence and \( E_{1}/M_{1} \) ratio

The relevant energy dependence and \( E_{1}/M_{1} \) ratio is not well known in the 3s region. Measurements of average keV neutron capture in Cu \( \{\text{Al68a}\} \) and Ni \( \{\text{AKS68}\} \) suggest \( n \sim 3 \), and in zinc \( n \sim 4 \) \( \{\text{AM70a}\} \). A value of \( R \sim 1 \) is obtained from threshold photonuclear experiments with the residual nuclei \( ^{53}\text{Cr}, ^{56}\text{Fe}, ^{57}\text{Fe} \) and \( ^{63}\text{Ni} \) \( \{\text{Ja74b}\} \). Both of these results are at variance with \( n = 5, R = 7 \) observed in the 4s region (section 1.1), and it is necessary to ensure that values of \( R \) and \( n \) are used which are consistent with the radiative width data available in the 3s region.

Using the keV \( \gamma \)-ray data for \( ^{56}\text{Fe} \) \( \{\text{Bi+73}\} \) and the s-wave and p-wave radiative widths obtained earlier, the average \( E_{1} \) and \( M_{1} \) reduced widths are found to be

\[ <k(E_{1})> = 1.5 \times 10^{-9} \ \text{MeV}^{-3} \]

\[ <k(M_{1})> = 4.8 \times 10^{-9} \ \text{MeV}^{-3} \]
where \( k(El) = \Gamma_{\lambda \mu}^{(El)} \left[ E_{\gamma \mu}^{3} \cdot D_{0, J} \cdot A_{\gamma} \right]^{-1} \), \( \ldots (8.6) \)

\( k(Ml) = \Gamma_{\lambda \mu}^{(Ml)} \left[ E_{\gamma \mu}^{3} \cdot D_{1, J} \right]^{-1} \), \( \ldots (8.7) \)

and the subscripts 0 and 1 designate s- and p-wave resonances. The El average is for 14 \( \gamma \)-ray intensities measured at thermal and 27 keV, and the Ml average is over 25 intensities at 1.17 keV, \( <36>, <52> \) and \( <72> \) keV.

While \( <k(El)> \) is similar to that obtained in the \((\gamma, n)\) reaction \( \{Ja74\} \), \( <k(Ml)> \) is only 25% of the corresponding value. Since a large number of \( l > 0 \) resonances were missed in the \((\gamma, n)\) measurement, it would seem that the average Ml reduced width so obtained is not representative of the overall population.

The ratio of reduced widths is

\[ <k(El)> = 0.3 \cdot <k(Ml)> \]

and this gives the result

\[ <\Gamma_{\lambda \mu}^{(El)}> = 7 \cdot <\Gamma_{\lambda \mu}^{(Ml)}> \]. \( \ldots (8.8) \)

This ratio is in excellent agreement with that found in the 4s region by Bollinger \( \{Bo73\} \).

**Variance**

The variance of the average statistical width is also of interest. The calculated partial width \( \Gamma_{\lambda \mu} \) is in fact the expectation value \( <\Gamma_{\lambda \mu}> \) for the distribution of \( \Gamma_{\lambda \mu} \) over many resonances (\( \lambda \)). The quantity

\[ x_{\mu} = \frac{\Gamma_{\lambda \mu}}{<\Gamma_{\lambda \mu}>} \]

is assumed to follow a \( \chi^2 \) distribution with one degree of freedom (\( v = 1 \)) with mean \( <x_{\mu}> = 1 \) and variance \( \nu_{\mu}^2 = 2 \). The variance of \( <\Gamma_{\lambda \mu}> \) is then

\[ \sigma_{\mu}^2 = 2 <\Gamma_{\lambda \mu}>^2 \]. \( \ldots (8.9) \)

The total width and variance of the statistical component for \( n \) partial widths are the sum of means and variances, respectively,
\[ <\Gamma^S_{\lambda Y}> = \sum_{\lambda=1}^{n} <\Gamma_{\lambda}\mu> \]

\[ \sigma^2_s = \sum_{\lambda=1}^{n} \sigma^2_{\lambda\mu} = 2 \sum <\Gamma_{\lambda}\mu>^2 \quad ...\quad (8.10) \]

The relative variance of the weighted sum of \(n\ \chi^2\ (\nu = 1)\) distributions is then

\[ \nu^2_s = \frac{\sigma^2_s}{<\Gamma^S_{\lambda Y}>^2 \nu_{\text{eff}}} \quad ...\quad (8.11) \]

where the subscript \(s\) denotes the statistical model and \(\nu_{\text{eff}}\) is the effective number of degrees of freedom. Normally \(\nu_{\text{eff}} \ll n\) as a result of the \(E\gamma^3\) weighting of the partial widths to the low lying discrete states.

**Normalisation**

Some authors have chosen (e.g. Jo77) to normalise the threshold values of \(<\Gamma_{\lambda Y}>\) to the extrapolation of the GDR. This procedure is acceptable only if the \((\gamma,n)\) experimental cross section below a few MeV above threshold faithfully reproduces the Lorentzian curve. This is not the case for light nuclides nor those near closed shells (Be73), nor is such a smooth effect predicted by p-h calculations (section 7.3).

The method adopted here is to normalise the p-wave statistical calculation to the observed p-wave radiative width, and then determine the statistical component of the s-wave width directly.

The ratio of statistical radiative widths for s- and p-wave resonances \((R^S = <\Gamma^S_{\lambda YS}>/<\Gamma^S_{\lambdaYP}>\) depends critically on the energy of the first positive parity state. Most nuclides in the 3s region have low lying p-states, to which El transitions can occur from s-wave resonances. M1 transitions to these states from p-wave resonances will therefore be weaker by the factor \(R\). The ratio of s- and p-wave statistical widths can therefore be compared with the experimental data after correcting \(<\Gamma_{\lambda Y}>\) for non-statistical components. It is assumed that the p-wave
width has $<\Gamma^S_\gamma > = <\Gamma_\gamma >$.

### 8.2 Partition of the $s$-wave Radiative Widths

The $s$- and $p$-wave radiative widths are partitioned as follows:

$$<\Gamma^S_\gamma > = <\Gamma_\gamma >$$

$$<\Gamma^S_\gamma > = R^S <\Gamma^S_\gamma >$$

$$<\Gamma_\gamma > = <\Gamma^V_\gamma > + <\Gamma^U_\gamma > + <\Gamma^S_\gamma > \quad \ldots(8.12)$$

with variance

$$\sigma^2 = \sigma^2_V + \sigma^2_U + \sigma^2_S + \sigma^2_{EX} \quad \ldots(8.12)$$

where

$$\sigma^2_S = \frac{2}{\nu_{\text{eff}}} <\Gamma^S_\gamma >^2$$

$$\sigma^2_V = 2 <\Gamma^V_\gamma >^2$$

and $\sigma^2_{EX}$ is the average experimental variance. The ratio $R^S$ is obtained from the statistical calculations described in the preceding section, and the valence width $<\Gamma^V_\gamma >$ is calculated following the procedure of section 1.3. It remains to obtain estimates of the doorway component $<\Gamma^U_\gamma >$ and its variance $\sigma^2_U$ which reduces the magnitude of the initial state correlation (equation 5.6).

$s$-wave radiative widths are taken from Table 4.1 and from the latest revision of ORELA results, in which the full Monte Carlo analysis of the prompt background correction is included (AM77). These results supersede those given in Table 4.1.

Conclusions based on equation 8.12 are qualitative only, because of the intrinsic uncertainties in the calculated quantities, as well as the experimental errors and large variances of the $s$- and $p$-wave radiative widths.

Results of calculations for a range of nuclides in the $3s$ region (assuming $R = 7$, $n = 3$ for the statistical model) are given in Table 8.1. In the main, these results are sensitive to the assumed values of $R$ and $n.$
### TABLE 8.1

**PARTITION OF s-WAVE RADIATIVE WIDTHS IN THE 3s REGION (eV)**

<table>
<thead>
<tr>
<th>Target</th>
<th>$J^\pi_\lambda$</th>
<th>$\langle \Gamma_{\gamma s} \rangle$</th>
<th>$\langle \sigma_{\text{ex}}^2 \rangle$</th>
<th>$\langle \Gamma_{\gamma p} \rangle$</th>
<th>$R_s^s$</th>
<th>$\langle \Gamma_{\gamma s}^s \rangle$</th>
<th>$\sigma_{s}^2$</th>
<th>$\langle \Gamma_{\gamma}^V \rangle$</th>
<th>$\sigma_{V}^2$</th>
<th>$\langle \Gamma_{\gamma}^U \rangle$</th>
<th>$\sigma_{U}^2$</th>
<th>$\rho_{I_C}$</th>
<th>$\rho_I$</th>
<th>S.D.</th>
</tr>
</thead>
<tbody>
<tr>
<td>$^{40}$Ca</td>
<td>$1/2^+$</td>
<td>1.5</td>
<td>0.12</td>
<td>0.90</td>
<td>0.36</td>
<td>0.92</td>
<td>0.33</td>
<td>0.02</td>
<td>0.52</td>
<td>0.54</td>
<td>0.6</td>
<td>0.16</td>
<td>0.4</td>
<td>0.2</td>
</tr>
<tr>
<td>$^{45}$Sc</td>
<td>$3^-$</td>
<td>0.84</td>
<td>0.01</td>
<td>0.21</td>
<td>0.50</td>
<td>1.08</td>
<td>0.54</td>
<td>0.006</td>
<td>0.10</td>
<td>0.02</td>
<td>0.2</td>
<td>0.17</td>
<td>0.1</td>
<td>0.04</td>
</tr>
<tr>
<td>$^{46}$Ti</td>
<td>$1/2^+$</td>
<td>1.2</td>
<td>0.08</td>
<td>0.36</td>
<td>0.60</td>
<td>0.94</td>
<td>0.56</td>
<td>0.03</td>
<td>0.53</td>
<td>0.57</td>
<td>0.1</td>
<td>-</td>
<td>1.0</td>
<td>0.43</td>
</tr>
<tr>
<td>$^{48}$Ti</td>
<td>$1/2^+$</td>
<td>1.4</td>
<td>0.11</td>
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<td>1.26</td>
<td>0.42</td>
<td>0.03</td>
<td>0.83</td>
<td>1.4</td>
<td>0.2</td>
<td>-</td>
<td>0.8</td>
<td>0.82</td>
</tr>
<tr>
<td>$^{53}$Cr</td>
<td>$1^-_2^-{1.4_2^4}$</td>
<td>1.4</td>
<td>0.42</td>
<td>0.64</td>
<td>1.0</td>
<td>1.83</td>
<td>1.63</td>
<td>0.26</td>
<td>0.22</td>
<td>0.10</td>
<td>-</td>
<td>-</td>
<td>0.39</td>
<td>0.31</td>
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<tr>
<td>$^{56}$Fe</td>
<td>$1/2^+$</td>
<td>3.1</td>
<td>1.3</td>
<td>6.8</td>
<td>0.45</td>
<td>2.6</td>
<td>1.2</td>
<td>0.25</td>
<td>1.8</td>
<td>6.7</td>
<td>-</td>
<td>-</td>
<td>1.0</td>
<td>0.79</td>
</tr>
<tr>
<td>$^{59}$Fe</td>
<td>$1/2^+$</td>
<td>1.5</td>
<td>0.1</td>
<td>0.64</td>
<td>0.6</td>
<td>3.5</td>
<td>2.1</td>
<td>0.84</td>
<td>0.28</td>
<td>0.16</td>
<td>-</td>
<td>-</td>
<td>0.50</td>
<td>0.52</td>
</tr>
<tr>
<td>$^{57}$Fe</td>
<td>$0^-$</td>
<td>2.2</td>
<td>0.71</td>
<td>1.96</td>
<td>0.52</td>
<td>1.7</td>
<td>0.88</td>
<td>0.17</td>
<td>0.17</td>
<td>0.06</td>
<td>1.1</td>
<td>1.0</td>
<td>0.02</td>
<td>0.57</td>
</tr>
<tr>
<td>$^{59}$Fe</td>
<td>$1^-$</td>
<td>1.7</td>
<td>0.30</td>
<td>1.69</td>
<td>0.62</td>
<td>2.8</td>
<td>1.53</td>
<td>0.22</td>
<td>0.54</td>
<td>0.58</td>
<td>-</td>
<td>0.64</td>
<td>0.58</td>
<td>0.66</td>
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<tr>
<td>$^{60}$Ni</td>
<td>$1/2^+$</td>
<td>1.4</td>
<td>0.02</td>
<td>0.04</td>
<td>0.23</td>
<td>1.72</td>
<td>0.40</td>
<td>0.03</td>
<td>0.65</td>
<td>0.85</td>
<td>0.4</td>
<td>-</td>
<td>0.65</td>
<td>0.71</td>
</tr>
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</table>
Exceptions are $^{40}$Ca, $^{45}$Sc and $^{46}$Ti where low lying states of mixed parity occur and $R^S$ has a value of unity.

Equations 8.12 and 8.13 provide independent evidence for the existence of the doorway component for a number of nuclides. When such evidence is absent, the appropriate columns in Table 8.1 are left blank.

Results for $^{40}$Ca and $^{45}$Sc, because of their insensitivity to the statistical model assumptions, provide strong evidence for the doorway component. In the case of $^{45}$Sc the observed variance of the s-wave radiative widths for a large number of resonances can only be explained by invoking a doorway component with enhanced transitions to a small number of $l_n = 3$ states, as described in section 7.5.

The magnitude and variance of the doorway components in $^{40}$Ca and $^{45}$Sc reduce the expected initial state correlation, yielding good agreement with the measured values.

The valence model accounts for the observed variances in $^{46}$Ti, $^{48}$Ti and $^{54}$Fe, but the valence variance greatly exceeds the reported value in $^{60}$Ni. A reduction in the valence estimate for $^{60}$Ni would require an increase in $<\Gamma^U_\gamma>$, and a lower calculated correlation. With the exception of $^{46}$Ti, good agreement between the expected and observed correlations is obtained, while a larger value for $<\Gamma^U_\gamma>$ would improve the agreement in $^{46}$Ti.

In $^{53}$Cr, the large experimental variance dominates the observed variance and an uncorrelated component is not required to account for the data.

The assumption that $<\Gamma^S_\gamma>$ = $<\Gamma_\gamma>$ may also cause discrepancies, as observed in $^{53}$Cr and $^{56}$Fe. For $^{56}$Fe, the value $<\Gamma_\gamma>$ = 0.3 eV reported in this thesis, has been superseded by elastic scattering measurements at ORNL by F. Perey et al. (private communication), which gave $l_n$ and $J^\pi$ estimates. The new value $<\Gamma_\gamma>$ = 0.6 eV cannot be equal to $<\Gamma^S_\gamma>$ since the calculated statistical width for s-wave resonances (2.1 eV) exceeds the observed value (1.5 eV). As noted in section 6.1, p-wave spectra are essentially non-statistical in character and the p-wave statistical assumption is invalid. A value
If the same argument holds for p-wave resonances elsewhere in the 3s region, then the magnitudes and variances of the doorway component given in Table 8.1 will be lower limits.

An interesting spin dependence is apparent in $^{57}$Fe. Low lying states in $^{58}$Fe have spins $0^+, 1^+, 2^+$, and only states with $1^+$ can be reached by El transition from $0^-$ resonances, whereas all states are accessible from $1^-$ resonances. Consequently, statistical and valence components are larger for $1^-$ resonances, yet the observed radiative widths are smaller \cite{Al+77c}. The doorway component is therefore quite large for $0^-$ resonances, but is observed only through its variance in the $1^-$ resonances. The sample size for $0^-$ resonances is small (note the large error on the correlation), and possibly the measured widths may not be a representative sample from the parent population. Alternatively, the doorway component may suffer a spin dependence.

**8.3 FINAL STATE CORRELATIONS**

The magnitude of the uncorrelated component is often comparable to or greater than the average valence width. It is therefore expected that the $\gamma$-ray energy dependence of the transitions to $l_n = 1$ final states from the doorway part of the resonance wave functions will strongly influence the observed energy dependence. Since the p-h doorway transitions are dependent on the final state spectroscopic strengths, correlations with those strengths would also be expected.

Kopecky \cite{Ko73} varied the $\gamma$-ray energy dependence to obtain optimised final state correlations for thermal capture in the 3s region. The same approach is applied to average $\gamma$-ray intensities observed in the keV neutron energy region, where p-wave resonances also make a substantial contribution to the capture spectrum. The optimised final state correlation coefficients in $^{54}$Fe and $^{56}$Fe are consistently obtained for a value of the exponent $n$ in equation 8.5, in the range zero to unity for both
**TABLE 8.2**

OPTIMISED FINAL STATE CORRELATIONS FOR keV CAPTURE IN THE 3s REGION

<table>
<thead>
<tr>
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<th>$\rho_F$ (n=0)</th>
<th>$\rho_F$ (n=3)</th>
<th>$\rho_F$ (n=5)</th>
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<tr>
<td>$^{42}$Ca</td>
<td>0.67</td>
<td>0.23</td>
<td>0.11</td>
</tr>
<tr>
<td>$^{44}$Ca</td>
<td>0.91</td>
<td>0.50</td>
<td>0.06</td>
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<tr>
<td>$^{48}$Ti</td>
<td>0.79</td>
<td>0.75</td>
<td>0.71</td>
</tr>
<tr>
<td>$^{54}$Fe</td>
<td>0.90</td>
<td>0.25</td>
<td>0.02</td>
</tr>
<tr>
<td>$^{56}$Fe</td>
<td>0.45</td>
<td></td>
<td></td>
</tr>
<tr>
<td>$^{58}$Ni</td>
<td>0.98</td>
<td>0.75</td>
<td>-0.15</td>
</tr>
<tr>
<td>$^{60}$Ni</td>
<td>0.77</td>
<td>0.51</td>
<td>0.18</td>
</tr>
<tr>
<td>$^{64}$Zn</td>
<td>0.42</td>
<td>0.34</td>
<td>0.13</td>
</tr>
<tr>
<td>$^{66}$Zn</td>
<td>0.56</td>
<td>0.50</td>
<td>0.39</td>
</tr>
</tbody>
</table>
and p-wave resonances. Since the valence model requires a maximum
correlation for n = 2, and the statistical model predicts n = 3 and a
zero correlation (independent of n), it is evidence that an additional
capture mechanism is present. This must be the uncorrelated component
which is needed to reduce the initial state correlation and to account
for the observed magnitude and variance of s-wave radiative widths.

Similar results are observed for average capture spectra in $^{42}\text{Ca}$,
$^{44}\text{Ca}$, $^{48}\text{Ti}$, $^{58}\text{Ni}$, $^{60}\text{Ni}$, $^{64}\text{Zn}$ and $^{66}\text{Zn}$. Averaged γ-ray transitions
(Bi+73) from a mix of s- and p-wave resonances in all these nuclides
have maximum final state correlations for n = 0-1 (Table 8.2). It
therefore appears that both s- and p-wave capture are influenced by
doorway effects.

The existence of an energy dependence in the distribution of final
state spectroscopic strengths has not been included in the thermal and
resonance final state correlations. Generally, the lowest lying state
have the largest spectroscopic strengths and, as a result, the correla­
tion with γ-ray energy $\rho((2J+1)\theta^2, E^n_\text{Y})$ is large (-0.6-0.9) for most cases.
However, this correlation is relatively insensitive to the exponent n,
save that n is positive.

The observed energy dependences (n = 0-1) are therefore significant
since for keV capture most correlation coefficients are sensitive to
the exponent n. Since the dipole phase space factor is proportional to
$E^3_\text{Y}$ and $\sigma = E^3_\text{Y} |<0>|^2$, the dipole matrix element for doorway transitions
will be proportional to $-E^{-1}_\text{Y}$. This result is similar to that expected
for potential capture.
CHAPTER 9
CONCLUSIONS

The neutron capture mechanism has been investigated in a series of γ-ray and cross section measurements using NaI and Ge(Li) spectrometers and the total energy detector, respectively.

The contribution of resonance scattered neutrons to capture γ-ray yields in the total energy detector has been studied and the effective empirical efficiency for neutron detection obtained as a function of scattered neutron energy. This neutron sensitivity function has then been used in Monte Carlo analyses to fit prompt backgrounds and resonance capture yields simultaneously. In this way, it has been possible to account for observed asymmetries in s-wave capture resonances as well as the background structure between resonances.

Using this method of analysis, the interference between resonant and non-resonant capture cross sections has been investigated in $^{54}$Fe. An upper limit has been set for the magnitude of the non-resonant cross section which provides a new benchmark for calculations of direct capture.

The resonance capture data are found to exhibit moderate initial and final state width correlations in the 2p, 3s and 3p regions. The magnitude of the total valence radiative widths are found to be a significant fraction of the non-statistical widths, i.e. about one half and two thirds in the 3s and 3p regions, respectively. In the 4s region, the valence process is generally insignificant, yet moderate initial state correlations still occur. It is apparent that valence processes are exempt from elevation of El strength up into the GDR. This result holds for s- and p-wave capture and $^{40}$Ca results indicate that d-wave capture is also exempt.

A detailed investigation of neutron capture in $^{54}$Fe and $^{56}$Fe has shown the dominance of valence capture in the former, and the presence of non-statistical effects in the latter. Resonance-resonance inter-
ference is observed in $^{54}$Fe, which is found to be a direct consequence of the valence model. In $^{56}$Fe, capture $\gamma$-ray spectra for s-, p- and d-wave resonances are found to be quite similar, but essentially non-valence, and are interpreted in terms of a 2p-1h interaction.

A similar explanation is advanced to explain the anomalous spectra in $^{45}$Sc and $^{139}$La, where intense transitions occur to the low lying $\ell_n = 3$ states. Resonance capture cross section measurements, however, do not exhibit initial state correlations.

A detailed investigation of capture in the 3s region has provided evidence for significant doorway contributions to the observed s-wave radiative widths, their variance and correlation with the reduced neutron widths. Since statistical and valence variances can be estimated accurately, increased attention needs to be placed on the measurement and interpretation of the variance of radiative widths.

Overall, the correlation data can be explained if varying numbers of El doorway states are present in all mass regions within a few MeV of threshold. In the 3s and 3p regions these doorway states appear to be responsible for reducing the correlations expected from valence capture, while they may be the cause of the observed correlations in the 4s region. Results for Sc and La indicate that since non-statistical spectra occur in the absence of initial state correlations, many doorway states may contribute to resonance capture, particularly for odd-even target nuclides.

These states would have large amplitudes for El p-h states which have unperturbed energies in this range, and carry a few per cent of the total El strength. Superimposed on this background of radiative doorway states, valence effects, being exempt from GDR depletion, can dominate near closed neutron shells where maxima occur in the resonance spacing, neutron strength function and binding energies of bound single particle states. Isolated particle-vibrator states can occur in the elastic
neutron scattering channel, but in general these appear to carry little intrinsic E1 strength in capture reactions.

Only limited theoretical support for this model is found. It is evident that fragments of E1 strength can escape from the GDR and that a number of doorway states may exist in the threshold region with widths of 0.5 to 3 MeV. Further realistic calculations are required to determine in detail the energies, widths and wave functions of these states in regions near the closed neutron shells. Complementary γ-ray measurements are also needed for neutron energies up to at least 1 MeV to provide information on the energy dependence of the capture mechanism.
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